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FOREWORD

The publications of 14th International Congress on Acoustics are bound in five volumes, four volumes of Proceedings and one volume of Programme. The Proceedings contain the text of 838 papers presented on the Congress. The Programme contains abstracts and technical programmes of all the papers, and general information. The Congress is sponsored by the International Union of Pure and Applied Physics (IUPAP), and organized by the Acoustic Society of China and the Institute of Acoustics, Academia Sinica. The Congress is held in Beijing, China from 3rd through 10th September 1992.

The Proceedings reflect the recent scientific and technical achievements made by about 1700 authors around the world. Twelve plenary lectures review the state of arts of some widely interested subjects in acoustics. The rest research papers are arranged according to the subfields of acoustics as follows:

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The papers in each subfield are divided into a number of subjects and then numbered in the subjects. The papers were presented, during the Congress, orally in 97 Sessions and by posters in two. Those papers presented in the Poster Sessions had been displayed for three days during the Congress and are classified by subfields only and then numbered.
The 14th International Congress on Acoustics has been a great event in the field of acoustics, and well received by the acousticians all over the world. The publication of the Proceedings helps to disperse the scientific information in acoustics cumulated in the last few years and to arouse the interest of scientific community for further development of the subject.

The editor would like to thank Professor DAI Genhua and Dr TIAN Jing for their contribution in editing the publications.

Li Peizi
*Editor*
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AP  Poster Papers
ELASTIC NONLINEARITIES IN SINGLE CRYSTAL GALLIUM ARSENIDE BETWEEN 77 AND 300 K

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Third order elastic (TOE) constants are necessary for a fundamental description of many physical properties of solids resulting from lattice nonlinearity. The properties in question include thermal expansion, heat conduction, temperature dependence of the specific heat, temperature and pressure dependence of the elastic constants, difference between the adiabatic and isothermal elastic constants, phonon viscosity, thermal attenuation of acoustic waves, shock deformation of materials, acoustically induced static stress and strain, etc. Previously \(^1\), we have shown that measurement of ultrasonic harmonic generation gives information about combinations of TOE constants from helium temperatures to at least 350° K. Since no single technique gives all six TOE constants of cubic crystals it was necessary to introduce a second technique in order to plot all six of the TOE constants of the semiconductor Gallium Arsenide (GaAs) over a wide temperature range. The purpose of this report is to describe the method used because it makes optimum use of theory and experiment to arrive at data which would not be available otherwise. The technique should be applicable without modification to the evaluation of all TOE constants of all zincblende structure compounds. The data on GaAs verify its validity.

Measurement of Harmonic Generation

The propagation of a finite amplitude ultrasonic wave along the three principal directions in a cubic lattice is described by

\[
\frac{\partial^2 U}{\partial t^2} = K_2 \frac{\partial^2 U}{\partial x^2} + (3 K_2 + K_3) \frac{\partial^2 U}{\partial y^2} \frac{\partial^2 U}{\partial z^2} .
\]

(1)

Assuming an initially sinusoidal wave allows one to write a solution in the form

\[
U = A_1 \sin (k_1 x - \omega t) + A_2 \cos (k_2 y - \omega t) + ...
\]

(2)

where

\[
A_2 = -3 K_3 \frac{K_1}{K_2} A_1^2 k^2 x .
\]

(3)

Defining the nonlinearity parameter

\[
\beta = \frac{A_2}{A_1} \frac{k^2}{x}
\]

(4)

allows one to write the nonlinearity parameters in terms of measured quantities

\[
\beta = \beta_0 + \beta_110 + \beta_{111}
\]

(5)

where \( k = \frac{2\pi}{\lambda} \) is the propagation constant and \( x \) is the sample length. In terms of TOE constants the nonlinearity parameters along the principal directions are

\[
\begin{align*}
\beta_{100} &= -3 + \frac{C_{111}}{C_{11}}, \\
\beta_{110} &= -3 + \frac{C_{111} + 3C_{112} + 12C_{166}}{2C_{11} + C_{12} + 2C_{44}}, \\
\beta_{111} &= -3 + \frac{C_{111} + 6C_{112} + 12C_{144} + 24C_{166} + 2C_{123} + 16C_{444}}{3C_{11} + 2C_{12} + 4C_{44}}
\end{align*}
\]

(6, a, b, c)

By measuring the amplitudes \( A_2, A_1 \) the frequency (from which \( k \) can be calculated), and the sample length \( x \) in the principal directions one can evaluate the nonlinearity parameters. The results for GaAs are shown in Fig. 1 plotted between 77 K and room temperature.

![Fig. 1. Temperature dependence of nonlinearity parameters of GaAs.](attachment:image1)

To evaluate the individual TOE constants from these data, additional information is needed. The Keating theory is a three-parameter theory from which one can calculate all six TOE constants. Thus, in principle these data plus the Keating theory would be sufficient to isolate all six TOE constants of GaAs. We chose to provide additional experimental data, however.

Measurement of Pressure Variation of Sound Velocity

The most accurate evaluation of TOE constants of cubic crystals at room temperature appears to come from the combination of harmonic generation data with those taken by use of pressure variation of ultrasonic wave velocity.\(^4\) A pressure bomb was used to take pressure variation data at room temperature with the same GaAs samples. A plot of the normalized frequency (essentially sound velocity) as a function of pressure is given in Fig. 2. The curves all are straight lines except for the longitudinal wave in the [111] direction (labelled [111][111] in Fig. 2). The slopes of these curves can be used to calculate combinations of TOE constants at room temperature.

![Fig. 2. Variation of normalized frequency (ultrasonic wave velocity) with pressure in GaAs.](attachment:image2)

Results

The results of the two sets of measurements can be combined to isolate all six TOE constants at room temperature. The results have been evaluated for GaAs and compared with room temperature values of other researchers. Having these values one now is able to use the Keating model along with the harmonic generation data to calculate the temperature dependence of each TOE constant. Results of the values of all six TOE constants of GaAs between room temperature and liquid nitrogen temperature are given in Figs. 3 and 4 in which the curves are least squares fits of the data with a fifth order polynomial. Most of the
TOE constants are linear functions of temperature; however, both $C_{112}$ and $C_{123}$ exhibit remarkable temperature variations and emphasize the importance of being able to measure these fundamental parameters over a range of temperatures.

![Graph showing temperature dependence of $C_{111}$, $C_{166}$, and $C_{456}$](image1)

![Graph showing temperature dependence of $C_{112}$, $C_{123}$, and $C_{144}$](image2)

We have examined the strong Cauchy relations over the available temperature range by comparing the measured quantities $C_{112} + 4C_{155}$ and $C_{123} + 6C_{144} + 8C_{456}$ with $\frac{3}{2}C_{111}$ and find that there is a tendency for better agreement with the TOE constant Cauchy relations as $0K$ is approached. Since both the Keating model and the Cauchy relations strictly should apply only at $0K$, this may indicate that the lattice interaction in GaAs is basically of the central type, and the deviations from the Cauchy relations are caused by thermal effects. An analogous observation has been made for germanium and silicon.

**Summary**

All six TOE constants of GaAs have been evaluated between room temperature and liquid nitrogen temperatures. Two of the TOE constants, $C_{112}$ and $C_{123}$, exhibit considerable variation, (at least a factor of 10 for $C_{123}$). The other four are almost linear functions of temperature.

The technique we have used can be applied directly to evaluation of the TOE constants of all zincblende structure compounds. Other structures or samples in which interstitials or dislocations are prominent require further analysis.

**References**


**Acknowledgement**

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INTERFACE ACOUSTIC NONLINEARITY
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I. THE PHYSICS OF NONLINEAR INTERFACE PROBLEM

The self- and/or cross-interactions of bulk acoustic waves have long been studied, but only the interactions complying with phase match condition could be accumulated to a significant amount and observed.[1] At the interface, however, all these waves due to self- and cross-interactions would participate the boundary conditions. In addition, the fundamental incident, reflected and deflected waves have the second harmonic stresses with them as well. In order to satisfy the boundary conditions, one has to have the free propagating waves radiated from the interface. This phenomenon characterizes the nonlinear interface problem, which differs from ordinary bulk nonlinear effects, but depends also on the second- and third-order elastic constants of both materials.

In some special cases over the critical incident angle, the reflected and/or deflected waves are of the evanescent property, and the power flow in a wide incident beam will squeezed into a narrow layer at the interface, the high power density of which causes a strong nonlinear effect just as it happens at the nonlinear surface wave case.

II. REFLECTION OBSERVATION

As shown in Fig. 1, the incident SV wave may cause the occurrence of an reflected SV harmonic wave. For the incident angle which is far above the critical one, only SV harmonic will propagate in the bulk and, thus, can be measured.(2)

The materials were glass ($\rho = 3.54$ g/cm$^3$, $v_L = 4.16$ km/s, $v_T = 2.50$ km/s) and iron ($\rho = 7.82$ g/cm$^3$, $v_L = 6.07$ km/s, $v_T = 3.30$ km/s). With the modulated pulse input (8.3 MHz), a double frequency reflective wave which obeys the square law was detected by the transducer. The signal dropped significantly if the iron was removed. Fig. 2 shows the relation of the received harmonic amplitude versus the incident angle, which indicates a maximum generation efficiency around the critical angle of the generalized Rayleigh wave generation (53° calculated).

The physical explanation of the peak is that the generated leaky wave radiates its harmonic component back into the media while it propagates along the interface. In the angle range where most reflected or deflected waves are evanescent ones, the high energy density occurs within the thin layer near the interface, hence, a strong nonlinearity was examined.

III. CONVOLUTION EFFECT

On the basis of the above experimentation, a convolution effect on the interface was studied(3). As shown in Fig. 3, two transducers were fed with the same modulated rectangular pulse of a primary power so that two beams of the SV incident waves with the same frequency will generate a driving wave whose wave vector is normal to the interface. Because of the boundary condition, a double frequency longitudinal wave will ejected from the interface, and its generation efficiency reveals the convolution of two incident waves, i.e. the received wave has a double frequency and a triangle waveform.

![System setup](image)

![Echoes (upper trace) & received signals (lower trace)](image)

Fig. 3 Experiment of the convolution effect
Such a harmonic signal was accessed (7.7 MHz was the primary frequency). It satisfied the square law and its maximum amplitude increased if the pulse width increased. As the pulse width exceeded the interaction zone, the shape of the harmonic became a trapezoid and the maximum amplitude reached a constant. The correlation between the convolution signal and the incident angle is outlined in Fig. 4. The maximum peak arose almost at the same angle as the case of interface reflection (Fig. 2).

At the incident angle of 51°, the bilinear factor C was measured as 106.6 dbm, and the figure of merit 4×10^-7 W/Hz. These values are smaller than the value of Y-Z lithium niobate surface wave device. The reason is that the SAM device uses the surface wave, which has a higher energy density near the boundary than the bulk wave as was used in this experiment.

![Graph](image)

**Fig. 4 Convolution signal vs. incident angle**

**Fig. 5 Theoretical expectation of the reflected SV harmonic**

(The generated harmonic has been normalized by the wave vector & the square of the amplitude of the incident primary SV wave.)

**IV. THEORY OF THE INTERFACE REFLECTION**

In an infinite, uniform and isotropic material, the second-order approximation will give us a linearly accumulation factor $\alpha$ in case of self- and cross-interactions under the phase match condition. As well known, $\alpha$ could be determined by the second- and third-order elastic constants of the material. Once an oblique bulk wave incidence on an interface is involved, the nonlinear motion equation must be expressed in a two dimensional system; and for a plane wave solution, there exist two constants $\alpha$ and $\alpha'$ for the linear increases both the in propagation direction and within the wave front. $\alpha'$ keeps the same meaning, and $\alpha'$ is related with the interface.

For a general case with an incident $L$, SV- and/or SH-type waves, we have written out all the self- and cross-interaction terms contributed by the incident, reflected and deflected waves. As mentioned above, all the second harmonic or sum frequency displacements and stresses at the interface can be expressed in terms of the material parameters and the incident amplitude. (We will not write out all of them here because of the text space.) We take it for granted that the boundary conditions in the nonlinear case is still the continuity of displacements and stresses everywhere, that is, continuity of every harmonics. Including $\alpha'$ (some displacements and stresses are increasing along the interface) into the boundary conditions, one will double the equations of the boundary conditions with the $\alpha'$ variables. Finally, the problem is solvable and the solution is unique. The detailed theory and solving process will be shown in another paper.

**V. CONCLUSION**

This paper presents a comprehensive theory to analyze the problem of a general nonlinear boundary, and some related experimental results have been acquired. The discovery of convolution effect of the interface may lead to the realization of nonlinear bulk acoustic wave device. And the simplified version of this theory has been applied to the analysis of the SAW nonlinearity[5][6].

**REFERENCE**

THE ORIGIN OF NONLINEAR RESPONSE IN ROCK

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INTRODUCTION

Seismic velocities generally increase with depth in the earth. The increases in moduli with depth (and thus with pressure) implied by these observations are much greater than those measured for other materials, such as gases, liquids, metals and intact ceramics. Increases in the elastic constants that govern wave propagation are caused by closure of the open, compliant porosity present near the earth's surface where overburden pressures are relatively low. Microscopic examination of representative specimens show that porosity with low aspect ratio (crack width over length) is a common feature of the microstructure. Waves of finite amplitude should be altered by nonlinear effects, because independent estimates of the nonlinearity parameter yield values of order 10\textsuperscript{-2} or greater. We are exploiting this phenomenon to form parametric beams in rock having low divergence and improved penetrating power (P.A. Johnson et al., this volume), and therefore need to identify and characterize the microstructural features that control the strength of the nonlinear interaction to optimize source performance.

The purpose of this paper is to summarize experimental results that are distinctive for rocks and then examine potential physical mechanisms which might plausibly explain the observations. These mechanisms will provide the basis for computational models to be developed in order to infer physical properties from measurements of nonlinear response and to simulate wave propagation effects. Mechanisms causing the nonlinear response of rocks are not well understood at present. Experiments indicate that different mechanisms come into play depending on the amplitude of the disturbance. At the lowest strains of interest here (10\textsuperscript{-4} to 10\textsuperscript{-3}) closure of long thick cracks with high compliance strongly influences deviations from linear response. At higher strains, above \textasciitilde10\textsuperscript{-5}, sliding along internal surfaces occurs over a range of scales producing a hysteretic nonlinearity. Throughout the entire strain range, movement of pore fluids that are often present in near surface rocks can cause nonlinear strain response through local changes in the crack configuration. A related and potentially important effect is the motion of the contact line that develops between the liquid and gas phases in partially saturated material.

MECHANISMS OF THE NONLINEAR RESPONSE

Crack closure

Laboratory measurements of ultrasonic velocities under hydrostatic pressure clearly show the effect of closing microfractures. Typical behavior is illustrated by the data of Figure 1. (Bonner and Wanamaker, 1990) which show measurements of compressional and shear velocities along with the computed bulk modulus for a fine grained sandstone known to contain a large concentration of long, thin cracks. Measurements were made using contact pulse transmission in a fluid pressure medium under isothermal conditions. The pressure derivative of the bulk modulus, dK/dp, computed from the linear fit shown in Figure 1 is approximately 4.4x10\textsuperscript{3}.

Figure 1. Ultrasonic velocities and bulk modulus for Nugget sandstone at elevated hydrostatic pressure.

The sandstone stiffens as the softening effect of cracks is eliminated by applying normal stress. Various idealized models, which usually simplify the geometry of the cracks, have been proposed to explain this effect. Budiansky and O'Connell (1976) give a method for calculating modulus reductions resulting from introducing thin cracks into an intact solid. The infinitesimal stress-strain relations depend on crack density. For example, the pressure-volume relation for elliptical cracks takes the form:

\[ dP = \left[ \frac{16K}{9} \left( \frac{1-V_s^2}{1-2V_p} \right)^\frac{1}{2} \right] dV \]

where \( K \) is the bulk modulus of the intact solid, \( \varepsilon \) is a crack density parameter, which equals the product of number density and a shape factor, and \( V_p \) and \( V_s \) are the effective bulk modulus and Poisson's ratio of the cracked solid which are also functions of the crack density. Expressions of this type predict nonlinear response for any finite disturbance because the crack density parameter is no longer constant. When pressure increases, the crack density decreases as cracks close, and the effective Poisson's ratio of the composite also changes. Both of these effects combine to introduce a strong hardening nonlinearity into the stress-strain relation.

Microscopic sliding

The clearest evidence for understanding nonlinearities associated with sliding comes from experiments performed on macrofractures. Fractured granite samples show several indications of strong nonlinear response at low frequencies, including load dependent modulus, amplitude dependent attenuation, hysteretic loading history and the transfer of energy to higher frequencies (Bonner and Wanamaker, 1991). The hysteretic nonlinearity can be either of the hardening or softening type, depending on stress state. Different stress/strain laws apply for normal and shear deformation. Although the stress singularity that is inherent to the tip of a sharp crack insures that partial slip occurs for any finite amplitude disturbance, sliding becomes dominant only for strains greater than \textasciitilde10\textsuperscript{-6}. Mavko's (1979) model of sliding on a rough crack explains the approximately linear increase of amplitude dependent attenuation observed for strains > 10\textsuperscript{-6}. The theoretical result is scale invariant; e.g., it applies to cracks from grain scale to discrete discontinuities. Experiments show that nonlinear attenuation in granites increases when microcrack density is increased by appropriate heat treatment or cyclic fatigue (Bonner and Wanamaker, 1990).

Movement of pore fluids

The change of compressional velocity for partially saturated volcanic rock in the direction perpendicular to uniaxial loading (Bonner and Wanamaker, 1990) is plotted in Figure 2.
Figure 2. Travel time increase with uniaxial load for partially saturated Butte lapilli tuff.

The observed increase in travel time is greater than that expected for the increase in path length calculated for the Poisson effect. The sensitivity of velocity to stress, which is approximately two orders of magnitude higher than usual (Figure 1), and the change of slope with small changes of saturation suggest that additional nonlinearity is associated with partial saturation. The cause of the apparent softening nonlinearity is not known, but it is clear that at least at larger loads, velocities become increasingly anisotropic as cracks oriented perpendicular to the loading direction close. The softening shown in Figure 2 may occur as cracks parallel to the loading direction are inflated by fluid expelled from cracks oriented perpendicular to the loading direction. This behavior is consistent with that predicted by expressions such as Eq. (1), although expressions for linear compressibilities in an anisotropic cracked solid would apply in this case. It is also worth noting that the movement of the contact line that develops between liquid and gas in a partially saturated solid is nonlinear. Contact angle hysteresis and stick slip occur as the contact line reverses direction. Motion of the contact line will contribute to the deformation of the saturated solid, and if the contact lines move by a series of arrested chain reactions as described by Jansons (1985), a softening nonlinearity such as that observed will occur.

CONCLUSION

A variety of mechanisms can contribute to the observed nonlinear stress-strain response of rock, depending on the amplitude of the exciting stress, and the fluid content of the pore space. At strains of order $10^{-7}$, crack closure causes a hardening nonlinearity. For strains approaching $10^{-6}$, frictional sliding on rough surfaces becomes increasingly important. The nonlinearity can be of the hardening or softening type, depending on the relative size of normal and shear stresses at the crack. Movement of pore fluids in response to nonhydrostatic deformation may cause a softening nonlinearity in particular directions for partially saturated rock. Movements of the contact lines that develop between liquid and gas in partially saturated rock may contribute an additional softening nonlinearity.

ACKNOWLEDGEMENTS

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ELASTIC ANHARMONICITY IN COMPOSITE PLATES

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INTRODUCTION

Elastic properties of piezoelectric composite materials have been mostly studied in the linear regime [1,2]. In the low frequency range it has been found that vibrational properties can be understood in terms of acoustic wave resonances due to successive reflections on boundaries between two different materials [2]. The eigenmodes of composite structures have been optically measured and good agreement has been found between experimental and calculated patterns, obtained by using a previously developed model both for periodic [3] and aperiodic composites [4]. While the linear properties of composite piezoelectric structures are well understood, there are relatively much less information about their nonlinear properties. The aim of this paper is the study of the response of such systems in resonant regime, when subject to high dynamical stress.

EXPERIMENT AND RESULTS

We experimented with periodical composite samples prepared by slicing in parallel stripes of width \( l_2 = 0.9 \) mm and \( l_2 = 1 \) mm thin piezoelectric ceramic plates (PZT5), of dimensions 26x14x0.4 mm. The stripes of width \( l_2 \) have been removed and substituted by a soft epoxy resin. (The parameters of the two materials are: \( \rho_1 = 7.65 \times 10^3 \) kg/m\(^3\), \( \rho_2 = 1.17 \times 10^3 \) kg/m\(^3\), \( v_1 = 3200 \) m/s, \( v_2 = 1800 \) m/s; \( v_1 \) e \( v_2 \) are effective velocities of \( S_0 \) Lamb mode [2]). Aluminum electrodes have been sputtered on the plate surfaces which are perpendicular on z-axis (see the right inset in Fig. 1).

The linear response of the plate at low level of excitation is reported in Fig. 1. A number of resonant modes are evidenced, that corresponds to resonances of the lowest Lamb modes [5] along the largest size of the plate. The sample was then linearly excited at these frequencies and the vibration patterns measured through a sensitive acoustooptical device. Some of the frequencies have been selected and the vibration amplitude of the plate at its center point, at these frequencies, was detected with the acoustooptical device, whose output was analyzed with a HP4194A spectrum analyzer. At low applied voltage, a typical spectrum presents the usual ±ω components.

Fig. 1 Experimental admittance curve versus frequency, with the first part of the spectrum enlarged in the left inset. The marked peak corresponds to the measurements presented in Figs. 2 and 3. (The two highest peaks correspond to symmetrical stopband edge resonances of the periodical plate).

By increasing the voltage, nonlinear effects appear in the spectrum, as harmonic (Fig. 2a) and subharmonic (Fig. 2b) modes; (The central frequency \( \omega_0 = 70 \) MHz)

Fig. 2 Frequency spectrum of the periodical sample excited at \( \omega = 250 \) kHz; voltage amplitude is (a) \( V = 5 \) V; (b) \( V = 26 \) V. The numbers identify the order of harmonics and subharmonics.
is the frequency of acoustooptical modulator and $\omega = 250$ kHz is a resonant mode of the sample; subharmonic generation is a threshold phenomenon [6], that suddenly appears at a certain value of the fundamental mode amplitude which corresponds in our experiment at the voltage value $V = 22$ V (Fig. 3). The amplitudes of the harmonic waves increase with power laws of the fundamental mode, as can be seen in Fig. 3. Similar effects have been obtained also with an aperiodic composite sample having the two materials distributed according to a four generation Cantor sequence [4] (Fig. 4). It can be observed from Fig. 2b and 4 that sum frequencies like $(\omega + \omega/2)$ are generated, too [7].

![Graph showing voltage amplitude versus frequency](image)

**Fig. 3** Harmonics and $1/2$ subharmonic amplitude versus voltage amplitude for the periodical sample excited at $\omega = 250$ kHz.

A theoretical model, based on a perturbative method, was used to calculate the threshold amplitude of the fundamental mode, in the approximation of considering it as a pure longitudinal mode, and the corresponding voltage [7]. The considered structure was an elastically nonlinear plate, where the piezoelectric material, which is highly nonlinear ($C_{111} \approx 100 C_{11}$), acts as an input gate for the energy flow. The nonlinearity in the resonator couples normal modes of the resonating structure, so that a net energy flow can be established from the driving frequency to higher (harmonics) or lower (subharmonics) modes. The stability condition for the subharmonic imposes a threshold condition for the fundamental mode amplitude. The threshold value is found to depend on strain overlapping of fundamental and subharmonic modes, as well as on frequency matching between subharmonic and an eigenmode of the structure (details about model and calculations are presented elsewhere [7]). A good agreement has been obtained between the calculated and experimental values of the threshold amplitude.

**CONCLUSION**

Nonlinear elastic properties of piezoelectric composite materials have been experimentally studied. Harmonics and subharmonics generation has been obtained both in periodic and aperiodic structures. Frequencies of the subharmonics has been found to be close to eigenmodes of the structures.

**ACKNOWLEDGEMENTS**

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ELECTRIC POTENTIAL IN PIEZOELECTRIC MEDIUM AND ITS INFLUENCE ON MEASUREMENT OF ULTRASONIC NONLINEARITY PARAMETER

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INTRODUCTION

In the experiment of second harmonic generation (SHG), the capacitive detector is used to measure the absolute amplitude of acoustic wave. [1] The assembly of the detector and sample can be simplified as shown by Fig. 1. Usually the capacitive detector is mechanical displacement sensitive. When the sample is piezoelectric, however, it has been observed that capacitive detector gives output even there is no DC-bias applied to it. [1] In the present paper, the origin of the no DC-bias output is discussed and its influence on the measurement of ultrasonic nonlinearity parameter is estimated. The calculation is compared with experiment. Although the analysis is done only for longitudinal wave along Z-axis of crystalline LiNbO₃, the procedure of the analysis can be easily extended to any piezoelectric medium as long as there exists a piezoelectric-stiffened wave in certain direction.

ELECTRIC FIELD IN AIRGAP

Crystal LiNbO₃ has symmetry of 3m. The longitudinal wave along its crystallographic Z-axis is piezoelectric-stiffened. Without loss of generality one dimensional problem is treated here because most of SHG experiments are performed for pure longitudinal wave direction. Hence the coupling equation of particle displacement u and electric potential φ can be expressed as:

$$\rho_0 \frac{\partial^2 u}{\partial t^2} = \frac{1}{c_{33}} \left( \frac{\partial^2 u}{\partial x^2} + \frac{\partial^2 \phi}{\partial y^2} + \frac{\partial^2 \phi}{\partial z^2} \right)$$

and

$$\frac{\partial^2 \phi}{\partial x^2} - \frac{\partial^2 \phi}{\partial y^2} + \frac{\partial^2 \phi}{\partial z^2} = 0$$

(1)

(2)

Moreover, in addition to incident and reflected waves there should be a reflected evanescent wave at in the piezoelectric medium in order to fulfill the boundary conditions. [3] φ₀ satisfies Eqs. (1) and (2) with the solution of u=0, and φ₀=0.

OPEN-CIRCUIT

In open-circuit case, the bottom surface of sample is unmetallized. The solution for reflected evanescent wave is

$$\phi_e = A_e (x-h)$$

(4)

where h is penetration depth of ϕ, i.e. ϕ=0 when x=h. The solution of one-dimensional Laplacian equation is

$$\phi_e A_e (x-d) \phi_e$$

(5)

which also satisfies that Aₖ₆=0 when r=0. A₆ is the potential developed at load Y. Use of boundary conditions gives:

$$-j k c \bar{A}_e + j k c \bar{A}_e B_e e_i A_e = 0$$

(6a)

$$\frac{\bar{A}_e}{\bar{A}_e} \frac{\bar{A}_e}{\bar{A}_e} = \frac{\bar{A}_e}{\bar{A}_e} \bar{A}_e = d$$

(6b)

$$e_i A_e = -e_i A_e$$

(6c)

In addition, the circuit equation

$$\phi_e + \frac{\partial (L \phi_e)}{\partial t} = 0$$

should be applied. Here, Aₖ₆ and Bₖ₆ are amplitude of incident and reflected waves, respectively, s is top surface area of detector button. By solving Eqs(6) the following results are obtained:

$$\frac{B_e}{A_e} = 1 - \frac{2 k_0^2 \phi^2}{(k_0 d_a)^2 + (k_0 d_a)^2}$$

(7)

$$\phi_e = \frac{2 j \omega}{\frac{(k_0 d_a)^2 + (k_0 d_a)^2}{k_0^2 \phi}}$$

(8)

Here

$$\theta = \frac{1}{\omega} \left( \frac{k d_a}{k_0^2 - k d_a} \right), x = 2 \omega \left( \frac{k^2 y}{d_a} \right) \frac{y}{y_0}$$

$$x = \frac{k^2 y}{d_a} - 2 \omega \left( \frac{k^2 y}{d_a} \right) \frac{y}{y_0}$$

$$k_0^2 = \frac{y_0}{y_0}$$

$$k_0^2 = \frac{y_0}{y_0}$$

$$k_0^2 = \frac{y_0}{y_0}$$

L is the electromechanical coupling factor.

SHORT-CIRCUIT

In the case of short-circuit, the bottom surface of sample is coated with good conductor film and grounded.
In the airgap electric potential is equal to zero because the grounded bottom surface shields electric field in the sample. Under this condition, the measurement of ultrasonic nonlinearity parameter $\beta$ for piezoelectric sample is the same as for nonpiezoelectric one.

THE EFFECT OF PIEZOELECTRIC POTENTIAL

In the case of open-circuit, the vibration of bottom surface caused by incident acoustic wave cannot be detected by capacitive detector no matter whether there is DC-bias applied to it or not because the circuit is open. Obviously, the output of the capacitive detector is caused by penetrating airgap potential. The detector becomes electric charge-sensitive in this case. It can be seen that output $q_o$ of capacitive detector is proportional to amplitude $A_o$ of incident acoustic wave.

In order to estimated the influence of penetrating potential on the result of measurement of ultrasonic nonlinearity parameter $\beta$ of a crystal $Z_{\text{LmB}_6O_{12}}$ sample is measured using the apparatus and procedure in I when its bottom surface has the different metalization extent i.e. the metal film coated on the surface has different conductivity. Table 1 gives measured results. The measured value of $\beta$ is 2.95 when DC-resistance of the film is very small (indicated by o in the table), which is in agreement with the previous one. The measured value of $\beta$ decreases with the metalization extent of the bottom surface becoming worse. Obviously, this is due to the effect of penetrating potential.

In a practical experiment the metalization extent of the bottom surface may be between completely short and open circuit. "Thin coated films" in Table I is corresponding to this situation. The capacitive detector becomes sensitive to both mechanical displacement and electric charge. The experimental manifestation of this fact is that the capacitive detector will give output without DC-bias applied to it, but the output increases markedly when DC-bias is applied. In the present experiment, it is found that the output with DC-bias is 1.4 times that without DC-bias for fundamental and 1.9 times for the second harmonic. If the calibration procedure in (i) is used to calculate amplitude of acoustic wave, the calculate value will be 1.71 times as high as that it should have for fundamental ($A_f$) and 1.53 times for second harmonic ($A_2$) because the procedure in (i) takes only mechanical displacement contribution to the output into account. Since ultrasonic nonlinearity parameter $\beta$ is proportional to $A_f/A_2^2$, the ratio of $\beta^2$ measured under "thin coating" to $\beta^2$ measured under 0-resistance is 0.52 which is in good agreement with the result indicated in table 1, where the ratio is about 0.5. Therefore, the electric potential in piezoelectric medium could affect measured result of ultrasonic nonlinearity parameter if metalization extent of bottom surface is not good enough to completely screen it. In other hand, the fact shows that the capacitive detector can also be used to detect the electric potential. Especially, the nonlinear acoustic field generated by nonlinearity of the medium is accompanied by corresponding nonlinear electric field which would be easily detected by capacitive detector if the bottom surface of sample is open-circuited. Moreover, experiment shows that the capacitive detector is more sensitive to electric charge than to mechanical displacement. Hence, a new approach to determine nonlinear constants of piezoelectric medium through direct detection of nonlinear electric potential may be presented.

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![Figure 1. Simplified assembly of capacitive detector and sample.](image)

<table>
<thead>
<tr>
<th>DC resistance of coated film</th>
<th>$\beta$</th>
<th>$C_{333}$</th>
<th>$(10^{11} \text{ N/m}^2)$</th>
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<td>-10.43</td>
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LABORATORY AND FIELD OBSERVATION OF PARAMETRIC BEAM FORMATION IN ROCK

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INTRODUCTION

The purpose of this paper is to show how nonlinear elastic waves generated within a material can be made to interfere and produce a difference frequency signal f, at distances where the primary signals have disappeared due to attenuation. In addition, we show results from low frequency attenuation studies using a torsional oscillator illustrating that a material can be made more nonlinear by inducing additional microcracks [1,2]. Both amplitude dependent attenuation and the strength of the nonlinear interaction are thought to result from the presence of microcracks. Therefore, if strain dependent attenuation can be increased by oscillating the sample in torsion at high amplitude, the relation between nonlinear attenuation and microcrack density is firmly established because new microcracks are produced by fatigue damage. Higher microcrack density should then improve parametric array formation by increasing the elastic nonlinearity. Lastly, preliminary field experimental results are shown which indicate that a parametric array can be used in imaging geologic strata.

EXPERIMENTAL METHOD AND APPARATUS

Experimental configurations for the experiments are described in this section. The experiments illustrating parametric array formation took place in a 183 cm-long sample of Berea sandstone. For the torsional oscillator measurements a 20 mm-long rod of Sierra White granite was employed. Because of their elastic uniformity both rock samples are standard materials in rock physics experiments.

Parametric Array Measurements

The two primary-wave signals are electronically summed inside a function generator, amplified, and fed into a single transducer. The transducer inputs two approximately planar, collinear waves propagating with separate frequencies, whose interaction creates the difference frequency signal. The detected signal is preamplified and recorded on a digital oscilloscope, and then relayed to a SUN IPC for processing.

Torsional Oscillator Measurements

The torsional oscillator experiments were designed to quantify the departure from linearity as a function of strain. In these experiments, a sample is oscillated in torsion over a range of strains from $10^{-7}$ to $10^{-4}$ at a frequency of 1 Hz (the torsional oscillator operates between 0.1 and 100 Hz). Shear attenuation is a direct function of the phase angle between torque and twist as a function of frequency and amplitude. For a complete description of the experimental configuration see Bonner and Wannamaker [1].

Field Experiments

In the summer of 1990 we began conducting Vibroseis® experiments to determine whether or not it may be feasible to image reflectors with a parametric array composed of Vibroseis® sources. In the largest experiment conducted to date, six Vibroseis® sources were employed. Rather than fix the primary wave frequencies the sources were swept as follows. Simultaneously three of the sources were swept upward in frequency while three were swept downward. Thus the nonlinearly-generated $f_\phi$ signal was the running difference frequency between the two sets of Vibroseis® sources. The $f_\phi$ signal between the two simultaneous sweeps was then cross-correlated with the received signal.

RESULTS

Parametric Array Measurement

Fig. 1 illustrates the strong parametric generation of a difference frequency signal in sandstone. The figure shows an example of the measured spectrum when the primary wave frequencies were fixed at 555 and 559 kHz, respectively, and thus the difference frequency was 44 kHz. It is remarkable that because of frequency dependent attenuation there is no trace of the primary wave signals while the difference frequency and the second harmonic of the difference frequency remain strong, despite the fact that the conversion efficiency between primary waves and the difference wave is on the order of 1% [2]. For this sample the specific dissipation $Q$, which is inversely proportional to attenuation, is approximately 70.

Figure 1. Parametric array result.

Torsional Oscillator Measurement

In Fig. 2 the phase angle between the torque and twist is shown for the granite as a function of shear strain amplitude, before cycling (dots) and after $10^7$ cycles (solid triangles) at a strain of $3\times10^{-5}$. The departure from linear to nonlinear elastic behavior is shown to take place at strains of greater than approximately $6\times10^{-5}$ for the sample of Sierra White granite. The strain sensitivity of the attenuation increases dramatically as a function of higher strains.

Field Results

Fig. 3 shows preliminary results from an experiment conducted in an active oil field in west Texas (USA). In this case, one set of Vibroseis® sources was swept from 90 to 50 Hz while the other set was swept from 50 to 90 Hz. The difference frequency was then 40-4-40 Hz, advantageously outside the band of the primary frequencies. A known reflector determined from the standard reflection seismic method was located along with the surface wave as shown by the arrows in the figure.

Figure 2. Torsional oscillator result.
Figure 3. Difference frequency seismic reflection profile from west Texas.

DISCUSSION

An encouraging result from our current work, shown in Fig. 1, is that producing a directed, nonlinear-wave source at the difference frequency may well be possible in the earth. Careful measurements were conducted using the identical experimental configuration but without the rock (transducers were face-to-face) to be certain that the difference frequency was created in the rock and not in the associated electronics and/or transducers.

The torsional oscillator experiment shows that introducing additional cracks into rock increases the nonlinearity of the material significantly. This work implies that damaging the near-source region should enhance parametric array formation which may be useful in Earth imaging studies. The possibility of using a fractured ceramic for creation of a parametrically-derived difference frequency signal outside of the material under study and then injecting it into the material is intriguing.

The Vibroseis® result indicates that it is possible to image at f. However, further research must be conducted in order to determine whether or not the majority of the difference frequency signal is created at the source or in the earth (Vibroseis® sources are notoriously nonlinear). This is important from the standpoint of directionality of the source: L created at the source is nearly omnidirectional while that created in the earth would be directional. We note that the using the standard reflection method this reflector was imaged with better signal/noise; the fact that a reflector was imaged with the difference frequency at all was an important indication of future possibilities. A further question to be addressed is whether or not deep reflectors can be imaged that may not otherwise be possible by study of parametric formation below standard frequencies where Vibroseis® sources operate i.e., below about 5 Hz.

CONCLUSIONS

Most notably, when two collinear, primary pressure waves were simultaneously injected into a 183-cm long sample at ultrasonic frequencies, we detected strong difference frequency signals across the sample but found that the higher-frequency primary waves had been entirely attenuated. We also showed that the nonlinearity of the material could be increased by inducing additional cracks. We conclude that the possibility of creating a low frequency, directional source by nonlinear elastic wave mixing is promising for laboratory and field applications.

ACKNOWLEDGEMENTS

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SECOND ORDER ACOUSTIC NONLINEARITIES

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The so-called nonlinearity parameter which is the combination of second and third order elastic constants has been extensively studied for crystalline solids [1] and has been related to microscopic behavior of the solid, e.g., a harmonic behavior of the interatomic potential as well as to macroscopic quantities e.g., thermal expansion, residual stress and material hardness. There are several indications that higher than third order elastic constants could play a role in material characterization, e.g., fourth and fifth order elastic constants enter in the temperature dependence of velocity as functions of stress [2]. There are practically no reported results on fourth order elastic constants available. The stress dependence of the velocity when carried out up to second order offers some possibilities for obtaining second order nonlinearities, and hence perhaps fourth order elastic constants. The objective of this paper is to report some results that correlate stress dependence of second order elastic constants to third and fourth order elastic constants and to demonstrate the possibility to measure second order nonlinearities.

THEORETICAL CONSIDERATIONS

The sound velocity in a stressed solid may be expressed

\[ C(\sigma) = C_0 + C_1 \sigma + C_2 \sigma^2 + \ldots \]  

(1)

where \( C_0 \) is the velocity in the unstrained material. \( C_1 \) is actually the first order acousto-elastic constant and is a combination of the second and third order elastic constants that can be obtained. \( C_2 \) may be called a second order acousto-elastic constant and it may be given as the combination of second, third, and fourth order elastic constants and related to second order nonlinearity parameters. Depending on the direction of \( \sigma \) with respect to the wave propagation and polarization direction, both \( C_1 \) and \( C_2 \) have different forms. Using finite deformation theory [3] we have calculated the second order effective elastic constants for isotropic and cubic media subjected to either uniaxial compression or shear stress [4]. In case of an applied simple shear stress for example the stress dependent second order elastic constant \( C_{11}(\sigma) \) may be given as

\[ C_{11}(\sigma) = C_{11}(0) + \frac{1}{2}(\sigma^2/C_{12}) [4C_{12} + 8C_{44} + C_{112} + C_{110}] \]  

(2)

Notice that because of the symmetry due to applied shear stress only second order nonlinear term is present (stress square dependence). On figure 1 the calculated stress dependent quas longitudinal wave velocity is plotted for fused quartz as function of pure shear stress.

Polarization Technique [5]

When an isotropic solid is prestressed under uniaxial tension or compression a slight anisotropy is introduced in the material. As a result of this anisotropy, the velocity of the shear wave will depend on the polarization direction with respect to the applied stress. This phenomenon is known as "acoustic birefringence." By designating the stress dependent shear wave velocity with parallel polarization to the applied stress \( C^{(0)} \) and the shear wave velocity with perpendicular polarization to the applied stress by \( C^{(0)} \), Eq. (1) may be written as

\[ C^{(0)} = C_0^{(0)} + C_1^{(0)} \sigma + C_2^{(0)} \sigma^2 \]  

(3)

and

\[ C^{(\perp)} = C_0^{(\perp)} + C_1^{(\perp)} \sigma + C_2^{(\perp)} \sigma^2 \]  

(4)

In Eqs. (3) and (4) it is assumed that the shear velocities are not the same at both polarization directions at zero stress, i.e. there is some texture in the material. Now consider an ultrasonic shear wave polarized in an arbitrary angle \( \theta \) with respect to the axis of the applied stress.

The received birefringent signal may be written as

\[ u = \text{cos}^2 \theta \text{ exp} \text{i} \omega 2D/(C^{(0)} + \text{sin}^2 \theta \text{ exp} \text{i} \omega 2D/C^{(\perp)}) \]  

(5)

where \( D \) is the sample thickness. Here, the amplitude variation of the two signals is expressed only in terms of the direction of polarization. In order to measure the weak second order effect at moderate stress levels, an optimal angle for shear wave polarization is selected where the stronger first order effect can be reduced or eliminated. At this angle the velocity may be written as

\[ C_0 = C_{\perp 0} + C_{\parallel 0} \sigma^2 \]  

(6)

and the required angle is

\[ \tan \theta = -C_{\perp 0}/C_{\parallel 0} \]  

(7)

Equation (7) implies that the required condition to find a polarization angle \( \theta \) is that the slope of the two stress dependent velocities has to be of opposite sign. The zeroth and second order coefficients from equation (6) may be given as

\[ C_{\perp 0} = C_{\perp 0} \cos^2 \theta + C_{\perp 0} \sin^2 \theta \]  

(8)

and

\[ C_{\parallel 0} = C_{\perp 0} \cos^2 \theta + C_{\perp 0} \sin^2 \theta \]  

(9)

The experimental objective is to find an angle of polarization predicted by Eq. (7). In our experimental system an ultrasonic contact transducer at 5 MHz (either longitudinal or shear) is coupled to the aluminum block. The transducer is operated in pulse-echo mode and the received echoes are displayed on a LeCroy 9400 digital oscilloscope. The aluminum block was loaded in uniaxial compression and tension ranging from -19 KN to 11 KN in increments of 1 KN. The
ultrasonic velocity was measured as a function of applied stress.

In Fig. 4 shear wave velocity for the case of parallel and perpendicular polarization is given as the function of applied stress. The corresponding first order acousto-elastic constants are obtained as

$$C'_1 = -0.459 \text{ m s}^{-1} \text{ kst}^{-1}$$  \hspace{1cm} (10)$$
and

$$C''_1 = 0.129 \text{ m s}^{-1} \text{ kst}^{-1}$$  \hspace{1cm} (11)$$

From Fig. 4 it should be noticed that at zero stress the two shear velocities are not the same indicating some texture in the material. In order to measure the second order acousto-elastic coefficient $C_{o2}$ the polarization angle was determined from Eq. (9) by substituting the measured values of $C'$ and $C''$. The polarization angle at which the first order effect is minimized for this aluminum is 20°. The shear wave velocity with 20° polarization angle was measured as the applied stress in the aluminum sample. The experimental result is plotted in Fig. 5. A second order polynomial "best fit" routine was used to determine the relationship between the stress dependent shear wave velocity and the applied stress resulting in the following equation:

$$C_{o2}(0) = 3174 - 0.001555\sigma$$  \hspace{1cm} (12)$$

where 3174 m/sec is the zero stress shear velocity with 20° polarization angle relative to the applied stress. The second order acousto-elastic coefficient $C_{o2} = 0.001555 \text{ m s}^{-1} \text{ kst}^{-1}$. The coefficient $C_{o2}$ is a combination of second, third, and fourth order elastic constants in the form predicted by equation (1).

CONCLUSION

In this paper we have introduced the concept of second order nonlinearity by considering the acousto-elastic effect up to second order. When a small amplitude ultrasonic wave propagates in a material which is under a pure shear stress the velocity will depend on the square of the applied stress. The coefficient will include second, third, and fourth order elastic constants of the unstressed material. By minimizing the effect of the shear wave "birefringence" (used to measure first order nonlinearities) due to stress induced anistropy at an optimized polarization angle, second order nonlinearity was measured in aluminum.

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FOURTH ORDER ELASTIC MODULI OF DIAMOND STRUCTURE MATERIALS

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1. INTRODUCTION

One of the most successful schemes for interpreting the elastic properties of diamond structure materials is the Keating model. In the formulation of the strain energy in this model, the quadratic contribution to the energy requires only two adjustable parameters, a bond stretching force constant $\alpha$, and an angle bending constant $\beta$, while the cubic contribution contains only three force constants, a bond-stretching constant $\gamma$, an angle bending constant $\delta$, and a mixed bond-stretching angle-bending constant $\epsilon$. Indeed, the comparison between experimental results and calculated values shows a gratifying agreement for the second order (SOEM) and third order (TOEM) elastic moduli.

Examining the numerical values of the various Keating model parameters for diamond, Si, and Ge, it is immediately obvious that the ratio of the bond-bend to the bond-stretch parameters decreases as one moves to higher orders. This is shown in Table I, where the ratios $\beta/\alpha$, $\delta/\gamma$, and $\epsilon/\gamma$ for diamond, Si, and Ge are presented. Assuming that this trend continues to the higher order force constants as well, it is reasonable to expect that the fourth order force constants will be dominated by the bond-stretch parameter, $\chi$. This is also supported by the fact, that at 0 K the TOEM of Si and Ge tend to obey the Cauchy relations better than at higher temperatures. The above assumption will be utilized in order to evaluate the fourth order elastic moduli (FOEM) of the diamond structure elements with the help of the Keating model.

2. CALCULATIONS

Following the procedure devised by Keating, the elastic moduli, up to fourth order, may be expressed in terms of $\alpha$, $\beta$, $\gamma$, $\delta$, $\epsilon$, and $\chi$. The expressions thus obtained for the SOEM and TOEM are identical to those presented by Keating. The values of the 11 independent FOEM in terms of the various force constants are shown in Eq. (1).

\[ C_{1111} = 6\gamma - 6\delta + 16\chi + 6C_{1111} = -54\epsilon + 16\chi \]
\[ C_{1112} = 3\gamma - 3\delta + 16\chi - 3C_{1112} = -3\epsilon + 16\chi \]
\[ C_{1122} = 2\gamma - 2\delta + 16\chi - 2C_{1122} = -2\epsilon + 16\chi \]
\[ C_{1123} = \gamma + 3\delta + 16\chi - C_{1123} = 3\epsilon + 16\chi \]
\[ C_{1222} = 4\gamma(5\delta - 4) + 6\delta(7\chi + 4) - \epsilon(2\chi^2 + 2\zeta - 1) + 16\chi(\zeta - 1)^3 + 2\zeta c_{123} \]
\[ C_{1223} = 4\gamma(5\delta - 4) + 6\delta(7\chi + 4) - \epsilon(2\chi^2 + 2\zeta - 1) + 16\chi(\zeta - 1)^3 + 2\zeta c_{123} \]

3. DISCUSSION

In order to evaluate the FOEM presented in Eq. (1),
one experimental measurement is required, since only one fourth order force constant, \( \chi \), is assumed in the present formulation. The most straightforward procedure would be to utilize high pressure lattice constant data, readily available from measurements in the diamond anvil cell. Unfortunately, both Si and Ge undergo a phase transition at relatively low pressures of 12.5 and 8 GPa respectively, hence the equation of state cannot be used for the evaluation of the FOEM. Another readily available experimental datum is the volume dependence of the Grüneisen constant, i.e. the second order Grüneisen constant,

\[
q = \left( \frac{\partial \ln \gamma_0}{\partial V} \right)_T = -B_T^T \left( \frac{\partial \ln \gamma_0}{\partial P} \right)_T
\]

(3)

where, \( V \) is the volume, \( P \) the pressure, \( \gamma_0 \) the Grüneisen constant, \( B_T^T \) the isothermal bulk modulus. Such data may be obtained from adiabatic measurements,\(^6\) pressure dependence of the thermal conductivity,\(^7\) or pressure variation of the thermal expansion,\(^8\) shock wave investigations,\(^10\), thermodynamic considerations via the Grüneisen-Anderson constant,\(^11\), as well as various theoretical models.\(^12\) Also some theoretical considerations\(^13\)-\(^15\) indicate that \( q = 1 \). Assuming the validity of the continuum anisotropic model,\(^16\) and the high temperature limit (a temperature of the order of the Debye temperature or higher), \( q \) may be expressed in terms of the SOEM, TOEM and FOEM.\(^17\)

\[
q = \frac{B_T^T}{18N\gamma_0} \sum_{p,N} \frac{\delta_{a_mN}}{w_p(N)} u_p(N) u_p(U,U_s - \delta_v)
\]

\[
+ \delta_{i_i(\text{stv})} + 2c_{i \text{stv}N} N_j N_l - \frac{1}{w_p(N)} + \frac{1}{2c_{i \text{stv}N} N_j N_l} (U,U_s - \delta_v) + c_{i \text{stv}N} u_p(N) u_p(U,U_s - \delta_v)
\]

(4)

Lower case Roman subscripts are Cartesian indices (i,j,.. = 1 to 3), the Einstein convention of summation over repeated indices being implied. \( N \) and \( U \) are the propagation and polarization vectors, \( N \) the number of propagation directions, \( \delta_v \) the Kroncker delta, and \( u_p(N) \) is given by

\[
u_p(N) = c_{i j k N} N_j N_k U_i
\]

(5)

The summation is to be carried out over the irreducible part of the Brillouin zone, i.e. in our case over the spherical triangle whose apexes are the [100], [110] and [111] directions.

Since all FOEM depend linearly on \( \chi \), they may be expressed as

\[
C_{i j k l m n r} = C_{i j k l m n r} + \chi C_{i j k l m n r}
\]

(6)

where \( C_{i j k l m n r} \) and \( C^{\prime}_{i j k l m n r} \) depend only on known parameters, and can therefore be evaluated. Inserting eq. (6) into (4), we obtain

\[
q = q' + \chi q''
\]

(7)

\( q' \) and \( q'' \) depend on the SOEM, TOEM, \( C_{i j k l m n r} \) and the propagation direction, and may be calculated. Hence, once a value for \( q \) has been assumed, \( \chi \) may be computed, its value plugged back into eq. (4), thus deriving the values of the FOEM. In Table II the calculated values of \( \chi \) for various sets of \( q \) are displayed, and Table III presents the values of the FOEM for the same \( q \) values. Comparing the magnitudes of the third order bond-stretch force

<table>
<thead>
<tr>
<th>Table II. Values of the fourth order force constant ( \chi ) for various ( q ) values. (Units are GPa.)</th>
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</thead>
<tbody>
<tr>
<td>Diamond</td>
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<tr>
<td>( q=1 )</td>
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<td>445</td>
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\*Reference 11 \*Reference 8

<table>
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<tr>
<th>Table III. Values of the FOEM for different ( q ) values. (Units are GPa.)</th>
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<tr>
<td>Diamond</td>
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<tr>
<td>( q=1 )</td>
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constant \( \gamma \) with its fourth order counterpart \( \chi \), it is obvious that \( \chi \approx \gamma / 3 \), which is contrary to the intuitive expectation that the force constants should increase with ascending order. Examining Table III, and comparing the FOEM presented there with the analogous quantities for the alkali halides,\(^3\) it is evident that the two sets are very different in character. From Table III, it is also evident that except for \( C_{i111} \), the FOEM are not strongly dependent on the choice of \( q \); thus, one may expect the values of the FOEM in Table III to be a reasonable estimate for the actual magnitude of the FOEM.

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APPLICATION OF NONLINEAR ACOUSTICS TO NONDESTRUCTIVE MATERIALS CHARACTERIZATION AND EVALUATION

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As a consequence of Murnaghan's nonlinear theory or elasticity, two acoustic effects can be predicted. The first is a linear dependence of the acoustic sound velocity on the applied stress, the so-called acoustoelastic effect,7 in which the propriety of a notch in a structure as an external load is applied. The second effect is found practical applications such as in bolt testing devices.5 More recently, applications to the analysis of complex stress states have been studied. An uniformly inhomogeneous stress field is produced due to the presence of a notch or crack in a structure as an external load is applied. Kino et al.4 used ultrasonic waves to map out such a stress field. The change in velocity is determined by the sum of the two principal stresses. On the other hand, using two shear waves with orthogonal polarizations provided information on the differences in the principal stresses or the maximum shear stress, t, a measurement technique known as acoustoelastic birefringence.6 In both cases,4,5 such measurements yield a "stress intensity factor" which provides information on the driving force on the crack.

The second effect predicted is acoustic harmonic generation in which case energy, oscillating at the frequency of the fundamental wave, is pumped into the second and higher harmonics.6 The strength of the effect depends on the driving (fundamental) amplitude, frequency, the propagation distance of the waves, and the second and third order elastic constants.7 Effects of reflections from specimen boundaries on second harmonic generation have to be taken into account, however.8,9

Alloying affects the second and third order elastic constants, and thus both the acoustoelastic effect and harmonic generation will change, as has been observed10 by hydrogen additions to a niobium single crystal. In this case, the differences of the third order elastic constants were found to be roughly 5-10 times larger than those of the second order elastic constants. The effect of carbon on the acoustoelastic effect of four different steels was determined by Heyman et al.11 The carbon was located in interstitial sites as the hydrogen in niobium. In contrast, metallic additions in high strength aluminum alloys are on substitutional sites. By thermal aging after quenching from the solid solution phase, second phase particles precipitate and grow. The "acoustoelastic constant" as well as the "nonlinearity parameter" change12 strongly with volume fraction of second phase present.

A large number of phase transformations occur by spontaneous or diffusionless transformations connected with a large change in the second order elastic constants (as well as a large attenuation effect). One would expect large nonlinear effects, since the atomic potential is heavily distorted.13,14 Hikata et al.15 introduced the idea that the displacement due to the bowing-out of dislocations is a contributor to harmonic generation. Therefore, the use of motion by dislocations does provide for nonlinear acoustic effects due to the strong deviation from elasticity as plasticity arises. Similarly, Richardson15 pointed out that an unbonded interface, subjected to a sufficiently intense incident acoustic wave, acts as a harmonic generator. Experiments on fatigue cracks16 using acoustic surface waves confirmed this prediction. The generation is most efficient when the fracture surfaces touch lightly and increases with the growth of the cracks. If the cracks are completely closed, due to a compressive stress, or fully open, due to a tension stress, harmonic generation disappears as one would expect from this model.15

In summary, the influence of microstructural changes provides a rich field of research opportunities in nonlinear acoustics. From a nondestructive materials characterization point of view, measuring harmonic generation has advantages over techniques using the acoustoelastic effect since stress does not have to be applied. However, an absolute determination of the wave amplitude is necessary. To date, a number of techniques have been employed to measure strain17 and displacement amplitudes of ultrasonic waves. The capacitive detector being the most common. However, due to the experimental difficulties encountered in using the capacitive detector, this technique has not gained popularity outside the laboratory setting.

EXPERIMENTAL PROCEDURE

It has been shown that contact transducers with high conversion efficiency can be used for absolute displacement amplitude measurements if they are first calibrated using a broadband pulse-echo technique.20,21 However, the transducer/ couplant arrangement must be reciprocal, since the calibration signal passes through the receiving transducer twice. Obviously, a linear operation of the system is a necessary condition for this reciprocity to hold. The system, shown in Fig. 1, is based on a 33MHz microcomputer communicating with a bus with an oscilloscope. A function generator tuned to the fundamental frequency supplies the c.w. to a gated amplifier, providing the high-power monochromatic toneburst for the harmonic generation measurement. The transducers are spring loaded and use LiNO3 wafers coupled to the sample with a light oil. The transducers are mounted on movable "stages" that support the sample and allow removal of either transducer without disturbing the sample or the other transducer.

Figure 1. Harmonic generation system.

During calibration the "transmitting" transducer is not in place. In this case a low level broadband pulse (to ensure linearity) is
produced by a pulser/receiver and fed to the "receiving" transducer. By determining the input voltage and current and those due to the first back wall echo, a "calibration constant" is obtained that is used in the actual measurements to convert the output current to a displacement. The first, \(|A_1|\), and second harmonic, \(|A_2|\), amplitudes are determined independently with details, particularly on the conversion of acoustic to electrical power and vice versa, given in Ref. 21. An independent test, we are now measuring the fundamental amplitude on the receiver side of the specimen directly using a broadband laser interferometer. Preliminary results are encouraging and will be presented elsewhere.

RESULTS AND CONCLUSIONS

Evaluation of the system's performance was achieved on fused silica based on its lack of microstructural contributions to the nonlinearity and its low attenuation. As expected \(|A_2|\), the absolute value of the second harmonic amplitude increases linearly with \(|A_1|^2\), or the square of the fundamental amplitude, yielding a "nonlinearity parameter" \(\beta\):

\[
\beta = \frac{\nu^2 |A_2|}{\omega_1^2 |A_1|^2}
\]

where \(\nu\) is the longitudinal ultrasonic velocity, \(\omega_1 = 2\nu_0\) (\(\nu_0\) = fundamental frequency) and \(x\) is the propagation distance. In this present case, we obtained a value of \(|\beta| = 12.4 \pm 0.2\), comparing very favorably with other results, at least above an \(|A_1|\) of about 0.5 mm (-5\(\AA\)). As shown in Fig. 2 below.

Figure 2. \(\beta\) versus \(|A_1|\).

this amplitude there is some leakage of harmonically distorted fundamental c.w. from the signal generator unit into the receiver-side electronics. This leakage is considerably reduced in metallic specimens, such as Al6061-T6, where we obtained a \(|\beta| = 4.5 \pm 0.2\), again comparing favorably with a literature value.21 We note that the Al6061-T6 results have been obtained on two different runs, between which the sample/transducers combination was totally disassembled and recalibrated. Measurements of the \(x\) and \(\omega_1\) dependence of harmonic generation are in progress.

Based on these preliminary results, it is believed that a viable new method to measure materials' nonlinearities will promote a wider acceptance for materials characterization and nondestructive evaluation.

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MATERIALS CHARACTERIZATION USING NONLINEAR ACOUSTICS

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Nonlinear acoustics is a broad topic touching on many aspects of solid state science and technology including nonlinear dynamics, and lattice anharmonicity, thermoelasticity, and nondestructive materials characterization. Nonlinear acoustical phenomena are manifested in many ways in material systems and can be studied, for example, using the methods of acoustoelasticity, ultrasonic harmonic generation techniques, or, as we have recently shown, from a consideration of the stationary (static or "dc") waveform generated by the material nonlinearity in the presence of a driving acoustical wave. The emphasis in this paper is placed on the relationship of the stationary waveform to the thermal expansivity of the material.

NONLINEAR STATIONARY WAVES

We consider the propagation of an elastic wave in a lossless semi-infinite solid of arbitrary crystalline symmetry. The nonlinear equations of motion along a given propagation direction may be transformed into the form [1]

$$\frac{\partial^2 \xi}{\partial \tau^2} = c_s^2 \left( 1 - \frac{\partial \xi}{\partial \tau} \right) \frac{\partial^2 \xi}{\partial \eta^2}$$

(1)

where $\xi = \xi(\eta)$ is a mode index representing a wave of polarization $\eta = 1,2,3$ and direction of propagation $N$, $\sigma$ is the Lagrangian coordinate transformed such that it is always along the direction of wave propagation, $\tau$ is time, $c_s$ is the particle displacement for mode $\sigma$, $c_s$ is the "linear" wave speed, and $\beta_2$ is the modal nonlinearity parameter of the solid.

Although the harmonic wave solution to Eq.(1) is well known [2-4], a nontrivial, stationary (static or "dc") solution also exists that provides useful and sometimes unique information regarding material nonlinearity. Cantrell [1] obtained a stationary solution to Eq.(1) in the form (the acoustic radiation-induced static strain)

$$\frac{\partial \xi}{\partial \tau} = \frac{\beta_2}{4 \mu_s} \xi^2$$

(2)

where $\xi^2$ is the energy density of the propagating acoustic wave and $\mu_s = \rho c_s^2$ (mass density).

THERMODYNAMICS AND ACOUSTIC FIELDS

Eq.(2) predicts the existence of an acoustic radiation-induced static strain ("dc" level) in the crystalline solid for each mode of acoustic wave propagation; the strain is either dilative (positive) or contractive (negative) depending on the sign of the acoustic nonlinearity parameter for that particular mode. The experimental confirmation of such modal static strains [5,6], together with a recent model showing the relationship of the acoustic nonlinearity parameters to anharmonic lattice parameters [7], suggest that the static strains associated with thermal vibrations in the crystal lattice may be related directly to the thermal expansivity.

Such a connection can be obtained formally by considering the crystalline solid to consist of a large number $N_0$ of incoherent acoustic radiation sources which may be identified with the vibrating lattice particles of the crystal. It is useful to consider the Earnshaw [8] particle velocity solution to the nonlinear wave Eq.(1) which we write in the form

$$\frac{\partial \xi}{\partial \tau} = A \sin(\alpha_0 \eta + \phi_0), \quad \phi_2 = \kappa \beta_2 \frac{\partial \xi}{\partial \tau}.$$  

(3)

where $\kappa = \omega/\omega_c$ is the wave number. The energy density of the propagating nonlinear acoustic wave described by Eq.(3) is

$$E^2 = (1/2)\rho_0 \beta_2^2.$$  

In order to apply the above results to a system of incoherent radiation sources we must randomize the nonlinear acoustic field represented by Eq.(3). Cantrell [9] has performed such a randomization based on the methods of stochastic electrodynamics [10,11] by assuming: (a) there exists a fluctuating, nonlinear acoustic radiation field in the crystal at the absolute zero of temperature having an average energy per unit mass $E_0 = p_0\alpha^2 E_0$, (b) this zero-point radiation field fluctuates randomly as if it were produced by a large number of incoherent sources, and (c) at a finite temperature $T$ the averaged total energy per unit mass $\bar{E} = p_0\alpha^2 \bar{E}$, is composed of the sum of the zero-point field $E_0$ and a stochastically independent, temperature-dependent nonlinear acoustic radiation field $\bar{E}_T = p_0\alpha^2 \bar{E}_T$. The randomization leads to an expression of the Helmholtz free energy per unit mass $F$ as [8]

$$F = \sum \left( \frac{<E^2>}{k_B T} + k_B T \ln (1 - \exp(-<E^2>/k_B T)) \right)$$

(4)

where the angular brackets denote an average over random phases of the incoherent radiation field. Eq.(4) emphasizes the dependence of the thermodynamic state functions on the nonlinear acoustic modal energies $<E^2>$.

The thermal expansivity of a solid can be obtained from the Helmholtz free energy according to the expression

$$\alpha = \frac{S_{ijrs} \rho_0}{\sigma^{ijrs}} \frac{\partial F}{\partial \sigma_{ijrs}} = \sum S_{ijrs} C_{ij} \frac{G_p}{G_p^*} \beta_s$$

(5)

where $S_{ijrs}$ are the compliance coefficients of the material. The last equality in Eq.(5) follows from substituting Eq.(4) into Eq.(5). The $C_{ij}$ are the modal heat capacities and $G_p^*$ are constants determined by the modal directions.

It is clear from Eq.(5) that the thermal expansivity may be viewed as a temperature-dependent weighted average of the modal nonlinearity parameters of the crystal. The nonlinearity parameters may be positive or negative for a given mode, although for cubic crystals the nonlinearity parameters along the pure mode propagation directions always have been found to be positive. The thermal expansivity for most crystals is positive for most temperature ranges, although for some crystals, such as germanium and silicon, the thermal expansivity is found to be negative over a limited, relatively narrow, low temperature range.

The swing from positive to negative thermal expansivity is particularly pronounced in the amorphous solid vitreous silica (glass). The structure of vitreous silica is crystalline quartz-like up to a radius of a few atomic diameters from a given atomic site [12]; thereafter, the structure is characterized as amorphous giving rise to isotropic, pure mode, acoustical properties. Vitreous silica is known to have a large negative thermal expansivity at low temperatures where long wavelength vibrational modes dominate the dynamical properties. At higher temperatures, where the expansivity is positive, the short wavelength vibrational modes become more populated and the lattice dynamics is dominated by the local
quartz-like structure having positive nonlinearity parameters along the pure mode propagation directions. Although Eq.(5) strictly applies only to crystals, the implications of the equation may be extended with considerable caution to more complex structures. Thus, Eq.(5) suggests that the sign of the thermal expansivity of glass at high and low temperatures is reflected in the sign of the nonlinearity parameters appropriate to the atomic structure "seen" by the dominant lattice vibrations at a given temperature. This would imply that the sign of the nonlinearity parameter for long wavelength acoustic propagation is negative in vitreous silica. The negative value of the acoustic nonlinearity parameter has, indeed, been confirmed in vitreous silica (Suprasil) using 30MHz tonebursts [13].

ENVELOPE SOLITONS IN GLASS

Eq.(2) predicts that the static strain in materials possessing a negative value of the acoustic nonlinearity parameter, as Suprasil vitreous silica, has a polarity opposite to that of materials possessing a positive value of $\beta$. The predicted contractive nature of the acoustic radiation-induced static strain of vitreous silica has been confirmed experimentally using 30MHz acoustic tonebursts [5]. The contractive pulse itself, however, was found to give rise to the appearance of secondary peaks in the static (stationary) wave profile. It was also found that concomitant with the appearance of the secondary peaks is an amplitude-dependent velocity dispersion. It is well-known that the atoms in vitreous silica can sit in one of two closely spaced relative spatial positions giving rise to an effective double-well interatomic potential [14]. This two-level configuration is the apparent origin of a sufficiently strong relaxation to produce the observed velocity dispersion.

We consider theoretically the effects of dispersion on the generation and propagation of the static displacement waveforms by replacing the phase term $\frac{k}{2c} \frac{\partial u}{\partial t}$ in Eq.(3) by a more general "modulation" factor $\Phi_0(x,t)$, containing dispersive as well as nonlinear components. We assume that in general $A = A(x,t)$, and, hence, that $\hat{E}$ is not constant. We regard the resulting equation as representing the modulation of a continuous (carrier) waveform and postulate that the "local" frequency and wave number varies as $\omega(x,t) = \omega_0 + \Phi(x,t)$ and $k(x,t) = k_0 + \Phi_0(x,t)$, respectively. We also assume that the dispersion relation depends upon the wave amplitude as well as the wave number and expand the relation in a Taylor series as

$$\omega = \omega_0 + \frac{\partial \omega}{\partial k} k_0 + \frac{1}{2} \frac{\partial^2 \omega}{\partial k^2} k_0^2 + \beta_A A^2 + \cdots$$  \hspace{1cm} (6)

where $\omega_0 = \Phi_0(x,t)$, $\omega_0 = \frac{\partial \omega}{\partial k}$, $\beta_A = \frac{\partial^2 \omega}{\partial k^2}$, and $\beta_A = \frac{\partial^2 \Phi_0}{\partial k^2}$ is a measure of the nonlinearity, not to be confused with the nonlinearity parameter $\beta$ used above. Substituting the "local" frequency and wave number expressions into Eq.(6), we get

$$\frac{\partial^2 \Phi_0}{\partial k^2} + \frac{\partial^2 \Phi_0}{\partial A^2} + \omega_0^2 A^2 + \beta_A A^2 - \frac{\partial^2 \Phi_0}{\partial k^2} A^2 = 0$$  \hspace{1cm} (7)

Since Eq.(7) is a single equation with two unknowns, $A$ and $\Phi$, it is necessary to make another equation involving the same unknowns. An appropriate expression is the energy conservation equation:

$$\frac{\delta^2 A^2}{\partial t^2} + \frac{\partial}{\partial k} (c_A A^2) = 0$$  \hspace{1cm} (8)

where $c_A$ is the group velocity in the linear approximation. A special solution to Eqs.(7) and (8) is the solitary wave (soliton) form given by

$$A(x,t)^2 - A_0^2 = \frac{1}{2} \left[ \frac{\beta_A}{c_A} \right]$$  \hspace{1cm} (9)

where $A_0$ is the amplitude. Thus, Eq.(9) indicates that the initial energy distribution $(-A_0^2)$ evolves into a series of solitons. Since from Eq.(2) the acoustic static strain is directly proportional to the average energy density, the equations predict the initial radiation-induced static strain also to evolve into a series of envelope solitons.

CONCLUSION

The theoretical prediction and experimental confirmation of nonlinear stationary (static strain) waves is seen to provide useful fundamental information about crystalline solids that can be linked directly to their thermoelastic properties. Specifically, we have shown by randomizing the nonlinear acoustic radiation sources in the crystal that the thermodynamic state functions, hence thermal expansivity, may be expressed in terms of a sum over nonlinear acoustic modal energies. In more complicated structures, such as vitreous silica, the exact link between acoustic nonlinearity and thermal expansivity is likely to be more complicated than that given by Eq.(5), but the equation is still instructive and, perhaps, directive as well. Eq.(5), after all, does imply the existence of a negative nonlinearity parameter for fused silica at ultrasonic frequencies that is experimentally confirmed. Although dispersive effects are not directly incorporated into or specifically predicted by Eq.(5), such effects may be expected when the structural features of the material changes dramatically with scale as does glass. The effect of dispersion in glass on the static waves is seen to be quite substantial. The dispersion serves to catalyze the evolution of the initial static displacement signal profile into a series of envelope solitons. The mere existence of acoustic envelope solitons is itself a matter of some importance with deep-seated implications to the thermo-statistical physics of certain solids.

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CONSTANT FREQUENCY ACOUSTIC PULSE PHASE-LOCKED LOOP METHOD FOR MEASURING MATERIAL NONLINEARITY

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INTRODUCTION

Material nonlinearity manifests itself in various ways. One manifestation is the well known fact that an ultrasonic wave distorts as it propagates through the material. As a consequence of this the second harmonic amplitude increases as the propagation path increases. Equation (1) gives $\beta$, the nonlinearity parameter, in terms of measurable quantities,

$$\beta = \lim_{A_1 \to 0} \frac{B' \cdot k}{A_1^2 k_2}$$

where $A_2$ is the amplitude of the second harmonic, $A_1$ is the amplitude of the fundamental, $k$ is the propagation constant, $2\pi\lambda$, and $\lambda$ is the length of the propagation path. The nonlinearity parameter is directly related to combinations of second order and third order elastic constants.

Measurements of $B'$ from single crystals have a weak dependency on $A_1$. However, other materials show a stronger amplitude dependence. In engineering materials $B'$ shows a relatively large variation with fundamental amplitude. We have experimentally examined this variation as a function of fatigue level in the aluminum alloy, Al 2024 T4, and show the results in Fig. 1.

![Amplitude vs. Fatigue](image)

**Amplitude (Angstroms)**

Fig. 1. A plot of nonlinearity parameter of virgin Al 2024 and of fatigued Al 2024 vs. fundamental wave amplitude.

Figure 1 shows the dependence of $B'$ on the fundamental amplitude for an unfatigued and a fatigued sample of this alloy, subjected to 10,000 cycles of loading from 0 to 40 kilo pounds per square inch. A comparison of $B'$ for the two samples shows (a) a better than one hundred percent increase in $B'$ with increasing level of fatigue, (b) a transition to a higher value for $B'$ at higher acoustic drive amplitudes, and (c) a shift to the left, for the fatigue sample, of the acoustic drive amplitude at the transition takes place. These results suggest the existence of a strong microstructural feature in Al 2024 T4 that can exist in one of two states.

In order to investigate further the possibility of two state microstructures in Al 2024 T4 we have considered another manifestation of material nonlinearity based on the fact that a change in the pressure and temperature of a material results in a change of the sound velocity that is fundamentally dependent on the direction of propagation and polarization state of the ultrasonic wave. This dependence is also expressed in terms of combinations of second and higher order elastic constants. To pursue the investigation, we have developed an apparatus to measure the temperature dependence of the pressure derivative of ultrasonic velocities in materials.

THE CONSTANT FREQUENCY PULSED PHASE-LOCKED LOOP

For measurement of the change in the natural velocity we developed a new instrument based on pulsed phase-locked loop technology because of its inherent sensitivity to small velocity changes associated with modest pressure changes, thereby giving improved resolution with small pressure increments. A diagram of the constant frequency pulsed phase-locked loop (CFPLL) circuit that we developed for the pressure derivative system is shown below.

![Diagram of CFPLL circuit](image)

**Fig. 2.** A block diagram of the constant frequency pulsed phase-locked loop (CFPLL). Unlabeled triangles are buffers.

For an understanding of the CFPLL operation consider a constant frequency oscillator and two signal paths. Along the first path (the experimental measurement path) an acoustic signal is generated and traverses the sample. The second path, which includes a voltage controlled phase shifter (VCPS), is the reference path whose output is used for phase comparison with the first path. The control voltage to the VCPS is automatically changed until the output of the detector is zero volts (indicating quadrature between the signals). A calibration procedure that uses a line stretcher in the reference path permits the conversion of the control voltage change to a phase shift change between the paths.

Consider the block diagram of the instrument shown in Fig. 2. The output from the synthesizer is split to path 1 (measurement) and path 2 (reference). The splitter also feeds timing circuits for synchronizing all of the timing pulses. Path 1 includes progression from the frequency source to the gate which forms an adjustable width toneburst. After passing through the drive amplifier it is coupled to the transducer through the coupling / decoupling network (C/D).

The acoustic toneburst travels through the sample, is reflected at the opposite end, and impinges upon the transducer as an echo. Any change in the propagation conditions (e.g. velocity changes or path length changes) produces an associated phase change in Path 1.

After toneburst conversion to an electrical signal, the C/D routes the signal to the preamplifier and to an input of the phase detector.
Path 2 is the reference signal path from the frequency source to the phase detector. The input to the VCPS comes through a buffer from the power splitter. The VCPS output passes through a line stretcher (for system calibration) and buffer to the phase detector.

Phase comparison of the two paths is made by the phase detector, whose output voltage is one half the product of the input voltage amplitudes times the cosine of the phase differences between the two signals\(^2\). The output of the phase detector is passed to the sample and hold circuit, which selects the desired portion of the phase signal. The portion of the phase signal chosen for measurement is selected by a timing pulse to the sample and hold circuit whose output is passed to the integrator (I).

The loop control circuit, (the sample and hold, the integrator, the phase set point potentiometer, and the adder (V)) provides the control voltage for the VCPS. The control voltage is a sum of two voltages one of which comes from the integrator and the other from the phase set point potentiometer (PSPP). The PSPP sets the nominal phase shift about which the system operates. When the phase detector output reaches null, the integrator output voltage stabilizes thus indicating quadrature between paths 1 and 2. The phase shift control switch (PSCS) is normally in the "locked" state during measurement. The "unlocked" position is used to initially set the phase near quadrature.

The timing control section forms all of the timing signals for the various sections. The origin of all timing signals are determined by counting down from the synthesizer. Timing sequences are adjusted with an oscilloscope.

**MEASUREMENT TECHNIQUE**

The phase shift control switch is placed in the "unlocked" position. Using an oscilloscope, the frequency of the synthesizer is adjusted so that the phase comparison output voltage of the desired echo is relatively "fast topped" and stable. The PSPP is adjusted until the phase comparison output voltage corresponding to the desired echo is approximately zero volts. The control point is adjusted so that the phase signal is sampled well into its latter half. The PSCS is then placed in the "locked" position for data taking.

Calibration of the VCPS network is accomplished by keeping the system locked in the date-taking mode. With the experimental apparatus adjustments unchanged the calibrated line stretcher is adjusted to introduce a known phase shift into the calibration path. This results in a change in the phase shift control voltage to the VCPS. The change in phase shift control voltage is recorded and a change of output voltage for a corresponding phase shift is thus determined and is used to calculate the phase shift changes.

**MEASUREMENTS OF ALUMINUM 2024 - T4**

Room temperature measurements of the changes, \( \Delta w \), and the natural velocity, \( w \), were made with a system similar to that above described above as a function of applied hydrostatic pressure, \( p \), making certain that the sample temperature \( T \) was held constant during the run. The sample temperature was raised slightly and the measurements were repeated. The pressure limit for both of these was 1.38 MPa (200 psig). Fig. 3 is a plot of the

$$\frac{\Delta w}{p_{\text{at}} w}$$

as a function of pressure, calculated for room temperature. We observe constant values for pressures below 0.8 MPa and for pressures above 1.1 MPa. The value changes by approximately an order of magnitude within pressure range from 0.85 MPa to 1.025 MPa. The transition zone width is approximately 185 kPa. Sound pressure levels in the sample are estimated to be quite low for this system. As in Fig. 1 there is evidence of a two-state system.

**DISCUSSION**

Explanation of the experimental results shown in Figs. 1 and 3 is under investigation. Active consideration of dislocation interaction mechanisms are being explored because of the generation of large numbers of dislocations during the fatigue process reported in single cryogalts\(^3\). Two dislocation dipole interactions have shown promising results. Edge dislocation of opposite polarity can form dipoles which can exist in one of two equilibrium states. In one state the dislocations form a stable pair at an angle of 45° with respect to the Burgers vector, \( \vec{b} \), while the other is 135° with respect to \( \vec{b} \). The transition between the two states occurs by slip in the direction of \( \pm \vec{b} \). The external force necessary to activate a transition is proportional to a shear stress \( \sigma_y \). Within a polycrystalline material a finite amplitude compressive wave of an amplitude similar to those in Fig. 1 is sufficient to provide the resolved shear stress necessary to activate the transition. This mechanism can explain a two state system such as that observed in Fig. 1.

Perpendicular to \( \vec{b} \) are two equilibrium states for the edge dislocation dipole. These occur at -45° and +45° to \( \vec{b} \). The transition from one state to the other (climb) requires a tensile stress. The tensile stress required to activate the transition has been estimated from consideration of the variation of the chemical potential due to changes in vacancy concentration at the dislocation sites. The calculated value of 1.4 MPa is reasonably close to the experimentally assessed stress of approximately 1 MPa determined from Fig. 3. Further study of dislocation dynamics using the CFPPLL are under way.

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EXPERIMENTAL INVESTIGATION OF THE NONLINEARITY PARAMETER B/A OF WATER ENHANCED BY TRAPPED BUBBLES

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INTRODUCTION

It has been known for some years that the presence of bubbles may enhance nonlinear parameter B/A of a medium. Two distinct physical mechanisms are believed to be involved. The first one is due to bubble nonlinear oscillations; the second is driven by an acoustic wave near resonance frequency. Although a systematic theory in this regard is still needed, some theories estimate that the nonlinear effect of oscillations of spherical bubbles near resonance may increase the effective nonlinear parameter B/A for a bubble liquid by two orders of magnitude over that of either the gas or the liquid phase. The second mechanism is due to static properties of a bubbly medium; it can be explained, for example, by a mixture law for the acoustic nonlinear parameter B/A developed by Apel.1 For the second mechanism, theory predicts that the effective nonlinear parameter B/A of a bubbly medium is a strong function of the gas volume fraction. Xg, B/A quickly reaches the maximum value of about 9x10^7 from water value of about 5, then falls to air value of about 0.4. The second mechanism dominates at frequencies much lower than the resonance frequency of the bubbles.

An experimental measurement of B/A of a bubbly liquid to compare with the theories would be meaningful. The main reason that such a measurement has not yet been done, in our view, is due to the difficulty of establishing such a system. Spherical bubbles, in general, are quite unstable. Large bubbles will rise in a fluid due to the buoyant force. Small bubbles, on the other hand, may dissolve due to the high Laplace pressure caused by surface tension. A technique which uses hydrophobic polycarbonate thin membranes containing randomly spaced, fairly uniform cylindrical pores immersed in water to trap bubbles has been developed.24 Unlike free spherical bubbles, the cylindrical gas bubbles trapped by the membrane are quite stable. We have used these membranes to obtain information on frequency-response curves of trapped gas bubbles.24 Thickness of the membranes used in this experiment is 13 μm; the pore diameter is fairly uniform and equal to 4 μm; pore density is about 720/mm². The experimental results indicate that, when a hydrophobic micropore membrane is immersed in a liquid, air partially fills the pores. Upon exposure to ultrasound, the trapped bubbles' surfaces oscillate very much like a drumhead. The resonance frequency of the two-dimensional trapped bubble ensemble is a function of pressure amplitude of the incident ultrasound; the higher the pressure amplitude, the lower the resonance frequency. The shifting of the resonance frequency is found to be primarily due to the net displacement of the air-water-membrane triple-phase line and the air loss of the trapped bubbles. Those effects are caused by the radiation pressure and microbubble production induced by the incident ultrasound. Another observation relevant to this study was that the membranes without pores are essentially transparent to ultrasound in MHz frequency range.

The primary purpose of this study is to measure the effective B/A of a three-dimensional aqueous system which contains randomly distributed uniform-sized stable cylindrical bubbles. Such a system is achieved in this study by immersing into water a stack of many identical membranes with equal minimal water pressure at the nearest neighbors to trap bubbles. As pointed out earlier, the membrane material has been proven to be acoustically transparent at MHz frequencies. Furthermore, it was measured that density of the membrane is 1.06 g/cm³, the speed of sound in the membrane is 1.5x10^3 m/s, thus, as far as the acoustic characteristic impedance is concerned, the membrane material matches water well. The membrane, besides holding the cylindrical bubbles, is very much like water for our experimental purpose.

EXPERIMENTAL METHOD

The experimental arrangement used in this study is shown in Fig. 1. The experiments were carried out in a water tank (dimension 56 x 36 x 30 cm³). The tank was filled with tap water and left standing for days in order to let large bubbles leave the liquid. In Fig. 1, the transmitter T is a broadband (in frequency range of 0.5 to 3 MHz) ceramic (PZT) transducer of 1.3 cm diameter (Staveley Sensors Inc., East Harford, CT, USA). The receiver receiver of 0.5 mm diameter active sensing area (NTR systems, Inc., Seattle, WA, USA), which has essentially a flat response curve in the frequency range of 1 to 10 MHz. During the measurements, T was driven by RF toneburst; the duration of the toneburst was chosen long enough to approximate cw conditions, yet short enough to prevent standing waves. Between the two transducers, a stack of the membranes (MM) was installed perpendicularly to the axis of the incident beam; the incident sound beam emitted by T impinged normally on the membranes; the trapped bubbles oscillate along the beam axis. The transmitted signal was received by R and amplified by a broadband amplifier and then sent to a computer. Fast Fourier transform (FFT) of the gain of close shot was done by the computer through a software program (Aasyst Software Technologies, Inc., Rochester, NY, USA) and the results were displayed in graphical form by a plotter. From the FFT analysis of the signal, the amplitudes of the fundamental and the second harmonic were determined.

Experimental Results

Measurements show that the waveform of the transmitted signal suffers significant deformation after passing through several membranes when the frequency of the incident wave was about equal to the resonance frequency of the trapped bubbles. During this measurement, the amplitude of the incident wave was 5.3x10^-3 Pa and the frequency was 0.95 MHz, which is close to the resonance frequency of the trapped bubbles at the amplitude of 5.3x10^-3 Pa.

Fig. 2 contains plots of FFT results for the above measurements. Fig. 2(a) is the FFT result for the case of no membranes; essentially it is a pure sinusoidal signal. The identical FFT result was achieved if several no-pore membranes were placed between the two transducers; it shows that water and no-pore membranes are alike for our purpose. Fig. 2(b), (c), and (d) are, respectively, for case of one, two, and five membranes. The second and the third harmonic appears. One observation of this figure is that the absolute amplitude of each harmonic component decreases as the number of membranes increases; meanwhile the ratio of the amplitude of the second harmonic to that of the fundamental increases as the waveform becomes more deformed.

It is essential to keep two nearest membranes some distance apart; plugging membranes directly on top of another may inhibit the water-air interface of the trapped bubble oscillating freely under influence of ultrasound. It may also cause large air pockets trapped by membranes. Measurements show that the spacing between two nearest neighbor membranes and that between two transducers do not have any significant effect on the wave deformation. For instance, as the spacing between the two nearest membranes reduced from 1 cm to 1 mm, there was no significant change of the measured amplitudes of the fundamental and the second harmonic. According to our opinion, that is because B/A for water and the attenuation coefficient of water are much smaller than those for the system which contains the ensembles of the trapped bubbles oscillating at resonance.

An expression for the second harmonic magnitude, p2(x), as a function of source distance, x, was derived by Fubini3 for a homogeneous system and extended by Thuras et al.4 based on the assumption that the linear attenuation coefficients of the fundamental and the second harmonic amplitudes are mutually independent and the rate of change with propagation distance of the second harmonic amplitude is the sum of changes caused by nonlinear generation of the second harmonic and by its attenuation. This expression of p2(x) is given by

\[ p_2(x) = (1 + B/A) \frac{1}{2}(P_0^2 - p_0^2) \exp(-\alpha_1 x) - \exp(-\alpha_2 x) \]

where p0(x) is the magnitude of the acoustic pressure of the fundamental at the source, p0 is the density of the medium, c is the speed of sound of infinitesimal waves of the medium, f is the fundamental frequency, \(\alpha_1\) and \(\alpha_2\) are linear attenuation coefficients of the fundamental and the
second harmonic respectively. When \((a_0 - 2a)x\) is small, Eq. (1) can be simplified as

\[
p_2(x) = [(2 + B/A)\psi(2a)\psi(\psi\psi\psi) + \lambda]e^{-\lambda x}.
\]  

(2)

This approximation introduces an error of about 1% when \((a_0 - 2a)x\) is 0.5.\(^7\) For our cases both \(a_0\) and \(a_0\) are large, but \(x\) is small. The condition of \((a_0 - 2a)x\) being small can be satisfied. But both equations of (1) and (2) are for a homogeneous system; due to the presence of bubbles, the system is obviously not homogeneous any more. On the other hand, this is the only existing theory we may apply. The uncertainty of both \(ao\) and \(ao\) prevented us assessing them individually within reasonable accuracy. Fortunately, the necessity of knowing \(ao\) and \(ao\) can be eliminated and \(B/A\) can be determined by plotting \(\ln(p_2(x)/p_0^2(x))\) versus \(x\) and extrapolating to \(x = 0\), i.e.,

\[
B/A = 2\ln[p_2(x)/p_0(x)] - \ln[p_2^0(x)/p_0^0(x)] - 2.
\]  

(3)

Fig. 3 is a plot of \(\ln[p_2(x)/p_0^2(x)]\) versus \(x\). The crosses are the experimental data and the solid line is the best fit of the data. The effective \(B/A\) for this case was calculated to be \(2.3 \times 10^2\). In the calculation, \(\rho_0\) was 1.05 g/cm\(^3\) and \(\sigma_0\), the surface tension of the membrane, was 1.5 \times 10^2 m/s/s. In fig. 5, the water path was not included in \(x\). Since minimal water path is needed for the system, the value calculated is the asymptotic maximum value of the effective \(B/A\) of the system.

**CONCLUSION**

The experimental results show that the dramatic enhancement of the effective nonlinear parameter \(B/A\) is primarily due to the nonlinear resonance oscillation of the trapped bubbles. The measured effective \(B/A\) for water which contains the ensemble of trapped cylindrical bubbles oscillating at resonance is of the magnitude of \(10^3\). The details of discussion could be found in our recent work\(^6\).

REFLECTION AND REFRACTION OF FINITE-AMPLITUDE SOUND WAVE ON A PLANE INTERFACE BETWEEN TWO SOLID MEDIA

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INTRODUCTION
In recent decades, there was a rapid development of nonlinear acoustics in infinite-extension elastic solids. Especially, to 2nd order approximation, the nonlinearity of them can be described by the third-order elasticity (TOE). However, the theory in this area can not provide enough relations to definite all of them, for example, Munson’s third TOE n was not included and other methods have to be invoked [1, 2]. On the other hand, the extension of a solid is limited and many experiments are always done at its surface, thus, the boundary effects on nonlinear waves will be essential. In Ref. [3] reflection for 2nd harmonics in isotropic solids were reported by means of the theory [4] when the incidence wave was P and SV waves. It was shown that some effects are accumulative. Suppose the contribution of the solutions of a homogeneous wave to be negligible, Ref. [5] considered nonlinear reflections of P incidence on a plane half space. It is worth noting that the solutions of the inhomogeneous wave equation may have accumulative (secular) terms, which are proportional to the coordinates and at least a variable occurs in the expressions of boundary conditions of relevant homogeneous wave equation so that the integral constants are hardly to be determined [6, 7]. In order to overcome this difficulty, Ref. [8] gave a group of accumulative solutions to satisfy the boundary conditions. In application of the theories [4, 8], Refs. [9, 10] investigated nonlinear bulk and surface sound waves. In this paper, reflections and refractions at a plane boundary will be investigated when the incidence are P, SV, and SH waves, respectively, and the solutions of the boundary surface and of Q waves are analysed.

WAVE EQUATIONS AND BOUNDARY CONDITIONS
By a selection of the coordinates, the vibrations of P and SV waves can be in (x, z) plane and the one of SH wave is in the direction of its normal. Applying the theory [4], wave equations of the second-order bulk harmonics \( \mathbf{q}_{1}, \mathbf{q}_{2}, \mathbf{q}_{3} \) (SH waves, respectively) can be obtained [9, 10]

\[
\begin{align*}
\frac{\partial \mathbf{q}_{1}}{\partial t} &= \nabla \times \mathbf{F}_{1} - \mathbf{F}_{0}, \\
\frac{\partial \mathbf{q}_{2}}{\partial t} &= \nabla \times \mathbf{F}_{2} - \mathbf{F}_{0}, \\
\frac{\partial \mathbf{q}_{3}}{\partial t} &= \nabla \times \mathbf{F}_{3} - \mathbf{F}_{0}
\end{align*}
\]

where

\[
\begin{align*}
\mathbf{F}_{1} &= \frac{-1}{2\sqrt{\pi}} \int \left[ \ln \left( \frac{1}{(\lambda + 2\mu)} + \frac{1}{(\lambda - 2\mu)} \right) \right] \nabla \psi_{1} \cdot d\mathbf{r}, \\
\mathbf{F}_{2} &= \frac{-1}{2\sqrt{\pi}} \int \left[ \ln \left( \frac{1}{(\lambda + 2\mu)} + \frac{1}{(\lambda - 2\mu)} \right) \right] \nabla \psi_{2} \cdot d\mathbf{r}, \\
\mathbf{F}_{3} &= \frac{-1}{2\sqrt{\pi}} \int \left[ \ln \left( \frac{1}{(\lambda + 2\mu)} + \frac{1}{(\lambda - 2\mu)} \right) \right] \nabla \psi_{3} \cdot d\mathbf{r}
\end{align*}
\]

The definitions of rest signs can be found in [4, 9, 10]. At the boundary surface, z = 0, the stresses and the displacements must be continuous, respectively, i.e.

\[
\begin{align*}
\mathbf{q}_{1} &\big|_{z=0} = \mathbf{q}_{1}^{'}, & \mathbf{q}_{2} &\big|_{z=0} = \mathbf{q}_{2}^{'}, & \mathbf{q}_{3} &\big|_{z=0} = \mathbf{q}_{3}^{'}
\end{align*}
\]

SOLUTIONS FOR P INCIDENT WAVE
In the first maximum (z = 0), the primary incidence of P wave can be denoted as

\[
\mathbf{q}_{1}^{'} = A_{0} \exp(3k_{x}z + ik_{z}x) \mathbf{e}_{x} + B_{0} \exp(-3k_{x}z - ik_{z}x) \mathbf{e}_{x}
\]

and a primary SV reflection is

\[
\mathbf{q}_{2}^{'} = \mathbf{e}_{y} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} - k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{x} + \mathbf{e}_{x} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} + k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{y}
\]

Similarly, in second medium (z > 0) one has

\[
\mathbf{q}_{1}^{'} = \mathbf{e}_{y} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} - k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{x} + \mathbf{e}_{x} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} + k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{y}
\]

where

\[
\begin{align*}
k_{x}^{2} + k_{z}^{2} - k_{s}^{2} = \omega^{2}/c_{s}^{2}, & \quad k_{x}^{2} + k_{z}^{2} = \omega^{2}/c_{s}^{2}, \\
k_{x}^{2} + k_{z}^{2} - k_{s}^{2} = \omega^{2}/c_{s}^{2}, & \quad k_{x}^{2} + k_{z}^{2} = \omega^{2}/c_{s}^{2},
\end{align*}
\]

\(c_{s1} \) and \( c_{s2} (1 = 1, 2) \) are the sound speeds for P and SV waves in two media, respectively. Substituting (7-11) into the eqs. (3, 4) and (10) of Ref. [4], we can obtain \( T_{ij} \) and \( \rho_{p} \), subsequently, \( \mathbf{L} \) and \( \mathbf{E} \) as well as means of eq. (5) [cf. Refs. [9, 10]]. Thus, eqs. (1-2) can be

\[
\begin{align*}
\psi_{1} &= -E_{0}^{i}, & \psi_{2} &= -E_{0}^{i}, \\
\psi_{1} &= -E_{0}^{i}, & \psi_{2} &= -E_{0}^{i},
\end{align*}
\]

Similarly, the solutions in z > 0 can easily be obtained as long as let \( h_{5} = 0, h_{0} = h_{0}, h_{1} = h_{1}, \) and \( h_{4} = h_{4}, h_{5} = h_{5} \) in eqs. (15-16). Now, we discuss the homogeneous solution. In general, it has two arbitrary constants, which will be defined by boundary conditions. Obviously, there are eight constants in two media and only four conditions can be available. Suppose the secondary field is caused by the primary incidence, as did in Ref. [8], we can reduce them to four (for two solids) or three (one of them is a fluid). Thus, they are denoted as

\[
\begin{align*}
\psi_{1h} &= h_{0} E_{0}, & \psi_{2h} &= h_{1} E_{0}, \\
\psi_{1h} &= h_{0} E_{0}, & \psi_{2h} &= h_{1} E_{0},
\end{align*}
\]

which are the special solutions and \( h_{0}, h_{1}, h_{4} \) and \( h_{5} \) can be determined uniquely. For example, if a P wave is vertically incident onto a free surface, the displacements of the second-order waves can be

\[
\begin{align*}
u_{z} &= \frac{1}{2} (z + (z + 2) \cos(\omega \cdot z)), \\
\nu_{z} &= \frac{1}{2} (z + (z + 2) \cos(\omega \cdot z))
\end{align*}
\]

Obviously, at the surface z = 0, the nonaccumulative and the contributions of the homogeneous solutions can not be negligible.

SV INCIDENCE
In this case, the primary wave can be denoted as

\[
\begin{align*}
\mathbf{q}_{1}^{'} &= \mathbf{e}_{y} - \mathbf{e}_{y} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} - k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{x} + \mathbf{e}_{x} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} + k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{y}
\end{align*}
\]

Similarly, we obtain the special solutions i.e.

\[
\begin{align*}
\mathbf{q}_{1h} &= \mathbf{e}_{y} - \mathbf{e}_{y} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} - k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{x} + \mathbf{e}_{x} \left[ (\omega - k_{z}c_{s}^{2}) \mathbf{e}_{z} + k_{z} \mathbf{e}_{x} \right] \mathbf{e}_{y}
\end{align*}
\]

\[
\begin{align*}
\psi_{1h} &= h_{0} E_{0}, & \psi_{2h} &= h_{1} E_{0}, \\
\psi_{1h} &= h_{0} E_{0}, & \psi_{2h} &= h_{1} E_{0},
\end{align*}
\]

\( \mathbf{q}_{1h} \cdot \mathbf{h}_{1} \) can be given while \( A_{0}^{2} \mathbf{B}_{0} = 0, A_{0}^{2} \mathbf{B}_{0} = 0, \) and
\[ k_x = k_y = k_z = k_1 = 1 \] are taken. By a similar procedure, the homogeneous solutions and a complete solution can be given. Fig. 2 shows all of the waves. The results tell us that the third TOE n appears not in the expressions whether the incidence is P or SV waves, thus, a SH incidence has to be invoked.

**SH Incidence**

In case of SH incidence, the primary waves are

\[ \psi^H_{01} = B_1 e^{iB_2} \quad \psi^H_{11} = (B_1 e^{iB_2}) k_y / k_z \]  

and the special solutions in

\[ \tilde{\psi}^{H}_{01} = \left( k_{y} / 4 (k_{x}^{2} + k_{z}^{2}) \right) e^{i(k_{x} x + k_{z} z)} \]

\[ \tilde{\psi}^{H}_{11} = \left( k_{y} / 4 (k_{x}^{2} + k_{z}^{2}) \right) e^{i(k_{x} x + k_{z} z)} \]

(22)

The special solutions in \( \pi \neq 0 \) can be similarly given, thus the homogeneous and complete ones can be obtained. Fig. 3 shows the results. A fact is excitable that the third TOE n is included both in \( Q_{SH} \) and in the homogeneous solutions so that it can be obtained by an experiment.

**Appendix**

\[ \alpha = \alpha_0 \left( \frac{3k^2}{2} + 2k_x^2 + 4k_y^2 + 2k_z^2 \right) \lambda / 4(\lambda^2 + \alpha^2) \]

\[ \alpha_0 = \left( \lambda^2 + \alpha^2 + \alpha^2 \right) / 4(\lambda^2 + \alpha^2) \]

\[ \alpha_1 = \left( \lambda^2 + \alpha^2 + \alpha^2 \right) / 4(\lambda^2 + \alpha^2) \]

\[ \alpha_2 = \left( \lambda^2 + \alpha^2 + \alpha^2 \right) / 4(\lambda^2 + \alpha^2) \]

(23)

\[ \beta_{13} = N_{13} k_x k_y k_z \]

\[ \beta_{14} = N_{14} \]

(24)

\[ \beta_{15} = \left( -N_{15} k_z \right) / \left( k_y k_z \right) \]

\[ \beta_{23} = -4k_x^2 k_y k_z \]

\[ \beta_{24} = -2k_x k_y k_z \]

\[ \beta_{25} = -2k_x k_y k_z \]

\[ \beta_{34} = 2k_x k_y k_z \]

\[ \beta_{35} = 2k_x k_y k_z \]

\[ \beta_{45} = 2k_x k_y k_z \]

(25)

\[ n_{01} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{11} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{21} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{31} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{02} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{12} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{22} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

\[ n_{32} = \left( 1 + 2 \pi / k \right) / 2 \lambda \]

(26)

Let

\[ M = \left( \lambda^2 + \alpha^2 + \alpha^2 \right) / \left( \lambda^2 + \alpha^2 + \alpha^2 \right) \]

\[ N = \left( \lambda^2 + \alpha^2 + \alpha^2 \right) / \left( \lambda^2 + \alpha^2 + \alpha^2 \right) \]

\[ \beta_{13} = N_{13} k_x k_y k_z \]

\[ \beta_{14} = N_{14} \]

(27)

\[ \beta_{15} = \left( -N_{15} k_z \right) / \left( k_y k_z \right) \]

\[ \beta_{23} = -4k_x^2 k_y k_z \]

\[ \beta_{24} = -2k_x k_y k_z \]

\[ \beta_{25} = -2k_x k_y k_z \]

\[ \beta_{34} = 2k_x k_y k_z \]

\[ \beta_{35} = 2k_x k_y k_z \]

\[ \beta_{45} = 2k_x k_y k_z \]

(28)

\[ \tilde{\psi}^{H}_{01} = \left( k_{y} / 4 (k_{x}^{2} + k_{z}^{2}) \right) e^{i(k_{x} x + k_{z} z)} \]

\[ \tilde{\psi}^{H}_{11} = \left( k_{y} / 4 (k_{x}^{2} + k_{z}^{2}) \right) e^{i(k_{x} x + k_{z} z)} \]

(29)

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FOURTH-ORDER ELASTIC CONSTANTS AND THE VELOCITY OF ELASTIC WAVES IN STRESSED SOLIDS

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INTRODUCTION

The higher-order elastic constants are of interest in the study all of the anharmonic phenomena of solids. A number of investigators discussed the pressure dependence of velocity of elastic waves in a stressed solid in terms of second- and third-order elastic constants and measured third- and order elastic constants for various materials.

So far very few quantitative results on the fourth-order elastic constants were published. The theoretical consideration and experimental measurement of fourth-order elastic constants were presented only for the case of hydrostatic compression. However, such a technique can evaluate some combinations of the fourth-order elastic constants which are not sufficient to determine all of independent fourth-order elastic constants for solid material. Hence, an analysis for the case of a small amplitude elastic wave propagating through a solid which is simultaneously subjected to either an uniaxial compression or shear stress was undertaken by using the finite deformation theory.

I. FOURTH-ORDER ELASTIC CONSTANTS

According to Bruggen's formal thermodynamic definition, the fourth-order adiabatic elastic constant is given by:

$$\begin{align*}
C_{ijklmn}^\prime & = \rho_0 s^4 U / \varepsilon_{ij} \varepsilon_{kl} \varepsilon_{mn} a_{iv} \varepsilon_{uv} a_{uv},
\end{align*}$$

where $\rho_0$ is the density in the unstrained state, $U$ the internal energy per unit mass, $\varepsilon_{ij}$ the component of the Lagrangian strain, $a_{iv}$ the entropy, and the derivatives of $U$ evaluated at zero stress. Because of the rotational invariance of $U$, the number of independent components of $C_{ijklmn}^\prime$ decreases to 126 for a triclinic crystal. As the symmetry of the material increases, the number of independent constants is reduced. By using the Voigt notation, for 43m, 432 and m3m point group of the cubic symmetry, the eleven independent fourth-order elastic constants are:

$$C_{1111}, C_{1122}, C_{1212}, C_{1222}, C_{1233}, C_{1313}, C_{1323}, C_{1333}, C_{2233}, C_{3322}, C_{3333}.$$ 

In the case of isotropic solids, seven further relations among these eleven constants were derived by Powell and Krishnamurti. Thus, the number of independent fourth-order elastic constants for an isotropic solid reduces itself to four only. They may be denoted by $C_{1111}, C_{1122}, C_{1212}$ and $C_{2233}.

II. SUMMARY OF FINITE DEFORMATION THEORY

A. Lagrangian-Strain

Let $X_1$ be the Cartesian coordinates of a particle in the unstrained state. The coordinates of the same particle after deformation are denoted by $X_2$. Then, Lagrangian strain tensor is defined:

$$\eta = 1/2 (X' \cdot J - \delta),$$

where $J$ is the Jacobian of the coordinates of a point with the transformation coefficient,

$$J = \delta X_1 / a X_i',$$

where $X'$ is the transpose of $X$, and $\delta$ the Kronecker delta.

B. Stress-Strain Relation

In the finite deformation theory, the relationship between the strain components and stress components is expressed as:

$$T_{ij} = \rho s^4 (\varepsilon_{ij} / \varepsilon_{kl} \varepsilon_{mn} a_{iv} \varepsilon_{uv} a_{uv}),$$

where $T_{ij}$ is the component of the stress tensor, $\rho$ the density in the strained state.

III. SECOND-ORDER EFFECTIVE ELASTIC CONSTANTS OF A STRESSED SOLID

There are different ways to define the second-order effective elastic constants of a stressed solid. As second derivatives of the internal energy $U$ with respect to the components of Lagrangian strain which evaluated at initial statical stress:

$$C_{ijkl}^\prime = \rho s^2 U / \varepsilon_{ij} \varepsilon_{kl} a_{iv} a_{uv}.$$ 

The second-order effective elastic constants for a stressed solid can be expanded in terms of the second-, third-, and fourth-order elastic constants evaluated at zero stress:

$$C_{ijkl}^\prime = (\rho / \rho_0) s^2 (C_{ijkl} + C_{ijkl} a_{iv} a_{uv} + 1/2 C_{ijkl} a_{iv} a_{uv} + ...).$$

2. As the coefficients in a linearized equation of motion which governs the propagation of a small amplitude elastic waves in the stressed solids:

$$S_{ij}^\prime = T_{ij} \delta a + C_{ijkl}^\prime,$$

where $T_{ij}$ is the initial statical stress.

A. Case of Uniaxial Compression

We consider an uniaxial compression deformation given by:

$$X_1 = (1 - \beta) X_1,$$
$$X_2 = (1 - \beta) X_1,$$
$$X_3 = (1 - \beta) X_1,$$

where $\beta$ and $\gamma$ are the deformation coefficients, the homogeneous uniaxial pressure $P$ is applied along the $X_1$ axis. From Eq.(3), Jacobian follows

$$J = \begin{vmatrix}
1 & - \beta & 0 \\
0 & 1 & 0 \\
0 & 0 & 1 - \eta
\end{vmatrix},$$

and hence from Eq.(2)

$$- \beta + \beta^2 / 2 = 0,$$
$$\eta = \begin{vmatrix}
- \alpha + \beta^2 / 2 \\
0 & 0 & - \alpha + \beta^2 / 2
\end{vmatrix}.$$

We calculated Eq.(6) with the aid of Eqs.(8-10). After carrying out the sums involving repeated indices and taking the Voigt notation, the second-order effective elastic constants for a cubic or isotropic material evaluated at the homogeneous uniaxial compression are written as follows:

$$C_{ijkl}^\prime = C_{ij} + a(C_{11} - C_{ij}) - 2/3 a^2 C_{ij} + C_{ij} + C_{ij} + a^2 / 2 C_{ij},$$

$$C_{ijkl} = C_{ij} + a(C_{11} - C_{ij}) - 2/3 a^2 C_{ij} + C_{ij} + C_{ij} + a^2 / 2 C_{ij}.$$
\[ C'_{n} = C_{n} + \alpha \left( C_{n} + C_{m} \right) + \beta \left( 3C_{n} - 2C_{m} + 3C_{ms} + 3C_{ns} \right) + \gamma (\beta C_{n} - C_{ms} + 3C_{ns}) + \delta (\beta C_{n} - C_{ms} + 3C_{ns}), \]

\[ C_{m} = 2C_{s} + 2C_{m} \]

and all others zero. We also obtained,

\[ \gamma = \frac{T_{12}}{(2C_{m})}. \]

It is seen from results listed in Eq.(15) that the application of a statical pure shear stress to a cubic or isotropic material results in changing cubic symmetry or isotropy into monoclinic symmetry.

From above-obtained second–order effective elastic constants \( C_{n} \), we can calculate the combination of \( S_{101} \) which is equal to \( \rho V^2 (V \text{ is the phase velocity}) \) for a given elastic mode. For certain pure mode direction of a stressed isotropic solid, the values of \( \rho V^2 \) are listed in Table I.

**Table I. Values of \( \rho V^2 \) for particular plane waves in an isotropic solid under uniaxial and shear stress.**

<table>
<thead>
<tr>
<th>Type of stress</th>
<th>Wave Vector Direction</th>
<th>Mode Direction</th>
<th>Displacement Direction</th>
<th>( \rho V^2 )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Uniaxial</td>
<td>( \perp ) to stress</td>
<td>L</td>
<td>( \parallel ) to k</td>
<td>( C'_{n} ) (Eq.11)</td>
</tr>
<tr>
<td>pressure</td>
<td>( \perp ) to stress</td>
<td>S</td>
<td>( \parallel ) to stress</td>
<td>( C'_{m} ) (Eq.11)</td>
</tr>
<tr>
<td>pure shear</td>
<td>( \perp ) to stress</td>
<td>L</td>
<td>( \parallel ) to k</td>
<td>( C'_{n} ) (Eq.15)</td>
</tr>
<tr>
<td>stress</td>
<td>( \perp ) to stress</td>
<td>S</td>
<td>( \parallel ) to k</td>
<td>( C'_{m} ) (Eq.15)</td>
</tr>
</tbody>
</table>

Thus, for the cases of uniaxial and shear stressed isotropic solid, we can obtain five combinations of the fourth–order elastic constants which are sufficient to determine four independent fourth–order elastic constants.

**REFERENCES**

PHOTOACOUSTIC PULSE DETECTIONS IN TRANSPARENT AND IN OPAQUE MEDIA

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Photoacoustic pulse detections have been studied and utilized for various materials with a wide range of optical absorbilities, from highly transparent materials to highly opaque materials. This paper reviews some photoacoustic signal generation mechanisms for the various cases, as well as the corresponding applications.

PHOTOACOUSTICS IN A TRANSPARENT MEDIUM

In a transparent medium, the PA source is of the shape of a long cylinder, with a source radius R (i.e., effective radius at the end of the laser pulse of duration t, being determined by the larger of the incident laser-beam radius R and v, which is the product of the acoustic velocity in the medium and the laser pulse duration). This type of PA generation has been treated by numerous authors, both rigorously and semi-rigorously. Examples of some of these studies are given in Refs. 1-9. A simple treatment emphasizing physical insights was recently given in Ref [10], which gives the PA pressure amplitude P at a radial distance r from the axis of the PA cylindrical source as:

\[ P(r) = \frac{\beta E R^2}{2 \pi c_c R^2 + 1/2} \]

where \( \beta \) is the expansion coefficient, \( c_c \) is the specific heat at constant pressure, \( E \) is the laser pulse energy transmitted into the medium, and \( v \) is the acoustic velocity. The above equation indicates the types of applications possible utilizing photoacoustic generation in a transparent medium:

1. Since the PA amplitude is proportional to the absorption coefficient \( \alpha \), absorption spectroscopy can be performed by measuring the PA signal amplitude as a function of excitation wavelength. However, this assumes that the de-excitation probability is constant over the wavelength range.
2. The PA amplitude goes linearly with the excitation laser energy E, hence, better signal-to-noise ratio can generally be obtainable by increasing the laser energy. This does not violate energy conservation as long as the acoustic energy produced is small compared to the absorbed optical energy.
3. The signal is proportional to the "PA efficiency coefficient" \( \beta c_c \). Thus, PA conversion can be optimized by choosing a medium with large PA efficiency coefficient. Furthermore, physical or chemical changes in the medium often cause this coefficient to vary, and hence becomes detectable by PA monitoring.

PHOTOACOUSTICS IN AN OPAQUE MEDIUM

When an excitation laser pulse is incident on an opaque sample, PA generation is possible via the following mechanisms:

1. Thermoelastic generation in the irradiated sample.
2. Thermoelastic generation in a medium adjacent to the sample.
3. Surface desorption or ablation of the sample.
4. Plasma breakdown at the sample surface.

In all these cases, the PA source is a thin layer at the surface, of thickness being larger of the optical penetration depth or the thermal diffusion length. This thin PA source at the surface can produce a large variety of PA waves: longitudinal, shear, or surface acoustic waves, depending on the boundary conditions at the surface as well as on the observation direction; many different kinds of applications are thus possible:

1. PA spectroscopy of opaque materials.
2. Nondestructive materials evaluation and acoustic or thermal-wave imaging.
4. Monitoring of laser desorption or ablative [13, 14].
5. Generation of mechanical motions or movements of microscopic objects [15].

CONCLUSIONS

This paper provides an overview of PA generations in optically transparent samples and in opaque samples. Diverse applications are possible, including excitation spectroscopy, sensing of thermoelastic properties and fluid flows, as well as producing microscopic motions like ejection of submicron particles from a surface.

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References

LOCAL OPTOACOUSTIC PHENOMENA

Consider low frequency acoustic wave produced by earth’s surface transient heating by modulated sunlight intensity. This wave can be measured by "local" receivers, i.e. by infrasound pressure transducers or their networks. Light intensity modulation can be a result of cloud motion or solar eclipse.

During solar eclipse local area of earth’s surface is subjected to a time-dependent heat and cooling action. The adjacent gas layer can form, for example, expansion or contraction thus producing low frequency acoustic pulse. It is possible to describe this pressure perturbation using the steady state piston radiator approach [5].

It should be said that transient heat sources analysis for eclipse conditions requires definite innovations. First of all, one must consider thermal wave radiation into unoccupied gas volume. Secondly, the important condition of homogeneous ground layer depth being much larger than thermal wavelength \( \lambda \), should be met (thermal wave is coming from thin upper part of the ground layer, adjacent to gas boundary). Then we obtain a one-dimensional case surface time-dependent component of the surface temperature spectrum (relative to ambient):

\[
T(\omega) = \frac{B \Omega(\omega)}{4(L^2 + 2B + 2B)} + \frac{3}{4(L^2 + 2B + 2B)} + \frac{3}{4(L^2 - 4 - 4B)}.
\]

where \( \lambda_{\infty} = \frac{\lambda_{0,\infty}}{2\omega} \) - thermal wavelengths in the air and Earth layer, consequently, \( \beta_{\infty} \) - heat diffusivities, \( k_{\infty} \) - heat conductivities, \( I(\omega) \) - sunlight intensity modulation spectrum, \( B = \alpha I \), where \( \alpha \) - light absorption coefficient. By the distance \( l \) from the surface transient component of temperature goes down from its maximum to zero, so the characteristic temperature spectrum of the gas layer may be estimated as \( T_{\text{layer}}(\omega) = 1/T_0 \).

Using the one-dimensional piston radiator approach, one the spectrum of subsurface pressure pulse in the air,

\[
p(\omega) = p_{\infty}.c.\frac{T_{\text{layer}}(\omega)}{T_0}.
\]

where \( p_{\infty}c \) - air density and sound velocity consequently, \( T_0 \) - the initial surface temperature. Examination of (1),(2) shows that at high frequencies, for i.e., \( B \ll 1 \), \( p(\omega) \) goes down proportionally to \( 1/T_0 \). At low frequencies, \( p(\omega) \) - related pulse magnitude doesn’t depend on \( \alpha \) or \( \lambda \) and with convenient transformation we find for peak pressure.
\[ p_m = \frac{1}{\sqrt{2}} \rho_a c_a \sqrt{\rho_a} \cdot \frac{I_0}{c_b [\kappa + 1]} \]

where \( I_0 \) peak intensity sunlight modulated pulse. This expression is valid for pulses of sufficiently large width \( \tau \gg \tau_{\text{exp}} / \beta_a \). For eclipse conditions we have \( \tau = 200 \ldots 500 \text{ sec} \). For typical values of earth and air parameters we obtain \( p_{m[Pa]} = 10^{-6} I_0 \) in W/sq.m. When sunlight intensity goes down from the value of 1400 W/sq.m to zero in several minutes the pulse of rarefraction being 0.01 Pa (near field value) is radiated.

The characteristic acoustic wavelength and the size of shielded area are of the same order, so the radiated pulse will affect nearby destructive interference. Thus, necessary acoustic measurements should be conducted in the cooling region of the moon shadow. We suppose that most adequate measurements could be provided by infrasound pressure transducer immersed in a shallow closed pond thus minimizing low frequency noise caused by air flow around the transducer. According to (6) low frequency noise level in such ponds is extremely small, 30-40 dB re 1µPa/HZ and the lower limit of detected magnitude constitutes 1-10 µPa. In the course of “eclipse listening” device design one should take precautions against heating and cooling of transducer by surrounding water due to sufficient temperature dependence of piezoceramics sensitivity. As a convenient measure we can propose to place transducer in additional sand reservoir and, finally, interference caused by cloud motion sunlight modulation may be minimized by using transducer network.

Acoustic pulse of sufficiently higher frequencies occurs when intense light of bolide trace flush is absorbed by earth’s surface (7). It may explain the fact that an observer of a bolide can hear concomitant sound without any delay. We can use following expression for rough estimation of pulse magnitude in the vicinity of 1-10 kHz frequency band:

\[ p_m = \rho_a c_a \frac{x}{\tau} \sqrt{\rho_a} \tau \Delta T \]

where \( x \) - thermal expansion coefficient, \( \tau = 10^{-5} \text{sec}^{-1} \) - flush length, \( \Delta T \) - surface temperature growth due to absorption of light pulse having peak intensity \( I_0 \), \( \Delta T = \frac{I_0}{c_p \rho_a} \sqrt{\rho_a} \tau \Delta T \geq 3 \times 10^{-5} \text{deg} \). With \( x = 1/300 \text{ deg}^{-1} \) and \( \beta_a = 0.1 \) sq.cm/sec we obtain \( p_m = 3 \times 10^{-5} \text{Pa} \), that exceeds the lower limit of detection of human ear in the kHz frequency band.

The values given in the report may be regarded as approximate. Detailed calculations are in progress now.

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LASER GENERATION OF INTENSE ACOUSTICAL PULSES

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INTRODUCTION

The optical generation of high intensity acoustical pulses by an incident on a surface of condensed media laser beam is considered. The experimental data on a sound generation by a laser pulse in vaporization regime are presented. Theoretical estimation of focusing optically exited pulses are given in comparison with the experimental data on a generation of high intensity sound pulses in a process of optical breakdown.

EXPERIMENTS

We have investigated experimentally the optical generation of sound in water both in a regime of vaporization and in a regime of optical breakdown.

The CO2-laser was used in the experiments on vaporization regime of sound generation. The experimental set-up is presented schematically in the Fig.1.

![Diagram of experiment setup](image)

Fig.1

Laser pulse energy was varied in a range 10-1000 J, the duration of optical pulse is 5 μs. Light beam is focused on a water surface, the diameter of light beam on it is about 0.5 cm.

The acoustical pulse was registered by a PZT-hydrophone with a frequency band about 7 MHz. Its sensitivity is about 0.5 μV/Pa. Signals after amplification was recorded by the oscilloscope. One example of such a oscillogram is presented on Fig.2.

The profile of optical pulse, which generate sound is given on Fig.2a. The oscillogram of sound pulse and its frequency spectrum are given on Fig.2b and Fig.2c respectively. The signal was registered at a distance 100 cm from the source in the far field zone of radiating array produced laser beam action on the free surface of liquid, its peak amplitude is $P_m = 10^4$ Pa, duration on a level -6.0 dB, $\tau = 10 \mu s$, intensity is $I = 0.1$ W/cm². Based on this data theoretical estimation gives at a distance $r = 1$ cm $P_m = 10^5$Pa, $I = 10$ W/cm² representatively, which are in agreement with data (1).

The peak pressure of acoustical signal grows up with the increasing of light intensity in a laser pulse in a proportion $p^\alpha$, where $\alpha = 0.6 - 1$. The optically generated sound pulse can be acoustically focused by a sound lens or in the case of optical generation of sound in a transparent container with a focusing wall. Consider an focusing system with aperture $\alpha$ and focal distance $F$. The coefficient of amplification of such a system is

$$K = \frac{P_{1f}}{P_m}$$

(1)

where $K = \frac{IF}{\alpha^2 \sin \alpha}$.

(2)

where $\alpha$ is angular aperture of a system, $P_m$, $P_{1f}$ - amplitudes of sound on the system respectively. For a reasonable values $\alpha = 10^\circ$, $F = 30$ cm, $\alpha = 10$ cm.

This means, that by acoustical focusing of a laser generated sound pulses one can obtain in the focal region of a system acoustical field of intensity $I = 10^4$ W/cm², $P_m = 10$ Pa.

Let us compare this results with the case of sound generation by focusing beam of light (2,3,4). At a distances $r = 1$ cm from a source in a case of extended laser spark we obtained $P_m = 0.3$ Pa, $I = 10$ W/cm² (4), while for the full developed spark both theoretical and experimental data corresponds to $p = 10$ Pa, $I = 10^4$ W/cm².

CONCLUSION

Experimental investigation of optical generation of sound in vaporization regime and in the regime of optical breakdown and corresponding theoretical estimation of sound field characteristics in cases of optically generation focusing sound beam and generation of sound by focused laser beam demonstrate that in the space scale of 1 cm the average intensity of sound fields in both cases are close together in the order of magnitude.

However, at smaller distances the intensity of a sound field produced at optical breakdown can sufficiently exceeds that of field generated in a vaporization regime.
profile of sound pulse

Fig. 2a

spectrum of sound pulse

Fig. 2b

LASER PULSE

Fig. 2c

REFERENCES

HIGHLY SENSITIVE DETECTORS FOR CAPILLARY ZONE ELECTROPHORESIS USING PHOTOTHERMALLY INDUCED CAPILLARY VIBRATION.

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INTRODUCTION

For its extremely high separation efficiency, the theoretical plate number presently reaching near \(10^6\), capillary zone electrophoresis (CZE) is a powerful separation method especially for biological materials (1). The method has another merit that sampling volume can be reduced to less than nano liter. Nevertheless, conventional absorbance detectors are not sensitive enough for such small sampling volume according to its short optical path. The authors have already proposed another high sensitive detection method named capillary vibration induced by laser (CVL). The details of the signal generation and detection mechanisms were described in the literature (2,3), however, its lower limit of detection was in the order of \(10^{-5}\) in absorbance. This CVL method was applied to a detector for CZE, and ultramicro amounts of nonderivatized amino acids of femtomole level were detected (3).

As in the previous report, the capillary vibration was caused by local tension fluctuation of the capillary irradiated by intensity modulated excitation laser beam. Liquid samples in the capillary tube absorbed this laser beam, and generated heat periodically according to photothermal effect, which made the capillary tension fluctuated locally. This vibration emitted an acoustic wave, and was detected by an optical beam deflection (OBD) method (4).

In this report, the capillary vibration was directly detected by a piezoelectric transducer (PZT). The energy conversion process in this direct detection method does not include an energy transfer from the mechanical vibration energy to acoustic wave energy, this method is expected to be more sensitive than the OBD detection. Furthermore, in the PZT detection, alignment of the probe beam is not required.

EXPERIMENTAL SECTION

One of the holders which supported the capillary tube (50 \(\mu\)m i. d., 150 \(\mu\)m o. d.) was replaced by a PZT disc (20 mm in diameter, 2 mm thick) as shown in Fig. 1. The excitation beam was an argon ion laser beam of 488 nm wavelength and 100 mW output power. The beam intensity was modulated by a light chopper. The beam was focused with a 50 mm focal length lens on the capillary tube, which was tension applied with a hanging weight. The capillary vibrates like a string between two supports, as the result of the light absorption by liquid. The mechanical vibration of the capillary was directly detected by the PZT detector. The liquid samples were aqueous solutions of sunset yellow dye. The sample concentrations were adjusted step wise dilution with distilled water to be in the range of \(1.7 \times 10^{-3}\)cm\(^{-1}\) to \(1.7 \times 10^{-2}\)cm\(^{-1}\) absorption coefficient.

![Fig. 1. Experimental arrangement of the direct detection system of the CVL with a PZT detector.](image)

RESULTS AND DISCUSSION

For sunset yellow dye solution of 5.0 \(X 10^{-3}\)cm\(^{-1}\), which corresponds to absorbance of 2.5 \(X 10^{-5}\) for 50 \(\mu\)m capillary i. d, Stable signal was obtained, and was confirmed to be due to the mechanical energy of vibration of the capillary was directly transferred to the PZT disc.

The dependence of signal magnitudes on dye concentration was measured. For 50 \(\mu\)m i. d. capillary, absorbances of the prepared samples corresponded to \(8.4 \times 10^{-6}\) to \(8.4 \times 10^{-5}\) Abs..
The signal magnitudes showed linearities for the absorbances of the sample. The detection limit ($S/N=2$) of absorbance was calculated to be $8.5 \times 10^{-7}$ for the 50 um capillary. Considering that the detection limit reported previously in the OBD method was $1.5 \times 10^{-5}$ for the same capillary (3), the PZT detection was proved to be at least one order of magnitude more sensitive than the conventional OBD method.

Comparing the energy conversion process of the two detection methods of CVL, energy loss in the PZT detection process is smaller than the other one. Energy conversions from optical to thermal, and to mechanic energy of vibration by photothermal effect are common for both methods. However, the mechanical energy is directly converted to electric one in the PZT detection, while there are two processes until the final conversion at the detector, acoustic emission of vibration and interaction between acoustic wave and the probe beam, in the OBD detection. Therefore, total energy loss in the PZT detecting process was considered to be smaller than the OBD detection, and this resulted in providing superior sensitivity.

As a PZT detector is very sensitive to electric field because it is electrically equivalent as combination of a capacitor and a resistor, base-line stability and noise due to the separation current and potential of electrophoresis is suspected. However, so far as the experimental result of applying this method as a CZE detector, shown in Fig. 2, no serious effect of separation current was observed. In this experiment the sample was riboflavin of 2-3 hundreds fg and separation conditions were same as the previous report. In addition to the superior sensitivity, the direct detection method does not need optical alignment for probing. This may result in favorable reproducibility and simpler structure of the instrumentation.

Fig. 2. Electropherogram of riboflavin obtained with the direct detection method.

LITERATURE CITED

LASER AS AN UNDERWATER SOUND SOURCE AND ITS APPLICATION FOR OCEANOGRAPHY

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I. INTRODUCTION

The interaction of a pulsed light with air to generate acoustic signals in air had been discovered in the last century[1] and this phenomenon has been found to be the principle of a microanalysis method[2]. Since the discovery of laser, the opto-acoustic sources have been attracted great attention and devoted to study the mechanisms of the sound generation.

The perfectest one of mechanisms is the thermo-acoustic mechanism which applies direct heating of the acoustic medium to produce a controlled local thermal expansion which in turn generates sound signals. The thermal energy is delivered by a laser beam to the water without mechanical contact. The thermal acoustic sources in principle can transmit the sound with frequencies from several hertz to several tens megahertz and with the preset waveform and can possess various directionalities. A moving thermal acoustic source may be obtained, causing the position of the thermal sound source to change as a function of time. The insufficiency of the thermal sound source is the conversion efficiency to be too low to be useful at sonar frequencies at the levels of laser energy that are considered to be practicable.

For the thermoacoustic mechanism, the acoustic medium does not undergo phase change, while for the evaporation mechanism and the breakdown mechanism, the acoustic medium undergoes serious phase change. These mechanisms have been found experimentally to be more efficient, but they are not well understood.

The optoacoustic source by the laser-induced evaporation and breakdown can transmit sound signals with sufficient levels, but their frequencies, waveforms and directionalities can not be preset, furthermore their processes of sound generating are unstable. Thus the potential applications of these sources are constrained due to such shortcomings.

In our experiments, a very intensive laser pulse has been used, the phenomenon of laser-induced selffocusing is found and the sound signals due to the selffocusing is received. The characters of the sound signals due to the selffocusing are presented and the potential applications for oceanography are reviewed in this article.

II. EXPERIMENT

1. Equipment

The lasers used are the short pulse laser and the ultrashort pulse YAG laser with the pulse length 40ns and 20ps respectively. The maximum pulse energy is approximately 20mJ and 1mJ respectively. The laser beam was directed vertically down and focused below the water surface by one of the converging lens with different focal length, in order to produce the different spot size at the water surface. The ultrashort pulse laser can pulse a train of pulses with the maximum energy greater than 8mJ. The photodiodes used were the optical detector to monitor the energy of the laser pulse. The acoustical detectors included a NV-10 hydrophone, a shock probe and a composite PZT probe with different frequency response respectively. The block diagram of acoustic data collection and analysis system are shown in Fig. 1.

The acoustic data were recorded on a TES 2232 digital oscilloscope.

2. Measurements

The initial experiments were performed by using a short pulse laser and the acoustic signals were received by the composite PZT probe whose upper frequency limit was lowest. A ultrashort pulse laser was used during the subsequent measurements and a NV-10 hydrophone, whose upper frequency limit was highest, was adopted to monitor the acoustic signals generated in the water. In both cases, there was a luminous segment beneath the water surface elongated to the focus of the converging len. It is interesting to point out that the laser pulse length in the water is approximately 6m for the 20ps laser pulse. It is well known that once the optical breakdown occurred, the induced plasma absorbed the energy of the laser. The sequential breakdowns along the luminous segment infer the fact that the optical breakdown occurs at the tail of the laser pulse. It suggested that the optical breakdown is induced by the selffocusing effect[4]. A horn-shaped laser pulse due to the block diagram of acoustic data collection and analysis system. The optical breakdown in the water at its tail (the squeezed part). Certainly, the selffocusing process may be the appropriate mechanism to interpret the sequential breakdowns.

It is obvious that the optoacoustic sources by sequential breakdowns due to the selffocusing have the directionalities, which can be varied by adjusting the length of the sequential breakdowns.

The measurements show that the acoustical spectra of the laser-induced sound are alike for the laser pulses with different pulse length from several tenths microsecond up to several tens picosecond (Fig. 2). It turned out, the model to illustrate the mechanism of generating sound in the water by optical breakdown is closely dependent on the plasma formation. Once the breakdown is initiated, the characteristic of the breakdown-induced sound is
III. THE POTENTIAL APPLICATION OF LASER-INDUCED SOUND FOR OCEANOGRAPHY

Optoacoustic source appeared to be an original tool for remote acoustic sensing in oceanology. Some useful applications of the acoustic technique such as bathymetry and profilometry etc. had been reported. But the sound generated by thermal mechanism has not the sufficient intensity to match the practical purposes, so the thermoacoustic waves can not be used as an effective information carriers in the hydrosphere. However, the sound based upon the thermoacoustic conversion can be utilized to diagnose the natural ocean waves, since the surface wave motion greatly effects the characteristics of the laser-induced source. Furthermore, it is easily to realize simultaneous many-point diagnosing wave motion[5].

The optoacoustic source due to the laser-induced selffocusing in the water has been found to be more intense for practical uses at same frequencies. Meanwhile this source bears the directionality which is the essentials for practical uses, but it is not understood well. A new hypothesized optoacoustic generation process - selffocusing process - has to be studied more thoroughly and detailedly.

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SOME SPECIAL FEATURES OF NONPROPAGATING SOLITONS AND THEIR TRANSITION TO CHAOS

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I. INTRODUCTION

The renewed interest in Faraday's oscillating water tank experiment of 1831 has led to several recent important works[1,2] concerning either one of the two dominant phenomena, the soliton and chaos, in contemporary nonlinear physics but none of these works involves the transition of one to the other. The present paper describes a further experimental investigation of the hydrodynamic nonpropagating solitons[11] discovered in a small rectangular trough with the aforementioned transition as one of the special features. We found that this experiment provides interesting phenomena not yet fully explored in the preceding paper.

II. EXPERIMENT

The experimental setup is extremely simple. A long rectangular trough of water (with \(a=39\text{cm}, b=2.45\text{cm}, h=2\text{cm}\)) was driven vertically (i.e., along the z-direction) by a vibrator at a frequency about twice that of the lowest normal mode (0,1), governed by the equation

\[
\frac{d^2 \phi}{dt^2} + c^2 \frac{d\phi}{dx} + \gamma \frac{d^2 \phi}{dx^2} = 0
\]

where \(c\) is the phase velocity, \(\gamma\) and \(p\) the integer. By adjusting the amplitude and frequency of the parametric drive in the "stability range", these authors observed some highly localized pulses in the x direction, but a sloshing motion in the y direction. These pulses correctly have been called nonpropagating solitons.

By repeating this experiment with a trough of similar dimension (\(a=38, b=2.7, h=2\text{ cm}\)), the present authors supply the following finding or results: (1) the effect of the trough dimensions, the driving frequency, and the amplitude on the formation of the solitons, (2) the more detailed observation of multisoliton collisions and, in particular, (3) the condition and behavior necessary for the transition of this solitary system to chaos via bifurcations. Some tentative explanations about the physical mechanism as well as some unsolved problems are proposed.

Our experiments show that the range of stable soliton generation depends strongly on the length of the trough. For example, at \(h=38\text{ cm}, b=2.7\text{ cm},\) even if the driving amplitude and frequency (10 Hz) are appropriately set in the soliton range, the usual case is that irregular Stokes waves on the water surface appear after the drive is turned on. As reported[11], this symbolizes the competition of various modes. The stable soliton select (0, 1) or (0, 1) only being generated by essentially discouraging the unwanted longitudinal modes "by rocking the resonator" or "with the help of a hand-held paddle". However, the spontaneous generation of these solitons is made possible at decreasing \(h\) (the shape of a single soliton remains unchanged until \(h=36\text{ cm})),\) comparable to its extension as shown in Fig. 1. The explanation for this is that the number of longitudinal modes excited varies inversely with \(h\). (The gap between the wavevectors of any two neighboring longitudinal modes \(\Delta k = \pi/k\). Those \(p, q\)'s with \(p \leq h\) are all smaller than \(k\), and many easily excited.) From the above experiments we may conclude that: (1) the length of the trough, which was assumed to be immaterial to the soliton formation both in experiment[17] and theory[13], should be reconsidered, and (2) for a rectangular trough of any special geometry, the most stable soliton always comes from the dominant mode as a result of the mode competition. Hence, the stability region defined as "the range of drive amplitude and frequency in which individual solitons are observed without hysteresis" depends strongly on (1) and (2).

It was also reported that these standing solitons could become moving solitons (along the x direction) in various ways. The interaction of two neighboring solitons seems to have more attention because one of the important requirements for solitary waves to be classified as solitons is that they have to maintain their identity after interaction.[11] These solitons behave somewhat differently from KdV solitons. It is observed that when two solitons of opposite polarity (i.e., they slosh in the y direction 180° out of phase) would repel as a certain distance is reached; those of like polarity would attract (as shown by the upper arrows in Fig. 2(a)), "pass through" each other, "exchange places" (lower arrows) and then repeat the process indefinitely. However, our findings show that the collision of a like polarity pair will be inelastic, resulting sometimes in an overlapping stable soliton shaped like Fig. 2(b), which eventually turns out to be a single soliton. An elastic collision as described with the overlapping one like Fig. 2(c) can occur only if the system is fed with an additional excitation energy \(\Delta E\). This bears the semblance of Bohr's liquid drop model of "symmetrical fission",[14] a nuclear model but of classical origin. It may be stated most simply as "a normally stable liquid drop can be made to break up if a mechanical vibration of enough energy is set up in it". The mode after collision of a like polarity pair also varies with \(h\), e.g. for \(h=20\text{ cm}\), it is a clear-cut separation as indicated by the lower arrow of Fig. 2(a), but for \(h=6\text{ cm}\), the overlapping soliton does not die out instantly [Fig. 2(d)]

There were several theoretical papers dealing with the transition from the solitary system to chaos as exemplified by the one dimensional perturbed KdV and the damped driven sine-Gordon[15] solitons. However, we have observed some very interesting phenomena in this matter by carefully investigating the amplitude variations in the Stokes wave mode, and driving at a fixed frequency (~10.06 Hz) and inserting parallel electrodes to condense resistance variations as the liquid mode changes. In the range of the amplitude \(A, 0.65\text{ mm}\) up to \(1.2\text{ mm}\), we obtain a series of wave amplitude \(A\) versus time responses employing FFT spectrograms before and after the soliton formation. For example, Fig. 3(a) shows the \(A\)'s vs. t response for \(A\sim 0.76\text{ mm}\). The corresponding FFT spectrum [Fig. 3(b)] depicts essentially two prominent lines with one near one-half of the drive frequency (5 Hz) and the other with 2 Hz. However, the lines obey (1/2)^n, where n=2 is masked by other fractional subharmonics. This periodic doubling seems to be a general feature in the whole range (0.65 mm to 1.2 mm) before the onset of chaos. What interests us most is that at certain frequencies, even a certain subharmonic from a hand-held paddle causes the soliton immediately to form. A typical FFT spectrum of the soliton so generated in this range is shown in Fig. 4 (A=0.8).
With a further increase of \( A \) up to 1.3 m, no stable solitons can be formed by any means. The \( A' \) vs \( t \) response pattern (Fig. 5(a)) shows no constant period. The corresponding FFT continuous spectrum (Fig. 5(b)) with the complicated wave pattern and temporal fluctuations shows the onset of chaos.

Let us go back to the single solitons for a long enough. For simplicity, if we illustrate, as follows, a normalized semiempirical formula based on the theory of multiple scale expansion valid for the solitons exited at (0,1) mode

\[
\begin{align*}
z = & \text{sech}(0.65x)[b - \cos ky \cos wt + c - \cos 2ky \cos 2wt \\
+ & d - \cos 3ky \cos 3wt + \ldots] \\
\end{align*}
\]

where \( k = \sqrt{1}, \) and \( b > c > d. \) This equation represents the shape of an envelope soliton that is localized in the \( x \) direction and sloshing in the \( y \) direction as already stated. The space and time dependence of \( z \) has been given in some detail.\(^{[2]} \) The inclusion of small amplitude higher harmonics and the omission of negligible subharmonics are in agreement with Fig. 4.

### III. CONCLUSION

From what we have observed so far, the problem is further complicated by the phenomena in connection to collision dynamics and the transition of this solitary system to chaos via bifurcations. As shown by Eq.(2), the trough is a strictly two-dimensional system. The response of a pair of electrodes also varies with its position. The set of data presented in this paper for the above transition is obtained at the trough (node) of the Stokes wave. A comparison with the set obtained at the wave peak (antinode) shows also the predominance of 1/5 subharmonic (Fig. 3(b)), but overall spectral distributions are different in finer details insofar as the bifurcation diagram is concerned. But the threshold as well as the general trend is invariant in this regard. Further experiments with a view to generalizing the picture of the transition, computer simulations, and theoretical calculations are underway. We are just entering into the second phase of the present study, and would also regard this problem as one that may provide more linkage between the classical fluid dynamics or acoustics and modern physics as repeatedly claimed\(^{[2]} \) the further we go.

### References


THE PERIODICAL PASSING THROUGH OF A PAIR OF SOLITARY WAVES

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In case of travelling wave solution of solitons group velocity $v_s$ is constant. Two solitons with different $v_s$ can not meet again after passing—through about each other. However it has been found that two solitary waves with same polarity can pass through and oscillate periodically about each other in rectangular water trough undergoing vertical parametrical excitation. Mao and Wu et al found their group velocities are equal and opposite and vary periodically with time, reach the maximum value when two solitary waves complete overlap and decrease to zero gradually after separate, then come back and move to each other pass—through and separate again. This sequence repeats infinitely. Period of passing—through is not the same in different experimental conditions, but about 10 seconds usually. Until now, there is no reasonable theoretical explanation of these phenomena.

As usual, we take the coordinate system with the origin on the undisturbed free surface of fluid, the z coordinate positive up and the x,y coordinates along the length and the width of the water trough, respectively, and consider the irrotational motion of incompressible fluid in a long trough with width b and depth d at rest in a gravitational field. Then the velocity potential $\Phi(x,y,z,t)$ satisfies following equations

\[ \begin{align*}
\Phi_{xx} + \Phi_{yy} + \Phi_{zz} &= 0, \\
\Phi_{z} - \eta &\Phi_{z} = r - \eta, \\
\Phi_{x} + \frac{1}{2}(\Phi_{x}^2 + \Phi_{y}^2) + \Phi_{z} + \eta + \eta_{t} + 2\beta \Phi &= 0, \\
\Phi_{x_{t}} &= 0, \\
\Phi_{z_{t}} &= 0.
\end{align*} \]

(1)

Where $\eta(x,y,t)$ is the disturbance of the free surface relative to the rest of fluid, $L = 4\omega_{s}^2 A_{s} \cos \omega_{t} t$ is the acceleration of vertical excitation exerted on water trough. The suffixes of $\Phi$ and $\eta$ denote the partial differentiation with respect to corresponding quantities. We suggest that $\Phi, \eta, \Phi_{x}, \Phi_{y}, \eta_{x}, \eta_{y}, f, \omega_{s} A_{s}, \Delta \omega = \omega - \omega_{s}$ are first order of $\varepsilon$ and $\Phi, \eta, \eta_{x}, \eta_{y}$ are second order of $\varepsilon$. Where $g$ is the gravitational acceleration, $\omega_{s}$ half the driving frequency, $A_{s}$ the driving amplitude, $\omega$ the frequency of (0,1) mode of the water trough and $\beta$ the damping factor. We expand $\Phi$ and $\eta$ in series of form

\[ \begin{align*}
\Phi = \Phi^{(1)} + \varepsilon \Phi^{(2)} + \varepsilon^{2} \Phi^{(3)}, \\
\eta = \eta^{(1)} + \varepsilon \eta^{(2)}.
\end{align*} \]

(2)

In the approximation of second degree, let the secular term equal zero

\[ \begin{align*}
\phi^{(1)}_{xx} + f \omega \phi^{(1)}_{z} e^{-\Delta \omega t} = 0, \\
\phi^{(1)}_{x} + \Delta \omega \phi^{(1)}_{x} + f \omega \phi^{(1)}_{y} = 0.
\end{align*} \]

(3)

Introducing $\phi^{(1)} = \phi_{x} e^{-\Delta \omega t}, f = f_{x} + f_{y} + if_{z}$ and putting $if_{x} - if_{y} \phi^{(1)}_{x} = i \beta \phi^{(1)}_{x}$ one has

\[ \begin{align*}
\phi_{x_{t}} + \Delta \omega \phi_{x} + f \omega \phi^{(1)}_{y} &= 0, \\
\phi_{y_{t}} &= \psi^{(1)}_{t} e^{-\Delta \omega t} - \Delta \omega \phi_{y} - \frac{\Delta \omega}{f_{x}} \phi^{(1)}_{x} e^{-\Delta \omega t} - \frac{1}{2} \phi^{(1)}_{y} e^{-\Delta \omega t}.
\end{align*} \]

(4)

The periodical solution of eq (2) is

\[ \begin{align*}
\phi_{x} = \psi^{(1)} e^{-\Delta \omega t} - \frac{\Delta \omega}{f_{x}} \phi^{(1)}_{x} e^{-\Delta \omega t} - \frac{1}{2} \phi^{(1)}_{y} e^{-\Delta \omega t}.
\end{align*} \]

(5)

Formula (3) is in agreement with the experimental value of the period of passing through as shown in Fig 1.

We suggest $\psi$ is a function of $\eta$ and $\theta$, where $\theta(t + T_{1}) = \theta(t)$. It should be determined experimentally and in general $\theta = \sum \theta_{k} \sin \omega_{k} t$. For simplicity we put $\theta = \frac{\theta_{0}}{\Omega} \sin \Omega t, \Omega = \frac{2 \pi}{T_{1}}$. In the third order of approximation the nonlinear schrodinger equation of $\psi(x,\theta)$ is obtained

\[ \begin{align*}
\begin{align*} 
\psi_{x} &= \frac{\theta_{0}}{4} \psi + \delta_{x} \psi_{xx} + \gamma_{0} |\psi|^{2} \psi.
\end{align*}
\end{align*} \]

(6)

By use of the solution of two solitons of eq (4)

\[ \begin{align*}
\eta_{1}^{(1)}(1 + \frac{\Delta \omega}{f_{x}}) \cos \omega_{s} t \{ \cos \alpha_{+} \sech \beta_{+} + \cos \alpha_{-} \sech \beta_{-} \} - D e^{-\beta_{-}} \cos \alpha_{-} + e^{-\beta_{+}} \cos \alpha_{+} + a_{+} \} \Delta_{1} \sech \beta_{+} \sech \beta_{-} \}
\end{align*} \]

(7)

\[ \begin{align*}
\eta_{2}^{(1)}(1 + \frac{\Delta \omega}{f_{x}}) \sin \omega_{s} t + \{ \sin \alpha_{-} \sech \beta_{-} + \sin \beta_{-} \sech \beta_{+} \} - D e^{-\beta_{-}} \sin \alpha_{-} + e^{-\beta_{+}} \sin \alpha_{+} + a_{+} \} (\Delta_{1} \sech \beta_{+} \sech \beta_{-} \}
\end{align*} \]

(8)

\[ \begin{align*}
\Delta_{1} = 1 - \sech \beta_{-} \sech \beta_{+} \{ \frac{1}{4} \frac{\alpha_{+}^{2} + \alpha_{-}^{2}}{2 \alpha_{+} + 2 \alpha_{-}} e^{-2 \alpha_{+} \alpha_{-}} \}
\end{align*} \]
+ \frac{\xi^2}{2(\zeta^2 + \xi^2)} \left[ \left( \zeta^2 - \xi^2 \right) \cos \frac{4\pi}{\delta e} x - 2\xi \sin \frac{4\pi}{\delta e} x \right] \right] 
- 2\xi \left( \sin \frac{4\pi}{\delta e} x + \frac{\xi}{\delta e} \right) \right) 
\sigma = \Omega + 4(\zeta^2 - \xi^2) \theta + \frac{2\xi}{\delta e} (x \pm x_0) - \sigma + \frac{\xi^2}{4} \theta 
\beta = \frac{2\xi}{\delta e} (x \pm x_0) - 8\xi \xi \theta 

By the expressions of \( \beta \), the group velocities of the peaks of two solitary waves can be determined

\[ V_{xz} = \pm 4\xi \sqrt{\frac{\delta e}{\delta_0}} \theta \]  
where \( \theta = \frac{\delta e}{\delta_0} \cos \Omega t \). From eq. (6) the group velocities of two solitary waves are equal and opposite and vary periodically with time. Their period is \( T = \frac{2\pi}{\Omega} \). The positions of two peaks are \( x \pm x_0 = \pm 4\xi \sqrt{\frac{\delta e}{\delta_0}} \theta \) and the maximum distance of separation is \( 2D + 2x_0 = 8\xi \sqrt{\frac{\delta e}{\delta_0}} \theta / \Omega \).

Besides the above theory the theory of bound state is also given. We considered the free irrotational motion of an incompressible inviscid fluid in a long trough with length L and write \( \eta_1 \) in the form

\[ \eta_1 = \frac{\omega}{g} (A(Z + Y) + c.c.) \]
\[ Z = e^{i\omega t - \omega_0 t}, \quad Y = e^{i\omega t - \omega_0 t}, \quad k = \frac{\pi}{b} \]

Assuming \( \frac{\omega^2 - \omega_0^2}{\omega_0^2} \sim O(\epsilon^3) \) where \( \omega_0^2 = gkT \)

\( (k\delta)(1 - \frac{\lambda \epsilon^2}{g}) \), and \( \lambda \) is the surface tension of the fluid, we obtain, to the order \( \epsilon^3 \)

\[ 2i\omega_3 A_{1,3} + C^2 A_{2,4} + N|A|^2 A - k^2 (1 - T)BA \]
\[ + (\omega_0^2 - \omega_2^2)A = 0 \]  

For the pulse-type initial condition

\[ A(x, t = 0) = a_0 \text{sech}(\sqrt{1/3}N / (2C^2))^{1/2} a_0 x \]

where \( a_0 = g a_0 / (4\omega_0) \)

and \( a_0 \) is the maximum initial amplitude, the eq. (8) yields a solution of bound state of solitary-wave pulses

\[ A(x, t) = 6a_0 (A_1(x, t) / A_2(x, t))e^{-\theta t} \]

with the recurrent period

\[ T = \frac{144\omega^3}{g^2 a_0^2 N} \]

Some periodic process appears in numerical calculation but one cannot get the group velocity of two solitary waves. The sophisticated study is in progress.

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EFFECTS OF SURFACE-TENSION ON WU'S SOLITON AND KINK SOLITON

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The discovery of the so-called nonpropagating hydrodynamic soliton by Wu et al. was reported in 1984.1 In Wu’s experiment, the water was poured into a Plexiglas channel 38 cm long and b = 2.54 cm wide to a depth of d = 2.00 cm, which was placed on a loudspeaker whose cone is driven at a frequency 2ν = 2ωb/2π ≈ 100 Hz in the vertical direction. Above a certain frequency threshold, there may be a solitary surface wave excited at half the drive frequency, which is localized longitudinally and “shoeshoe” transversely. Before long, the nonlinear equation and its nonpropagating soliton solution were derived by Larrasa and Putterman5 and Miles6 in different methods. In a Cartesian coordinate system with its x, y, and z axes parallel to the longitudinal, transverse, and vertical directions of channel, respectively, and the still free water surface as z = 0, the leading term of the solitary surface wave is:7,8

\[ \xi(x, y, t) = u(x, t) \cos(\pi y/b) e^{i\omega t} + c.c. \]

where c.c. denotes the complex conjugate of the previous term. When the viscosity, the surface-tension and the drive are neglected as in Ref. 2,

\[ u = \frac{2}{B} (\omega_1^2 - \omega^2)^{1/2} \tanh[(\omega_1^2 - \omega^2)/2] \]

where, ω1 (natural frequency), c and B > 0 are constants which depend upon the width and depth of water. In Ref. 3, the expression of the term of surface-tension (the right side of Eq. (8)) is too simple so its high order effects on the soliton were omitted.

Six years later, Denardo et al.4 reported their observation of a kink soliton in the similar experiment as in Ref. 1 except the width b and depth d of water being 5.71 cm and 1.00 cm, respectively, and here B < 0. If the viscosity, surface-tension and drive were not taken into account as in Ref. 4, then

\[ u = \frac{2}{B} (\omega_1^2 - \omega^2)^{1/2} \tanh[(\omega_1^2 - \omega^2)/2] \]

In this paper, we’ll derive the effects of surface-tension on both the nonpropagation and kink soliton. If the motion of water is supposed to be rotationless as in Ref. 2-4 and the viscosity is neglected for the moment, the velocity potential φ(x, y, z, t) and the surface displacement of water ξ(x, y, t) will satisfy the following hydrodynamical equations and boundary conditions

\[ \phi_x + \phi_y + \phi_z = 0 \quad -d \leq z \leq \xi(x, y, t) \]  
\[ \phi_y = 0 \quad y = 0, b \]  
\[ \phi_z = 0 \quad z = -d \]  
\[ \phi_x = \xi + \phi_x, \phi_y + \phi_z = \xi(x, y, t) \]  
\[ \phi_x = (\omega^2 + \omega_1^2) \xi + \frac{1}{2} (\phi_x^2 + \phi_y^2 + \phi_z^2) \]  
\[ \alpha \xi_{xx} + (1 + \xi^2) \xi_{yy} + (1 + \xi^2) \xi_{zz} - 2 \xi_{xy} \xi_{xz} = \xi(x, y, t) \]  

where g is the gravitation acceleration, α is the ratio of the surface-tension coefficient of water to its density, Z0 = a cos(2ωt) denotes the drive, Z0z denotes its second order derivative with respect to time t, the subscripts express the partial derivative with respect to relevant variables. According to the standard multiple scales method, we introduce the variable of multiple scales \( \chi_j = e^{i\omega t}, \) j = 0, 1, 2, ... and the asymptotic expansions \( \phi = \sum_{j=1}^{\infty} \chi_j \phi_j, \) \( \xi = \sum_{j=1}^{\infty} \chi_j \xi_j. \) Substituting them into Eqs. (4)-(8), and comparing the coefficients of each power of e, we then obtain an infinite system of linear equations which can be solved step by step according to the method of Ref. 5. The main results up to third order are listed as follows

\[ \xi = (u \cos(\pi y/b)) e^{i\omega t} + c.c. \]

\[ + \frac{k(1 + \alpha)^3}{2T} \frac{(3 - T^2)}{(3 - \sigma)(3 - T^2)} u \cos(2ky) e^{i\omega t} + c.c. \]

\[ + \frac{k(1 + \alpha)^3}{2T(1 + 4\sigma)} |u|^2 \cos(2ky) \]

\[ - \frac{k(1 + \alpha)^3}{2T} |u|^3 + O(e^4) \]

(9)

here the local amplitude u satisfies the following nonlinear equation

\[ 2i \omega_1 u + (\omega_1^2 - \omega^2 + 2i \omega_1 \sigma) u - \epsilon^2 u_{\omega_1} - B |u|^2 u - ru^* = 0 \]

(10)

where \( u^* \) denotes the complex conjugate of u, \( s \) is a parameter describing the size of the viscosity,9 and

\[ \omega_1 = [gkT(1 + \sigma)]^{-1/3} \quad (natural \ frequency) \]

\[ k = \pi / b, \quad T = \tanh(kd), \quad \sigma = \alpha k^2 / g \]

\[ \epsilon^2 \xi = \frac{2k^2}{\omega_1^3} [1 + \sigma] [T + kT(1 - T^2)] = \frac{\omega \partial \omega}{k \partial k} \]

\[ r = \frac{2kT}{\omega_1^2} \]

\[ B = \frac{k^2g}{8T} (1 + \sigma) \]

\[ x^2 [(T^2 - 9)(1 - T^2) + \sigma(T^2 - 11)(3 - T^2) + \sigma \omega_1^2] \]

\[ + \frac{2(1 + \sigma)^3}{1 + 4\sigma} (1 + T^2)^2 + 4(1 + \sigma)T^2(T^2 - 1) \]

(15)

In the derivations of (9)-(15), it has been supposed that \( s = O(\epsilon^3) \) and \( \omega_1 - \omega = O(\epsilon^4). \)

The type of the soliton solution of Eq. (10) is determined by the sign of constant B which depends upon the width and depth of water. In the case of nonpropagating soliton, \( B = 2.54 \text{cm}, d = 2 \text{cm}, B > 0 \), it can be checked directly that Eq. (10) has the Wu’s soliton

\[ u = Ae^{-i\pi b \xi}(x/L) \]

(16)

where

\[ A = \{ \frac{2}{B} [g \omega_2 - \omega^2 + (\omega^2 - 4\omega_1^2 \sigma^2)]^{1/3} \}

\[ L = \cos(\varphi - \omega_1^2 \sigma^2)^{1/3} \]

\[ \cos \theta = \frac{1}{\sqrt{2}} \left[ 1 - \frac{(4\omega_1^2 \sigma^2)^{1/3}}{\omega^2} \right] \]

if \( \omega^2 - 4\omega_1^2 \sigma^2 \geq 0 \) and \( \omega_1^2 - \omega^2 + (\omega_1^2 - 4\omega_1^2 \sigma^2)^{1/3} > 0 \). While in the case of kink soliton, \( b = 5.71 \text{cm}, d = 1.00 \text{cm}, B < 0, \)

...
\[ u = A e^{-\frac{t}{L}} \tanh (\frac{x}{L}) \]  

(20)

where

\[ A = \left( \frac{1}{|B|} \right) \left[ (\omega_0^2 - \omega_1^2) + (r^2 - 4\omega_1^2s^2)^{1/3} \right]^{1/3} \]  

(21)

\[ L = c \left( \frac{1}{2} \left[ (\omega_0^2 - \omega_1^2) + (r^2 - 4\omega_1^2s^2)^{1/3} \right] ^{1/3} \right) \]  

(22)

\[ \cos \theta = \frac{1}{\sqrt{2}} \left[ 1 + (1 - 4\omega_1^2s^2)^{1/3} \right] \]  

(23)

on conditions that \( r^2 - 4\omega_1^2s^2 > 0 \) and \( \omega_0^2 - \omega_1^2 + (r^2 - 4\omega_1^2s^2)^{1/3} > 0 \). One can easily see from (16)-(23) that when \( r = s = \sigma = 0, B = k^2g(6T^4 - 5T^2 + 16 - 9/T^2)/9T, \) and (16) and (20) are consistent with (2) and (3) respectively.

To estimate the effects of surface-tension on the solitons, we take notice of \( \omega_1, r, c, \) and \( B \) being functions of \( \sigma \). For water, \( \alpha = 73\text{ cm}^2\text{sec}^{-1}, \gamma = \omega_1^2/\rho = 0.114 \) (for the case of Wu's soliton) and \( \sigma = 0.0225 \) (for the case of kink soliton), which are much smaller than unit, so the above quantities can be expanded into Taylor series of \( \sigma \) up to first order for a rough estimation: \( \omega_1 \approx \omega_0(1 + \sigma/2), r \approx r_0(1 - \sigma), c \approx c_0(1 + \sigma/2), B \approx B_0 + \sigma B_1, \) here the subscript zero denotes relevant quantities for \( \sigma = 0 \), while \( B_1 = B_0^2/2 \left( 4T^2 - 21T^2 + 7 - 97T^{-2} - 27T^{-4} \right) \) is the value of derivative of \( B \) with respect to \( \sigma \) at \( \sigma = 0 \). Substituting them into Eqs. (17), (18) and (20), (21), and neglecting the viscosity, one can easily obtain

\[ \frac{A^{(N)}}{A_0^{(N)}} = 1 + \frac{1}{2} \left[ -\frac{B_1}{B_0} + \omega_0^2 - r_0 \left( A_0^{(N)} \right)^{-1} \right] \]  

(24)

\[ \frac{L^{(N)}}{L_0^{(N)}} = 1 + \frac{1}{2} \left[ -\frac{B_1}{B_0} + \omega_0^2 - r_0 \left( L_0^{(N)} \right)^{-1} \right] \]  

(25)

and

\[ \frac{A^{(K)}}{A_0^{(K)}} = 1 + \frac{1}{2} \left[ \frac{B_1}{B_0} + \omega_0^2 + r_0 \left( A_0^{(K)} \right)^{-1} \right] \]  

(26)

\[ \frac{L^{(K)}}{L_0^{(K)}} = 1 + \frac{1}{2} \left[ \frac{B_1}{B_0} + \omega_0^2 + r_0 \left( L_0^{(K)} \right)^{-1} \right] \]  

(27)

respectively. In the case of Wu's soliton, \( \sigma = 0.114, \omega_1^2 = 1196 \text{ sec}^{-2}, \omega_0^2 = 363\text{cm}^2\text{sec}^{-2}, B_0 = 1774\text{cm}^{-2}\text{sec}^{-1}, B_1 = -11155 \text{ cm}^{-2}\text{sec}^{-2}, A_0^{(N)} \approx 2\text{cm}, L_0^{(N)} \approx 1\text{cm}, \) thus from Eqs. (24) and (25), we obtain \( A^{(N)}/A_0^{(N)} \approx 1.38, \) \( L^{(N)}/L_0^{(N)} \approx 0.88, \) i.e., the height of Wu's soliton is increased while the breadth is decreased by the surface-tension. On the contrary, surface-tension makes the height of kink soliton decrease and its breadth increase, which can be seen obviously from Eqs. (26) and (27).

The modification of surface-tension on Wu's soliton is of practical significance in studying the so-called equal-wave-response-curves. The experimental curves obtained by Wu et al. are drawn in Fig. 1, in which the full-line curves are equal-wave-response-curves and the number associated with each curve is the peak height in centimeters of the solitons above the equilibrium level of the water. The theoretical curves drawn in Fig. 2 are obtained from Eq. (17) in which the boundary dashed lines are given by \( \omega_1 - \omega = O(c^2), a = O(c^2), \) and \( r = 2x_1c, \) and the number beside the lower end of each curve is the peak height of the Wu's soliton without considering the surface-tension while the number beside the upper end is the peak height with the surface-tension being taken into account. Obviously, the modification of surface-tension makes the theoretical curves more consistent with the experiment.

Fig. 1 Experimental equal-wave-response-curves (after Wu et al.)

Fig. 2 Theoretical equal-wave-response-curves.

REFERENCES

DISCOVERY OF NONPROPAGATING SOLITONS - A TRANSITION FROM AN EXTENDED STATE TO A LOCALIZED STATE

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INTRODUCTION

In 1984, we made the first observation of a nonpropagating hydrodynamic soliton. The soliton was observed in a rectangular water trough that was continuously and parametrically excited by vertical oscillation. It is a self-trapped, highly localized, and stationary transverse surface-water-wave. Unlike other traveling solitons at the water surface (KG or envelope type), this local-mode excitation is stationary in the length (the longest dimension of the trough) direction. In the length, it has the form of the hyperbolic secant function, a characteristic soliton solution of the nonlinear Schrödinger equation, the disturbance is also characterized by a sloshing motion back and forth across the width of the trough at half the frequency of the drive.

WHY DOES A SOLITON FORM?

The formation of the nonpropagating, hydrodynamic soliton is based on the idea of the cutoff frequency. The cutoff frequency is the frequency below which waves of a particular mode cease to propagate down a waveguide and decay exponentially in space. This concept is a result of linear wave theory that applies to small amplitude cases.

A long narrow water trough can be considered as a waveguide for surface waves in water. The linear cutoff frequencies, \( f_{cm} \), corresponding to (m, n) mode are given by

\[
f_{cm} = \left( \frac{c}{2L} \right) \sqrt{(m^2 + n^2)},
\]

where \( m \) and \( n \) are integers, \( L \) is the length, \( W \) is the width of the trough, and \( c \) is the phase velocity of the surface-water-wave. Here \( c \) changes with wavelength, \( k \), and is given as

\[
c = \frac{g(k + Th/2)}{\tanh(Th/2)},
\]

where \( g \) is gravitational acceleration, \( T \) is the surface tension of a fluid, \( h \) is the fluid depth, and \( T \) is the fluid density. For deep water cases (\( kT > 1 \), \( \tanh(Th/2) = 1 \). Therefore, \( f_{cm} \) can be written as

\[
f_{cm} = \left( \frac{1}{2\pi} \right) \sqrt{k_{cm}^2 + (nW)^2},
\]

where \( k_{cm} \) is given by

\[
k_{cm} = \pi \sqrt{\frac{mL^2}{2}} + (nW)^{1/2}.
\]

For the nonlinear case, corresponding to finite amplitudes, the cutoff frequencies are amplitude dependent. According to experimental measurements, 1 the higher the amplitude, the lower the cutoff frequency. Thus, if a localized disturbance with the characteristic soliton profile of a hyperbolic secant function occurs in the waveguide, the amplitude at the center of the soliton is higher than that at the edges, so the cutoff frequency at the center, which we call the nonlinear cutoff frequency, is lower than the linear cutoff frequency. Now, if the direct drive frequency (refer to the next section) is below the linear cutoff frequency and above the nonlinear cutoff frequency, the disturbance will propagate away from the center. When it gets to both edges, it can not get out but bounces back toward the center, exhibiting a phenomenon known as self-focusing. The theories were worked out by Larraza and Puttermann 4 and Miles 6 independently.

EXPERIMENT

There are two different drives that could generate the solitons. Both drives should be capable of producing waveguide displacements of several millimeters.

The first is the direct drive which is characterized by the fact that the frequency of the drive is the same as that of the wave generated. The waveguide is driven horizontally along its width direction. The main components needed to study a soliton using the direct drive are a Plexiglas waveguide with inside dimensions of 15 cm long \( \times \) 2.5 cm wide \( \times \) 6 cm high, a 15-cm-diam stereo speaker connected to a low-frequency oscillator through an amplifier, two 15-cm metal rods, a thin aluminum tubing. The detail of this drive is referred to Fig. 1.

The second type of drive is the parametric drive in which the waveguide is vertically driven by a stereo speaker. In this case, the frequency of the driving oscillator is twice the frequency of the waves generated. The set-up of this drive is illustrated in Fig. 2.

In any case, the depth of water should be about 2 cm to ensure that the deep water condition is satisfied. Several drops of surfactant (Kodak Photo-Flo) is added to reduce the pinning effect at the walls of the container.

Using the second type of drive as an example, the trough was driven up and down at a frequency little greater than twice the linear cutoff frequency of a particular mode determined by Eq. (3), a standing wave was formed. Then, we gradually decreased the frequency, the amplitude of the standing wave increased. When the frequency was slightly below twice that of the linear cutoff frequency, the above-mentioned self-focusing phenomenon appeared. As the frequency was reduced a little bit further, a single disturbance with the hyperbolic secant profile in the length direction formed. Once the soliton formed, it appeared to be very stable. The profile of the soliton was measured using a video camera.

We were able to generate solitons of (0, 1) and (0, 2) modes; they are illustrated in Fig. 3 and Fig. 4, respectively. The specific frequencies used for the solitons are indicated in the captions. The position of the disturbance formed usually was at the center of the trough with a total length of about 5 cm; the remaining part of the water surface in the 15 cm long trough was tranquil, although the whole trough was driven uniformly.

Another important aspect of these solitons is that they can pass through another similar soliton and remain unchanged. At higher amplitudes, two solitons can be made to form side by side. These solitons may move toward and pass through each other. The detailed description can be found in Ref. 1 and 5.

Fig. 1 A schematic of the direct drive apparatus. The long length of the waveguide is perpendicular to the plane of the schematic.

Fig. 2 A schematic of the parametric drive apparatus. The short length of the waveguide is perpendicular to the plane of the schematic.

Fig. 3 A three-dimensional plot of a $(0,1)$ mode soliton that formed when the waveguide was parametrically driven at 10.2 Hz. The plot is based on curve fit to measurements taken from photographs. The $x$ axis is along the length of the waveguide; the $y$ axis is along the width of the waveguide, and the $z$ axis represents the amplitude of the wave. The soliton was at the center of the trough, and extended about 5 cm. (a) A plot of $z = 	ext{sech}(x/1.12) \cdot (2.8 \exp (-1.1y) - 0.70)$, a mathematic expression for the $(0,1)$ mode soliton; (b) A plot of the same soliton at a half-period later.

Fig. 4 A three-dimensional plot of a $(0,2)$ mode soliton that formed when the waveguide was parametrically driven at 19.1 Hz. This plot is based on curve fit to measurements taken from photographs. The $x$ axis is along the length of the waveguide, the $y$ axis is along the width of the waveguide, and the $z$ axis represents the amplitude of the wave. (a) A plot of $z = \text{sech}(x/(0.53 \cos(\pi x/1.27)) + 0.14 \cos(2\pi y/1.27) + 0.007)$, a mathematic expression of the $(0,2)$ mode soliton; (b) A plot of $z = \text{sech}(x/(0.43 \cos(\pi x/1.27) - 0.004 \cos(2\pi y/1.27))$, the same soliton a half-period later.
Dynamical Behaviour of a Nonpropagating Soliton under a Periodically Modulated Oscillation

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We consider a fluid in a small long rectangular channel of length \( l \), breadth \( b \) and quiescent depth \( d \), which is subjected vertically to a slowly modulated simple harmonic oscillation of frequency \( 2\omega_0 \), such as
\[
Z_a = a(t) \cos 2\omega_0 t, \quad a(t) = a_0 + a_1 \sin \Omega t, \quad \tau = \varepsilon^2 \omega t
\]
where \( \Omega \) is dimensionless modulation frequency and \( \varepsilon \) is a small parameter. When half of the basic driving frequency, \( 2\omega_0 \) is near the nature frequency of the lowest cross-mode, a cross-wave at frequency \( \omega \) with slow longitudinally and time varying amplitude is induced. We can extend the theory of Miles(1984) and Lararra & Puttermann(1984) to describe this problem as following.

\[
i u_x + \beta u + u_{xx} + 2|u|^2 u + i\mu_1 u + \gamma(\tau) u' = 0
\]
where
\[
X = \varepsilon x/\sqrt{K}, \quad u = \frac{\sqrt{N}}{N^2} A, \quad (3)
\]
\[
\alpha = \bar{\delta}/\varepsilon^2, \quad \beta = \frac{\alpha_1 - \alpha_0}{2\varepsilon^2 \omega \omega_0}, \quad \gamma(\tau) = \frac{a(t) \omega^2}{\varepsilon^3}, \quad (4)
\]
\[
\omega_1 = \sqrt{\bar{g} \varepsilon k T}, \quad k = \pi b, \quad T = \tanh(k d), \quad (5)
\]
\[
K = \frac{1}{4k^2 T} (T + k d \text{sech} k d), \quad (6)
\]
\[
N = \frac{k^2}{64 T} \left( 6T^2 - 5T^4 + 16T^2 - 9 \right), \quad (7)
\]

the asterisk * denotes complex conjugate, \( A \) is the amplitude of the dominant lowest cross-mode, \( \omega_1 \) is the natural frequency of the mode, \( \bar{\delta} \) is the damping ratio which is same as \( \delta \) in Miles(1984) and \( \bar{g} \) is the gravitational acceleration. If the \( \alpha \) and \( \gamma \) is small, (2) is expected to have a dynamical nonpropagating soliton solution. By using the inverse scattering perturbation technique (Lamb 1980), we can obtain a single nonpropagating soliton solution
\[
u = 2\rho e^{-4\delta} \text{sech} Z, \quad Z = 2\rho (X - \xi_0) \quad (8)
\]
\[
\rho_1 = -2 (\alpha + \gamma \sin 2\delta) \rho, \quad (9a)
\]
\[
\delta_1 = -\beta - \gamma \cos 2\delta - 4\rho^2. \quad (9b)
\]

Numerical solutions of (9a,b) show that the amplitude of the soliton can vary periodically at 1,2,3,4,6,8,12 at and so on times of modulation period and chaos as well (to see Fig.1(a-g) and Fig.2).

The anticipated experiment was also carried out afterwards, in which the phenomena of multiple-periodic and chaotic amplitude-modulated motion of the nonpropagating soliton were observed clearly. We use the frequency specters to distinguish the dynamical characters of the motion of the soliton to see Fig.3 and Fig.4). The phenomena were also recorded by a video-camera. The experiment shows qualitative agreement with the theoretical prediction.

REFERENCES


Fig. 1. Phase portraits \((\rho, \dot{\rho})\) and Poincaré maps of some typical attractors for \(\alpha = 0.136, \beta = -0.8582\) and \(\gamma_i = 1\).
(a) periodic limit cycle at 1 time of \(\Omega, \Omega = 3.34\) and \(\gamma_i = 0.41\).
(b) periodic limit cycle at 2 times of \(\Omega, \Omega = 3.3\) and \(\gamma_i = 0.48\).
(c) periodic limit cycle at 3 times of \(\Omega, \Omega = 3.34\) and \(\gamma_i = 0.42\).
(d) periodic limit cycle at 4 times of \(\Omega, \Omega = 3.3\) and \(\gamma_i = 0.49\).
(e) periodic limit cycle at 6 times of \(\Omega, \Omega = 3.34\) and \(\gamma_i = 0.43\).
(f) strange attractor, \(\Omega = 3.3\) and \(\gamma_i = 0.5\).

Fig. 2. Detailed Poincaré maps of the case in Fig.1(f), where data are selected at \(\tau = m2\pi / \Omega, 250 < m < 4000\).

Fig. 3. The experimental result of the amplitude-modulated motion of the soliton at 6 times of the driving modulation period, when \(2f = 10.20\) Hz, \(d_0 = 1.066\) mm, \(F = 1.62\) Hz and \(\gamma_i = 1.158\).
(a) the elevation spectrum of amplitude.
(b) the wave elevation at the centre of the soliton and near the wall.

Fig. 4. The experimental result of the chaotically amplitude-modulated motion of the soliton, when \(2f = 10.20\) Hz, \(d_0 = 1.077\) mm, \(F = 1.11\) Hz and \(\gamma_i = 1.15\).
(a) the elevation spectrum of amplitude.
(b) the wave elevation at the centre of the soliton and near the wall.
PARAMETRICALLY EXCITED MULTISOLITONS AND THEIR INTERACTIONS

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I. INTRODUCTION

In 1984, Wu et al. [11] observed a nonpropagating soliton in a parametrically driven long trough of water. Two such solitons will oscillate slowly along the trough length with a frequency of much less than the Faraday resonant frequency $f_0 = f_0/2$. In this process, the solitons overlap completely at their collision instants and hence they are indistinguishable. [3]. Further properties of the solitons, including the trough dimension dependence of the stability region (SR) and the transition to chaos, were investigated by Wei et al. [3, 4].

We focus our attention on the behaviours of multisolitons. Our apparatus is essentially same as Wu et al. [11]. The used troughs of finite length L and width W are filled with the distilled water of height h. Our detailed investigations have revealed many interesting phenomena not explored by others. Hereafter, a N-soliton state will be represented by a symbol $S(S_1 S_2 ... S_N)$, where the first soliton $S_1$ is denoted by sign "$+$" while the i-th $S_i$ (i=1, 2, ..., N) either by "$+$" or "$-$" depending on whether it vibrate in phase (same polarity) or 180° out of phase (opposite polarity) with the first one.

II. EXPERIMENTS

Bound State

First consider $S(++)$, the like pair. Depicted in Fig.1 is its SR in the driving amplitude $A_0$ and frequency $f_0$ plane and that of $S(++)$. Generally longitudinal direction with a variable period from 5 to 15 second. Therefore, the energy of the state should be greater than that of a single soliton but less than the sum of two independent ones. This is why we call $S(++)$ as the bound-state (or bound-pair in multisoliton states). (4) Moreover, especially in short troughs, there exists a critical value $A_c = A_c$ (A1+A2+A3) across which the complete overlapping type of interaction transits to that of partial overlapping and thus the solitons become completely indistinguishable (Fig.2). This fact clearly shows that the two solitons undergo the periodic reflections rather than exchange places in the course of interaction. This is verified as follows.

![Fig.2 Motion of S(++)](image)

(a) The two solitons are at their maximal separation and start attracting; (b) they reach their minimal separation and begin repelling. ($L_x = 20 \times 2.5 \times 2 \text{ cm}^2$, $A_0 = 0.88 \text{ mm}$ and $f_0 = 10.9 \text{ Hz}$).

Mirror Effect

If we vertically dip a smooth, rigid and thin board into water at the symmetric center of $S(++)$ so as to partition the fluid, we have surprising found that the motion of each soliton is unaltered. It says that the action of one soliton on the other is equal to that of the partition board. Conversely, any boundary (the partition board or one of the end-walls) has the "mirror effect" on a soliton state. Thus we expect that a soliton in the vicinity of a boundary may oscillate slowly as if there existed a virtual soliton on the other side, so that the real and virtual solitons would form what we will call the virtual bound-pair, as is verified below.

We generate a soliton near one end-wall and fill some absorbent at the other end-wall so as to enlarge the SRs. At some very small $A_0$, the soliton is attached to the wall and only half its envelope appears. As $A_0$ is increased, it becomes oscillating in the vicinity of the wall with a frequency $f_0 (\sim 0.1 \text{ Hz}) \ll f$ (Fig.3). Further increase of $A_0$ will diminish the motion and the soliton will settle down away from the wall ($\sim 0$ cm).

Of course, mirror effect strictly hold only in the ideal case, i.e., the boundaries are smooth and rigid, and the fluid-surface tension is negligible. It results in the slowing of the solitonic modulation. Due to the smoothness, the fluid surface is nearly perpendicular to the boundaries.

Applying the mirror effect to both end-walls $x=0$ and $L$, we can draw that a soliton state in a trough of finite length L is equivalent to an infinite system of period 2L through symmetric inversions at $x=\frac{L}{2}$. For example, $S(++)$ is
equivalent to
\[ S(\ldots \uparrow \uparrow \uparrow \downarrow \downarrow \downarrow \downarrow \downarrow \ldots) \] (1)

Fig. 3 Motion of a soliton near one end. The soliton is (a) attracted and (b) repelled by the wall alternately. (LxWxh=20x3x2 cm\(^3\), A\(_0\)=0.85 mm, and f\(_0\)=9.8 Hz).

**General Multisoliton State**

It was once observed\(^1\) that the unlike pair \(S(\uparrow \uparrow \downarrow)\) become standing after a repulsion process. But our observation shows that, due to the interaction with the end-walls, the two solitons oscillate slowly in the same direction with a frequency \(f_s\) about 0.1 Hz (Fig.4). The same motion also occurs in \(S(\downarrow \downarrow \uparrow \uparrow)\) and \(S(\uparrow \uparrow \uparrow \uparrow \uparrow \uparrow \uparrow \uparrow)\) etc. which have the same equivalent form (1). Thus we classify these states to the same pattern (Pattern 1), of which A N-soliton state \(S(N=2n, n=1, 2, \ldots)\) is composed of \((n-1)\) real bound-pairs and two virtual pairs; neighbouring pairs must oscillate oppositely in the longitudinal direction with a frequency \(f_s<<f\) as well as in the vertical direction (see Fig.5). We have also observed other two stable patterns: \(S(\uparrow \uparrow \downarrow \uparrow \downarrow), S(\downarrow \downarrow \uparrow \uparrow \downarrow \uparrow \downarrow)\) etc. (Pattern 2) and \(S(\uparrow \uparrow \downarrow \uparrow \downarrow \uparrow \downarrow \uparrow \downarrow)\), \(S(\downarrow \downarrow \uparrow \downarrow \uparrow \downarrow \uparrow \downarrow \uparrow \downarrow)\), etc. (Pattern 3). There is no difference among these patterns but that a N-soliton state of Pattern 2 \((N=2n, n=1, 2, \ldots)\) is composed of only a real bound-pair, while a N-soliton state of Pattern 3 \((N=2n+1, n=1, 2, \ldots)\) of a real and one virtual pair.

Fig. 4 Motion of \(S(\uparrow \uparrow \downarrow)\). (a) The left (right) soliton is attracted (repelled) by the left (right) wall, and (b) vice versa. (LxWxh=20x3x2 cm\(^3\), A\(_0\)=0.85 mm and f\(_0\)=9.7 Hz).

Fig. 5 Structure and motion of Pattern 1.

Fig. 6 Motion of \(S(\uparrow \uparrow \downarrow \uparrow \downarrow)\). One bound-pair are repelled each other and the other pair are attracted, and vice versa. (LxWxh=34x2.5x2 cm\(^3\), A\(_0\)=0.92 cm and f\(_0\)=10.9 Hz).

The states like \(S(\uparrow \uparrow \downarrow \uparrow \downarrow, S(\uparrow \uparrow \downarrow \uparrow \downarrow' \uparrow \downarrow'), etc. are also observable but less stable. It seems that a phase-reversed soliton between two bound-pairs would weaken the stability.

The most unstable states are those containing more than two solitons of like polarity that are adjacent each other. They exist only in a short and indefinite time after their formations. Perhaps, such an instability is attributed to the impossibility of forming bound-pairs among these solitons. According to the mirror effect, the very thing would also happen in Pattern 2 and 3. However, we should note that the effects of the surface tension and the friction, no matter how small but inevitable in practice, may force the solitons, e.g. the left in \(S(\uparrow \uparrow \downarrow)\), "choose" its real neighbor as its partner to form a bound-pair. Thus we are sure that it is the these non-ideal effects that stabilize \(S(\uparrow \uparrow \downarrow)\). Indeed, if we dip a thin board in the middle of the first and second solitons of \(S(\uparrow \uparrow \downarrow)\), a bound-pair is formed between the second and third solitons and the solitons become stable.

**III. DISCUSSIONS**

It is fascinating for the mirror effect to exist in our nonlinear system. Due to its existence, the finite-boundary-value problem is translated into a periodic one, which is significant in analytic treatment. The special patterns of the structures and motions of multisolitons exhibit the other particle-like characters than those of well-known solitons. An analytic understanding is underway.

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**REFERENCES**

SUBHARMONIC SEQUENCES IN PARAMETRICALLY
EXCITED ANNULUS RESONATOR

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§1. Introduction

The study was stimulated by Keolian et al's observation of subharmonic sequences in
annulus resonator with one-dimensional property and subjected to a vertical oscillation [1].
These authors observed rich spectrum of subharmonic instabilities which do not generally
follow the universal period-doubling sequences for long period subharmonics. These subharmonic
sequences may be expressed by the fraction f/m where f is the vertically
exciting frequency and m is prime numbers or their integer times. For typical preturbulent
series, the value of m observed at various drive amplitudes and frequencies (without the rigid reflektor) are

m=1, 2, 4, 12, 14, 16, 18, 20, 22, 24, 28, 35

It has been well known that a nonlinear system can generate subharmonics besides
superharmonics. Lord Rayleigh noted that two classes of driven systems possess subharmonic
response [2]: driven nonlinear oscillator described by a Duffing equation and parametrical
excitation described by a Mathieu equation. Feigenbaum has recently shown that a
universal period-doubling sequence arises quite generally in nonlinear one-dimensional
map, appearing as a precursor to chaotic behaviour [3].

In this paper, we try to provide a theoretical model for understanding Keolian et al's experimental observation.

§2. Hamiltonian formulation

We consider irrotational gravity waves in an inviscid fluid that fills a rigid annulus trough to a quiescent depth h. Let r = (r, θ) and z be horizontal and vertical
coordinates, respectively, in a reference frame fixed in the trough. The velocity potential Ψ(r, z, t) and free-surface displacement η(r, t) of the fluid must obey the canonical equations [4]

\[ \frac{\partial \eta}{\partial t} = \frac{\partial \psi}{\partial \eta}, \quad \frac{\partial \psi}{\partial t} = -i\partial \psi/\partial \eta \]  

\[ H = \int_{\mathbb{R}^2} \left[ \left( \frac{\delta^2}{\delta z^2} + \frac{\delta^2}{\delta r^2} \right) \psi^2 + \frac{1}{2} \frac{\delta^2}{\delta r^2} \eta^2 \right] \left( 1 + \gamma \cos 2\sigma t \right) \eta^2 \right] \frac{dr}{r^2} \]

where \( \psi(r, z, t) \) is the gravitational acceleration, \( \nabla \) is the two-dimension

gradient operator, \( 2\sigma \) and \( \gamma \) are the exciting frequency and amplitude, and \( R_1, R_2 \) are the inside- and outside-radius of annulus resonator. We can completely determine the flow of fluid if only \( \eta \) and \( \psi \) are given [5].

By means of the Fourier transformation

\[ \eta(k) = \frac{1}{2\pi} \int \eta(r) e^{-ik \cdot r} dr, \]

\[ \psi(k) = \frac{1}{2\pi} \int \psi(r) e^{-ik \cdot r} dr \]

and the canonical transformation

\[ \eta(k) = \sqrt{2}/2 \left( T(k) / \omega(k) \right)^{1/2} \left( a(k) + a^*(k) \right), \]

\[ \psi(k) = -i \sqrt{2}/2 \left( \omega(k) / T(k) \right)^{1/2} \left( a(k) - a^*(k) \right) \]

where \( \omega(k) = |g| k |T(k)|^{1/2} \) and \( T(k) = \text{tanh}(|k|h) \)

equation (2.1) is expressed

\[ \partial a(k)/\partial t = -i \partial \mathcal{H}/\partial a^*(k) \]  

and the Hamiltonian \( \mathcal{H} \) become a series of \( a(k), a^*(k) \).

Assuming the exciting amplitude \( \gamma = 0(\varepsilon^2) \)
where \( \varepsilon \) is a small parameter, we can expand \( a(k) \) according to the multiple-scale method as

\[ a(k) = \varepsilon A(a(k)) + \varepsilon^2 f(k, t) \text{exp}(-i \omega(k)t) \]

where \( \varepsilon = 2 \alpha \) and \( \omega^2(k) - \omega^2(k_0) = 0(\varepsilon^2) \). Substituting (2.5) into (2.4),
etuating terms like order of \( \varepsilon \) and then solving \( f \), we obtain

\[ \{ \varepsilon A(k) + \omega(k)^2 \} / \omega(k_0)^2 \]

\[ \times \{ \varepsilon A(-(k) + \varepsilon^2 f(k, t) \text{exp}(i \omega(k)t) \}

\[ + \omega(k_0)^2 (\omega(k) - \omega(k_0)) \delta(k_0 - k) \delta(k_0 - k_0) \}

\[ + \omega(k_0)^2 (\omega(k) - \omega(k_0)) \}

\[ = 0 \]  

where weak linear damping has been incorporated and \( a \) is the damping rate, and

\[ T(k, k_1, k_2, k_3) = -0.2, 0.2, -0.2, 0.2, -0.2, 0, 0.2, 0, -0.2, 0, 0.2, -0.2 \]

\[ \Delta = 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0 \]

\[ \text{ and } \]

\[ T(k, k_1, k_2, k_3) = -0.2, 0.2, -0.2, 0.2, -0.2, 0, 0.2, 0, -0.2, 0, 0.2, -0.2 \]

\[ \Delta = 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0 \]

\[ \text{ and } \]

\[ T(k, k_1, k_2, k_3) = -0.2, 0.2, -0.2, 0.2, -0.2, 0, 0.2, 0, -0.2, 0, 0.2, -0.2 \]

\[ \Delta = 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0 \]

\[ \text{ and } \]

\[ T(k, k_1, k_2, k_3) = -0.2, 0.2, -0.2, 0.2, -0.2, 0, 0.2, 0, -0.2, 0, 0.2, -0.2 \]

\[ \Delta = 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0, 0.2, 0 \]
and comparing the result with (2.3), we obtain

\[ a(k) \text{sn}_{m_l-1,l,m_l}^0(\tau) \exp(i\Delta \omega t) \tag{3.5} \]
in Koelion et al's case.

§4. Numerical analysis

Substituting \( A(k), \omega(k), \omega(t) \) and \( k \) in (2.6) with \( A_{ml}(\tau) \) of (3.5), \( \omega_{ml} \) and \( (k_{ml}, l) \) respectively, we obtain the discretized version of this equation

\[ i(\theta-\omega)(A_{ml}^{(l+1)} - A_{ml}^{(l)}) = \frac{1}{4} \gamma^2 A_{ml} \tag{4.1} \]

where \( l \) and the summation indices \( i,j,n \) are taken over all dominant modes in (3.5). In order to analyze the time-evolution of these dominant modes, numerical integrations of the set of differential equations (4.1) were carried out using an Adams-Bashforth routine with a round-off error of less than 10^{-6}. We examined the modulational evolution of two dominant modes both excited parametrically and determined the power spectrum of N-points runs of \( H \) through a standard fast-Fourier-transformation routine. We find that the spectrum, as illustrated in Fig.1, is in qualitatively agreement with the experimental observation.

\[ \uparrow \text{This work was supported by CYNS Foundation.} \]

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\[ \text{exciting freq. 10.2 Hz} \]
\[ \text{exciting ampl. 0.08 cm} \]
\[ \text{damping rate 0.24} \]

FIG. 1. The subharmonic sequence of excited gravity waves, as given by (4.1).
THERMOACOUSTIC REFRIGERATION RESEARCH at NPS

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INTRODUCTION

Observation of the production of sound by thermal gradients, such as the Rijke tube [1] or Taconis oscillations in liquid helium, is relatively common. The reverse process – thermoacoustic heat pumping – is far less well known and was the first intentional demonstration of a new class of intrinsically irreversible heat engines [2].

In thermoacoustic engines, the irreversibility caused by imperfect (diffusive) thermal contact between the acoustically oscillating working fluid and a second thermodynamic medium provides the required phasing for a heat engine. This “natural phasing” has produced heat engines which require no moving parts other than oscillations of the working fluid.

A Simple Lagrangian Model of the Heat Pumping Process

A schematic diagram of the Space Thermoacoustic Refrigerator (STAR) flown on Space Shuttle Mission STS-42 is shown in Figure 1. The driver assembly at the top, consisting of a highly modified commercial electrodynamic loudspeaker epoxied to a bellows, sets up the standing wave within the gas-filled tube. Its frequency is chosen so that the driver excites the fundamental $(\lambda/4)$ resonance of the tube. Below the driver unit and hot heat exchanger there is a stack of plates (the “stack”) whose spacing is chosen to be a few thermal penetration depths.

\[ \delta_x = \frac{\sqrt{\frac{\lambda}{k\rho p}}}{} \]

This length scale is crucial to understanding the performance of the thermoacoustic cycle since the diffusive heat transport between the gas and the “stack” is only significant within this region. Accordingly, the stack and the spacing between its plates are central to the thermoacoustic cycle.

In Figure 2, a small portion of the stack has been magnified and a parcel of gas undergoing acoustic oscillations is shown. The four steps in the cycle are represented by the four boxes shown as moving in a rectangular path for clarity; in reality the parcels simply oscillate sinusoidally back and forth. As the fluid oscillates along the plate, it undergoes changes in temperature because of the adiabatic compression and expansion of the gas from the pressure variations which accompany the standing sound wave.

![Figure 1. Space Thermoacoustic Refrigerator](image)

![Figure 2. Thermoacoustic Heat Transport](image)

The thermodynamic cycle can be considered as consisting of two reversible adiabatic steps and two irreversible isobaric (constant pressure) steps. The plate is assumed to have a mean temperature, $T_m$, and a temperature gradient, $\nabla T$, referenced to the mean position, $x = 0$.

In the first step of this four-step cycle, the gas parcel is transported along the plate by a distance $2x_1$ and is heated by adiabatic compression from a temperature of $T_m + x_1 \nabla T$ to $T_m + x_1 \nabla T + 2T_1$, represented as T+ in Figure 2. The adiabatic gas law provides the relationship between the change in gas pressure, $p_1$, and the associated change in temperature, $T_1$. Work in the form of sound was done on the gas parcel which, in its present location, now has a temperature higher than that of the plate.

In the second step, the warmer gas parcel transfers an amount of heat, $dQ_{hot}$, to the plate by thermal conduction at constant pressure. In the third step, the gas is transported back along the plate to position $-x_1$ and is cooled by adiabatic expansion. This temperature is lower than the original temperature at location $-x_1$, so in the fourth step the gas parcel adsorbs an amount of heat, $dQ_{cold}$, from the plate thereby raising the parcel temperature back to its original value, $T_m - x_1 \nabla T$.

The net effect of this process is that the system has completed a cycle which has returned the parcel to its original state and an amount of heat, $dQ_{cold}$, has been transferred to the system. This is a temperature gradient by work done in the form of sound. If this point that no mechanical devices were used to provide the proper phasing between the mechanical motion and the thermal effects. In a thermoacoustic engine, the proper phasing occurs naturally since the sinusoidal temperature oscillation of the gas parcel is phase shifted (in time) with respect to its motion as a result of heat diffusion through the thermal boundary layer.
OTHER THERMOACOUSTIC REFRIGERATORS AT NPS

The STAR is the first in a series of spacecraft cryocoolers now under development at the Naval Postgraduate School. Projects in progress include the third generation cryocooler, TAR-3, and the Thermoacoustic Life Sciences Refrigerator, TALSR, which is designed to pump 200W at 4°C and 120 W at -22°C.

The requirements of spacecraft cryocoolers (temperature spans of 100°F – 200°F) are opposite those of residential refrigerator/freezers and air conditioners. These applications require modest temperature spans (AT = 25 - 45°F), but heat pumping powers on the order of hundreds of watts for refrigerators and thousands of watts for air conditioners. Designs capable of pumping about 2 kW of heat across a 30°C temperature span with a COP = 3-4 (about 30% of Carnot performance) including electroacoustic conversion efficiencies have been completed.

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PIEZOE XCITA TION OF RAYLEIGH- TYPE ACOUSTIC WAVES INDUCED BY NEAR-SURFACE ABSORPTION OF LASER PULSES
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The highest efficiency of light conversion into Rayleigh surface acoustic wave (SAW) has been achieved by laser action on a photoconductive piezocrystal, with the simultaneous application of an external electric field. Theoretical analysis supports the opinion of the authors of ref. 1 that SAW excitation was primarily carried out by the external piezoeffect in the nonstationary electric fields induced by spatial separation of optically generated charge carriers.

In our experiments high-resistance CdS samples were used. A constant external voltage, 0.3 kV ≤ U ≤ 3 kV, was applied to the electrodes deposited on the front surface of the crystal at the distance of 0.5 cm one from another. To realize near-surface absorption of light, we irradiated CdS by the third harmonic of pulsed YAG:Nd<sup>3+</sup> laser (hv = 3.5 eV - energy of optical quantum, τ<sub>L</sub> = 10 ns - laser pulse duration, α = 2.5 × 10<sup>8</sup> cm<sup>-2</sup> - coefficient of light absorption). The laser beam was focused with a cylindrical lens into a strip (dimensions X<sub>S</sub> × Y<sub>S</sub> ~ 3 × 3 mm<sup>2</sup>) parallel to electrodes (Fig. 1).

Fig. 2 represents typical profiles of SAW pulse registered by the large bandwidth piezoelectric transducer. The profile depicted in Fig. 2 (a) was obtained in the case when the external electric field was directed opposite to the SAW propagation vector. The inversion of the SAW profile caused by the reversal of voltage polarity provides the first direct experimental evidence for a piezoelectric origin of the observed optoacoustic (OA) conversion. The SAW signal amplitude A scales linearly with external electrical field (A ~ U) (Fig. 3). Fig. 4 shows that the photoexcited SAW amplitude saturates when the light intensity I reaches values 100 W cm<sup>-2</sup> ≤ I ≤ 300 W cm<sup>-2</sup>. We have additionally confirmed that the saturation effect persists up to I ≤ 10<sup>4</sup> W cm<sup>-2</sup>.

Theory describing the saturation of Rayleigh SAW piezoexcited as a result of weak (bulk) light absorption has been developed in ref. 2. It predicts the saturation effect for the intensities I ≥ I<sub>L</sub><sup>10</sup>. L<sub>L</sub><sup>10</sup> occurs when the concentration of free carriers generated in the process of interband light absorption becomes so high that the external electric field fails to separate spatially all the electrons and holes. The amplitude of the nonstationary electric field saturates at the levels of the external external electric field of magnitude E. The critical intensity L<sub>L</sub><sup>10</sup> can be estimated using the condition that the field of the "capacitor" (formed in the process of charge separation) totally compensates the external field. The crude estimate of the maximum concentration of electron-hole (EH) plasma n<sub>e</sub> can be obtained considering E ~ (U/a) and relative dielectric constant ε ~ 10

\[
(n_e^+ cm^3)^{3/2} \approx 1.0 \times 10^{11} (U/eV)
\]

In our experiment the absorption region is diffusively thin (√D<sub>n</sub> ≪ α<sup>1</sup>), D - coefficient of ambipolar diffusion. Then neglecting diffusion along the surface (√D<sub>n</sub><sqrt>ε<sub>e</sub>) and EH recombination:

\[
(n_e/cm^3) \approx (1-R)ln/√D<sub>n</sub>hν \approx 1.4 \times 10^{14} (1/W cm^2)
\]

Here the following values of the parameters have been adopted, R = 0.2 is the light reflection coefficient, D ≈ 1 cm<sup>2</sup> s<sup>-1</sup>. Comparison of Eqs. (1) and (2) shows that theoretical estimate does not correlate with experiment (Fig. 4). It should be noted that similar analysis yields correct order-of-magnitude estimates for the critical intensities in experiments using bulk absorption of laser radiation.<sup>1</sup> In the present experiment the SAW signal increases with laser pulse intensity as the result of the growth in size of the space region where piezoresources of the acoustic wave are located. This growth is initiated by the diffusion of photogenerated EH plasma into the bulk of the sample.

Theoretical analysis revealed that a limitation in size of the space region over which the electric field is "swept out" may be caused by the development of a nonlinear recombination process. When the time of bimolecular recombination τ<sub>B</sub> = (Bn)<sup>-1</sup> (B - recombination constant) becomes less than laser pulse duration τ<sub>L</sub>

\[
(Bn)^{-1} \approx \tau<sub>L</sub>
\]

then a quasistationary distribution of the nonequilibrium plasma concentration takes place. This concentration can be derived from equation

\[
D<sub>n</sub> - Bn^2 = 0
\]

and the boundary condition at x = 0 D · n<sub>E</sub> = -(1-R)I(hν)^1/2(x/x<sub>S</sub>)

In Eq. (5) the function ψ describes the surface distribution of laser intensity. The explicit solution of the problem Eqs. (4), (5) can be written in the form

\[
n(x) = n_0 \tilde{\psi}^2 \tilde{\psi} \tilde{\psi}^2 [1 - \tilde{\psi}^2]^2 D<sub>n</sub> (n_0) [1/2]^{-2}
\]

where n<sub>E</sub> = [(3/2BD)](1/2)(1-R)/hν]^{-1/3}

Using Eq. (6) one can determine the characteristic dimensions x<sub>n</sub> and x<sub>e</sub> of the screened region of space.

Considering ψ = exp [-2(x/x<sub>S</sub>)]<sup>2</sup>, for I ≫ I<sub>L</sub><sup>10</sup> we find

\[
x_e \approx \sqrt{3/2} ln (n_0/n_e)/[x_S/2]
\]

Theoretical Eq. (7) describes an extremely mild dependence of x<sub>n</sub> on laser intensity (x<sub>n</sub> ~ [ln(1)]^1/2). It is important that x<sub>n</sub> considerably exceeds the radius of the laser beam x<sub>l</sub>/2. For B = 10<sup>-8</sup> - 10<sup>-6</sup> cm/s<sup>-4</sup>, Eq. (6) yields the estimate (n_e/cm^3) ~ (3-1.5) × 10<sup>14</sup> (1/100 W cm<sup>-2</sup>). It means that in the intensity range 100 W cm<sup>-2</sup> ≤ I ≤ 10<sup>4</sup> W cm<sup>-2</sup> and U ~ 1 kV the relation x<sub>n</sub> ~ (4-5) (x<sub>e</sub>/2) is fulfilled. Consequently, the characteristic duration of a SAW pulse may be of the order of the sound propagation time across the screened
part of the crystal $\tau_0 \sim 2x_0/c_0 \sim 700-900 \text{ ns}$ ($c_0 \approx 1.7 \times 10^3 \text{ cm/s}$) is the velocity of the Rayleigh type SAW on $Z$ - cut (cds). This estimate correlates with the experimental results (Fig. 2). It is extremely important that in accordance with Eq. (6) the plasma concentration at the distance larger than the recombination length ($x > \sqrt{D_{n0}(n_0)}$) does not depend on the laser intensity at all; $n \approx 6D/Be^q$. It is worth emphasizing that this situation is specific to nonlinear EH recombination. According to considered equations one can show that for $L \gg L_{n0}$ the inequality $x > \sqrt{D_{n0}(n_0)}$ is valid. That is why fulfillment of Eq. (3) is sufficient to bound the acoustically active volume.

With the help of Eq. (3) and Eq. (6) we obtained the following estimate of the critical intensity; $L_{n0}^{(c)} \sim (500-50)$ $\text{ W/cm}^2$. Taking into account the uncertainty of the recombination constant $B$, one can consider this result to correlate well with the experiment (Fig. 4).

In conclusion we have observed saturation of SAW piezoexcited by laser induced screening of external electric field in CdS crystal. The experimental estimates support the hypothesis that the saturation effect is connected with the influence of nonlinear recombination on the photogenerated EH plasma evolution ($\tau_0 \leq \tau_0$ for $L \gg L_{n0}$).

References
RELATIONSHIP BETWEEN THE MOLECULAR CONSTANT, THERMO-AcouSTIC PARAMETERS AND INTER-MOLECULAR VOLUME EXPANSIVITY OF LIQUIDS

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INTRODUCTION

There are a number of useful acoustic and anharmonic parameters which are interrelated. The anharmonic microscropic isobaric $\Pi$, isothermal $\Pi'$ and isochoric $\Pi''$ Gruneisen parameters (as measures of anharmonicity of molecular vibrations) are treated identical [1] to the corresponding acoustical parameters $\lambda$, $\lambda'$ and $\lambda''$ for a liquid. The parameters $\Pi'$, $\Pi''$, $\lambda$ and $\lambda''$ expressing respectively the isochoric temperature derivative of sound speed, microscopic isothermal Gruneisen parameter, isothermal volume derivative of thermodynamic Gruneisen parameter, fractional free volume (as a measure of disorder due to increased mobility of the molecules) are found to be important factors which contribute significantly to the thermo-acoustic, anharmonic and non-linear properties of liquids [2-4].

Using the expressions obtained [2-4] for the anharmonic and thermo-acoustical parameters, the parameters $\Pi$, $\Pi'$, $\lambda$, $\lambda''$, $\lambda$ and the dimensionless parameter $A^*$ are interrelated and expressed in terms of the molecular constant $n$ (as a measure of the ratio of the internal pressure to cohesive pressure) of the liquid. Also, the ratio ($\Pi/n$) of the thermodynamic to microscopic isobaric Gruneisen parameter [1,4] is utilized to measure the intermolecular vibration contributions to heat capacity and volume expansivity of the liquid, furnishing understanding of the significance of microscopic Gruneisen parameter and fractional free volume through molecular constant, in describing the anharmonic behaviour with regard to molecular order and intermolecular interactions in liquids, has been developed. The treatment has the distinct advantage that the anharmonic parameters can be evaluated from the thermo-acoustic data of liquids.

THEORY

The thermodynamic Gruneisen parameter $\Pi$ (used for structural study of liquids) for a normal mode model of liquid state [1,4,5] and the microscopic isobaric Gruneisen parameter $\Pi$ for a liquid [1,4] can be expressed as

$$\Pi = \lambda V B / C_v = L / T C_v = M C_v^{\Pi} / C_p$$

$$= (\gamma - 1) / \lambda T$$

(1)

$$\Pi' = -\frac{1}{2}(\rho \ln B / \lambda T) + 1$$

(2)

in which $V$, $C_v$, $M$, $L$, $\lambda$, $B$, $B_s$ are respectively the molar volume, sound speed, molecular weight, internal molal latent heat of vaporization, volume expansivity, isothermal and adiabatic bulk modulus, and $\rho = (C_p/C_v)$ is the heat capacity ratio, $C_p$, $C_v$ are isobaric and isochoric heat capacity of the liquid at absolute temperature $T$ and pressure $P$.

To establish agreement between eqs. (1) and (2) and generalize the suggestions [1,4], the intermolecular contribution to various quantities, using eqs. (1) and (2), can be written as

$$\Pi = \Pi = \frac{L}{T} C_v = \lambda / B V / C_v$$

$$= M C_v^{\Pi} / C_p = (\gamma - 1) / \lambda T$$

(3)

where $\lambda$, $\gamma = (\Pi / \Pi')$, $\Pi$, $\Pi'$, $C_v$, $C_p$, respectively represent intermolecular contribution to $\lambda$, $\gamma$, $\Pi$, $\Pi'$, $C_v$, $C_p$ for the liquid.

Eqs. (1) and (3) transform into the relations given by

$$C_v / C_v = \lambda / C_p / C_v = (\gamma - 1) / \lambda T$$

$$= (\Pi / \Pi') = \lambda / B / C_v / \lambda / C_v$$

(4)

$$\lambda / \lambda = (\gamma - 1) / \gamma = (B / B_v)$$

$$= (\Pi / \Pi') / (1 - \gamma - (1 - \Pi / \Pi'))$$

(5)

Using eq. (2) and the expression obtained [2] for the molecular constant $n$ as a measure of the isochoric temperature derivative of internal pressure), the parameters $\Pi'$, $\Pi''$, $\lambda$, $\lambda'$, and $\lambda''$ may be related to $n$ as

$$K'' = -K' = K = (d ln C_v / d T_v) / \lambda = 1 + (1 - n) \gamma / \lambda$$

(6)

$$n' = (\Pi / \Pi') (1 - n) \gamma / \lambda T$$

(7)

$$f = (\Pi'' + 1) = k + 1 + k' + k''$$

(8)

$$A^* = 1 + (\Pi' / \Pi' + 1)$$

(9)

$$= 1 + f / \Pi'$$

in which $V'$, $C'_v$, $M$, $L'$, $\lambda'$, $B'$, $B_s'$ are respectively the molar volume, sound speed, molecular weight, internal molal latent heat of vaporization, volume expansivity, isothermal and adiabatic bulk modulus, and $\rho = (C_p/C_v)$ is the heat capacity ratio, $C_p$, $C_v$ are isobaric and isochoric heat capacity of the liquid at absolute temperature $T$ and pressure $P$.

Eqs. (1) and (10) establish the interrelationships between various thermo-acous-
tic and anharmonic parameters and can be deter-
mined from the thermo-acoustic data of li-
guids.

DISCUSSION

Eqs. (1), (2), (4) and (5) are of great
interest once these can be utilised to
measure the intermolecular vibration contrib-
utions $C_{v,k}$, $\alpha$, $B_3$, $C_p$, $\alpha$
to the respective quantities from the knowl-
edge of thermo-acoustic data on $C_{v,k}$, $\alpha$, and
$(dC/d\omega)^*$ for liquids. Eqs. (4) and (5) give
general, correct form of the relationship
between intermolecular volume expansivity,
heat capacity and the Gruneisen parameter
for technical purposes. The present results
of eq. (5) confirm the analogous proposal
[6-9] that the bulk modulus and volume expansi-
vity of a polymer glass are determined by
the interchain forces.

For liquids, it has been observed
[2,4,10] that $\eta > \eta_0$ which implies that a
certain fraction of the modes in a liquid
has a finite and common fractional frequency
dependence on volume at constant temperature
and at constant pressure and the remaining
modes are unaffected by volume. Consequently,
the value of $(\eta_0/\eta)$ for a liquid is more
than unity, as is evident from eqs. (4) and
(5), so that the intermolecular contribution
to $C_p$ are expected much higher by the factor
$(\eta_0/\eta)$ than that to $W$ in the liquid. Thus
the volume dependent, anharmonic acoustic
modes represent only a fraction of the
normal mode molecular vibrations that con-
tribute to $C_v$, $C_p$, $\alpha$ of the liquid. This
result agrees with similar suggestions [4,10] that
not all the normal mode contributions to $C_p$
also contribute to the volume change of
sound speed thereby exhibiting molecular
rotation in liquids with strong intermolec-
ular interactions and anharmonicity effects.

Eq. (6) shows that $K^*$ can be evaluated
from the experimental data on $n$ and $\alpha$
available in literature [11]. Eq. (6)
imparts a value of $K^*$ less than unity in
agreement with experimental value [2,10].
Using these calculated values of $K^*$ and
those of $X$ from experimental data on $C_{v,k}$, $\alpha$
and $(dC/d\omega)^*$, the parameters $\eta'$, $f$, $A^*$ and
$\lambda$ can then be evaluated using eqs. (7) - (10)
for the liquids. Eqs. (7) and (8) show that
the fluorocarbon fluids having higher values
of $\eta'$ are strongly anharmonic as compared to
other liquids [4,12] having higher values
of $f$, similar to that observed for polymers
[3].

Eq. (9) shows that at lower temperatures,
close to absolute zero, the substance would
tend to be ordered, exhibiting little ther-
mal expansion and fractional free volume,
thereby making the parameter $A^*$ a constant
equal to unity. Eqs. (6) - (10) demonstrate
the significance of $n$, $\eta'$ and $f$ in describing
thermo-acoustic properties and anharmonic
behaviour with regard to molecular order
and inter-molecular interactions in liquids.

The present treatment offers a convenient
means for establishing relationship
between the thermo-acoustic and anharmonic
and correlating with fractional
free volume and Gruneisen parameter through
molecular constant as a useful parameter for
investigating several, thermo-acoustic
properties of liquids.

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PHOTOACOUSTIC PIEZOELECTRIC IMAGING OF SUBSURFACE STRUCTURES

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INTRODUCTION

Photothermal techniques are a non-destructive tool for obtaining information about the variations of the optical and thermal properties of a sample, also as a function of depth. They are able to detect the presence of surface and subsurface inhomogeneities.

A large number of studies has concerned the applications of thermal wave imaging in metals, semiconductors, ceramics, polymers. A review is given in [1]. These analyses are based on examining the changes of signal amplitude and phase, as a consequence of differences in material properties which influence the generation, propagation and detection of thermal waves. These properties include the surface optical characteristics (optical absorption and scattering), which determine the initial distribution of the thermal waves generated after the absorption of the incident energy. But the most important parameter is the thermal diffusivity of the material, by which the deep propagation is controlled. The presence of thermal boundaries causes thermal wave scattering, reflection and refraction.

In the particular case of a strongly absorbing material, thermal waves are generated near the surface and decay while propagating inside the material. The wave amplitude decreases by a factor 1/e in one "thermal diffusion length" \( w = (\alpha/\sigma \cdot \tau) \), where \( \alpha \) and \( \tau \) are the thermal diffusivity of the sample and the modulation frequency, respectively. This relationship indicates that the thermal wave propagation distance in a sample can be changed by changing the modulation frequency. For common materials of interest in thermal wave imaging, such as metals, semiconductors and ceramics, the thermal diffusion length varies within a range from a few microns to one millimeter. For instance, in the case of crystalline silicon, \( w = 1.5 \) mm at 1000 Hz, \( w = 15 \) \( \mu \)m. This range of thermal diffusion length values corresponds to the range of depth values along which it is possible to get information about the structural conditions of the sample.

The thermal-wave image is obtained by detecting the signal at one of the sample surfaces, as a function of the position of the point of excitation and/or detection. Several methods have been developed in last years for the detection of thermal waves, each one with its own advantages and disadvantages. Some methods require the physical contact with the sample (gas-microphone, piezoelectric and piroelectric detection); with other methods the detection is performed remotely (radiometric, optical beam deflection -laser-, photothermal displacement and modulated reflectance detection).

This paper reports some experimental results obtained by a photoacoustic piezoelectric imaging technique. The principle of piezoelectric detection is that the absorption-induced heating of the sample causes a thermal expansion. A transducer attached to the sample converts the thermal stresses and strains into a measurable voltage. A complete discussion of this type of detection was performed in [2] and in analysis of its depth-profiling capabilities in [3].

With respect to other detection techniques, piezoelectric method has the disadvantage to require the contact with the sample; moreover in some cases piezoelectric signal is difficult to interpret because it depends not only on optical and thermal properties but also on local thermoelastic properties and boundary conditions. The main advantage is the possibility to operate at high modulation frequencies (even of MHz) and thus to get a better image resolution (shorter diffusion length). Another advantage is to be an easy-to-implement and strong technique, which can be used also under operative conditions. Piezoelectric imaging is a back detection method, which therefore has the advantage to allow an imaging depth in the material approximately two times higher than from surface detection. The difference in the involved depth range was shown in [4], by means of a direct comparison with gas-microphone technique.

EXPERIMENTAL MEASUREMENTS AND DISCUSSION

The experimental apparatus was similar to that described by other authors. The detector was a piezoceramic disk (PZT 4 of Vernitron). The PZT was joined to the sample in two different ways, the first one by gluing the detector with a cyanoacrylate adhesive and the other one fastening mechanically the sample against the detector by means of a quartz window, into which the incident radiation enters. The results obtained with the two arrangements were substantially equal; the mechanical coupling had the advantage that the sample could be changed more easily. The beam from an argon-ion laser (wavelength 514 nm) was AO modulated and focussed on the sample surface, with a power in the range 0.1-0.8 W. Owing to the high impedance of PZT transducers, a charge amplifier was used before monitoring the signal with a two-phase lock-in analyzer. Scanning was performed by translating the sample stage with a stepping motor. The system was completely automatized and controlled by a PC.

The first piezoelectric imaging was made on an Al sample with a subsurface rectangular slot, 2 mm wide and with depth gradually increasing from the surface. Fig.1 reports the results, in amplitude and phase, obtained at a modulation frequency of 20 KHz in a scan area of 5x5 mm² and with a scan step of 100um.
The first consideration is the different resolution of the two images: the amplitude is influenced by the optical surface structure, in this case the raw surface finishing. The phase is not affected by the surface structure and clearly shows the increasing depth of the subsurface slot. At 20 Hz the thermal diffusion length in Al is nearly 1100 μm. In the considered field of view the ratio between the slot depth d and w varies between 1.3 and 1.8. The corresponding phase shifts are 60° and 10°. This confirms, as it was demonstrated by other authors [3], that the depth profiling capabilities of piezoelectric imaging depend on the ratio d/w.

The same considerations are confirmed by the results obtained on another Al sample with two subsurface holes, having a diameter of 1 mm and ending 200μm and 500μm from the surface. The phase images, at modulation frequencies of 8Hz, 20Hz, 160Hz, 320Hz are shown in Fig.2, for a scan area of 2x3 mm². The holes are visible in all images. In the investigated range of thermal diffusion length values, from 1700μm to 280μm, the less deep hole is always clearly visible, while the image of the deeper hole degrades when increasing frequency. The resolution at the boundaries depends on the lateral diffusion of the thermal wave and increases when decreasing the thermal diffusion length.

Another sample investigated was a semiconductor power device consisting of a n-type silicon wafer with p-diffused zones, bonded to a metal base which serves as a heat sink. If the wafer is not well bonded, the power device can undergo failure. Unbonded zones can be detected by photoacoustic microscopy, as a strong thermal discontinuity is involved. Fig.3 shows the phase image of the device, at 160Hz, and a wide unbonded zone is visible below the wafer, which is 360μm thick. In the figure also p-zones are evident. The possibility to detect these zones, which are approximately 50μm deep, cannot be explained on the basis of different thermal properties (the diffusion process does not cause appreciable lattice disorder), nor can be only interpreted as due to variations of the optical absorption coefficient. The same problem arises in interpreting the image of Fig.4, where the device boundaries are visible with a high resolution. Owing to the fact that the image obtained by PZT detection concerns different material properties, optical, thermal and elastic, a possible mechanism of generation could be connected to the different elastic properties of p-zones with respect to the n-type wafer [5].

CONCLUSION

In this paper it has been confirmed that photoacoustic piezoelectric imaging allows to get information on surface and subsurface properties in metals and semiconductors. This technique, although the physical contact with the sample is required, is the easiest to implement among photothermal techniques, also under real operating conditions. Because the piezoelectric signal is determined not only by optical and thermal properties, but also by elastic properties, the results in some cases are more difficult to interpret.

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A TRANSIENT SOLUTION FOR THE AXIAL PRESSURE FIELD OF A SPARK-SOURCE LITHOTRIPTER

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Lithotripters use focused shock waves to disintegrate kidney stones and gallstones. One way to generate the shock wave is with a spark source at the near focus of an ellipsoidal mirror. The acoustic energy in the outgoing spherical wave produced by the spark is reflected from the mirror and concentrated at the far focus of the ellipse, where the stone is positioned. In this paper, a transient solution is presented for the pressure field along the axis of the mirror. Short wavelengths, and lossless, small signal propagation in a homogeneous fluid are assumed. The solution provides insight into the main effects due to diffraction by the mirror. A simple result is obtained for the pressure at the far focus.

THEORY

The following theory is an extension of the solution obtained by Cornet and Blackstock\(^1\) for the reflection of an N wave from a spherical mirror.

The surface of the mirror is defined by \(z^2/a^2 + (x^2 + y^2)/b^2 = 1\) for \(-a \leq z \leq a - d\) (see Fig. 1), where \(2a\) is the length of the major axis, \(2b\) the length of the minor axis, and \(d\) the depth of the mirror. Through the definition of the eccentricity \(\epsilon = \sqrt{1 - (b/a)^2}\), the position of the near focus (where the source is located) is written \(z_1 = -\epsilon a\), and the far focus (where the stone is located) is at \(z_2 = +\epsilon a\). The pressure \(p_0\) of the spherical wave generated by an omnidirectional source centered at \(x = z_1\) is described by

\[
p_1/p_0 = (r_0/r_1)f(t - r_1/c_0)
\]

where \(r_1 = \sqrt{x^2 + y^2 + (z - z_1)^2}\) is distance from the near focus, \(r_0\) characterizes the pressure at reference distance \(r_0\), \(f(t)\) is an arbitrary dimensionless function of time, and \(c_0\) is the sound speed.

We consider first the solution \(p_{33}\) for the reflected pressure field that is predicted by geometrical acoustics. With the assumption that \(p_{33} = p_3\) everywhere on the surface of the mirror, the solution for the reflected field is found to be

\[
p_{33}/p_0 = \frac{r_0}{r_2} \left(1 + \frac{4\epsilon \sin^2(\theta/2)}{(1 - \epsilon)^2}\right)^{-1} \left[1 + \frac{t + r_2 - 2a}{c_0}\right]
\]

where \(r_2 = \sqrt{x^2 + y^2 + (z - z_2)^2}\) is distance from the far focus, and \(\theta\) is the angle formed with the \(x\) axis (recall Fig. 1). Note that the reflected field becomes increasingly concentrated about the \(x\) axis as \(\epsilon \rightarrow 1\).

Equation (2) does not account for the effects of diffraction. We shall take diffraction into account by obtaining a solution of the Kirchhoff integral. If the characteristic wavelength of the incident field is small compared with the minimum radius of curvature of the mirror, Eq. (2) provides an accurate boundary condition for the reflected field on the surface of the mirror. In terms of the time \(T\) that characterizes the periodicity or duration of \(f(t)\), the restriction on using geometrical acoustics to obtain the boundary condition is nominally

\[(1 - \epsilon^2)^2a\pi c_0 T \gg 1\]

which is generally satisfied by lithotripsy pulses. The Kirchhoff integral for the reflected pressure \(p_3\) is thus

\[
p_3 = \frac{1}{4\pi} \int_S \left\{p_{33} \frac{\partial}{\partial n} - \frac{1}{c_0} \frac{R}{\partial n} \frac{\partial p_{33}}{\partial t} - \frac{R}{R} \frac{\partial p_{33}}{\partial n}\right\} dS
\]

where \(p_{33}\) is given by Eq. (2), \(S\) is the surface of the mirror, and all other notation is standard.

Along the \(x\) axis, Eq. (4) reduces to

\[
p_3 = H_c(z)f(r_1) + H_e(z)f(r_1) + \frac{c_0}{c_0} H_a(z,t) f(t - r_1)\]

We shall refer to the first term as the edge wave, the second term as the wave form, and (following the terminology used by Naive Tjetta and Tjetta\(^2\)) the third term as the wake, where

\[
H_c = \frac{(1 + \epsilon)}{1 - \epsilon} \frac{r_0}{a - \epsilon a - z}
\]

\[
H_e = \frac{(1 - \epsilon^2)(a + R_a)(a - d)}{2[a - (a - d)]}\frac{r_0}{a - (a - d) + (z - \epsilon R_e + \epsilon^2(a - d))}\frac{R}{R}
\]

\[
H_a = \frac{\epsilon(1 - \epsilon^2)(a + R_a) + \epsilon a + \epsilon^2 a^2}{c_0}\frac{R}{R}
\]

\[
R_e = \sqrt{(1 - \epsilon^2)(2a - d) + (z - a - d)^2}
\]

\[
R_a = \sqrt{(1 - \epsilon^2)(a - d)^2 + (a - z - d)^2}
\]

\[
\tau_c = t - (a + (2 - \epsilon)a)/c_0
\]

\[
\tau_e = t - [R_a + (1 - \epsilon)\alpha + \epsilon d]/c_0
\]

\[
t_1 = [z + (2 - \epsilon)a]/c_0
\]

\[
t_2 = [R_e + (1 - \epsilon)\alpha + \epsilon d]/c_0
\]

\[
t_3 = [z + (2 - \epsilon)a]/c_0
\]

At the far focus \((z = z_2)\), Eq. (5) reduces to

\[
H_e = (1 - \epsilon^2)(a + R_a)(a - d)/2[a - (a - d)]\frac{r_0}{a - (a - d) + (z - \epsilon R_e + \epsilon^2(a - d))}\frac{R}{R}
\]

\[
H_a = \frac{\epsilon(1 - \epsilon^2)(a + R_a) + \epsilon a + \epsilon^2 a^2}{c_0}\frac{R}{R}
\]
More detailed discussion of the theory will appear in a forthcoming paper.

COMPARISON OF THEORY WITH EXPERIMENT

For the calculations in Fig. 2 we considered a half ellipse \((d = a)\) with eccentricity \(\epsilon = 0.7\) (and therefore a far focus at \(x_0/a = 0.7\)). The source function \(f(t)\) shown in Fig. 2(a) has dimensionless duration \(c_0T/a = 0.03\), and we have used \(r_0 = a\) for the reference distance from the source. These parameters approximate those of an experiment performed in water by Müller\(^4\), whose measured waveforms (reproduced from Fig. 4 of Ref. 4) at three axial locations \((z/a = 0.47, 0.75,\) and 0.93) appear in the left column of Fig. 2, opposite the corresponding theoretical predictions based on Eq. (5). Müller used a Dornier XL1 lithotripter \((a = d = 11\ cm, b = 7.8\ cm)\, and the duration of his incident pulse was approximately \(T = 4\ \mu s\). The pairs of vertical dashed lines in Figs. 2(c)-2(e) identify the approximate beginnings and ends of the oscilloscope traces in the left column. The label \(C\) in the figures identifies the beginning of the center wave, \(E\) identifies the beginning of the edge wave, and \(W\) identifies the most pronounced contribution due to the wake.

At \(z/a = 0\), the center wave and edge wave are well separated, and the wake produces a slight negative pressure immediately following the center wave. At \(z/a = 0.47\) the center wave and edge wave are still clearly resolved. Comparison of Fig. 2(c) with the oscilloscope trace to the left indicates that the small negative pressure at the end of the measured waveform is evidently due to the wake rather than the edge wave. The oscilloscope trace appears to end prior to the arrival of the edge wave. At \(z/a = 0.75\), just beyond the far focus, the center wave and edge wave overlap, and the wake produces a large positive pressure. However, the predicted waveform in Fig. 2(d) appears backwards in comparison with the corresponding measured waveform. This reversal is probably due to nonlinear effects introduced by the large peak pressure, which is approximately 800 bar (80 MPa). Nonlinearity would cause point \(W\) in Fig. 2(d) to catch up with point \(E\) and thus produce a waveform more like that which was measured. Note also that the amplitude ratio of point \(W\) in Fig. 2(d) to point \(C\) in Fig. 2(c) matches the ratio of the corresponding measured peak pressures. At \(z/a = 0.93\), both the measured and predicted waveforms possess a horizontal plateau that follows the arrival of the edge wave \((E)\), leading to a spike produced by the wake \((W)\) in the middle of the waveform. Again, nonlinear effects would cause point \(W\) in the predicted waveform to advance in time relative to point \(E\). Although the predicted amplitude ratio of point \(W\) to point \(E\) matches the measured ratio, the predicted negative pressure due to the center wave is much greater than was measured. At \(z/a = 2\), the edge wave has separated from the center wave, and the effect of the wake is reduced.

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PROPOSITION OF SPHERICAL AND CYLINDRICAL SAW-TOOTH WAVES OF FINITE AMPLITUDE IN RELAXING MEDIUM

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INTRODUCTION

As known, an original sinusoidal excitation of finite amplitude distorts during propagation to form a weak discontinuity and finally reach a stable saw-tooth wave due to the nonlinearity of medium. At present, the intensity of ultrasonic pulse employed in imaging technique is so high that it is readily to reach or approach the stage of saw-tooth wave. Moreover, since the propagation takes place often in media with multiple relaxation mechanisms and under the focused field, it is of significant interest to study the propagation of saw-tooth of finite amplitude under the focused field and in relaxing medium. However, to deal with the nonlinear focused field of saw-tooth wave will encounter considerable mathematical difficulties. In this work we will ignore the diffraction and only restrict to provide information concerning the converging effect of one dimensional spherical and cylindrical saw-tooth waves of finite amplitude. To our knowledge, even this aspect of study has not appeared in literature yet. The author derived a Fourier series solution for a plane saw-tooth wave of finite amplitude based on an extended Burgers' equation containing relaxation processes [1]. In the present work we will extend the approach of Fourier series solution to the cases of spherical and cylindrical waves. However it should be noticed that previous studies related to the propagation of finite amplitude waves were mostly restricted to the case of outgoing progressive waves. It is evident that for the plane wave no difference in propagation property may be observed between the outgoing and incoming waves. But, for the spherical and cylindrical waves, the outgoing and incoming progressive waves actually represent the diverging and converging waves, respectively.

EXTENDED BURGERS' EQUATION IN STRETCHED COORDINATES

In the stretched coordinates, the equation of motion and equation of continuity in dissipative fluids can be expressed [2] as:

\[ \rho \frac{\partial \mathbf{v}}{\partial t} + \rho \mathbf{v} \cdot \nabla \mathbf{v} = -\nabla p + \mu \left( \nabla \mathbf{v} + \nabla \mathbf{v}^T \right) - \frac{\partial \mathbf{a}}{\partial t} \]  

(1)

and

\[ \nabla \cdot \mathbf{v} = 0 \]  

(2)

Here \( p \) and \( \mathbf{v} \) represent the sound pressure and particle velocity, respectively. \( \rho \) is the density of medium and its variation. \( r \) stands for the radial coordinate. \( a \) may be taken as 1 for spherical wave or 2 for cylindrical wave, \( b = (4/3) \left( \gamma + (\gamma - 1) \right) \). Here \( \gamma \) represents the shear viscosity of medium, \( \tau \) is the ratio of specific heats and \( \rho_0 \) is the Prandtl constant. If the relaxation effect is taken into account, the state equation of medium can be expressed as [2]

\[ p = c_0^2 p' + (a-1) \frac{c_0^2}{\rho_0} p' \frac{\partial \rho_0}{\partial t}^2 - \frac{c_0^2}{\rho_0} \int_0^t \frac{\partial p}{\partial t} e^{(t-T)/\tau} dt \]  

(3)

in which, \( \omega = (c_0^2 + 3u_0^2/c_0^2) \) is a parameter related to the relaxation process, \( \tau \) is the corresponding relaxation time, \( c_0 \) and \( c_0 \) represent the velocities of sound of small amplitude without relaxation and with "frozen" relaxation, respectively. \( a \) is the nonlinearity parameter of medium. As it is required to deal with both the cases of outgoing and incoming progressive waves, we must transform the coordinates to a set of those moving with waves, i.e. the particle velocity should be transformed from \( v(r,t) \) to \( v(\mu, \tau) \). Here \( \mu \) stands for the acoustic Mach number and is a small quantity. \( \tau = t - r/c \) represents the retarded time, where the minus sign is taken for diverging wave and plus sign is for the converging wave, \( \tau_0 \) is a reference distance. \( c_0 \) is the speed of sound, and \( \Delta \) denotes the dispersion quantity and also is a small quantity. By carrying out the change of variable for (1), (2) and (3) and introducing further dimensionless variables \( \eta = (\omega c_0^2 / c_0^2) \tau \) and \( \eta = (\tau - r/c_0) / \tau_0 \), \( v_0 \) stands for a reference particle velocity, where \( \eta \) is an angular frequency, the extended Burgers' equation in stretched coordinates can be derived with keeping terms up to second order approximation as

\[ \frac{\partial \mathbf{v}}{\partial \eta} + \mathbf{v} \cdot \nabla \mathbf{v} = \frac{\partial \mathbf{v}}{\partial \eta} + \frac{1}{2} (\mathbf{v} \cdot \nabla \mathbf{v}) \]  

(4)

Here \( \mathbf{v} \) is the velocity vector, \( \mathbf{v} \) is the strain rate vector, \( \eta = \theta(\mathbf{v}_0) / \theta(\mathbf{v}_0) \), \( \theta(\mathbf{v}_0) \) and \( \theta(\mathbf{v}_0) \) are the strain rate and the strain rate vector, respectively.

\[ \mathbf{J} = \frac{\partial \mathbf{v}}{\partial \eta} - \frac{\partial \mathbf{v}}{\partial \eta} \]  

(5)

It should be noticed that during the derivation of Eq(4) the approximation \( k \ll 1 \) was made, which means Eq(4) valid only at distances at least several wave lengths away from the spherical center or cylindrical central axis.

FOURIERS SERIES SOLUTION

By means of the similar approach taken by Blackstock [3] for solving the lossless Burgers' equation in stretched coordinates, we may let

\[ \mathbf{v} = (\mathbf{v}_0 + \mathbf{v}_1) \]  

(5)
and transform the spatial coordinate \( z \) by letting 
\[ f_z = (\xi \eta) \xi \eta. \]
Putting those relations into (4) gives 
\[ W = W_{w}e^{+i \phi} = (C_{w} e^{i \phi}) W \]
(6)

in which \( f = e^{i \eta} = \sqrt{w \eta} \) when \( a = 1 \)
and \( f = e^{\frac{1}{2}i \zeta} \) when \( a = 1/2 \). Now we may try a
solution for Eq (6) in the following forms
\[ W = E(t) Y = E(\kappa \pm \kappa E)^{n}(\kappa \mp g) \]
(7a)

for the outgoing wave, and
\[ W = E(t) Y = E(\kappa \mp \kappa E)^{n}(\kappa + g) \]
(7b)

for the incoming wave, respectively. Here \( E = E(t) \)
is related to absorption and should be determined.
Obviously, when \( f = 0 \), it should be \( E(t) = 1 \). Thus for
the incoming wave, we have
\[ \frac{1}{E} \frac{dE}{dt} = \pm \kappa \] (8)

Now, we may expand (7b) into a Fourier series as
\[ W = \sum_{n=-\infty}^{\infty} \frac{2}{\pi} \frac{E}{\kappa \mp \kappa E} \sin(n \kappa) \]
(9)

Substituting (8) into (9) and carrying out the
integration, we may compare terms for \( \sin(n \kappa) \)
and \( \cos(n \kappa) \) on both sides of the equality. Thus if the
solution exists, the following equations must be
fulfilled, separately.
\[ \frac{dE}{dt} = \frac{D_{p} E}{\kappa \mp \kappa E} \]
(10)

and
\[ A_{n} = A_{n} + \frac{n^{2} \sigma_{n} \sigma_{n+1}^{2}}{1 + n^{2} \omega_{n}^{2} \sigma_{n}^{2}} \]
(11)

in which \( A_{n} = \frac{dE}{dt} \) and \( B_{n} = \frac{dE}{dt} \) \( A_{n} \). Solving Eq (10) as the type of Bernoulli
equation gives
\[ E = E_{0} = e^{\frac{1}{2}(s-a) \tau} \frac{1}{1 + \frac{1}{2} (s-a) \tau} \]
(12)

for \( a = 1 \) the spherical wave, and
\[ E = E_{0} = e^{\frac{1}{2}(s-a) \tau} \frac{1}{1 + \frac{1}{2} (s-a) \tau} \]
(13)

for \( a = 1/2 \) the cylindrical wave. Here \( E(t) \) and 
\( E_{0} \) represent the Exponential-Integral function
and Degenerate-Hypergeometric function, re-
spectively [4]. Analogously, the solutions for
outgoing waves can be found (Here those expressions
are omitted for the limitation of space.)

ACOUSTIC SATURATION OF CONVERGING WAVES

As well known, the propagation of saw-tooth waves of
finite amplitude will lead to the appearance of acoustic saturation after certain distance. However
it is more difficult to reach the acoustic saturation
for the spherical and cylindrical outgoing waves than for the plane wave due to their diver-
gence of waves and slow accumulation of nonlinear
effect. However the picture should be reversed in
the case of converging waves. From (5) and (3) we
can obtain the amplitude of fundamental component
of converging wave.
\[ v_{la} = v_{x} \frac{\sigma_{n} \sigma_{n+1}^{2}}{1 + \frac{1}{2} \sigma_{n} \sigma_{n+1}^{2}} \]
(14)

when \( t \geq \tau > \tau \), (14) may be reduced to
\[ v_{la} = 2 \frac{c^{2}}{2} \phi_{o} \cos(\phi_{o} / \phi_{o}) \]
(15)

for the spherical wave, and
\[ v_{la} = 2 \left( \begin{array}{c}
\frac{\tau_{a}}{2} \\
\phi_{o}
\end{array} \right) \frac{c^{2}}{2} \phi_{o} \cos(\phi_{o} / \phi_{o}) \]
(16)

for the cylindrical wave, respectively. Both above
expressions indicate that the fundamental component
has kept independent of the peak of initial saw-
tooth wave \( v_{0} \) which means the acoustic saturation
has appeared in those conditions. We can choose -
\( \tau \geq \tau \) as a condition to determine the critical
distance of acoustic saturation \( r_{c} \). For example,
suppose the absorption is ignored, or let \( E(t) \) = 1
and \( \sigma_{n} \). we can estimate \( r_{c} = 0.443 r_{o} \) for spherical
wave and \( r_{c} = 0.33 r_{o} \) for cylindrical wave.
This work was supported by NSF of China

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THE PARAMETRIC ARRAY IN WAVEGUIDE

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INTRODUCTION

Investigations on the properties of parametric array generation under waveguide conditions are important in connection with the possibility of using parametric radiators to generate a field having a prescribed mode structure in ocean waveguide. Literature [1][2][3] studied the parametric sound excited by quasi-plane waves or point sources, respectively. In the present article, we apply the Green's function method to the study of the structure of the field of parametric array generated by a piston transducer in a waveguide and carry out the relevant experiments.

THEORY

Based on Westervelt's integral[4], the parametric source density, q(x,y,z,t), due to a primary acoustic pressure field p(x,y,z,t) is:

$$G(x',y',z') = \sum P(x',y',z') \Phi(x',y',z')$$

where $P_0$ is the static density, $c_0$ is the ambient sound speed, $\Phi$ is the nonlinear coefficient. Then the secondary pressure field is given by

$$P(x,y,z) = \sum G(x',y',z') G(x',y',z')$$

where $G(x,y,z)$ is the Green's function.

The geometry of the sound field of a piston transducer in free space is, as proposed in Ref.[5], indicated in Fig.1 and its numerical expressions are:

- in the nearfield, ($r > R$, $a_0 a_0$):
  $$P_1 = P_1 [\cos 2\pi (h - 1/2)f + \cos 2\pi (h + 1/2)f]$$

- in the farfield, ($r > R$, $e^0 y^2$):
  $$P_1 = \frac{a_0}{r} [\cos 2\pi (h - 1/2)f + \cos 2\pi (h + 1/2)f] D(\theta)$$

where $a_0$: radius of transducer; $R$: the Rayleigh length of the transducer; $\theta$: the angle of half-power; $D(\theta)$: the pattern of directivity; $A_{\Phi}$ and $f_0$ are the wavelength and frequency of primary wave, respectively.

![Fig. 1. The geometry of the sound field of a piston transducer.](image)

In a waveguide, we can use mirror method to modify eqs. (3) and (4) and substitute then into the eq. (1), then the source density of parametric array can be obtained.

The waveguide studied is supposed as an ideal one with the acoustically boundaries of which the upper one is compliant ($V_{\text{up}} = 1$), the bottom is a rigid one ($V_{\text{down}} = 1$). Thus, its Green's function is

$$G(x',y',z'|x,y,z) = \sum \sin(\nu x') \sin(\nu z') H(\nu (x - x') + (y - y')$$

where $\nu$ is the difference-frequency wavenumber. $D$ is the depth of water, $\nu_{x,s}$ and $\nu_x$ are the wavenumbers of normal modes in the horizontal and vertical directions, respectively. $H(\nu)$ is the first kind of Hankel function. To apply (5) to (2), the parametric sound field is rewritten as a fivefold integral. For the convenience of the numerical computation, some simplifying assumptions is invoked and the numerical value is obtained by use of computer. Fig. 2 and 4 show the theoretical value of the sound pressure of the parametric array, where the primary frequencies are 1000 Hz and 950 Hz, water depth are H=20 m, distances are L=2 m (nearfield) and L=2000 m (farfield), respectively.

EXPERIMENT

We performed a series of experiments in a tank. The schematic diagram of the experiment is shown in Fig. 5 and the experimental results are shown in Fig. 5, where the depth of Projector is 0.11 cm and its axis is oriented at an angle 0.11 (rad). By comparing it to the theoretical calculation as shown in Fig.2 under the same conditions, good agreement can be obtained.

It is necessary to mention that the size of the receiver, of which height is 3.0 cm and radius is 2.0 cm, has some influence on the structure of the sound field, so that the sharp peak is smoothed.

CONCLUSION

We now investigate the characteristics of the parametric sound in waveguide.

In the nearfield (within the range of Rayleigh length), a number of eigenmodes were generated. Since directivity of the parametric array is very poor there, the generation takes place almost at every directions. Furthermore, the influence of the waveguide's boundary on the propagation of parametric sound is insufficient, so that the eigenmodes generated are difficult to separated.

In the farfield, the parametric array has been developed sufficiently and the relative contribution of the array to the mode whose direction is close to the transducer's axis is greater than the others. So this mode increases with distance and gradually becomes dominant, then the waveform is almost composed of one or two normal modes.

In case of non-ideal bottom boundaries, the wavenumbers of normal modes in the horizontal and vertical directions, in general, are complex, it can be treated by the perturbation method if the image part of wavenumbers are much smaller than the real part. For a pulse wave, a similar procedure can be done as well by means of Fourier transform.

REFERENCE


**FIG. 2.** The theoretical result of the parametric array in a waveguide

**FIG. 3.** The experimental data of the parametric array in a waveguide

**FIG. 4.** The theoretical results of the parametric array in a waveguide (within the farfield). The projector's axis is oriented at the angles (a) 0.03 arc; (b) 0.09 arc; (c) 0.17 arc.

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**FIG. 5.** Schematic diagram of the experiment
DISTORTION OF THE LASER-GENERATED
FINITE-AMPLITUDE PHOTOACOUSTIC PULSES
DUE TO NONLINEARITY OF LIQUIDS

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INTRODUCTION

There are two methods of measuring acoustic nonlinearity parameter of liquids by means of
finite-amplitude acoustic waves. The first is the finite-amplitude harmonic wave method. The second is the
finite-amplitude acoustic wave method. Bozkho et al used the pulsed laser to generate the plane pulsed
acoustic waves with restricted beam in liquids. By measuring the variation of rise time of wavefront during propa-
gation, they obtained the nonlinearity parameter of the liquids used. However, in this case, the diffraction of
plane wave influenced the wavefront also. Sigrist et al. improved this method. They measured the rise times of
the wavefronts of two plane pulsed waves with different peak pressure amplitudes after propagating same dis-
tance in liquids. The acoustic pulses were generated by pulsed laser via thermoelastic mechanism also. They ob-
tained the nonlinearity parameter of liquids from the difference of rise times of two photoacoustic pulses. In this
case, two photoacoustic pulses must have the same rise time of wavefronts initially in order to ensure that the dif-
ference of rise times is caused by their different peak pressures. But this is difficult in experiment usually.

We have extended the case of plane pulsed waves to that of spherical. So, the diffraction effect can be avoided
and the strict reproducibility of pulse waveform is not required.

DISTORTION OF THE FINITE-AMPLITUDE
SPHERICAL PHOTOACOUSTIC PULSES

Due to the acoustic nonlinearity, the phase velocity
C of the finite-amplitude acoustic wave depends on the
particle velocity U in liquid. The phase velocity C for
plane waves in a lossless fluid is given by

\[ C = \frac{B}{A} + \frac{U}{2}, \]  

where \( C_0 \) is the phase velocity of the small-amplitude
acoustic wave; \( B / A \) is the nonlinearity ratio. If \( \rho_0 \) is the
density of liquid in the equilibrium state and \( P \) is the
sound pressure, then \( U = P / \rho_0 C_0 \).

For the plane pulsed acoustic wave with peak pressure \( P_m \), the propagation time \( t \) of the wave peak through
distance \( r \) is given by

\[ t = \frac{r}{C_0 + \frac{B}{A} + 2) \frac{P_m r}{2 \rho_0 C_0^3} \approx t_0 - \frac{(B}{A} + 2) \frac{P_m r}{2 \rho_0 C_0^3}, \]  

where \( t_0 = r / C_0 \) is the propagation time of the
small-amplitude wave through the same distance. So, the
decrement of the rise time \( \Delta t \) of the waveform of an
acoustic pulse is given by

\[ \Delta t = t_0 - t = \frac{B}{A} + 2) \frac{P_m r}{2 \rho_0 C_0^3}. \]  

From Eq.(3), it follows:

\[ \frac{B}{A} = 2(\frac{\rho_0 C_0^3}{P_m r} - 1). \]  

A finite-amplitude spherical acoustic pulse from a
point source may be considered approximately as a plane
pulsed wave in far field, but its peak pressure \( P_m \) is
inversely proportional to the propagation distance \( r \), i.e.
\( P_m = A_0 / r \). So, the velocity of the wave peak may be ap-
proximated by the following expression:

\[ C = \frac{B}{A} + (B + 2) \frac{A_0}{2 \rho_0 C_0^3}, \]  

and the propagation time of the wave peak from \( r_0 \) to \( r \) is
given by

\[ t = \int_{r_0}^{r} \frac{dr}{C_0 + \frac{B}{A} + 2) A_0}{2 \rho_0 C_0^3} \frac{a + C_0 r}{a + C_0 r_0}, \]  

where \( a = \frac{(B}{A} + 2) A_0}{2 \rho_0 C_0^3} \), \( \Delta t = \frac{r - r_0}{C_0} \).

The decrement \( \Delta t \) of the rise time of the waveform of the
spherical pulsed wave during propagation from \( r_0 \) to \( r \) is
given by

\[ \Delta t = \frac{a}{C_0^3} \frac{a + C_0 r}{a + C_0 r_0}. \]  

\( a / C_0 \) is usually much less than \( r_0 \) and \( r \). The
Eq.(7a) may be approximated by the following expression:
CRITICAL DISTANCE FOR THE FORMATION OF DISCONTINUITY

For a plane pulsed acoustic wave, if the rise time of its initial wavefront is $t_o$, the critical distance $r_{cr}$ for the formation of discontinuity may be obtained by putting $\Delta t = t_o$ in the expression (3). So, we have

$$r_{cr} = 2r_o \frac{\rho_o C_o^2}{(B/A + 2)P_m}$$  \hspace{1cm} (9)

For a spherical pulsed wave, if the rise time of its wavefront at distance $r_o$ from the source is $t_o$, the discontinuity of wavefront occurs when $\Delta t = t_o$ in Eq.(7b). So, the following expression for critical distance can be obtained:

$$\frac{(B/A + 2)A_o}{2 \rho_o C_o^2} r_{cr}^2 = 1$$  \hspace{1cm} (10)

It is interesting to evaluate the critical distance for a finite-amplitude spherical harmonic wave $P = P_m \sin(\omega t - k_o r)$ during propagation. Assuming that the pressure amplitude and the amplitude of the particle vibration velocity are $P_m = A_o/\tau_o$ and $U_{mo} = \rho_o C_o$ at $r_o$, the parameter $a$ in Eq.(7b) may be expressed as follows:

$$a = \frac{(B/A + 1)C_o}{2} r_o M$$  \hspace{1cm} (11)

where $M$ is the acoustic Mach number. At critical distance, $\Delta t = \pi/2a$. The following expression for $r_{cr}$ can be obtained from Eq.(7b):

$$\frac{2}{\pi} \beta k \rho_o M \ln \left( \frac{r_{cr}}{r_o} \right) = 1$$  \hspace{1cm} (12)

where $\beta = \left( \frac{B}{2A} + 1 \right) K_o = \frac{\omega}{C_0}$.

DISCUSSION

It can be seen from Eq.(7b) that the rise time of wavefront of a finite-amplitude spherical pulsed acoustic wave decreases logarithmically with propagation distance due to nonlinearity of the liquid. The nonlinearity ratio $B/A$ can be obtained by measuring the variation of the rise time of wavefront of one and the same spherical pulsed acoustic wave during propagation. So, the diffraction effect can be avoided and the strict reproducibility of pulse waveform is not required.

The factor $a/C_o^2$ in Eq.(7b) is usually less than $10^{-6}$. In order to increase the measuring accuracy, it is necessary to make use of a short acoustic pulse in experiment. The short photoacoustic waves with large amplitude generated by the high-intensive laser pulses are very appropriate to such measurement. By the way, the response time of the sound detector must be short enough.

The critical distance for the formation of discontinuity in the case of spherical harmonic wave is determined usually by assuming that the amplitude of second harmonic component is one half of the amplitude of the fundamental at $r_{cr}$. The expression for $r_{cr}$ obtained on this assumption is different from Eq.(12) by a factor $2/\pi$. But the critical distance determined by Eq.(12) is more close to the actual distance for the formation of saw-tooth wave.

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ENERGY DISSIPATION OF PULSED ULTRASOUND OF
FINITE AMPLITUDE

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INTRODUCTION

When finite-amplitude sound propagates in a fluid, the waveform is distorted and many harmonics are generated, which causes more energy dissipation than small-signal sound. If the waveform of the sound is an isolated pulse, its period elongates\(^{1,2}\) and the propagation curve of the sound energy is different from that of a continuous wave. Blackstock\(^{3}\) derived asymptotic absorption coefficients of plane waves for continuous sinusoidal, N wave and a shock wave with an exponential tail, neglecting ordinary absorption. The author proposed a direct numerical method to obtain the waveform at an arbitrary distance using an analytical solution of Burgers' equation which describes the propagation of plane sound waves of finite amplitude in a dissipative fluid.\(^{4}\) Since all of the calculation is carried out in the time domain, it is convenient to treat pulsed ultrasound which contains many harmonics. This method was applied to calculate the absorption of continuous sinusoid and a single-cycle sine wave taking ordinary absorption of the fluid into account.\(^{5}\) The results suggested that the absorption at a distance from the source falls into two groups if the ordinary absorption is not important.

In this report, the definition of absorption coefficient is examined so that it is applicable to pulsed ultrasound. Using the direct numerical method to obtain propagating waveforms of various pulsed source waveforms, the behavior of their absorption coefficients is discussed and compared with other theoretical results in a lossless fluid.

ABSORPTION COEFFICIENT OF FINITE-AMPLITUDE
PLANE SOUND

In the following discussion, only plane waves are considered. In general, absorption coefficient \(\alpha\) is defined using intensity \(I\) as

\[
\alpha = -\frac{\nabla I}{2I} = \frac{1}{2I} \frac{dI}{dx}.
\]

(1)

This definition brings ordinary exponential decay of intensity for small-signal sound. The intensity of sound is the rate of sound energy that passes through a unit area normal to the propagating direction, and this is defined by

\[
I = \int_0^T \rho u dt,
\]

(2)

where \(\rho\) is sound pressure, \(u\) is particle velocity, and \(T\) is a time interval to calculate the average. For periodic waves, the interval is an integral number of the period and \(I\) is determined uniquely. For pulsed ultrasound, \(I\) is significantly affected by the length of \(T\), and it is not appropriate to use intensity to define the absorption. Since a linear impedance relation holds even for finite-amplitude sound,\(^{3,4}\) the following variable \(W\) is introduced,

\[
W = \frac{I}{T} = \frac{1}{w} \left( \int_0^T w \right) \int_0^T \left( \frac{w}{d} \right) dt,
\]

(3)

where \(w\) is sound energy density and subscript "0" denotes the value at the source. This value thus corresponds to the normalized sound energy which passes through the unit area during time interval \(T\). If \(T\) covers all of the waveform of pulsed ultrasound, this value is definite. For periodic waves, the value corresponds to the normalized sound energy per period. Since \(T\) does not depend on \(x, f\) may be replaced by \(W\) in Eq. (1) to define absorption,

\[
\alpha = -\frac{1}{2W} \frac{dW}{dx}.
\]

(4)

NUMERICAL METHOD FOR ANALYTICAL SOLUTION

A plane sound wave of finite amplitude propagating in a dissipative fluid is expressed by Burgers' equation and its analytical solution of sound pressure is given by:

\[
P(x,t) = \frac{P}{\rho} \exp \left[ \int_{-\infty}^{x} \frac{1}{2D_0} \left( \frac{(T-x)}{c_0} \right)^2 - \frac{1}{2} \left( \frac{\rho c_0^2}{D_0} \right) dt' \right] d't'.
\]

(5)

where \(P\) is sound pressure normalized by the source amplitude, and \(T\) and \(x\) are nondimensional retarded time and distance, respectively. \(D_0\) is a nondimensional parameter that indicates the relative importance of dissipation caused by ordinary absorption to nonlinearity. If we use continuous sinusoid \(p(t)=p_0 \sin \omega t\) as a reference waveform, the above-mentioned variables are:

\[
P(x,t) = \frac{P}{p_0} \exp \left[ \int_{-\infty}^{x} \frac{1}{2D_0} \left( \frac{\rho c_0^2}{D_0} \right) \frac{x}{c_0} \right] d't'.
\]

(6)

\[
D_0 = \frac{\beta_x}{\beta_x} = \frac{\rho c_0^2}{D_0}.
\]

(7)

where \(\beta\) is the nonlinear parameter of the medium, \(D_0\) is the diffusivity of sound, \(x,\) is the shock formation distance in a lossless fluid and \(T\) is the Gouff\'s number.\(^{1,2}\) Substituting source waveform \(P(T,0)\) as a boundary condition in Eq. (5) and executing the indefinite integral analytically, numerical integration and the following numerical logarithmic differentiation gives the waveform of sound pressure at an arbitrary distance directly.

Pulsed ultrasound at the source is expressed as follows:

\[
P(T,0) = \int_a^b f(T) dT
\]

(8)

If indefinite integrals of the above are expressed by \(F(T), G(T)\) and \(H(T)\), respectively, the integral constants have to be chosen so that they are continuous at the boundaries: \(F(a) = G(a)\) and \(G(b) = H(b)\).

Further, additional constants have to be added to satisfy:

\[
-\int_T^{T+T_0} \int_0^T f(T) dt = 0.
\]

(9)

otherwise, the exponential function may overflow numerically. This operation of adding constants to the argument of the exponential function is justified by the form of the solution.

CONTINUOUS SINUSOID (C S)

POSITIVE SINGLE-CYCLE SINE WAVE (P S B)

NEGATIVE SINGLE-CYCLE SINE WAVE (N S B)

POSITIVE HALF-CYCLE SINE WAVE (P H B)

NEGATIVE HALF-CYCLE SINE WAVE (N H B)

POSITIVE TRAPEZOIDAL WAVE (P T)

NEGATIVE TRAPEZOIDAL WAVE (N T)

FIG. 1. Changes in waveform of seven kinds of waves for cases in which nonlinearity is dominant (Dp=0.03) and ordinary absorption is dominant (Dp=1).
NUMERICAL EXAMPLES

Figure 1 shows the changes in waveform of continuous sine-soid (CS), single-cycle sine wave beginning with positive sound pressure (PSS), single-cycle sine wave beginning with negative sound pressure (NSS), positive half-cycle sine wave (PHS), negative half-cycle sine wave (NHS), positive triangle wave (PT), and negative triangle wave (NT) at eight locations from X=0 to 20. At left are the cases in which ordinary absorption is less important and nonlinearity is dominant. For CS and NSS, the shock waves are formed inside the waveform and their propagations are confined by the propagation of the entire wave. On the other hand, for PSS, PHS, NHS, PT and NT, the shock waves formed can propagate freely and their periods elongate as they propagate. At right are the cases in which ordinary absorption is dominant. Strong shock waves are not formed and elongation of the period is observed even in NSS waveform.

Figures 2 and 3 show the propagation curves of normalized sound energy of various source waveforms. The abscissa is a nondimensional distance X. Figure 2 is the case in which nonlinearity is important and Fig.3 is that in which dissipation is important. Dotted curves are the propagation curves of sound energy without nonlinearity in Burgers' equation. It is seen from Fig.2 that the propagation curves fall into two groups: (CS, NSS) and (PSS, PHS, NHS, PT, NT). This classification does not depend on the polarity of onset sound pressure but to the confinement of propagation of shock waves. If the ordinary absorption is dominant as in Fig.3, the shock waves formed are weak, elongation of the waveform is apparent and the classification in Fig.2 does not apply. Solid curves above the dotted curve correspond to the elongation of the period, which causes decrease in fundamental frequency of the ordinary absorption.

ASYMPTOTIC ABSORPTION AT A DISTANCE

Once propagation curves of sound energy are obtained, absorption coefficients are calculated through Eq.(4). Blackstock did discuss the asymptotic absorption of finite-amplitude plane waves at a distance from the source neglecting ordinary absorption. According to his results, asymptotic absorption of continuous sinusoid is 1/x. For N wave, there is a discrepancy in the definition of absorption with the present. Recalculation of his results according to Eq.(4) gives 1/4x for N wave. Asymptotic absorption of a shock wave with an exponential tail is 1/x. Nakamura compared energy dissipation of repeated sawtooth wave and N wave neglecting ordinary absorption. When his results are substituted into Eq.(4), they also give 1/3 for sawtooth wave and 1/4 for N wave. These suggest that, if the ordinary absorption is not important, the asymptotic absorption at a distance from the source is 1/x or 1/4x according to the difference of behavior of the formed shock waves.

Figure 4 shows the comparison of nondimensional absorption αc of CS and NSS with the curve of 1/x when Dp=0.03, and Fig.5 compares PSS, PHS, NHS, PT and NT with 1/4x. It is seen that each absorption approaches the respective asymptotic curve.

CONCLUSIONS

In order to treat the absorption coefficient pertinent to pulsed ultrasound, the definition based on the sound energy that passes through a unit area should be used rather than that based on the intensity of sound. This definition is applicable to both periodic waves and pulsed waves.

Using numerical results of the analytical solution of Burgers' equation of various source waveforms propagating in a dissipative fluid, it has been shown that the absorption of the sound can be classified into two groups depending on whether the propagation of the formed shock waves is confined or not, if the ordinary absorption is not important compared to nonlinear attenuation. At a distance from the source, the absorption of the confined shock group is 1/x, whereas that of the free shock group is 1/4x.

REFERENCES
EFFECT OF TURBULENCE ON WAVEFORM
AND RISE TIME OF SONIC BOOMS:
A MODEL EXPERIMENT

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The ideal pressure signature of a sonic boom is an N wave. Important characteristics are peak overpressure $\Delta p$, shock rise time $\tau$ (defined as the time for the head shock to rise from $0.1\Delta p$ to $0.9\Delta p$), and duration $T$. The annoyance of the sonic boom to people grows with increased peak overpressure and decreased rise time. Field measurements show that the ground signature is rarely an ideal N wave. A wide variety of distorted N waves (Fig. 1), from spiked to rounded and even double-peaked, is seen. Moreover, observed rise times are two to five times larger than those calculated using molecular relaxation theory. Although it has long been accepted that waveform variability is due to atmospheric turbulence, the role of turbulence in increasing rise time is controversial. We report here the results of a model experiment to study the effect of turbulence on waveform and rise time of N waves.

EXPERIMENT

In our model experiment the N waves are produced by electrical sparks and the turbulence by a plane jet of air across which the N wave propagates. Typical values of $T$ and $\tau$ for spark-generated N waves are 10-30 $\mu$s and 0.5-2 $\mu$s, respectively, while the corresponding values for sonic booms are 100-300 ms and 1-10 ms. The scale-down factor for the model experiment is therefore of order 5,000 to 10,000. The same scale factor must relate the turbulence in the model experiment to that of the atmosphere. Since the atmospheric boundary layer has a thickness of about 1 km, an outer length scale of 100-200 m, and an inner length scale of order 10 mm, the model turbulent jet should have a thickness of 100-200 mm, an outer length scale of 10-20 mm, and an inner length scale of 0.001 mm. To approximate these properties, we used a plane jet that issued from a slit (length 250 mm, width variable but set at 27 mm for the measurements reported below) in the side wall of a settling chamber. The velocity at the slit was approximately 30 m/s (Reynolds number about 50,000). The jet spread so that 400 mm downstream its width was about 200 mm (the scaled thickness of the atmospheric boundary layer). The model N wave propagated across the jet at this point. Both spark and microphone were located in quiet air outside the jet (the total propagation path was about 400 mm). Although the sonic boom is a cylindrical wave whereas our spark-produced N waves are spherical, the difference was not important for our investigation. We also used a parabolic reflector to produce a planar N wave when desired, but all measurements reported here are for spherical waves.

Our receiver was a condenser microphone of very wide bandwidth, capable of measuring rise time as short as 0.5 $\mu$s. The design, construction, and calibration of this type of microphone is described by Wright.

RESULTS

The procedure was to make 100 measurements with no turbulence and then 100 measurements with the plane turbulent jet turned on. The rise time and peak overpressure of each individual waveform were computed and also average rise time $\tau_{ave}$, average peak overpressure $\Delta p_{ave}$, and standard deviation $\sigma$ for the 100 measurements. Figure 2 shows some typical waveforms for the spherical N wave after passage through the turbulent field. A wide variety of spiked and rounded N waves is shown. Clearly, the model experiment reproduces the waveform distortion effects seen in full scale sonic boom measurements. Figure 3 compares the data without turbulence to that with turbulence. Although variability of $\Delta p$ is increased enormously by the turbulence ($\sigma_{\Delta p}$ rises from 7 Pa to 61 Pa), $\Delta p_{ave}$ is not affected very much (190 Pa with turbulence, 166 Pa without). In the case of rise time, both average and variability are significantly increased. In particular, $\tau_{ave}$ changes from 0.82 $\mu$s (no turbulence) to 1.72 $\mu$s (turbulence present), a more than two-fold increase. The corresponding increase in $\sigma_{\tau}$ is from 0.05 $\mu$s to 0.96 $\mu$s. Even more important is that turbulence seems only to increase the rise time, never to decrease it. This can be seen from Fig. 4, which shows cumulative probability curves for both rise time and peak overpressure. Each curve compares the rise time and peak overpressure data with the average for no-turbulence N waves. It is seen that none of the turbulence rise time data is smaller than the no-turbulence value of $\tau_{ave}$. Our measurements tentatively confirm Pierce's prediction that turbulence only increases rise time.
DISCUSSION AND CONCLUSIONS

A model experiment has been done to study the effect of turbulence on the waveform and especially the rise time of spark-produced N waves. The distortion of the spark-produced N waves by the model turbulence (plane jet) is the same as the distortion of sonic booms by the atmosphere. An important result is that the turbulence always seems to be a shock thickening mechanism, as was proposed by Pierce. The rise time of each N wave, after passage through the turbulence, is increased when it is compared with the no-turbulence rise time. Later measurements with more intense turbulence have, however, produced a few cases in which turbulence causes a decrease in rise time. On a statistical average, the rise time is doubled after passage through the turbulence. The average peak overpressure is about the same as the no-turbulence average peak overpressure. Turbulence greatly increases the variability in rise time and peak overpressure. Future experimental work will include characterizing the strength and length scales of the turbulent field by hot wire anemometry and measuring the effect of acoustic nonlinearities (e.g., healing effect) on N wave propagation. [This work was sponsored by National Aeronautics and Space Administration.]

REFERENCES


ACOUSTIC SATURATION IN FLUIDS AT BIOMEDICAL FREQUENCIES AND INTENSITIES

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INTRODUCTION

The wide applications of ultrasound in medical diagnosis, therapy and surgery are always relied on the assumptions of linear acoustical theory. The hallmarks of this theory are that the received signal at a fixed point in direct proportional to the source amplitude and that the frequency of the source is the only frequency propagated in the medium. However, in recent ten years, many papers reported the generation of second Harmonics during propagation of ultrasonic wave in biological medium at biomedical frequencies and intensities. Also, the values of nonlinearity parameter of tissues and its dependence [1-4] in this paper another important nonlinear effect--acoustic saturation will be studied. The sound saturation has been studied in air and in the underwater sound frequency range [5,6]. Here, some experimental results of acoustic saturation at biomedical frequencies and intensities in the biological fluids, such as distilled water, ethylene glycol, aqueous solution of dextrose and human blood, etc., will be reported.

THEORETICAL ANALYZE

It is well known that according to the assumption of infinitesimal acoustics the amplitude of received signal at the fixed point in the medium increases linearly with the output of the source. However, in the case of high transmitted sound level the input-output intensity curve will gradually depart from the linear line and approach to a flattening. This phenomenon of sound saturation is one of the nonlinear effects of finite amplitude sound wave propagating in a medium.

Non-dimensional Burgers' equation used to describe the propagation of finite amplitude wave in fluids can be expressed as follows:

\[
\frac{\partial \mathbf{u}}{\partial t} + \frac{1}{\gamma} \frac{\partial \left( \mathbf{u}^2 \right)}{\partial x} = -\frac{\partial \mathbf{p}}{\partial x} - \frac{\epsilon}{\rho} \frac{\partial \mathbf{u}}{\partial x}
\]  

(1)

where \( \mathbf{u} \) is the particle velocity and its amplitude, \( \mathbf{u}_0 \) - partial velocity at the source, \( \lambda/\gamma \) is the distance from the source, \( \gamma=1+2A/\Delta A \) is the nonlinearity parameter, \( \Delta A / A \) is the acoustic wave number \( k \) is wave number, \( \mathbf{p} = \rho \mathbf{u} - \rho \mathbf{e} \) is attenuation coefficient of small amplitude wave, \( \mathbf{p} = \rho (1 - \omega^2 / \omega_0^2) \). Burgers' equation (1) has a known exact and quite complete solution but its form has thus far prevented any simple expression of it. So, usually for a concrete problem the procedure is to solve the nonlinear Burgers' equation (1) with the correct solution by sound attenuation. If boundary condition is \( \mathbf{u} = \text{sin}(\alpha \omega t) \) for \( x = 0 \), then the asymptotic form of solution for \( x = 0 \) and \( \omega = 1 \) is:

\[
\frac{u}{u_0} = \frac{2}{1 + \sum_{n=1}^{\infty} \frac{\sin \left( \frac{\omega t - kx}{\omega_0} \right)}{n}}
\]  

(2)

from (2) the pressure amplitude of fundamental frequency component (n = 1) can be obtained as follows:

\[
P_1 = \frac{2\rho c^2}{k}
\]  

(3)

Then, the sound saturation pressure amplitude is:

\[
P_{\text{sat}} = \frac{2\rho c^2}{k}
\]  

(4)

This means that saturation amplitude is in inverse proportion to the distance \( x \) and frequency \( \omega \), and is independent on the source intensity \( p_0 \). Considering the correct of attenuation of medium the pressure amplitude of fundamental frequency component can be written as:

\[
\frac{dP_1}{dx} = \frac{-0.2k}{2p_0 c^2} \sin \left( \frac{\omega t - kx}{\omega_0} \right)
\]  

(5)

The solution of Eq.(5) is given by:

\[
P_1 = \frac{2 p_0 c^2 e^{-\alpha x}}{1 + (1 + \alpha^2 x^2) / \omega_0^2}
\]  

(6)

then the saturation pressure can be expressed as:

\[
P_{\text{sat}} = \frac{2 p_0 c^2 e^{-\alpha x}}{\beta k(1 + \alpha^2 x^2)}
\]  

(7)

The formula (6) and (7) are the general expressions of sound saturation for finite amplitude wave of fundamental frequency component propagating in the lossy media.

EXPERIMENTS AND RESULTS

1. The experimental arrangement is similar to that described in ref[3]. The crystal source and receiver have the same resonance frequency of 4 MHz. The received signal can be measured by frequency spectroscope and by microcomputer. An oscilloscope is used to monitor the transmitted signal. The sound power of the source is measured by radiation pressure method. The sensitivity of the receiver is tested by self-reciprocity calibration.

2. The dependence of the received signal of 4 MHz at a fixed distance on the source intensity and the dependence of the acoustic saturation level on distance are measured by using this set-up. The experiments are carried out in the distilled water, ethylene glycol, aqueous solution of dextrose and human blood.

3. The wave-form distortions of finite amplitude sound at different position and at different source intensity are measured by using a wide-band PEF transducer. The process of harmonics generation will be analyzed simultaneously.

4. The experimental results of source-received intensity response curves at x = 5 cm, 10 cm, 15 cm in water are shown in Fig.1. Theory (solid line) and
experiments coincide quite well. Figure 2 shows the results for aqueous solution of ethylene glycol with concentration 30% and 100% at x=15cm. The curves for 40% aqueous solution of dextrose and for human blood at x=5cm are shown in Fig3. Using a wide-band transducer a set of pictures of distorted wave forms and harmonic components measured in the water at x=0 and x=10cm are presented in Fig 4(a) and 4(b) respectively.

DISCUSSION

1. As a result of the nonlinearity, the sine wave emitted by the source begins to distort, and a cumulative distortion in the propagating waveform consists of many harmonics. This means that the transmitted frequency is not the only frequency produced in the medium. In this case the nonlinear pumping of energy from the fundamental into the harmonics is to cause an extra-attenuation of sound in comparison with the linear theory. Therefore, the propagation distance of fundamental component may also be affected by the generation of harmonics.

2. The sound transmitting power in the fluids is limited by nonlinear effect of saturation. As the source intensity increases to a high level, amplitude of the received signal is limited and saturated. The existence of acoustic saturation is to limit the sound power deliverable to the medium at a given range and frequency.

3. Ultrasound is an useful tool in medical diagnosis and therapy. However, the nonlinear effect of finite amplitude ultrasound can accelerated losses and will set upper limit on the received sound power. Therefore, further study these problems is of great value to biomedicine.

REFERENCES

TRUSTY AND EXPERIMENTAL INVESTIGATION ON
FINITE-AMPLITUDE STANDING WAVE

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INTRODUCTION

In recent years, the researches on nonlinear acoustics of high-intensity sound have been developed rapidly because it is more and more important in contemporary technology and science. Aerodynamic noise emitted by rocket or jet engines, noise test of airframes and processing, and others, all involve finite-amplitude sound waves and, mostly standing waves. As for progressive, the theory of finite-amplitude waves has been well established, various properties and effects were investigated. But investigation on finite-amplitude standing wave was dissatisfactory. Although a number of theoretical and experimental papers[2-4] exist on investigations of resonant oscillations in closed tubes recently, various theories were proposed, which led to conflict results, none of them seems to be without question. One of the authors paid some attention to this situation for a number of years, a rigorous theory of one dimensional finite-amplitude standing waves has been proposed on the basis of the fundamental principles of hydrodynamics in a previous paper[1], and formulas of steady wave forms are derived. Experimental investigations on one dimension finite-amplitude standing wave have been carried out and interesting phenomena observed. It is the purpose of this paper to discuss these results.

A BRIEF REVIEW OF THE THEORY

The equations of Riemann simple waves can be written as

\[ r_s = (\gamma + 1) r_p = 0 \]  

\[ s_x = u \cdot (u - c_s) s_y = 0 \]  

Here, 2r-s and 2s-r, called Riemann invariants, u, p, \( p_c \) are the particle velocity, density and pressure, resp., \( \gamma = \frac{\partial p}{\partial \rho} / \frac{\partial p}{\partial \rho_c} \), \( \gamma \) the ratio of specific heats, \( \rho_c, \) the small signal sound velocity, c the velocity of the nonlinear sound wave.

In finite-amplitude standing wave, the principle of superposition in linear system no longer holds in nonlinear acoustics, and the parameters u, p, \( p_c \), each containing contributions from both incoming and outgoing waves, are no more directly related as in progressive wave. The r and s waves propagate in opposite directions. It should be noted that the quantity u is common to both equations, and it is the particle velocity of the resultant standing wave. As such u is nowhere a constant on the waveform, as in the case of progressive wave, and it produces a path difference equal to its time integral \( \int u dt \), instead of its time product. ut. The solution of Eqs. (1) and (2) are then

\[ r = \frac{1}{2} r' \cos(x - \left(c_s t - \frac{1}{2} c_s^2 t^2 \right)) \]  

The particle displacement may be found from the relation \( \varepsilon \), as

\[ \varepsilon = A_0 \cos \omega t \sin x \cdot \frac{1}{2} c_s t \]  

where \( A_0, \omega = \omega_0 \), being the first order solution. As \( A_0, \omega_0 \), the sound pressure is deduced from the relation \( \frac{d \varepsilon}{d p_c} / \rho_c \), which, after some manipulation, results

\[ p = p_0 \cos \omega_0 t \cos(x - \frac{1}{2} c_s t - \frac{1}{2} c_s^2 t^2) \]  

where \( p_0, \omega_0, c_s, u_0 \), with the value of given by Eq(5), this formula is correct to the third order. The theory can be applied to a one-dimensional model for the acoustic conditions inside a piston-driven, gas-filled pipe of length L and boundary conditions at the end and using \( \rho_0 \) at \( c_s \). The sound pressure can be expressed as:

\[ p = \frac{p_0}{\gamma} \cos \omega_0 t \frac{c_s}{16} \frac{3}{16} \frac{1}{\gamma - 1} \frac{\rho_0}{\gamma \rho_0 \sin^2 \alpha} \]  

The amplitudes are limited by the absorption at the closed end of the pipe, even the other dissipation is neglected. So Eq(7) is the theory basis of experimental investigation on one-dimensional finite-amplitude standing wave.

EXPERIMENTAL INVESTIGATION

Experimental Apparatus

The experimental scheme (Fig.1) consists of three relatively distinct parts: the signal- or driving system, the tube or waveguide system, the sound pressure measuring and data processing system. The electrical signal generated with Heterodyne Analyzer (BK 2010) is amplified by the power amplifier (Model SA-600) and used to drive a low-frequency high-power loudspeaker (16W, 100W) designed by Institute of Acoustics. The standing wave tube is a 255 cm. copper pipe with outside diameter of approximately 5.7 cm. and inside diameter of 4.5 cm. The pipe wall is thick enough to prevent the effect of its vibration. One end of the pipe is sealed with a flat stainless steel plug. The other end is joined to the loudspeaker through a cone reducer. Two microphones (BR 4136) are mounted at the end of pipe and the throat of the loudspeaker, resp., to pick up the acoustic signals. The driving current signal is measured by the voltage drop across a measuring resistance (120, 20 W). Both acoustic and current signals are fed into Dual Channel Analyzer(BK 2034), simultaneously, the driving voltage signal is introduced to High Resolution Analyzer(BK 2033), the acoustic signal from the throat of the loudspeaker is amplified with BK 2010 and then fed back to the Dual-Trace Oscilloscope(BR-6). The data displayed in analyzers are stored in the memory of computer. The distortion of the power amplifier (SA-600) is less than 0.5% at its rated power(300 W). The nonuniformity of frequency response of microphones is
less than 1 dB in the range from 20 Hz to 20 kHz. Both acoustic and driving frequency signals are analyzed by Dual Channel Analyzer (BK 2034) simultaneously as to discriminate the nonlinear distortion of driving system from that of standing wave field.

Results
The experiments were performed with sinusoidal wave in the range of 70 - 1000 Hz. The Fig 2 shows the magnitude of sound harmonic as a function of driving frequency at a constant fundamental frequency. Sound pressure level (SPL1) 146 dB. The curve of sound pressure level (SPL2) can be approximately described by a function Acos(At), A is a constant, depending the value of SPL1. According to the theoretical analysis and experimental results, the standing wave field in the closed tube can be divided into four regions in frequency domain: i.e., Steady Region(1), Transition Region(2), Resonance Region(3) and Antiresonance Region(4), and shown in Fig. 2. In this, the values of $\sin 2\theta$ is about unity, while region (3) and (4) are the case $\sin 2\theta$ approaching zero, and the values of $\sin \theta$ are near to zero and unity, resp. In region (2), $\sin \theta$ takes some values between zero and unity. The properties of standing wave in four regions are different and should be described with different formulas. In the region (1), called "steady region", the previous theories predict the magnitude of sound harmonic satisfactorily. Fig. 3 shows the sound pressure level at the end of pipe, i.e. $x=0$, the abscissas represents the fundamental frequency sound pressure level (SPL1) and ordinate indicates differences between SPL1 and SPL2. The fundamental frequency is 136 Hz. The results are the same for the other "steady region". The average differences between the theory and results are nearly within 1.5 dB. It may be said that the experiments agree well with the theory in the "steady regions", considering the effect of viscosity, experimental error and so on. In the other three regions, the properties of standing wave were also investigated by experiment in detail, some of results are as follows. In region (3), the nonlinear distortion is very large, and SPL1 may be as high as 185 dB, also shock waves can be easily observed in this region. The values of higher order harmonics are always smaller than lower order harmonics. When SPL1 is below about 155 dB, the changing of harmonics is linear with SPL1, i.e. SPL1 changes 1 dB, the values of 2nd, 3rd, 4th harmonics vary 2, 3, 4 dB, resp., while SPL1 gets higher, the harmonic variations are nonlinear with SPL1, e.g. SPL1 increase form 167 to 185 dB, the 2nd, 3rd harmonics only increase 1.1 and 1.2 dB, resp., instead of 2 and 3 dB. This is saturation phenomena of finite-amplitude standing wave. Fig. 4 shows results with the fundamental 294 Hz. Saturation phenomena of all the harmonics occurs at around SPL1 165 dB, this value of SPL1 considered as "Transition Point" between preshock and shock waves. In the region (4), where the fundamental is near antiresonance of the tube, SPL1 obtained is only around output of the loudspeaker, the relative value of SPL2 is about same in the region (3), but those of the harmonics higher than second are irregular, saturation and preshock formation are barely observed. Situations in region (2) are rather complicated, some of properties are similar to those in region (1), some in region (3) and (4). The detailed results and theories in these regions will be presented elsewhere.

Conclusion
The one-dimensional finite-amplitude standing wave field, on the basis of theoretical analysis and experimental investigations, can be divided into four regions in frequency domain, i.e., steady region, transition region, resonance region, and antiresonance region. It is shown that the experimental results agree well with the previous theory in the "steady region". In the other three regions, interesting phenomena were observed. Further work on experimental and theoretical investigations is in progress.

REFERENCES

Fig. 1. Diagram of apparatus and instrumentation

Fig. 2. Division of Standing Wave Field

Fig. 3. Experimental results in Steady Region

Fig. 4. Saturation of Standing Wave

LEAST-SQUARES ESTIMATES OF THE NONLINEAR CONSTANTS OF PIEZOELECTRIC CRYSTALS

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INTRODUCTION

The third-order nonlinearities in piezoelectric crystals are studied by means of various experiments. Past experience indicates that the theoretical relationship between m independent observations of an experimental quantity L and the third-order nonlinear constants in piezoelectric crystals can be recorded using a system of linear equations

\[ Y_i = \sum_{j=1}^{n} M_{ij} x_j, \quad \text{where } i = 1, 2, ..., m, \]

and where further

\[ Y_i = L_i - E_i. \]

In Eqs. (1) and (2) L_i represent individual observations of the experimental quantity L; x_j, j = 1, 2, ..., n, are the sought nonlinear material constants. The coefficients M_{ij} as well as the absolute terms E_i are functions of the linear constants of the piezoelectric crystal. They also depend on the geometry and other characteristic features of the particular crystal specimen as well as the experiment which produces the observations L_i. Their numerical values are fully calculable and known when system (1) is formulated.

The coefficients M_{ij}, forming matrix M of system (1) consist of columns referred to as vectors M_j, j = 1, 2, ..., n, each with m components M_{ij}, i = 1, 2, ..., m. Each nonlinear constant x_i in system (1) is thus associated with a column vector M_i.

To minimize the effect of random experimental errors in the observations L_i, the number m is made as large as possible, system (1) becomes overdetermined and the unknown nonlinear constants x_i are then sought in terms of their best estimates using the least-squares fit.

TRADITIONAL SOLUTION METHOD

Traditionally, all values of the nonlinear constants in system (1) which are already known from earlier work are substituted into the system and those nonlinear constants believed to be making a relatively small contribution are eliminated from it. This is done prior to the execution of the least-squares fit which is then made for the nonlinear constants remaining in the system and which concludes the process.

When the intended substitutions and eliminations are completed, the number of columns of system matrix M is reduced. The new system matrix is called N. It is assumed that matrix N always has a rank equal to the number of its columns. This is necessary and sufficient for the least-squares solutions that are to be discussed to actually exist. If the rank of matrix N is lower, its linearly dependent columns must be removed and the unknown nonlinear constants rearranged accordingly.

The substitutions and eliminations combined with the mechanical interpretation of the least-squares fit in terms of the nonlinear constants remaining in the system are a source of undesired effects and errors. Grouped and summarized in Table I, they are:

(1) distortion of the estimates,
(2) loss of potentially valuable estimates,
(3) change in the meaning of the estimates.

The impact on the standard errors is:
(4) misinterpreted meaning,
(5) distorted values.

Table I. Undesired effect and errors.

<table>
<thead>
<tr>
<th>least-squares estimates</th>
<th>substitution</th>
<th>elimination</th>
</tr>
</thead>
<tbody>
<tr>
<td>Case 1</td>
<td>(1)</td>
<td>(3)</td>
</tr>
<tr>
<td>Case 2</td>
<td>(2)</td>
<td>(2)</td>
</tr>
<tr>
<td>Case 3</td>
<td>(1) + (2)</td>
<td>(2) + (3)</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>standard errors</th>
<th>substitution</th>
<th>elimination</th>
</tr>
</thead>
<tbody>
<tr>
<td>Case 1</td>
<td>(4)</td>
<td>(4)</td>
</tr>
<tr>
<td>Case 2</td>
<td>(5)</td>
<td>(5)</td>
</tr>
<tr>
<td>Case 3</td>
<td>(4) + (5)</td>
<td>(4) + (5)</td>
</tr>
</tbody>
</table>

(1), (2), ..., are defined in the above text.

For any constant x_i in system (1) its corresponding vector M_i may be viewed as a sum of two perpendicular vectors

\[ M_i = M_i^v + M_i^w, \]

where M_i^w is a vector which can be expressed as a linear combination of the vectors forming matrix N while M_i^v is perpendicular to the space defined by the vectors of matrix N. The group of undesired effects and errors (Case 1, 2, 3, Table I) which occur if constant x_i is substituted for in (1) or eliminated from (1) depends on which of the two vectors, M_i^v and M_i^w, is different from zero: M_i^v (Case 1), M_i^w (Case 2), or both (Case 3).

Two 'real-life' examples of the undesired effects and errors are given in Table II. It lists the estimates and standard errors of the nonlinear constants of quartz determined using various field interactions with quartz and observed by means of the resonator method and the transit-time method. They illustrate Case 1 and 3, respectively. Most recent published results of this type can be found in [1] (long since corrected in [2]) and in [3]. Similar examples can be found in a number of papers published earlier.

There is nothing obvious in the results.
in Table II that would signal that they are incorrectly interpreted. To see this they must be compared with results which are formally correct.

Table II. Estimates and standard errors of the nonlinear constants of quartz afflicted by the undesired effects and errors of the traditional solution method.

<table>
<thead>
<tr>
<th>Resonator method</th>
<th>( f_{11} )</th>
<th>2.18 ± 0.05</th>
<th>( 1_{11} )</th>
<th>-3.28 ± 0.91</th>
</tr>
</thead>
<tbody>
<tr>
<td>( f_{113} )</td>
<td>-0.42 ± 0.07</td>
<td>( 1_{13} )</td>
<td>140 ± 118</td>
<td></td>
</tr>
<tr>
<td>( f_{114} )</td>
<td>0.19 ± 0.04</td>
<td>( 1_{13} + 21_{44} )</td>
<td>-8.98 ± 2.59</td>
<td></td>
</tr>
<tr>
<td>( f_{122} )</td>
<td>-1.14 ± 0.03</td>
<td>( 1_{14} )</td>
<td>-2.38 ± 0.54</td>
<td></td>
</tr>
<tr>
<td>( f_{124} )</td>
<td>0.76 ± 0.02</td>
<td>( 1_{14} + 21_{44} )</td>
<td>-12.40 ± 2.99</td>
<td></td>
</tr>
<tr>
<td>( f_{134} )</td>
<td>1.64 ± 0.03</td>
<td>( 1_{33} )</td>
<td>-8.39 ± 6.79</td>
<td></td>
</tr>
<tr>
<td>( f_{144} )</td>
<td>0.02 ± 0.03</td>
<td>( 1_{44} )</td>
<td>-4.41 ± 0.67</td>
<td></td>
</tr>
<tr>
<td>( f_{315} )</td>
<td>-0.79 ± 0.03</td>
<td>---</td>
<td>---</td>
<td>---</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Transit-time method</th>
<th>( f_{11} )</th>
<th>2.16 ± 0.06</th>
<th>( 1_{1} )</th>
<th>-4.44 ± 0.89</th>
</tr>
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<tbody>
<tr>
<td>( f_{113} )</td>
<td>-0.50 ± 0.06</td>
<td>( 1_{12} )</td>
<td>1.09 ± 0.88</td>
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</tr>
<tr>
<td>( f_{114} )</td>
<td>0.25 ± 0.05</td>
<td>( 1_{13} )</td>
<td>8.67 ± 2.50</td>
<td></td>
</tr>
<tr>
<td>( f_{122} )</td>
<td>-1.12 ± 0.05</td>
<td>( 1_{14} )</td>
<td>-2.38 ± 1.25</td>
<td></td>
</tr>
<tr>
<td>( f_{124} )</td>
<td>0.77 ± 0.02</td>
<td>( 1_{31} )</td>
<td>1.36 ± 2.12</td>
<td></td>
</tr>
<tr>
<td>( f_{134} )</td>
<td>1.64 ± 0.03</td>
<td>( 1_{33} )</td>
<td>-4.45 ± 4.69</td>
<td></td>
</tr>
<tr>
<td>( f_{144} )</td>
<td>0.08 ± 0.02</td>
<td>( 1_{41} )</td>
<td>-3.62 ± 1.85</td>
<td></td>
</tr>
<tr>
<td>( f_{315} )</td>
<td>-0.90 ± 0.03</td>
<td>( 1_{44} )</td>
<td>1.05 ± 1.49</td>
<td></td>
</tr>
</tbody>
</table>

\( f_{11} \) are electroelastic constants in N/(V.m), \( 1_{ij} \) are electrostrictive constants (dimensionless). The errors are standard errors. Given for room temperature, right-hand quartz and frame of reference according to IEEE Standard of 1978 [4].

FORMALLY CORRECT ESTIMATES AND STANDARD ERRORS

The undesired effects listed in Table I can be avoided or brought under control if the least-squares fit is executed for the linear system (1) in its original form i.e. prior to any intended substitutions or eliminations. The substitutions or eliminations - if still desired in the light of their consequences - can be made subsequently, with the same effectiveness and with the advantage of understanding their real impact.

Table III presents correctly interpreted results obtained by this calculation method. A comparison made with Table II discloses that a number of estimates have changed. One estimate has been added, another changed its meaning. The standard errors missing from Table III are those that are not properly calculable. Their counterparts in Table II are formally incorrect.

CONCLUSION

If nonlinear constants are substituted into or eliminated from linear systems solved by linear regression, then a mechanical interpretation of the results leads to false conclusions. The differences between the traditional (Table II) and the formally correct (Table III) interpretation are significant.

TABLE III. Formally correctly interpreted estimates of the nonlinear constants of quartz.

<table>
<thead>
<tr>
<th>Resonator method</th>
<th>( f_{11} )</th>
<th>2.18</th>
<th>( 1_{11} )</th>
<th>-3.28</th>
</tr>
</thead>
<tbody>
<tr>
<td>( f_{113} )</td>
<td>-0.42</td>
<td>( 1_{12} ) -4.889</td>
<td>( 1_{22} ) (-10^{22} )</td>
<td>140</td>
</tr>
<tr>
<td>( f_{114} )</td>
<td>0.19</td>
<td>( 1_{13} + 21_{44} )</td>
<td>-8.98</td>
<td>-2.38</td>
</tr>
<tr>
<td>( f_{122} )</td>
<td>-1.14</td>
<td>( 1_{14} )</td>
<td>-14</td>
<td>21 (-44 )</td>
</tr>
<tr>
<td>( f_{124} )</td>
<td>0.76</td>
<td>( 1_{31} + 21_{44} )</td>
<td>-12.40 ± 2.99</td>
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<tr>
<td>( f_{134} )</td>
<td>1.64</td>
<td>( 1_{33} )</td>
<td>-8.39 ± 6.79</td>
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<tr>
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<td>( 1_{41} )</td>
<td>-4.41 ± 0.67</td>
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<tr>
<td>( f_{315} )</td>
<td>-0.79 ± 0.03</td>
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<table>
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<tr>
<th>Transit-time method</th>
<th>( f_{11} )</th>
<th>2.14</th>
<th>( 1_{1} )</th>
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<tr>
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<td>( 1_{12} )</td>
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<td>( 1_{13} )</td>
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<td>( 1_{31} )</td>
<td>2.32 ± 2.05</td>
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<td>1.63</td>
<td>( 1_{33} )</td>
<td>-3.34 ± 4.39</td>
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<td>( f_{315} )</td>
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<td>( 1_{44} )</td>
<td>-0.74 ± 1.79</td>
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<td>( f_{11} ) (-10^{22} )</td>
<td>3.71</td>
<td></td>
<td></td>
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</table>

\( f_{11} \) is third-order permittivity in F/V. Other remarks are similar to those made for Table II.

nontrivial. This work further adds to the view that the use of the least-squares method must be complemented by a strategy which includes avoidance of trouble.

The nonlinear constants calculated in this paper are based on the same experimental data and nonlinear theory as used in [2,3]; other material constants of quartz needed here have been taken from [5,6,7].

ACKNOWLEDGEMENTS

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REFERENCES

SOLITARY WAVE IN COUPLING PROBLEM OF FLUID WITH SOLID

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1. INTRODUCTION

In recent years, nonlinear wave problems are taken seriously, particularly the investigation on solitary waves has got great development. As to the problem of fluid coupling with solid, the nonlinear wave problem has been studied and is widely present in many fields. For example, pipes transporting oil and tiny blood vessels in human body are all about the problems of fluid coupling with elastic wall. On the solitary waves in these problems, there have been a lot of studies so far. Johnson studied the wave propagation in an elastic tube full of liquid, found that the radial perturbation refers to the Korteweg de Vries equation with damping term and investigated the static state solutions. Lamb and Wang have discussed the propagating waves of tube’s cross section and the wall’s stress, the controlling equation changes to the KdV equation under the long-wave approximation. In Ref. 3, the nonlinear stress wave in an axial symmetric tube was studied by the singular perturbation method. In Ref. 4, another nonlinear wave equation satisfied by radial displacement of the tube wall has been obtained under the long and middle wave approximation assuming that the tube’s length is infinite, the tube wall is thin, the fluid flowing through it is incompressible and its viscosity can be neglected. Here we start with the dynamic basic equation, boundary conditions, and the symmetric condition satisfied by rotational motion of fluid, derive the nonlinear Schrödinger equation governing the tube wall’s radial displacement by the perturbation method of multiple scales, and obtain a modulated nonlinear solitary wave.

2. MODEL AND BASIC EQUATIONS

We discuss an infinite long elastic round pipe with the axial symmetry, its radius is R, and is full of fluid. Supposing that the pipe wall is thin, and its inertial force can be neglected, the fluid flowing in it is incompressible and its viscosity can be neglected, then the velocity potential of rotational motion of fluid satisfies the following equation and boundary conditions, symmetric condition:

\[ \nabla^2 \phi = 0 \quad (0 \leq r \leq \eta + R) \]

\[ \eta \phi + \frac{\eta}{2} \phi_r - \phi_r = 0 \quad r = \eta + R \]

\[ -\rho \phi - \frac{1}{2} \left( \frac{\eta}{r} \right)^2 \nabla \phi^3 = P \quad r = \eta + R \]

\[ P = E \eta \]

\[ r = \eta + R \]

\[ \phi = 0 \quad r = 0 \]

where \( \eta(x, r) \) denotes the radial displacement of the fluid, that is, the radial displacement of the pipe wall; \( x \) and \( r \) are co-ordinates along the length and radial directions of the pipe, respectively; \( P \) is the pressure of fluid forcing on the wall; \( E \) is the elastic coefficient of the pipe wall; \( \rho \) is the fluid density; \( \phi(x, r) \) is the velocity potential; \( \nabla \) is the gradient operator; The lower-right indexes express the derivatives with respect to relevant variables.

This is a complex nonlinear problem, we can study it by the perturbation method of multiple scales. First of all, we expand Eqs. (2), (3) and (4) into Taylor series around \( r = R \) up to the third order

\[ \eta = \eta_0 + \eta_\theta \psi_0 + \eta_\phi \psi_\phi + \frac{1}{2} \psi_0 \psi_\phi - \frac{1}{2} \psi_0 \psi_\phi = 0 \quad r = R \]

\[ \rho \phi + \rho \phi_\theta \psi_0 + \psi_\phi \psi_\phi + \frac{1}{2} \psi_0 \psi_\phi = 0 \quad r = \eta + R \]

We introduce a small parameter \( \epsilon \) and the multiple variables and asymptotic expansions as follows:

\[ t_0 = t \quad t_1 = \epsilon t \quad t_2 = \epsilon^2 t \]

\[ x_0 = x \quad x_1 = \epsilon x \quad x_2 = \epsilon^2 x \]

\[ \phi = \epsilon \phi_0 + \epsilon^2 \phi_1 + \epsilon^3 \phi_2 + \cdots \]

\[ \eta = \epsilon \eta_0 + \epsilon^2 \eta_1 + \epsilon^3 \eta_2 + \cdots \]

where \( \epsilon \) is a small parameter expressing the order of the relative size of the wave amplitude. Then from Eq. (8) we have

\[ \frac{\partial \phi_0}{\partial t} + \frac{\partial \phi_0}{\partial x} + \epsilon \frac{\partial \phi_1}{\partial t} + \epsilon \frac{\partial \phi_1}{\partial x} + \cdots \]

\[ \frac{\partial \phi_0}{\partial x} + \frac{\partial \phi_1}{\partial x} + \epsilon \frac{\partial \phi_2}{\partial x} + \cdots \]

\[ \frac{\partial \phi_0}{\partial \eta} + \frac{\partial \phi_1}{\partial \eta} + \epsilon \frac{\partial \phi_2}{\partial \eta} + \cdots \]

Substituting expressions (8)-(10) into (1) and (5)-(7) and equating the coefficients of corresponding power of \( \epsilon \), then we can obtain a series of approximate equations

\[ \phi_0^{(i)} + \phi_1^{(i)} = \alpha^{(i)} \quad r \leq \eta + R \]

\[ \phi_1^{(i)} = 0 \quad r = 0 \]

\[ \phi_0^{(i)} - \phi_0^{(i)} = \beta^{(i)} \quad r = R \]

\[ \phi_0^{(i)} + \delta \eta^{(i)} = \gamma^{(i)} \quad r = R \]

where \( i = 1, 2, 3, \ldots \)

\[ \alpha^{(0)} = 0 \quad \phi_0^{(1)} = -2 \phi_0^{(1)} \]

\[ \beta^{(0)} = 0 \quad \gamma^{(0)} = 0 \]

\[ \phi_0^{(1)} = 0 \quad \phi_0^{(2)} = \frac{1}{2} \phi_0^{(2)} - \phi_0^{(2)} \]

\[ \gamma^{(1)} = 0 \]

\[ \gamma^{(2)} = 0 \]

where \( \delta = \frac{R}{\epsilon} \) is a ratio of the tube’s elastic coefficient to the fluid density. We will replace \( x_0 \) and \( t_0 \) by \( x \) and \( t \) later, respectively.

3. APPROXIMATE SOLUTIONS

3.1 First Approximation
Putting \( i = 1 \) in Eqs. (11), (12), (18) and (19), we get the first approximation equations

\[
\phi^{(1)}_{s} + \phi^{(1)}_{v} = \alpha^{(1)} \quad 0 \leq r \leq \eta + R
\]  
(20)

\[
\phi^{(1)}_{v} = 0 \quad r = 0
\]  
(21)

\[
\phi^{(1)}_{r} + \delta \phi^{(1)}_{\theta} = \gamma^{(1)} - \delta \beta^{(1)} \quad r = R
\]  
(22)

\[
\gamma^{(1)} = \frac{1}{\delta} (\gamma^{(1)} - \phi^{(1)}) \quad r = R
\]  
(23)

where \( \alpha^{(1)} = 0, \beta^{(1)} = 0, \gamma^{(1)} = 0 \). The appropriate special solutions of Eqs. (20)-(23) are

\[
\phi^{(1)} = \phi^{(1)}_{s} + \phi^{(1)}_{r} \cosh(kr) \quad r + \text{c.c.}
\]  
(24)

\[
\eta^{(1)} = \eta^{(1)}_{s} + \text{c.c.} \quad \eta^{(1)} = -\delta \phi^{(1)}
\]  
(25)

\[
\theta = kz + \omega t \quad \omega = \delta k T \quad T = \text{tanh}(kR)
\]  
(26)

where \( \phi^{(1)} \) and \( \eta^{(1)} \) are functions of slow variables \( z, t, (j = 1, 2, 3, \ldots) \), c.c. represents the complex conjugate of the preceding term.

\[2.2 \quad \text{Second Approximation}\]

Putting \( i = 2 \) in Eqs. (11), (12), (18) and (19), we have

\[
\phi^{(2)}_{s} + \phi^{(2)}_{v} = \alpha^{(2)} \quad 0 \leq r \leq \eta + R
\]  
(27)

\[
\phi^{(2)}_{v} = 0 \quad r = 0
\]  
(28)

\[
\phi^{(2)}_{r} + \delta \phi^{(2)}_{\theta} = \gamma^{(2)} - \delta \beta^{(2)} \quad r = R
\]  
(29)

\[
\eta^{(2)} = \frac{1}{\delta} (\gamma^{(2)} - \phi^{(2)}) \quad r = R
\]  
(30)

where the expressions of \( \alpha^{(2)}, \beta^{(2)} \) and \( \gamma^{(2)} \) can be obtained from the first approximation results according to Eqs. (15b), (16b) and (17b). Solving Eqs. (27)-(30) and eliminating the secular term, we can easily obtain expressions of \( \phi^{(2)} \), \( \eta^{(2)} \), and

\[
2i\omega \phi^{(2)}_{v} + i\delta(kRT^2 - T - kR)\phi^{(2)}_{v} = 0
\]  
(31)

that is

\[
\phi^{(2)}_{v} = \phi^{(2)}(x_1 + v_1 t_1, x_2, t_2, \ldots)
\]  
(32a)

\[
v_1 = \frac{\delta}{\omega^2} [kR(1 - T^2) + T].
\]  
(32b)

\[2.3 \quad \text{Third Approximation}\]

The third approximation equations are

\[
\phi^{(3)}_{s} + \phi^{(3)}_{v} = \alpha^{(3)} \quad 0 \leq r \leq \eta + R
\]  
(33)

\[
\phi^{(3)}_{v} = 0 \quad r = 0
\]  
(34)

\[
\phi^{(3)}_{r} + \delta \phi^{(3)}_{\theta} = \gamma^{(3)} - \delta \beta^{(3)} \quad r = R
\]  
(35)

\[
\eta^{(3)} = \frac{1}{\delta} (\gamma^{(3)} - \phi^{(3)}) \quad r = R.
\]  
(36)

Similarly, solving equations, we have the third approximate solutions as follows:

\[
\phi^{(3)} = -\frac{1}{2} \phi^{(2)}_{v} - \frac{\delta}{\omega^2} \phi^{(2)}(x_2, t_2, \ldots) \cosh(kR) \quad r + \text{c.c.}
\]  
(37)

\[
2i\omega \phi^{(3)}_{v} + i\delta(kRT^2 - T - kR)\phi^{(3)}_{v} = 0
\]  
(38)

That is

\[
2i\omega (\phi^{(2)}_{v} - v_1 \phi^{(2)}_{v}) + (\phi^{(3)}_{v} + 2\omega kR \phi^{(3)}_{v}) - \delta R \phi^{(3)}_{v} - B_1 \phi^{(3)}_{v} \phi^{(2)}_{v} = 0
\]  
(39)

Substituting (25) into (40), we have

\[
2i\omega (\phi^{(2)}_{v} - v_1 \phi^{(2)}_{v}) + (\phi^{(3)}_{v} + 2\omega kR \phi^{(3)}_{v}) - B_1 \phi^{(3)}_{v} \phi^{(2)}_{v} = 0
\]  
(40)

where

\[
B_1 = \frac{k^4}{2} \left( 9T^3 - 12 + 13T^3 - 2T^3 \right) \frac{k^2 (1 - T^2)v_1 + 2\omega k^2}{v^2 - \delta R}
\]  
(41)

Setting \( u = c \phi^{(2)} \), \( u \) stands for the complex amplitude (envelope) of \( \phi^{(2)} \), which is the leading term of the radial displacement \( \eta \), then (41) changes to

\[
2i\omega (u - v_1 u) + (u^2 + 2\omega kR - \delta \delta) u = 0
\]  
(42)

that is, the radial displacement of the tube wall is

\[
\eta = c \phi^{(2)} + O(c^2) = c \phi^{(2)} + O(c^2)
\]  
(43)

\[
\text{Aesch}(x \phi^{(2)} + O(c^2)) = 0
\]  
(44)

This solitary wave propagates along the tube length direction, its amplitude \( A \), width \( D \) and propagating velocity are

\[
A = \frac{2}{B} (2\omega^2 - 2k^2 \omega_1 + k^2 \epsilon)^{1/2}
\]  
(45)

\[
D = \frac{1}{c} (2\omega^2 - 2k^2 \omega_1 + k^2 \epsilon)^{1/2}
\]  
(46)

\[
v = v_1 - \frac{k}{\omega^2} \epsilon
\]  
(47)

respectively, where \( v_1 \) and \( c \) are given in (32b) and (45), respectively.

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REFERENCES

MICRO-INHOMOGENEITY STRUCTURES AND HYSTERESIS EFFECTS IN CAVITATING LIQUID

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INTRODUCTION

According to results of a number of experimental and theoretical investigations of shock wave propagation, the free gas microbubbles are permanently present in distilled water and in water purified from mechanical impurities. The cavitation strength of liquid [1-5]. The mechanism of cavitation nuclei stabilization as an important factor of the cavitation process was analyzed by many authors.

The model proposed by Harvey [6] (air in cracks of solid uncontrollable particles) for solving the problem seemed to be ideal. The stabilization effect can be produced by the ion mechanism [7,8] or by the long molecules structure of surface-active substances [3,4,9]. It was found that the cavitation strength of water condensed in the atmosphere of different gases depends on their nature [10]. The stabilization model based on the relation between diffusion processes and physical gas adsorption was suggested in [10]. However, these models deal with a single particle and can’t explain some effects of a durable reaction of liquid on the bubble cavitation development.

COMBINATIVE STRUCTURES OF NUCLEI AND HYSTERESIS EFFECTS

A more detailed study of the cavitation processes as a result of underwater shock wave reflection from free surface carried out within the framework of experimental methods described in [11] has shown that the availability of single bubbles is deficient to explain the effects connected with an "original" memory of the cavitating medium. Fig.1 gives a comparison of the dynamics of the free surface of settled distilled water at the first reflection (curve 1) and the second identical shock wave loading (curve 2) carried out 1 min later. It is easily seen that in the first case the free surface displacement which is recorded by the capacitance gauge (CG, Fig.2) keeps during the time exceeding the shock wave duration (4 ms). This result points to the correlation between the dynamics of cavitation zone excited by reflection pulse and behavior of the free surface with their own characteristic time. Curve 2 does not possess an observed peculiarity. Such an effect can be connected with the change of nuclei structure or certain memory of cavitating liquids.

An analogous hysteresis is also observed at investigation of light scattering on inhomogeneities in liquids. The light of He-Ne laser (L1), scattered on microbubbles, was registered by a photomultiplier (PM) at an angle of 22° to the direction of a beam (LB) 1.5 mm in diameter transmitted at the depth of 3 mm under the free surface (Fig. 2). The dynamics of scattering indicatrix as a result of shock wave reflection has been studied. The incident shock wave with an amplitude of about a threshold one was generated by a membrane excited by pulse magnetic field. Fig. 3 shows the intensity of scattering light under primary (curve 1) and secondary shock loading of the same intensity in a few minutes (curve 2). The higher the shock wave amplitude the higher intensity of scattering light (curve 3) because of intense cavitation development. The return to the previous loading level shows that light scattering almost disappears (curve 4) and can be restored only in several tens of minutes.

The results of experiments on light scattering at static pressure point to the irreversible character of structure changes of some types of microinhomogeneities. In particular, at a decrease of pressure to 30-50 kPa and its subsequent increase to the initial level the initial indicatrix doesn’t restore.

The structure of microinhomogeneities was studied within the framework of static statement, on minisamples (thin layers, drops) of a distilled water, with the help of an immersion microscope and sensitive high resolution video apparatus (SONY CCD-V88).

It has been established that microinhomogeneities really have the complicated cluster structures which can be divided into 4 types (Fig. 4, A-D):

A) single solid microparticles,
B) free gas microbubbles which can contain solid particle of smaller size,
C) combinations of particles and bubbles (comparable in size),
D) formations which contain two or more bubbles adsorbed on the solid particle surface.

It is obvious that the combinitative structures of cavitation nuclei (B-D) will certainly react to intense loading and will be capable of varying under its action. A part of microbubbles can collapse and then restore in a rather long time period to the equilibrium 1-1.5 mm radius, “gathering” solute gas from water. The minimum radius, from which the bubble can yet grow, can be estimated on the basis of data on pressurization of water samples and on increase of its strength at pressure more than 3.5 MPa; it is about 0.04 mm. The bubble, probably, dissolves under the action of more strong compression. A structure peculiarity of the cavitation nuclei shown, in particular, in Fig. 4 D, is in that even gradual pressure decrease can result in the growth and subsequent of bubbles on the surface of solid particles thereby giving the change of light scattering.
CONCLUSION

The experimental researches of subtle effects of bubbly cavitation show that: a) hysteresis effects are observed in cavitating liquids, b) microinhomogeneities form combinative cluster structures capable of changing under static and shock loading. The process of structure reconstruction is relaxation, it is determined by characteristic time of diffusion processes at recovery of the equilibrium bubble and, in principle, can be irreversible.

REFERENCES


Fig.1. Free surface displacement h at the original (1) and second (2) loading.

Fig.2. Experimental set up: shock tube, L-laser, PM-photomultiplier, CG-capacity gauge of displacements, A-amplifier, PMF-pulse magnetic field source.

Fig.3. Intensity of scattering light at the different conditions of loading.

Fig.4. Combining structures of microinhomogeneities.
THE STUDY OF DYNAMICAL BEHAVIORS OF MULTI-SOLITON STATE IN RECTANGULAR TANK

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1. INTRODUCTION

The study of subharmonic wave of fluid in a vertically oscillating tank goes back to Faraday(1831). Recently the Faraday experiment induce once more great interest and remark[1,2]. In virtue of the developing of physics to nineties of 20 century, Prof.J.Krumhansl as president of American Physical Society in 1989 pointed out "today many scientists see nonlinear science as the most deeply important frontier for the fundamental understanding of nature"[3]. The solitons in nature including in water wave are the important practical example of nonlinearity and the unity of physics.

For the previous research of multi-soliton states, pioneer had pay attention to the nonlinear integrable equations as nonlinear schrodinger equation found N solitons solution of bound state by the inverse scattering transform in an infinite long channel or trough, and the solitons are propagating there. At times, they also took notice of that an initially symmetric wavepacket evolved in an asymmetric manner and broke up into "radiation" and a finite number of "solitons" that are bound together.

Now in this report, we devote ourself to study the multi-soliton state(NSS) in a finite length tank which as the forced standing and nonpropagating soliton if the soliton is in single. But they well interact and move in the multi-soliton situation.

2. EXPERIMENTAL RESULT

On the basis of previous analysis of dynamical behaviors of two solitons[4,5], we further study the excited state of multi-soliton in rectangular tank. For the two solitons having same parity we carefully observed that there is a threshold $a_c$ of excited signal amplitude on the excited plane of phase. If the signal was under the $a_c$, the behaviors of two solitons attract each other first, pass through and reach to some extreme positions, then they will repeat above process again and periodically, but by contrast if the signal was higher than the $a_c$, they will interact as elastic ball, turn over as they move closely enough and not pass through.

Under the situation of NSS, the excited energy of the system which are directly proportional to the number of solitons is increasing. The excited shaker wants to convey many energy as need, i.e. if the system working at the state having same frequency, the excited amplitude must be larger. The existence regime of NSS located higher than that of state of single soliton in the phase plane. But the regimes of the different number of solitons will have some overlap, which depend on the composite figure of solitons polarity state.

The stretch of every soliton was near same in the multi-soliton state. The relative phase of the water wave only have two favorable directions suited to the boundary condition of the vibrating tank, one is same polarity and the another is opposite polarity. The phase difference in between has not been observed in experiment. In order to describing the situation of soliton polarity, we introduce the symbol $T$ stands for one soliton have certain polarity, and $TT$ as the two soliton state of same polarity. The "TTTL" as the four soliton state which is composed of two opposite phase group of bi-soliton state of same polarity.

The production of multi-soliton state only if excite in their corresponding regime of existence, and must have right initial conditions (regular disturbance). Because at the same excited frequency and amplitude, many excited modes of surface water waves and their composition state will permit existence in the tank.

Based on the experimental observation of NSS, we have put in order of production of their composition mode according to whether are their easy to emerge.

Besides this, we had also sampled the time series of NSS at certain positions in the basin, drawn the time waveform and the FFT spectrum etc.

As view of the measured result, the states (LT),(TTLT),(TTTLT) are more stable among the states having even number soliton. There was formed easily (compare with the other soliton state of same number), and would return to normal state if they interfered with external disturbance, and the existence condition of this states are rather stability of all others.

The scenario of all the soliton with
same polarity in tank as the state (TT′) or (TTTT′) etc. are difficult to grow.

The scenario of neighboring solitons with inverse polarity as (TLTT) was poor stable than the configuration form of coupled pair of soliton. If single soliton emerges between pair soliton or a few of opposite polarity solitons, the stability of that state will be bad always.

We must emphasize that the stability feature denoting the life time of the state must long enough than that in which we can proceed the experimental measurement. Otherwise we will consider this state was instability or substability.

3. THE EXPLORATION OF THE COMPOSED RULE OF STATE

1. Two same polarity solitons made up the so call "bound pair", as we had discuss before[4], there is an attractive force between same polarity solitons and repelling force between inverse polarity soliton. The interactive force of soliton has potential energy in fluid field. Then the potential energy will decrease and located at lower energy level if two same polarity soliton are bound up in. There are the objective mechanism of bound pair in the water wave. But the two solitons having opposite polarity could not compose the similar "pair".

2. The relative phase of soliton state had only two favorable directions as shown( TLL ) queuing three solitons, if we wanted to turn the direction of middle soliton, the relative state of composition was maintained. The whole course of composition changes from beginning to end which will undergo the state of unfavorable direction, it might be hard to emerge. The change of state is forbidden.

3. The fluid field generally found itself in a state of disorder, and the disorder characteristic of the state will be remarkable with the raising temperature. The state of multi-soliton was an excited order state. If neighbor of the bound pair had more same polarity solitons, there will conflict with the disorder feature of fluid field.

4. In a limited container, the effect of basin wall on the water waves could not be negligible. In [4] there had explain the mapping effect of side wall's mirror image along the longitudinal direction and formed the virtual bound pair, which had been demonstrated by careful observation of experiments. Recently in [7] the experiment had shown that the viscous damping near the wall about two-fold larger than estimation value by theory (for circular cylinder basin) and near a order of magnitude higher(for rectangular basin). Due to this, once the movement figure near the wall had formed, they could not bear to broken suddenly.

In conclusion, all the description of the state composition rule of multi-soliton was considered phenomenologically. The dynamical system of fluid itself was comparatively complex, but physics would obtain succinct conclusion or structures from that. For example, the asymmetric equilibrium configuration lacking rotational could appear in a rotating fluid. We also observed the slowly and quickly, clockwise or counterclockwise forced standing solitary waves in a circular basin[8], which is the result of the symmetry breaking bifurcation[1]. So we can understand the unity of physics and plentiful phenomena and deep connotation of solitons study within nonlinear science.

Reference

THE NONLINEARITY OF THERMOACOUSTIC OSCILLATION

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I. INTRODUCTION

The history of the research of thermoacoustic effect is long, but there had been some difficulties for explanation of this effect. Not until the last quarter of the 20th century did it explain by famous Lord Rayleigh[1], who pointed out that "If heat be given to the air at the moment of greatest condensation, or be taken from it at the moment of greatest rarefaction, the vibration is destroyed." Since then the research on this effect was continued. J. Wheatley et al.'s[2] research on thermodynamic process appeared in the thermoacoustic effect revealed the physical essence of the conversion between thermal and acoustic energy in thermoacoustic effect. And this stimulated the research of thermoacoustic engine and its application, and a remarkable progress have been made in this respect.

But in research of passed decades people paid attention mainly to linear respect of the effect either theoretically or experimentally. In recent years people paid attention to nonline aspect of it, especially to soliton state[3] and nonlinear dynamics[4] of thermoacoustic oscillations. A. A. Atchley[5] recently observed nonlinear oscillation in a thermoacoustic prism mixer filled with 3 atm. of Helium.

The purpose of this paper is to give a preliminary results of our research on nonlinear phenomenon of thermoacoustic oscillations in a tube filled with natural air.

II. THE EXPERIMENTAL TECHNIQUE

The experimental set up is shown schematically in Fig. 1. It is a glass tube which closed at one end and open to air at the other end, and its length is 65 cm, the inner diameter is 5.5 cm. Inside of it is a bundle of small glass tube (called stack, it is know as second or secondary medium, the air the first or primary medium) which is in thermal contact with two heat exchangers (two copper grids). One end of the stack (near closed end of the tube) is held at higher temperature with electrical heater wound outside the large tube and the other end is maintained at ambient (room) temperature with room-temperature cooling water passing through a tube raped around the large tube. The length of the small glass tube is 8 cm, the inner diameter is 3 mm. Two electrical thermo-couples were fitted to two ends of the stack to measure the temperature of the two ends of the stack. The heating electric power was adjusted through a variac to control the temperature difference between the two ends of the stack. The acoustic pressure is measured with a capacitive microphone which sent the measured signal to oscilloscope and HP3562A spectrum analyzer.

III. THE RESULTS AND DISCUSSIONS

We adjust the position of the stack so that the distance between the stack and the closed end of the tube is 20 cm, and found that the threshold value for generation of sound oscillation in the tube is about 310°C, and the stable sound pressure level is higher, that is thermoacoustic energy conversion efficiency is higher under this condition.

By adjusting the heat power, when the temperature difference between two ends of the stack is larger than threshold value the audible sound is exited in the tube, and the level increasing gradually, and reached certain stable level eventually. The larger the temperature difference is, the higher the stable sound level reached is. The fundamental frequency of the sound is about 136 Hz, the pressure level at the open end of the tube is about 90~100 dB. The matter it is from the closed end of the tube, the higher the sound pressure level is. When the temperature difference between two ends of the stack is just higher than threshold, the sound wave is a pure sinusoidal wave, shown in Fig. 2 in which (a) is wave form, and (b) the corresponding frequency spectrum, it shows that the fundamental is 50 dB larger than the others. As the temperature difference is increased the wave form distortion is increased gradually. Fig. 3 shows the wave form and the corresponding spectrum for the temperature difference being 450°C. We see from Fig.3 that the wave form is non-symmetry for positive and negative half-cycles, the higher harmonic components appear in the spectrum.

You may ask that whether the generation of intrinsic high harmonics or the nonlinear distortion causes the generation of the high frequency components in the spectrum. We know that in a resonance tube the frequency of overtones depends upon dispersion relation ω=ω(k), while harmonic generation results in exact multiples of the fundamental. The all of high frequency components in the spectrum we measured in the experiment are multiples of the fundamental. But if we want to verify that the generation of high frequency components results from the nonlinearity we must verify that these components are not intrinsic high harmonics. But it is very difficult to find dispersion relation theoretically for such a complicated system with a stack of small glass tube in large tube and non-uniform distribution of temperature along the tube. We have measured the intrinsic frequency spectrum with loop-method experimentally, and found that for a tube with no stack of small glass tube in it and uniform distribution of temperature along the tube, the intrinsic frequency spectrum lines show a uniform distribution. But with a stack of glass tube in large tube the spectrum line show a non-uniform distribution whether the temperature distribution uniformly or not along the tube. Fig. 4 shows a spectrum for the heater wound in the tube and the temperature difference between two ends of the stack is just lower than the threshold valve, it shows that only the fundamental coincides with the peak in the oscillating spectrum. We say believe that there is no essential difference between intrinsic spectra for temperature difference being higher and lower than threshold valve, thus the fundamental in oscillation spectrum is the lowest harmonic of the system, and the higher components result from the nonlinearity not from the excitation of high harmonics.

But we couldn't determine the relative contributions from non-linear propagation and non-linear generation to the observed waveform distortion, nor such a determination will also require a dispersion relation of the system. This area is receiving attention now. Here we make a qualitative treatment for the problem according to
Coppens and Sandea's theory on finite amplitude standing waves. According to their theory the ratio of the amplitude of the second harmonic to that of the of the fundamental is given by

\[
\frac{R_2}{R_1} = \frac{1}{2} \left( \frac{1}{\sqrt{C_2}} \right) Q_2
\]  

(1)

Similarly, the amplitude of the third harmonic to that of the fundamental is given by

\[
\frac{R_3}{R_1} = \frac{1}{2} \left( \frac{1}{\sqrt{C_3}} \right) Q_3
\]  

(2)

where \( R_1, R_2, R_3 \) are the fundamental, second and third harmonic amplitude respectively, \( \rho \) is the ambient density of air, \( C_0 \) is infinitesimal amplitude sound speed, \( Q_2 = (\sqrt{\omega/4\pi C_2}) \cos \theta_2 \), \( Q_3 = (\sqrt{\omega/4\pi C_3}) \cos \theta_3 \), \( \theta_2 \) is the attenuation coefficient, \( \theta_2 = 0.414 \), \( \theta_3 = 0.732 \). Insert relevant experimental parameter into equations (1) and (2), we obtain

\[
\frac{R_2}{R_1} = -14 \text{ dB}, \quad \frac{R_3}{R_1} = -23 \text{ dB}
\]

comparing to Fig.3 shows a large difference. One of the reasons for this is that we consider here only the contribution of nonlinear propagation to waveform distortion. Further experiment and theoretical consideration, such as numerical calculation of a boundary problem for nonlinear coupled momentum and energy equations by the mode-truncation method are underway. We believe that these works are helpful for us to further understanding the nonlinear phenomenon of thermoacoustic effect.

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References

SPHERICAL WAVE PROPAGATION

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1. NEW FORMULA FOR SPHERICAL WAVES

Lamb/Sommerfeld’s integral formula [1], [2]

\[ \exp(-ik\sqrt{r^2 + z^2}) = \int_0^\infty J_0(\xi)\exp(-\xi\sqrt{r^2 + k^2})\xi d\xi / \sqrt{r^2 + k^2} \]  (1)

where \( k \) is the wavenumber, \( r \) and \( z \) are coordinates, has been used to investigate the spherical wave propagation for many decades. The integral in Eq.(1) is an improper integral, the integration is over infinite range and it has a singularity at \( \xi = k \). Since we have not found any discussion of the convergence of the integral in (1) in the literature, we shall show that the integral in Eq.(1) is convergent, and a new formula will be obtained as a consequence of the convergence analysis.

First we rewrite the improper integral in (1) as

\[ \lim_{\alpha, \beta \to \infty} \int_{\beta}^{\alpha} J_0(\rho \sqrt{\xi^2 + \rho^2}) e^{-i\rho \xi} d\xi = \int_0^{\infty} J_0(\rho \sqrt{\xi^2 + \rho^2}) e^{-i\rho \xi} d\xi, \]  (2)

where \( \rho = \sqrt{r^2 + k^2} \). We can eliminate the singularity by substituting \( \xi = \sqrt{r^2 + k^2} \) into Exp.(2), and when \( b \) tends to zero, Exp.(2) becomes

\[ \lim_{\alpha, \beta \to \infty} \int_{\beta}^{\alpha} J_0(\rho \sqrt{\xi^2 + \rho^2}) e^{-i\rho \xi} d\xi = \int_0^{\infty} J_0(\rho \sqrt{\xi^2 + \rho^2}) e^{-i\rho \xi} d\xi \]  (3)

The first integral of Exp.(3) is independent of the parameter \( \alpha \), and it is convergent over the finite range. By the Dirichlet test theorem, the second integral of Exp.(3) is convergent when the parameter \( \alpha \) tends to infinity. Hence the integral on the right side of Eq.(1) is convergent and we have also obtained a new expression for a spherical wave which avoids the singularity in Lamb’s integral. When written Exp.(3) in a compact form, we have

\[ \exp(-ik\sqrt{r^2 + z^2}) = \int_0^{\infty} J_0(\rho \sqrt{\xi^2 + \rho^2}) e^{-i\rho \xi} d\rho. \]  (4)

The integration path is taken from \( x = ik \) to \( x = i0 \) along an imaginary axis and from \( z = 0 \) to \( z = \infty \) along a real axis.

2. TWO FLUID HALF SPACES

We shall consider acoustical spherical waves in elastic media by using the classical theory of elastic wave propagation [3]. First, let us consider the problems of spherical wave reflection and transmission at a plane interface between two semi-infinite fluid media. In a cylindrical coordinate system, assuming that the point source, at \( z = h(> 0) \) on the \( z \)-axis, is in medium 1, the transmitted wave in medium 2, and the interface is at \( z = 0 \), the wave field is symmetrical about the \( z \)-axis and depends only on the coordinates \( z \) and \( r \). The boundary conditions require the discontinuities of the \( z \)-component displacements and the pressures at \( z = 0 \).

The classical solutions for this problem can be found in [3]. Using the method of transformation given above, we can obtain a new form of solution for the reflected wave, and the complete solution, in terms of displacement and potentials, can be written as

\[ \psi_1 = \exp(-ik\sqrt{r^2 + (z + h)^2}) / \sqrt{r^2 + (z + h)^2} \]
\[ \int_0^\infty \rho \frac{\partial P_2}{\partial z} e^{-i\rho \xi} d\xi + \int_0^\infty \rho P_2 e^{-i\rho \xi} d\xi. \]  (5)

\[ \psi_2 = \int_0^\infty \frac{2\rho_1}{\rho_2 + \rho_1} - 1 \int_0^\infty J_0(\rho_2 \xi) e^{-i\rho_2 \xi} d\xi, \]  (6)

where \( \rho_1 \) and \( \rho_2 \) are the densities, \( q_1 = \sqrt{r^2 + h^2} \) and \( q_2 = \sqrt{r^2 + k^2} \), and we have omitted the time factor \( \exp(i\omega t) \).

The boundary conditions guarantee the continuity of the energy flux in the \( z \)-direction for waves (5) and (6), but it is not easy to show, analytically, whether the energy flux of wave (5) or (6) is solenoidal. We shall numerically test solutions (5) and (6) for energy conservation. Consider now two plane surfaces, parallel to the interface, one is at \( z = h/2 \), and the other is at \( z = -h \). It is easy to show that the time average power crossing the plane, \( z = h/2 \), from the direct wave (the first term in (5)) is half of the total time average power generated by the point source. The law of energy conservation requires that the sum of the time average powers from the reflected wave (the second term in (5)) and the transmitted wave across these two surfaces must be equal to another half of the total power.

The time average power crossing an area of radius \( a \) on each plane is given by the integral of the time average energy flux over the area. The z-component of the time average energy flux, \( I_z \), is defined by the product of the pressure and the z-component particle velocity, and its time average can be expressed as

\[ <I_z> = \frac{1}{2} \rho c^2 \frac{\partial P}{\partial \theta} \frac{\partial \theta}{\partial \theta}, \]  (7)

where \( \theta \) indicates the complex conjugate.

Using the numerical integration routines in the NAG library, the time average powers were calculated, and are normalized to half of the total time average output power generated by the point source. The results are shown in Table (1).

The numerical integration becomes difficult at large values of

<table>
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Table 1: Calculation of time average powers crossing plane surfaces for the reflected wave and the transmitted wave. The kerosene-water interface is at \( z = 0 \), frequency 1000 Hz, source height 1m

\( r \) since the integrands are rapidly oscillatory. The sum of the reflected and transmitted wave powers at \( \rho = 1000 \) m is about 0.5974, which is in good agreement with the required result.

3. A SOLID PLATE IN A FLUID

Next, we consider spherical wave propagation in the case
of a solid in a fluid with a point source (at \( z = h \)) in the fluid. The boundary conditions at the fluid-solid interface (at \( z = 0 \)) or at the solid-fluid interface (at \( z = -l \)) require that the z-component displacements are continuous, the pressure in the fluid is equal to the normal stress in the solid, and the shear stress in the solid vanishes. The solutions for this problem, based on Lamb’s integral, can be found in [4] by Piquette. An equivalent, but new form of the solutions is given by

\[
\psi_1 = \frac{1}{\pi} \int_0^R \frac{1}{\sqrt{r^2 + h^2}} e^{-k(r - x)} \sin(kz) \, dz,
\]

\[
\psi_2 = \frac{1}{2\pi} \int_0^R \frac{1}{\sqrt{r^2 + k^2}} e^{-k(r + x)} \sin(kz) \, dz,
\]

\[
\psi_3 = \frac{1}{2\pi} \int_0^R \frac{1}{\sqrt{r^2 - k^2}} e^{-k(r - x)} \sin(kz) \, dz,
\]

\[
\psi_4 = \frac{1}{2\pi} \int_0^R \frac{1}{\sqrt{r^2 + k^2}} e^{k(r + x)} \sin(kz) \, dz.
\]

where \( q_1 = \sqrt{r^2 - k^2}, \quad q_2 = \sqrt{r^2 + k^2}, \quad q_3 = \sqrt{r^2 - k^2}, \quad q_4 = \sqrt{r^2 + k^2} \) and \( R, L_0, T_0, T_1 \), and \( \xi \) are the integrand coefficients determined by the boundary conditions. \( \psi_1 \) and \( \psi_2 \) represent the longitudinal and transverse wave potentials in the plate, respectively. The integrand coefficients are very complicated. Piquette did not present them in [4]. Our version of them can be found in [5].

Piquette has calculated the reflection coefficient at \( z = 0 \), and the transmission coefficient at \( z = -l \), both at the normal incidence angle, and his results show that in some cases the reflection coefficient or the transmission coefficient can be greater than one, however, his results have not been verified either by experiments or by other independent methods. Using the NAG routines, we have done similar calculations using Eqs.(8) and (11), and not found the "overpressures" predicted by Piquette. Our results are very similar to those for plane waves, and are shown in Table (2). It is likely that the "overpressures" are the results of computing error caused by the numerical integration around the singularities in the integrand coefficients.

<table>
<thead>
<tr>
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<th>aluminum</th>
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<td>5000Hz</td>
<td></td>
<td></td>
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</tr>
<tr>
<td>spherical wave</td>
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<td>0.7812</td>
<td>0.6526</td>
</tr>
</tbody>
</table>

Table 2: Comparison between spherical wave and plane wave. Calculation for the pressures above and below the centre of a metal plate in water, source height \( h = 2m \), the thickness of the plate \( t = 0.01m \).

Using NAG routines we have also calculated data using Eq.(8) for plotting pressure contours in a region between the point source and the fluid-solid interface for air-aluminium-air combination. The experimental tests were carried out to measure the pressure contours in the case of a point source in front of a disc. The experiment was conducted in an anechoic chamber. A circular disc made from aluminium was used. A spherical wave was generated from an open end of a brass tube which was coupled to a driver unit, and a single frequency sine wave signal was used.

The predicted pressure contours from the infinite surface model and the measured pressure contours in front of a disc are shown in Figs. (1) and (2). They are basically in agreement.

Figure 1: Predicted pressure contours, point source at \((0,1000)\), the first reflection surface at \( z = 0 \), plate thickness 8mm, frequency 1000Hz.

Figure 2: Measured pressure contours, point source at \((0,1000)\), the first reflection surface at \( z = 0 \), plate thickness 8mm, disc diameter 1200mm, frequency 1000Hz.

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A7-2

Visualization of the Sound Scattering from a Cylinder

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Cylinders with diameter 1.5 cm and length 3 cm and 8 cm respectively are immersed in the liquid, composed of a single liquid or two non-mixable liquids with layered interfaces. When the sound waves impinge on the cylinder normally or obliquely, the echo structure and the effects of the interface upon the echo structure are studied by schlieren method. The schlieren images show the effects caused by the discontinuity in the shadow zone. Using this fact, the locus of the helical surface wave induced by a sound wave obliquely incident on the cylinder can be described successfully. Therefore the three-dimensional scattering field can be illustrated with schlieren method and has not to use the stereoscopic visualization with complex appliance and complicated algorithm.

I. Scattering of a Cylinder for the Case of Normal Incidence

According to the theoretical prediction and the experimental result[1], it is shown that there are two kinds of scattering waves caused by the elastic cylinder immersed in water. The first one is the circumferential waves which are composed of wave, travelling in water along the surface of cylinder, i.e., Franz wave, and waves, travelling inner the cylinder along the boundary of the cylinder, i.e., whispering gallery waves. The second one is the waves obeying the geometrical laws, i.e., reflected waves and transmitted waves (the transmitted waves inner the cylinder travel with the longitudinal or transverse velocity), which pass through the cylinder directly or reflect several times in the cylinder undergo a mode transform (i.e., longitudinal to transverse or vice versa) or not, then project in water. The longitudinal and transverse velocity, obtained by experiments coincide with the known values.

II. Scattering of a Cylinder for the Case of Oblique Incidence

It is well known that the scattering of a cylinder can be visualized by schlieren method for the case of normal incidence. In this case, the sound wavefront of direct wave and scattering wave can be set in the plane of lights, i.e., the lights are perpendicular to sound rays.

For the case of oblique incidence, matters stand otherwise, because we can not set the wavefront in the plane of lights. So till lights, surface wave on the cylinder, scattered for the case of oblique incidence can not be visualized by schlieren method. However, it is found that the obstacle in the shadow zone of direct sound field can arouse scattering by primary scattering waves. If the size of obstacle is sufficient small, the second scattering wave can be assumed to be spherical wave, then this spherical wave can be visualized by schlieren method. Using this fact, the obstacle set on the surface of cylinder can be used as tracking mark to monitor the waves travelling along the surface, but whose wavefronts are not in the plane of lights. The travelling paths of the helical surface waves, caused by scattering of cylinder in the case of oblique incidence are thus reconstructed successfully, and the helical surface waves are examined in existence firstly by acousto-optic visualization.

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In spite of schlieren method can only visualize the two-dimensional image, but as the scattering wave of the obstacle is demonstrated in meantime, the travelling path of the surface wave in the three-dimensional space is tracked simultaneously, therefore the three-dimensional helical surface wave is reconstructed. So this method can be called the method of reconstructing by obstacle.

1. Comparison between the scattering wavefronts in the cases of normal and oblique incidence.

From Fig. 1a and 1b, we found that in the case of normal incidence the surface waves travel along the circle perpendicular to the axis of cylinder and radiate the scattering waves outwards along the tangents of the cross-section of cylinder. In this case, the wavefronts of reflected wave are in the plane of lights, so the image of reflected wave is clear. While for the case of oblique incidence, the surface waves travel along the helix and radiate scattering wave along the tangents of helix, here the angle between the tangents and the axis of cylinder is . But the wavefronts of reflected wave are no longer in the plane of lights, so their image is indistinct.

2. Comparison between the acoustic measurements and the schlieren visualization.

The second scattering wavefronts of the obstacle for the various incident angles at the same time after transmitting the incident sound pulse are shown in Fig. 2a and the time differences for the various incident angles are shown in Fig. 2b. The time delay due to various incident angles is same with what is obtained by the acoustic measurements[2]. So it is examined that when the incidences deviate from the normal incidence, the excited surface wave is no longer the Franz wave creeping round the circle, but is the helical surface wave travelling along the helix on the cylinder.

III. Scattering from Edges and Discontinuities in Shadow Zone.

The scattering from edges of cylinder is observed[3], where the sound waves impinge downward on the end-face of cylinder, and the incident wave, R is the wave reflected by the end-face and C is the echo caused by the edge.

It is evident that the acoustics echoes from a target are independent upon its shape in shadow zone, so the echo can not be caused from the shadow zone. However, if an obstacle (or the discontinuity) be set on the surface sitting in the shadow zone of the cylinder(fig. 4), an incident sound wave impinging downward on the edge of cylinder, the scattering wave caused by the edge travels along the cylindrical surface across the discontinuity (in shadow zone), then the second scattering is occurred. When the scattering from the edge is weakened, or disappears, the second scattering is weakened or disappears too. The same, such scattering waves can be used to demonstrate the waves, whose wavefront is not in the of lights.

IV. Target Scattering in a Layer of Two Non-mixable Liquids.

In order to visualize the sound field by the schlieren method, two kinds of transparent, non-mixable liquids are used. The upper layer is silicone
oil; the lower layer is distilled water. An aluminium cylinder is immersed in the lower layer and the images of wavefronts are obtained by schlieren method.

1. The scattering of target in two-layered medium is like what in single liquid, but the intensities of the scattering are weaker than that in single liquid (Fig. 5).

2. In two-layered medium, the target scattering in such condition is dependent upon the incident angle impinging on the surface and decreases as incident angle increases. The maximum scattering occurs at normal incidence. When \( c_2 < c_1 \) (where \( c_1 \) is the sound velocity of the upper layer and \( c_2 \) is the sound velocity of the lower layer), as the incident angle is equal to or greater than the critical angle, the scattering from the target will not occur (Fig. 6).

3. For the case of normal incidence, the echoes of the target consist of the echo caused by the direct incident wave and the re-scattering between the interface and the target (Fig. 7). If the incident wave is sufficiently intense, the multiple re-scattering will occur and the echo structures become more complicated.

4. When the sound wave is incident obliquely on the interface, the echoes of target mainly consist of the scattering by direct incident wave, the effect of re-scattering can be ignored.

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SCATTERING OF ACOUSTIC WAVE FROM AXISYMMETRIC ELASTIC BODIES

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1. INTRODUCTION

Acoustic wave radiation and scattering from submerged elastic structures are of great importance in underwater acoustics. Since the motion of the structures and the pressure of the surrounding fluid are coupled, wave propagation equations in elastic structures and fluid must be solved simultaneously. Analytic investigations have been limited to simple geometries such as spherical and cylindrical shells[1]. For bodies having arbitrary shapes, a variety of numerical methods[2] have been developed. In this paper, we present a simplified finite element(FE)/boundary element(BE) formulation for the radiation and scattering of acoustic wave from axisymmetric elastic bodies by means of the Fourier expansion method.

2. FE/BE FORMULATION

We consider an axisymmetric elastic shell submerged in an infinite fluid, where a plane wave \( P \) is incident with cone angle \( \theta \), from the symmetric axis as shown in Fig. 1. If we assume a harmonic wave in the form of \( \exp(i\omega t) \), the wave equation in the elastic body is given by

\[
\mu \nabla^2 \bar{U} + (\lambda + \mu) \nabla \cdot \bar{U} + \rho_\varepsilon \omega^2 \bar{U} = 0,
\]

where \( \lambda, \mu \) are the Lamé coefficients, \( \rho_\varepsilon \) is density of the body, and \( \bar{U} \) is displacement vector. In the fluid, the Helmholtz equation is satisfied

\[
\nabla^2 P + k^2 P = 0,
\]

where \( P \) is the total pressure (incident wave and scattered wave) and \( k \) is the wavenumber of the fluid.

At the surface, the displacement and pressure are related as

\[
\rho_\varepsilon \omega^2 \delta = \partial P/\partial n,
\]

where \( \delta \) is the normal component of the displacement and \( \rho \) is density of the fluid. Boundary condition for stresses is that the normal component of the stress is equal to the pressure, while tangential components must vanish. In addition, the scattered pressure must satisfy the radiation condition.

We first discretize the wave equation (1) in the shell using the FEM. In this study, we choose an isoparametric element originally developed for curved thick shells[3], which also has to vary linearly across the thickness direction. We use three nodal points in each element, in which nodal points are defined on the surface rather than on the mid-plane.

We expand the axial displacement \( U \) into Fourier series in terms of the azimuthal angle \( \phi \) as

\[
U(x, r, \phi) = \sum_{n=0}^{\infty} U_n(x, r) \cos n\phi,
\]

while radial and circumferential displacement \( V \) and \( W \) may be expanded in a similar manner. In each element \( U_n \) is interpolated in terms of the new curvilinear coordinates(see Ref. [3] for details).

We also expand pressure \( P \) and normal component of the displacement \( \delta \) on the surface as

\[
P = \sum_{n=0}^{\infty} P_n \cos n\phi, \quad \delta = \sum_{n=0}^{\infty} \delta_n \cos n\phi.
\]

We divide the generator into \( M \) nodes, and terminate the infinite summations of the Fourier expansions to \( N_F \) terms.

The FE formulation of Eq. (1) leads to the familiar equations of motion

\[
K_n \ddot{U}_n - \omega^2 M_n \dot{U}_n = -D_n \ddot{P}_n, \quad n = 0, 1, 2, \cdots,
\]

in which \( K_n, M_n \) are stiffness and mass matrices, and \( \ddot{U}_n, \ddot{P}_n \) are displacements and pressure vectors corresponding to the Fourier index \( n \). The matrix \( D_n \) is obtained from the work done by the pressure.

We next formulate the Helmholtz equation (2) in fluid by the BEM. We may rewrite Eq. (2) in the Helmholtz integral equations

\[
\gamma P(x) = P_0 e^{-ikx} \quad \text{in fluid}
\]

\[
+ \int_{S} [P(x') \frac{\partial}{\partial n'} G(x, x') - \frac{\partial P(x')}{\partial n'} G(x, x')] dS(x'),
\]

where \( k \) is the wavenumber for the Green's function \( G(x, x') \) is given by

\[
G(x, x') = e^{-ik|x-x'|}/4\pi|x-x'|.
\]

If \( x \) belongs to the surface, \( \gamma = 1/2 \) provided that the surface varies smoothly, while for exterior points, \( \gamma = 1 \).

After substituting Eqs. (5) into Eq. (7), we multiply \( \cos m\phi \) to both sides of Eq. (7) and integrate from 0 to \( 2\pi \) to find

\[
\int_{0}^{2\pi} P_0 (x) e^{-ikx} \cos m\phi R^2 \frac{\partial}{\partial n} d\phi
\]

\[
= \int_{0}^{2\pi} \left[ \Psi^n_m(x, z) - \omega^2 \Psi^n_m(x, z') \right] R^2 \frac{\partial}{\partial n} d\phi,
\]

in which \( \Gamma \) is the length along the generator and

\[
\epsilon_m = 2 \quad \text{for} \quad m = 0, \quad \epsilon_m = 1 \quad \text{for} \quad m = 1, 2, \cdots.
\]

The functions appearing in Eq. (9) are

\[
F_m(x) = P_0 \int_{0}^{2\pi} e^{-ikx} \cos m\phi d\phi,
\]

\[
\Psi^n_m(x, z) = \frac{\epsilon_m P_m}{4} \int_{0}^{2\pi} \cos m\phi R^2 \frac{\partial}{\partial n} e^{-ikR} du,
\]

\[
\Psi^n_m(x, z') = \frac{\epsilon_m \delta_m}{4} \int_{0}^{2\pi} \cos m\phi R^2 \frac{\partial}{\partial n} e^{-ikR} du.
\]

Integrals in Eqs. (12) and (13) have singular terms as \( R \) approaches zero. Evaluation of \( \Psi^n_m \) and \( \Psi^n_m \) can be performed. 

Figure 1. An axisymmetric shell
analytically by transforming them into the complete elliptic integrals of the first and second kind as shown in Ref. [4]. From the expressions in Eqs. (12), (13), we may rewrite Eqs. (9) in matrix equation

\[ \mathbf{A}_n \mathbf{F}_n + \rho_n \mathbf{B}_n \mathbf{U}_n = \mathbf{F}_n, \quad n = 0, 1, 2, \ldots, \]  

(14)
in which \( \mathbf{F}_n \) represents force vector due to the incident wave. Combining Eqs. (6) and (14) gives a linear system of fluid-structure interaction problem.

Scattered wave \( P_s \) at an arbitrary position \( x \) exterior to the body is given by

\[ P_s(x) = P(x) - P_0 e^{-i \omega x} \]  

(15)

\[ = \int_S |P(x')| \frac{\partial}{\partial n'} G(x, x') - \frac{\partial P(x')}{\partial n'} G(x, x') |dS(x'). \]

Evaluation of \( P_s \) in Eq. (15) may also be simplified after we expand Green’s function in terms of the spherical Bessel function and use the Fourier expansions Eqs. (5). By employing the asymptotic form of the spherical Bessel function when \( ks >> 1 \), we can write the scattered wave at the far field as

\[ P_s / P_0 = f(\theta, \phi) e^{-i \omega t} / s, \]  

(16)

where \( f(\theta, \phi) \) represents a scattering function and \( (s, \theta, \phi) \) is a spherical coordinate of \( x \). Backscattered wave is given if we take \( \theta = \pi - \theta \) and \( \phi = \pi + \phi \), where \( \theta \) and \( \phi \) are spherical angles of the incident wave (we chose \( \phi = 0 \)).

3. NUMERICAL EXAMPLES

As numerical examples, we compute the scattering at the far field from rigid spheroid subject to plane incident waves. Fig. 2 shows comparisons of scattering function \( |f| / \alpha \) of a rigid spheroid (major semi-axis = \( a \), minor semi-axis = \( 0.1a \)) with Tobocman’s result \([5]\) when \( ka = 1 \) (\( \theta = 0^\circ \) corresponds to the backscattering). The direction of incident wave is perpendicular to the symmetric axis. We next compare the backscattering functions \( |f| / \alpha |^2 \) of the elastic spherical shell with the analytic results\([1]\) in Fig. 3. The shells are surrounded by water, while vacuum is assumed inside the shell. Dimensions and material properties of the spherical shell have been chosen as: \( \lambda = 11.3 \times 10^{11} N/m^2, \mu = 7.4 \times 10^{11} N/m^2, \rho_s = 7840 \ kg/m^3 \), outer radius \( a = 2m \), thickness \( t = 2 \times 10^{-3} m \), Poisson’s ratio \( \nu = 0.3 \). Speed of sound and density of the water are \( C = 1500 m/s, \rho = 1000 \ kg/m^3 \).

![Figure 2. Scattering from a rigid spheroid when \( ka = 1(M = 33, N_F = 5) \).](image)

![Figure 3. Backscattering from an elastic spherical shell vs. wavenumber \( ka(M = 33, N_F = 0) \).](image)

4. CONCLUSIONS

We have presented a simplified BE/FE formulation for radiation and scattering of acoustic waves from axisymmetric elastic shells subject to non-axisymmetric boundary conditions. By employing the Fourier expansions of the dependent variables into the circumferential angle, we were able to reduce the dimension by one, which offers a significant savings of computations compared to the surface discretisation methods. The singular integrals encountered in the Helmholtz integral equation were performed analytically by transforming them into the complete elliptic integrals.

ACKNOWLEDGMENTS

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Sound Scattering of Plane Wave by a Transverse Isotropic Cylinder

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1. Introduction

The sound scattering of a plane wave by an infinite isotropic elastic cylindrical sample has been extensively studied. Faran's paper is the first investigation which dealt with the sound scattering of a plane wave by an elastic cylinder at the normal incidence. Flax and Li etc. have extended Faran’s work to the case including the oblique incidence. These theories enable the estimation of elastic coefficients of a cylindrical sample by measuring sound scattering from it. However, many of the engineering cylindrical samples are often transversely isotropic. There is no scattering theory on a transverse isotropic cylindrical sample.

In the present paper, we tried to solve the problem of sound scattering of a plane wave by a transverse isotropic cylindrical sample. The solutions of elastic vibration of an infinite transverse isotropic cylinder are necessary to solve this problem, and have been given by the authors. By using these solutions, the partial-wave scattering amplitude is derived from the boundary conditions. Both the scattering patterns and the backscattering responses are computed and compared for isotropic and transverse ones.

2. Theory of acoustic scattering by a transverse isotropic cylinder

Figure 1 shows the cylindrical coordinate and the direction of an incident plane wave. The z-axis is taken to be the axis of cylinder and makes an angle \( \pi/2 - \alpha \) with the incident plane wave.

(a) General case

In a liquid medium, the incident acoustic pressure can be expanded in the Rayleigh series form,

\[ P_0 = e^{-j\omega t} \sum_{n=0}^{\infty} e_n \left( -j \right)^n J_n(pr) \cos(n\theta) \]  

where \( q = k_0 \sin(\alpha), \rho = \rho_0 \cos(\alpha), k_0 = \omega/c, \) c is the sound speed of the liquid, \( e_0, \) is the Numann factor (\( e_0 = 1, \) if \( n = 0, \) \( e_n = 2, \) if \( n > 0, \)) and \( J_n(\cdot) \) is the first kind Bessel function of order \( n. \) The outgoing scattering wave is

\[ P_r = e^{-j\omega t} \sum_{n=0}^{\infty} S_n H_n^{(2)}(pr) \cos(n\theta) \]  

where \( H_n^{(2)} (\cdot) \) is the second kind Hankel function of order \( n \) and \( S_n \) is the partial wave scattering amplitude.

The partial wave scattering amplitudes can be calculated from the boundary conditions that the normal stress and the radial displacement must be continuous and that the tangential stresses must be free at the boundary of the cylinder. Since the vibration problem of an infinite transverse isotropic cylinder has been solved exactly by the authors, we can calculate the partial wave scattering amplitudes by using these solutions, Eqs. (1) and (2). We have the partial wave scattering amplitude as

\[ S_n = -j^n e_n \frac{\Phi_{J_n}(pa) - \rho a \Omega_n \Phi_{J_n}(pa)}{\Phi_{H_n^{(2)}}(pa) - \rho a \Omega_n \Phi_{H_n^{(2)}}(pa)} \]  

where \( j^n = -1, \) \( \Phi_n \) and \( \Omega_n \) are given in ref. (7), and primes denote the differentiation with respect to the argument. The partial wave scattering amplitude can be reduced to those of an isotropic cylinder if the transverse isotropic elastic constants are replaced with those of the isotropic material.

(b) Normal incidence

In this case, partial wave scattering amplitudes are given by

\[ \Omega_n = F_{13} F_{22} - F_{12} F_{23} \]

\[ \Phi_n = \rho (a \omega)^2 \frac{J_n(x_1)}{x_1^2} \left[ x_2 \frac{J_n(x_2)}{x_2^2} - x_3 \frac{J_n(x_3)}{x_3^2} \right] \]

and the partial wave scattering amplitudes given by these expressions are exactly the same as those of the isotropic cylinder, because the transverse isotropic cylinder shows the isotropic property in the x-y plane.
3. Numerical Results and Discussion

Numerical calculations were performed for a transverse isotropic cylinder immersed in water. The physical parameters of the materials are given in Table 1, where the elastic constant of the \( z \)-axis direction \( C_{33} \) is assumed to vary over the range \( C_{33}=\sigma C_{11} \) with \( \sigma = 1-10 \).

<table>
<thead>
<tr>
<th>Elastic parameters</th>
<th>( C_{11} )</th>
<th>( C_{12} )</th>
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<tbody>
<tr>
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<td>0.565</td>
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<table>
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<th>( C_{55} )</th>
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<table>
<thead>
<tr>
<th>Elastic parameters</th>
<th>( C_{66} )</th>
<th>( \rho )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Values</td>
<td>0.437</td>
<td>1.57</td>
</tr>
</tbody>
</table>

Table 1 Second order elastic constants (in \( 10^{11} \) dyn/cm\(^2\)) and density (in g/cm\(^3\)).

![Graph showing backscattering responses at the incident angle \( \alpha = 5.0 \): dashed line, \( C_{33} = 10.0C_{11} \); Solid line, \( C_{33} = C_{11} \).](image)

Fig.2 Backscattering responses at the incident angle \( \alpha = 5.0 \): dashed line, \( C_{33} = 10.0C_{11} \); Solid line, \( C_{33} = C_{11} \).

The backscattering responses and scattering patterns are plotted in Figs. 2 and 3 for both \( C_{33} = C_{11} \) and \( C_{33} = 10.0C_{11} \). The spectrum were computed at an interval of \( ka = 0.005 \) over the range from 0 to 8.0. Both the global pattern and the location of dips designated by arrows show significant differences between the curves of transverse and isotropic ones. These abrupt dips are caused by the modal resonances of the scatter as pointed out by Flax et al.\(^{8,9}\). The resonances of the cylinder occur at \( ka \) values at which \( \Omega_n = 0 (n = 1, 2, \ldots) \) and are proved to coincide with the dips.

![Graph showing scattering patterns at the incident angle \( \alpha = 5.0 \): dashed line, \( C_{33} = 10.0C_{11} \); Solid line, \( C_{33} = C_{11} \); \( ka = 5.0 \).](image)

Fig.3 Scattering patterns at the incident angle \( \alpha = 5.0 \): dashed line, \( C_{33} = 10.0C_{11} \); Solid line, \( C_{33} = C_{11} \); \( ka = 5.0 \).

4. Summary

A theoretical treatment of the sound scattering by a transverse isotropic cylinder is given. Numerical computations are performed for both the amplitude and phase change of the back-scattering spectrum in the region \( ka \leq 8.0 \). The sharp dips in the back-scattering spectrum are proven to be caused by the resonance of the scatter. The scattering pattern is also shown to vary sharply at the neighborhood of the resonance. Both the backscattering spectrum and the scattering pattern of a transverse isotropic cylinder were proved to differ from that of an isotropic cylinder at the oblique incidence of the sound wave.

Reference:

COHERENT AND INCOHERENT SCATTERING IN THE
ULTRASONIC INVESTIGATION OF INHOMOGENEOUS
MATERIALS

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INTRODUCTION

Many of the materials to which ultrasonic waves are applied
with an evaluative purpose are heterogeneous. They represent a
large fraction of the materials in our existence: concrete, geological
materials, foodstuffs, biological materials, human tissues and
particle or fibre reinforced composite materials. The tendency to
specialisation has probably militated against the development of
unified theoretical basis for the ultrasonic wave diagnostics of these
materials. However, there is a further source of potential
confusion, which is the inadequate development of a conceptual
structure for the wave propagation and measurement phenomena.
This contribution seeks to take some very first steps towards
improving this situation, and reducing the ambiguity of the
interpretation of our measurements.

ULTRASONIC PARAMETER MEASUREMENTS ON
BIOLOGICAL TISSUES

In the mid 1970s there was considerable concern about what
was perceived to be a potential "source of attenuation artefact" in
the measurement of attenuation in soft tissues (and other
inhomogeneous materials). Marcus and Carstensen [1] showed
briefly that if the wavefronts emerging from the specimen under
measurement were not planar, the voltage amplitude output from
the piezoelectric receiver would be artefactually low due to phase
cancellation of the wavefronts impinging upon the receiving
transducer surface. A potential solution to this problem was
identified by Busse, Miller and colleagues [2] using phase-
sensitive (cadmium sulphide) transducers. One of the results of
this has been the NCRP report [3] which states that "attenuation
coefficients obtained by phase-independent methods are the
preferred quantities to use in making predictions of intensity
reduction for ultrasound propagating in mammalian tissues".

There is certainly evidence [2,4] that the attenuation
coefficients of biological tissues measured with phase insensitive
receivers are lower than those measured with phase sensitive
receivers - and this atest to the non-planarity of the wavefronts
impinging on the receiving transducer. Nevertheless, the source of
this non-planarity remains essentially unknown. It may arise from
the diffractive nature of the sources used, from variations in
specimen thickness, or from scattering by the intrinsic
inhomogeneities in the tissue itself. If it is the last of these, it is
clear that, in principle, quantitative backscattering measurements
cannot reliably be made [5]. It is unfortunate that the only
systematic attempt to measure the fluctuations of the wavefronts
emerging from specimens of biological tissue, to identify their
source, proved inconclusive [6]. There is no doubt, though, that
scatterers can give rise to significant wavefront fluctuations when
other sources are removed [7].

INADEQUACY OF THE PLANE WAVE MODEL OF
ATTENUATION

The primary source of difficulties in this field appears to be
the usual attempt to straightjacket phenomena which involve spatial
wavefront variations, into the one-dimensional (plane) wave
equation which is used ubiquitousness to define attenuation. A more
flexible approach appears to be that which introduces the concepts
of coherent and incoherent scattering. Although the literature has
a number of references to these terms, it is not clear that the usage
of different authors is consistent. We adopt the approach of Foldy
[8] which is consistent with that of Morse and Ingard [9].

In discussing multiple scattering, Foldy introduces the
identity

$$\langle |\psi_0|^2 \rangle - \langle |\psi_\alpha|^2 \rangle = \langle |\psi_r|^2 \rangle = \langle \psi_0 \cdot \psi_\alpha \rangle$$

where $\psi_0$ is the incident wave, and $\psi_r$ the scattered wave. In
this he identifies the term $\langle \psi_\alpha \rangle$ as the coherent scattering, which adds
as a (plane wave) phase to the incident wave $\psi_0$. The last terms
on the right hand side are, together, identified as the incoherent
scattering.

It has been observed previously [10,11] that the term on the
left hand side of the above equation is the quantity that is measured
by a phase-insensitive receiver, while the first term on the right
hand side is the quantity that is measured with a phase-sensitive
(piezoelectric) receiver.

COHERENT, INCOHERENT AND TOTAL ATTENUATION

It would appear appropriate to move away from the concept
that piezoelectric receivers produce measurement "errors" and to
recognise, in the light of the above discussion, that they are actually
measuring a different physical quantity. Thus it is possible to
define a 'total' attenuation, measured with phase insensitive
devices, a 'coherent' attenuation measured with piezoelectric
devices, and the difference is the 'incoherent' attenuation.

Practical application of these concepts may be achieved in
one of two ways. Measurements can be made using both phase
sensitive and phase insensitive transducers. In this case, for reliable
results, it will be necessary to pay more attention to the detailed
calibration of receiving transducers than has generally been the case
to date. In particular the response of cadmium sulphide transducers
to non-planar wavefronts required careful evaluation.

Alternatively the wavefronts may be measured directly,
either by a scanned hydrophone [12,13] or by an optically scanned
hydrophone [14], and each of the terms in the scattering equation
evaluated. In this case, particular care needs to be taken with the
measurement of the phase of the ultrasonic waves, in order that the
wavefront disturbances are completely defined. This is a procedure
that requires high mechanical, thermal and electronic precision [12]
and a carefully calibrated hydrophone. Furthermore it should use a
hydrophone that is small compared to the wavelength. Whereas
this is relatively easy to achieve at low megahertz frequencies, it
becomes increasingly difficult as the frequency approaches 20
MHz. The use of a finite probe will introduce degradation of the
wavefront determination due to spatial averaging. The effects of
this may be minimised by deconvolving the probe response out of the
measured wavefront distributions, although for this it is
necessary to know the phase and amplitude directivity of the probe
[15]. It is vital, for correct interpretation of the results obtained,
to irradiate the specimen with waves that are closely planar. This can
only be achieved by detailed and direct measurement of the field
radiated by the transmitting transducer [16].
RESULTS AND CONCLUSIONS

The main data available in the literature appears to be that on the wavefronts emerging from suspensions of glass ballotini in silicone rubber [7] and from fresh beef liver [6]. As yet it is incompletely analysed, but the preliminary results are of some interest. For a 7.5% volume concentration of 500µm diameter glass spheres in silicone rubber, at 2 MHz, the ratio of coherent scattering intensity to incoherent scattering intensity appears to be as great as 3:2 [10]. For fresh beef liver, the incoherent attenuation appears to be only about 10% of the total [11].

It appears that this is an area of ultrasonic measurement science which deserves attention. In many circumstances, it will be found that the disruption of the wavefronts by the intrinsic material inhomogeneities is minimal, and an incident plane wave emerges as a plane wave. It is important, for unambiguous interpretation of the results to know that this is the case for a particular material and interrogating frequency.

If it is not the case, some detailed analysis will be needed, along the lines suggested here. The fluctuations can best be observed close to the scattering specimen, in the near field of the scattering volume. One of the major questions that remains for practical measurement procedures is identification of the highest spatial frequency that may be present in the fluctuations, and thus the size of finite receiver (if large piezoelectric or cadmium sulphide transducers are to be used) which will be large enough to give a statistically valid measurement. Some preliminary data analysis has been performed on the results from beef liver [11], but a more formal analysis is still required.

REFERENCES

THREE-PARAMETER BORN INVERSION OF ELASTIC WAVE SCATTERING IN A 3-D SOLID

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1. INTRODUCTION

Acoustic speed inversion under weak scattering has received a large amount of attention over the years. Recently, there has been an increasing interest in the inverse elastic wave scattering for the identification of objects defined by multiparameters in a 3-D elastic solid. The problem was investigated by Blackledge et al. for reconstructing 2-D configurations of the desity and the elastic parameters, and by Boyce and Keller for the case of 3-D elastic half-space with the source and geophone moving over the surface plane. More recently, Beylin and Zhu solved the problem both by applying generalized Radon transform but from different aspects. They showed it was sufficient to reconstruct the 3-D parameter configurations by using only 3-D observed data. However, in their methods, some proper approximations were made to implement the complicated inversion of the generalized Radon transform.

In this paper, a different Born inversion approach to inverse elastic scattering for density and two elastic parameters in three dimensions is presented. The inhomogeneous object is probed, from two orthogonal directions, by two sets of wideband plane waves, each contains P, SV, and SH waves. The scattered fields are measured over a closed surface of arbitrary shape surrounding the object. By calculating the plane-wave scattering amplitudes from the measured data and by choosing suitable polarizations of the incident plane waves, we uncouple the three unknown parameters in the Fourier transformation domain. Explicit solutions for the three parameters are obtained by combining the results obtained from the two illuminations in two directions, which, as x-ray tomography, has the form of filtered back-projection but with some weighting factors.

II. PROBLEM FORMULATION

The integral representation of the scattered elastic wave field within the Born approximation is given by (see Zhu)

\[ A' (K_s, k) = \int_\Gamma \left( u' (r) - \mu_0 (r) \nabla u' (r) - \mu_0 (r) \nabla (u' (r) \nabla) \right) dV, \]

where \( \rho_0, \lambda_0, \) and \( \mu_0 \) are the small fluctuations in the object \( V \) from the constant values of \( \rho_0, \lambda_0, \) and \( \mu_0 \) in the background medium, \( G(r,r') \) is the dyadic Green function, \( I \) is the unit dyadic, \( \nabla \) is the Laplacian, \( u' \) is the transverse of \( u_0 \), and where for simplicity the time factor \( \exp ( - i \omega t ) \) is omitted.

In the inverse problem, the scattered displacement and stress fields are measured over a closed surface \( S \) of arbitrary shape surrounding the object as shown in Fig. 1. Then, we introduce the plane-wave scattering amplitudes \( a \) from the observed data as follows:

\[ A' (K_s, k) = \int_\Gamma u' (r) - \sigma (u' (r) \cdot v' (r)) dS, \]

and

\[ \sigma (u) = \lambda_0 (V \cdot u) + \mu_0 (V u + u V), \]

\[ v' (r) = \frac{k}{\lambda_0 + 2 \mu_0} \exp (-i K_s \cdot r), \]

\[ K_s = \frac{\omega}{c_s}, \]

\[ K_s = \frac{t}{c_s}, \]

are plane \( P \) and \( S \) waves propagating in the \( k \) direction, where \( k, t \) are real unit vectors satisfying \( k \cdot t = 0 \). It can be shown that \( \sigma (u) \) represents the local-plane-wave amplitudes \( a \) of the scattered far-field in \( k (s = p) \) and \( t = (a, s) \) directions. The scalar counterpart of (2) was introduced by Devaney and Beylin for the single acoustic parameter inversion. Let

\[ [u_p', u_n', u_w'] = F (\omega \exp (- i K_s \cdot r), d_1 \exp (- i K_s \cdot r), d_2 \exp (- i K_s \cdot r) + r)] \]

be a set of incident waves in \( s \) direction with one longitude and two transverse waves, where \( (d_1, d_2, s) \) composes an orthogonal basis. For the incident \( P \) wave, we calculate the longitudinal scattering amplitude of the local-plane-wave with propagation vector \( k \), while for each of the two incident \( S \) waves, we calculate the transverse scattering amplitudes in \( t \) and \( d_2 \) directions, respectively, where \( (t, d_2, k) \) composes another orthogonal basis. Inserting the three incident fields of (4) into (1) respectively and then into (2), exchanging the order of volume and surface integrals, and applying vector Green’s theorem to the surface integral yields:

\[ A' (K_s, k) = - \frac{F (\omega)}{\lambda_0 + 2 \mu_0} \int_\Gamma \left( \int_\Gamma \left( \sigma (u') (r) \cdot (s - k) + \mu_0 (r) \cdot (s - k) \right) \exp (- i K_s \cdot r) dV, \]

\[ A' (K_s, k) = - \frac{F (\omega)}{\mu_0} \int_\Gamma \int_\Gamma \left( \sigma (u') (r) \cdot (s - k) + \mu_0 (r) \cdot (s - k) \right) \exp (- i K_s \cdot r) dV, \]

Clearly, for fixed \( \omega \) and \( k \), we obtain from above equations the linear combinations of the Fourier transforms of \( \rho_0, \lambda_0 \) and \( \mu_0 \) evaluated at \( K_s (k = s) \) for (5) and at \( K_s (k = s) \) for (6) respectively. If there existed three independent equations among (5) and (6) for all points in the wave-number space, the problem would be reduced to solving the equations for the Fourier transforms of the three parameters and taking the inverse transforms. Unfortunately, for some points, it is found that there are not as many as three independent equations in (5) and (6) and we have to seek a method to uncouple the parameters. This will be discussed in the next Sec.

III. RECONSTRUCTION OF THE THREE PARAMETER FUNCTIONS AND CONCLUSION

Let us first consider the coordinate transformation

\[ \xi = K_s (k = s), \]

\[ a = p_s, \]
It is easy to see that eq.17 builds one—one mapping between the 3-D Fourier transformed space $\zeta \in R^3$ and the continuous set of points $(K, \theta, \phi) \in (-\infty, \infty) \times (0, \pi) \times (0, 2\pi)$, where $\theta$ and $\phi$ are the spherical coordinates of $k$ as shown in Fig.2. The Jacobian of the transformation is found to be

$$\rho(k) = K_3^2 |(a-k)| \sin \theta.$$  

(8)

Next, we define the following 2 × 2 matrices:

$$[C] = \begin{bmatrix} (u \cdot d_1) & (u \cdot d_2) \\ (u \cdot d_3) & (a \cdot d_3) \end{bmatrix}$$  

(9a)

$$[E] = \left[\begin{array}{cc} d_1 & -d_1 \\ d_2 & -d_2 \end{array} \right]$$  

(9c)

$$[B] = \left[\begin{array}{cc} \mathcal{A}^{*}_{E} \mathcal{A}_{E} \end{array} \right]$$  

(9d)

Then, eq.7 can be rewritten as:

$$\left[\begin{array}{c} \tilde{\rho} \\ \tilde{\rho} \end{array}\right] = F^{-1}(\tilde{\rho}) \left(\begin{array}{c} \hat{\rho} \\ \hat{\rho} \end{array}\right) = \begin{bmatrix} \mu_1 & \mu_2 \\ \mu_2 & \mu_3 \end{bmatrix} \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} + \begin{bmatrix} \mu_2 \\ \mu_3 \end{bmatrix} D_{\omega}$$  

(10)

where a tilde denotes a 3-D Fourier transform. Note that there is a freedom in choosing the unit vector $t$ after $k$ is given, so we set

$$t_1 = (a \cdot k) \times \mu |(a \times k)|, \quad t_2 = k \times t_1, \quad d_2 = a \times d_1.$$  

(11)

Solving (10) for $\tilde{\rho}_1$ and $\tilde{\rho}_2$, we have

$$F^{-1}(\tilde{\rho}) = \begin{bmatrix} \mu_1 & \mu_2 \\ \mu_2 & \mu_3 \end{bmatrix} \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} = \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{11} = \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{22}.$$  

(12a)

$$F^{-1}(\tilde{\rho}) = \begin{bmatrix} \mu_1 & \mu_2 \\ \mu_2 & \mu_3 \end{bmatrix} \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} = \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{22}.$$  

(12b)

Clearly, $\tilde{\rho}_1$ and $\tilde{\rho}_2$ become uncertainty at $k = \pm s$, so we cannot determine the scattering functions by using only one incident direction. The requested data at $k = \pm s$ can be obtained by probing the scatter from another direction. Note from (7) that for different incident directions $s$, $\zeta$ represents different points in the wave—number space. Thus, to combine eqs.(12) for different $s$ at the same point in $\zeta$—space, we introduce a new variable $\eta$ by multiplying (12a,b) with $d(\zeta - \eta)/F(\omega)$ and integrate the result over the whole $\zeta$—space

$$\left[\begin{array}{c} \tilde{\rho}_1 \\ \tilde{\rho}_2 \end{array}\right] = \frac{1}{3} \int \left[\begin{array}{c} b_1 \\ b_2 \end{array}\right] d(\zeta - \eta)/F(\omega)$$  

(13)

where, $\zeta = 3-D$ Dirac delta, $\Omega = \sin \theta d \theta d \phi$, and where we have used (7) and following relation

$$K_3^2 [1 - (a \cdot k)] = K_3^2 [(a \cdot k)|s \times (a + k)| = K_3^2 |s - k|^2$$

(14)

In this way we can set $\eta$ in all equations like (13) for different $s$ at the same point. Now, choosing a second incident direction $s_2$ orthogonal to the first one $s_1$ and repeating the procedure of (5)–(13) we shall have a similar expression as (13). Taking inverse Fourier transform of the summation of these two expressions like (13) for the two incident directions, we finally obtain the reconstruction of $\mu_1$ and $\mu_2$

$$\left[\begin{array}{c} \tilde{\rho}_1 \\ \tilde{\rho}_2 \end{array}\right] = \frac{1}{3} \int \left[\begin{array}{c} b_1 \\ b_2 \end{array}\right] d(\zeta - \eta)/F(\omega)$$  

(15)

where

$$\begin{array}{c} b_1 \\ b_2 \end{array} = \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{22}$$  

(16a)

$$\begin{array}{c} b_1 \\ b_2 \end{array} = \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{22}$$  

(16b)

Then, from (12) and (5), the reconstruction of $k$ is given by

$$\tilde{\rho}_1 = \frac{1}{3} \int \left[\begin{array}{c} b_1 \\ b_2 \end{array}\right] d(\zeta - \eta)/F(\omega)$$

(17)

where

$$\begin{array}{c} b_1 \\ b_2 \end{array} = \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{22}$$

(18a)

$$\begin{array}{c} b_1 \\ b_2 \end{array} = \begin{bmatrix} \mu_1 \\ \mu_2 \end{bmatrix} + \frac{1}{2} \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} B_{11} + \frac{1}{2} \begin{bmatrix} 1 & -1 \\ 1 & 1 \end{bmatrix} B_{22}$$

(18b)

The results (15)–(18) have the form of filtered back—projection formula, except for the presence of the weighting factors $s \cdot (a-k)$ and $(1-|\eta| \cdot (s \times k))^2$.

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Fig.1 The plane wave illuminations in two orthogonal directions and the measuring surface.

Fig.2 The coordinate transformation.
FINITE ELEMENT--ARTIFICIAL BOUNDARY
METHOD FOR MANY-BODY SCATTERING
PROBLEM

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ABSTRACT

Scattering by many body is treated in near-field region with finite element--artificial transmitting boundary method, the solution in far-field region is computed with Kirchhoff's formula.

INTRODUCTION

The main feature of all many-body problems is that no one can solve them exactly. Some kind of approximation must be made to get an answer for any specific case.

Although it is much more complicated to deal with scattering by many bodies than scattering by single body, the scattering is an interesting subject and has been studied by many researchers with various methods, including analytic and numerical approaches.

Applications of analytic methods are restricted, because only a few scattering problems by regular shaped objects can be calculated with the methods. As developing of applied mathematics and electronic computer, numerical approaches are developed and are able to describe various shapes of boundaries, to treat layered media and inhomogeneity of medium, so that they are employed more and more widely in research and engineering.

Some numerical treatments, such as T-matrix method, boundary element method etc., have efficiently solved certain scattering problems. The applicability of these approaches, however, is confined by the complex computation processes.

In dealing with scattering in infinite space with a numerical treatment, the crux is to treat the radiation condition on the boundaries of a finite calculated region, that is to say, to take the boundaries of a finite discretised model as the infinite boundaries. Some processes to eliminate or decrease the effects of reflection from the finite boundaries have been put forward to overcome this difficulty. One of them is to use far enough boundaries. Consequently, large enough computer space and long enough computing time are required.

In this paper, multi--transmitting process is adopted to deal with artificial boundaries of a small, discretised region, so that the computer space and the computing time required are reduced significantly. The results show that fine transmitting effects on the artificial boundaries can be obtained. The idea of artificial transmitting boundary method is to model propagating of waves directly on artificial boundaries. Errors from this treatment are also waves propagating outward, therefore can be treated with the same process. This results in the proposal of the multi--transmitting process. Combining with finite element method, it can deal with many scattering problems in near-field region.

As an extension of our previous works[1]--[5], scattering by two circular cylinders is considered and computed with finite element--artificial transmitting boundary method.

DISCRETIZED MODEL, TRANSMITTING FORMULA AND FINITE ELEMENT EQUATION

The near--field region is discretised as shown in Fig.1. It is divided into boundary area and inhomogeneous area. The boundary area is the part between rectangles ABCD and EFGH, while the inhomogeneous area is the part inside the rectangle EFGH.

![Discretized model for scattering by two circular cylinders](image)

Fig.1 Discretized model for scattering by two circular cylinders

Let $\Psi(\tau)$ denote the potential without time dependence $\exp(-i\omega t)$. For the nodes on the boundary, the static transmitting formula is [6]:

$$\Psi = \sum_{i=1}^{N} (-1)^{i+1} C_{i} \Psi_{i-1} \exp(-i\omega c x / c) \tag{1}$$

where $N$ is transmitting order, $C_{i}$ is the binomial coefficient and $c$ is the acoustic speed in the water. For the internal nodes, the finite element equation is

$$-\omega^{2} [M][\Psi] + [K][\Psi] = [f] \tag{2}$$

where $[M]$ and $[K]$ are the mass and stiffness matrices, respectively. Solving the coupled equations (1) and (2), the numerical solutions for the discretised area can be obtained.

NUMERICAL RESULTS

We calculated the scattering field by soft, rigid and elastic objects for incident wave of various directions. When $k_0 = 3 \cdot 77 \cdot \omega / c = 0.5$, the real parts $Re(\Psi(\tau))$ of whole field of scattered wave are shown in Fig.2.
Employing these numerical solutions and Kirchhoff’s formula

$$\Psi(\bar{r}) = \left[ \Psi(\bar{r}_1) \frac{\partial G(\bar{r}/\bar{r}_1)}{\partial n} - \frac{\partial \Psi(\bar{r}_1)}{\partial n} G(\bar{r}/\bar{r}_1) \right] ds (3)$$

where L is an arbitrary integral path encircling the two cylinders, and $G(\bar{r}/\bar{r}_1)$ is the Green’s function of two dimensions, the solutions in far-field region and the extinction cross section etc. can be obtained. Fig. 3 shows the scattering pattern.

![Scattering patterns](image)

Fig. 3 Normalized far-field pressure-amplitude pattern for elastic cylinder Fig. 2(d).

CONCLUSIONS

In this paper a new method has been employed to deal with the scattering of an acoustic plane wave by two parallel cylinders. The new method has many features such as having clear physical views, simple processes, saving computing time and space etc. It is expected that this new method can be applied more and more widely.

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I. INTRODUCTION

The Singularity Expansion Method (SEM) has been developed and extensively used in the radar and sound scattering problems. The essence of the SEM is to expand the scattered field as a residual series of the complex frequency poles. In a previous article we have shown that the Resonance Scattering Theory (RST) can be derived from the SEM (Ref.1). The Generalized Singularity Expansion Method (GSEM) introduced in this paper will expand the scattered field of a solid elastic cylinder into a residual series of the complex wavenumber poles. After carrying out the Sommerfeld-Watson Transformation (SWT), each term of the series represents a contribution of the corresponding surface wave type. Hence the essence of the GSEM is to represent the scattered field by a sum of the contributions of various surface wave types.

II. GSEM AND SURFACE WAVE SCATTERING FORMULA

The total scattered field of a plane acoustic wave from a solid elastic cylinder can be separated into a rigid background $\rho^t$ and a pure elastic scattered field $\rho^s$:

$$\rho(r,\theta;x) = \rho^t(r,\theta;x) + \rho^s(r,\theta;x)$$

(1)

where

$$\rho^s(r,\theta;x) = \sum_{n=1}^{\infty} \sum_{m=1}^{2} a_n^m r^n e^{-j\alpha_n x} H_n^m(\alpha_n r) \cos \theta$$

(2)

$$\rho^t(r,\theta;x) = \sum_{n=1}^{\infty} \sum_{m=1}^{2} a_n^m r^n e^{-j\alpha_n x} H_n^m(\alpha_n r) \cos \theta$$

(3)

and all parameters are given in Ref.2. The function $S_n(x)$ is related with the frequency characteristic function $D_n(x)$ by

$$D_n(x) = -x^2 x H_n^{m+1}(x)D_n^m(x) S_n(x)$$

(4)

The real and imaginary of the function $S_n(x)$ are respectively

$$R_n(x) = \frac{D_n^{(1)}(x)}{D_n(x)} J_n(x) J_0(x) N_n(x) N_0(x)$$

(5)

$$I_n(x) = \frac{\alpha_n}{\text{tan}^{-1}[J_n(x)N_n(x)]}$$

(6)

$$\alpha_n = \tan^{-1}[J_n(x)N_n(x)]$$

(7)

Now the mode number $n$ is considered a variable in the complex $\nu$ plane. The function $I_n/S_n$ will have poles $\nu$. Using the Mittag-Leffler theorem the function $I_n/S_n$ is then expanded as

$$I_n(x) = \sum_{\nu} \frac{a_n}{\nu-n}$$

(8)

where $\nu$ satisfies $S_n(x)=0$ and $\nu$ is in the residue at the pole $\nu$. It is seen from Eq.(4) that $\nu$ are the wavenumbers of the pure elastic waves $\nu(x)\pm\xi(x)+\kappa(x)$, $\nu=1,2,\cdots$ (11)

in which $x$ is a real parameter. The dominant contributions to Eq.(8) come from those terms in which $n$ is near $\nu$. Neglecting the second term in Eq.(8) and substituting it into Eq.(3), an approximate expansion is obtained as

$$\rho^s(r,\theta;x) = \sum_{n=1}^{\infty} \sum_{m=1}^{2} a_n^m r^n e^{-j\alpha_n x} H_n^m(\alpha_n r) \cos \theta$$

(12)

For most solid objects immersed in water the liquid loading may be considered as light one. In this case, as in Ref.[3], it can be proved that $a_n=0$.

This shows that for the light liquid loading the radiation efficiency of the $J$-type surface wave can be represented by the imaginary of its pole. Using Eq.(13), the generalized singularity expansion (12) becomes

$$\rho^s(r,\theta;x) = \sum_{n=1}^{\infty} \sum_{m=1}^{2} a_n^m r^n e^{-j\alpha_n x} H_n^m(\alpha_n r) \cos \theta$$

(14)

III. THE CIRCUMFERENTIAL WAVE ANALYSIS

Eq.(12) can be transformed into a contour integral around a contour $c$ in the complex $\nu$ plane by SWT (Ref.[4])

$$\int_{c} d\nu \frac{\exp(i\nu x)}{\sin \nu v}$$

(15)

where the contour $c$ is shown in Fig.1. It is then deformed into a contour $c$.

![Fig.1](image1)

![Fig.2](image2)

The Eq.(12) can be transformed into

$$\rho^s(r,\theta;x) = \int_{\gamma} \frac{d\nu}{\sin \nu v} e^{i\nu x} \frac{e^{-i\alpha_n \nu}}{\nu - \nu}$$

(16)

Here only the elastic type poles lain in the first quadrant are involved. Using the identical relation

$$\cos \nu = \cos \nu (\nu - \nu) \exp(i\nu x) - \sin \nu \nu \exp(i\nu (\nu - \nu))$$

(17)

Eq.(16) can be split into two components

$$\rho^s(r,\theta;x) = \int_{\gamma} \cos \nu \frac{e^{-i\alpha_n \nu}}{\nu - \nu} d\nu$$

(18)

$$\rho^s(r,\theta;x) = \int_{\gamma} \frac{e^{-i\alpha_n \nu}}{\sin \nu v} d\nu$$

(19)

The first component, Eq.(18), is always convergent for each pole lain in the first quadrant. So the contour $c$ can be enclosed along a semicircle of infinitesimal radius in the upper half of the $\nu$ plane. Using the residue theorem gives

$$\rho^s(r,\theta;x) = 2\pi i \sum_{\nu} \frac{\cos \nu (\nu - \nu)}{\sin \nu v} e^{i\nu x} \frac{e^{-i\alpha_n \nu}}{\nu - \nu}$$

(20)
In the far field, taking the asymptotic form of the Hankel function and the expansion of the trigonometric function, the circumferential waves are then given by

\[
p_0(r, \theta, x) = \frac{2}{nkr} e^{i(kr-x)} \sum_{l=0}^{\infty} \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i}
\]  

(21)

where

\[
\phi_i = \alpha_i \theta_i, \quad \alpha_i = \alpha - k\theta_i
\]

(22)

Furthermore, using the Debye asymptotic form, Eq. (7) gives

\[
\alpha_i \approx - \frac{x \sin \gamma - y \cos \gamma}{n/4}
\]

(23)

where

\[
\gamma = \cos \theta
\]

(24)

If letting \( \gamma_i = \gamma_i \times \theta_i \), and assuming \( \gamma \approx \gamma_i \), Eq. (24) gives

\[
\theta_i = \cos \gamma_i, \quad \phi_i = \frac{x \sin \gamma_i}{n/4}
\]

(25)

Substituting Eqs. (23) and (24) into Eq. (20) and taking \( \alpha_i = \alpha \), the explicit expressions of the amplitudes and phases of the circumferential waves are given by

\[
p_0(r, \theta, x) = \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i}
\]

(26)

where

\[
\phi_i = \alpha_i \theta_i, \quad \alpha_i = -2k\pi \frac{l}{n} + \frac{1}{2}
\]

(27)

\[
\phi_i = \frac{x \sin \gamma_i}{n/4} + \frac{1}{2}
\]

(28)

Fig. 2 shows the geometrical meaning of the zero order (p=0) circumferential wave when \( \lambda = \gamma \). The wave is excited by an incident wave at the lower section of the cylinder. The total phase variation of the course from \( N \) to \( S \), with respect to the center of the cylinder, is

\[
\phi_{i-1} = \phi_{i-1} (2\pi + \phi_i) - 2x \sin \gamma_i
\]

(29)

Thus, the phase given by Eq. (28) can be rewritten as

\[
\phi_{i-1} = \frac{x \sin \gamma_i}{n/4} + \frac{1}{2}
\]

(30)

It must be emphasized that a constant phase \( Sn/4 \) is involved in the phase. Fig. 3 shows the propagating process of the zero-order circumferential wave when \( \lambda = \gamma \). It circumnavigates in the clockwise direction. When \( \theta < 2\gamma_i \), the wave circumnavigates along the surface over one circumference. When \( \theta > 2\gamma_i \), the wave circumnavigates less than one circumference.

\[
\theta < 2\gamma_i
\]

Fig. 3

\[
\theta > 2\gamma_i
\]

IV. THE GEOMETRICAL SCATTERING WAVES

The second component of the scattering field, Eq. (19), can be discussed for two situations:

(1) When \( \theta < 2\gamma_i \), the residual method is still valid because the integral is convergent. The Eq. (19) is then given by

\[
p_1(r, \theta, x) = \frac{2n}{nkr} e^{i(kr-x)} \sum_{l=0}^{\infty} \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i} H_{l+1}(kr)
\]

(31)

Taking some approximations gives

\[
p_1(r, \theta, x) = \frac{2n}{nkr} e^{i(kr-x)} \sum_{l=0}^{\infty} \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i}
\]

(32)

Adding this wave to the Eq. (20), the formalism of the circumferential waves is not varied, provided \( \theta > \gamma \).

(2) When \( \theta > 2\gamma_i \), the Eqs. (31) and (32) are not valid. In this case, the contour \( c \) can not be enclosed with a semi-circle of infinite radius. But if \( x > 1 \), the Eq. (19) may be evaluated by the steepest descent procedure. In the far field, the integral is

\[
p_1(r, \theta, x) = \frac{2n}{nkr} e^{i(kr-x)} \sum_{l=0}^{\infty} \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i}
\]

(33)

where

\[
A_l(v) = \frac{1}{v} \left( \frac{\sqrt{v^2 - 4 \cos \gamma_i^2}}{2} \right)
\]

(34)

\[
W_l(v) = \frac{\sin \gamma_i}{\sqrt{v^2 - 4 \cos \gamma_i^2}}
\]

(35)

It is easy to find that the saddle point \( v \) satisfies

\[
2v = \sqrt{v^2 - 4 \cos \gamma_i^2}
\]

(36)

Performing the steepest descent procedure gives

\[
p_1(r, \theta, x) = 2n \left( \frac{v}{nkr} \right) e^{i(kr-x)} \sum_{l=0}^{\infty} \frac{2n}{l+1} \sin \phi_i e^{i(l+1) \theta_i}
\]

(37)

where

\[
f_l(v) = \left( \frac{\sqrt{v^2 - 4 \cos \gamma_i^2}}{2} \right)
\]

(38)

This is a geometrical scattering wave. The wave has incident and reradiating angle \( \alpha_{i-1} = (\pi - \Delta) / 2 \) and circumnavigates along the surface through an angle \( \Delta \) as shown by solid line in Fig. 4. By comparison, the incident and reflected angle of the specular reflected wave are \( a_i+ (\pi - \Delta) / 2 \), as shown by dotted line in Fig. 4.

In fact, the steepest descent procedure can also be used to the case of \( \theta < 2\gamma_i \). So the second component of the pure elastic scattering wave can be represented by an unified Eq. (37) if the restriction \( \theta < 2\gamma_i \) is canceled.

Fig. 4

REFERENCES

MID-FREQUENCY ENHANCEMENT OF THE BACKSCATTERING OF TONE BURSTS BY THIN SPHERICAL SHELLS: RAY APPROXIMATION AND EXACT FOURIER SYNTHESIS

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This paper concerns a large contribution to the backscattering by thin spherical shells in water [1,2]. The contribution is caused by a subsonic guided wave described as a Snell's wave by some authors and designated here by \( l = a_0 \) to distinguish it from the wave designated by \( a_0 \) that becomes a leaky Lamb wave at high frequencies. Let the \( l \)th classification of a guided wave have a phase velocity \( c_l \) along the outer surface of the shell. When \( c_l \) exceeds the speed of sound \( c \) in the surrounding water, the wave is supersonic and usually leaks or radiates sound. To facilitate the extension of ray methods [3,4] to thin shells, it is necessary to include subsonic guided waves where \( c_l < c \). Figure 1 shows the modified ray picture [1,5] where the coupling is through an evanescent region having a thickness \( (b - a) \) where \( b = a_0/c_l \).

The contribution of a subsonic guided wave to the backscattering by thin spherical shells is especially significant when \( kh = 1 \) where \( h = a - b \) is the shell thickness and \( k = \omega c \) so that \( ka = a_0h \).

\[
\textbf{Figure 1} \text{ Ray diagram for a guided wave whose phase velocity along the outer surface of the shell is subsonic with respect to the surroundings.}
\]

Consider the example of an empty stainless steel 304 shell with \( a_0h = 40 \). The relevant guided wave properties are computed (based on the Watson transform) by finding the complex \( r \) roots of \( D_m(la) = 0 \) where \( D_m \) is the denominator of the \( m \)th term of the exact partial wave series for the form function [4]. The resulting properties are shown in Fig. 2. The wave labeled with long and short dashes, designated by \( l = a_0 \), is subsonic throughout the region of interest but displays a noticeable rise in the radiation damping parameter \( \beta_l \) as \( c_l \) approaches \( c \). This rise in coupling appears to be associated with the decrease in the thickness of the evanescent region near coincidence [1] and causes an enhanced backscattering. The contribution of the \( m \)th circumnavigation of the \( a_0 \) wave to the form function \( f \) has the magnitude

\[
|m_l| = 8\pi \beta_i (c/c_l) \exp(-\pi \beta_i - 2\pi m \beta_i), \tag{1}
\]

where for \( m = 0 \), the wave has traveled only around the backside of the sphere as shown in Fig. 1. The solid curve in Fig. 3 gives \( |m_l| \) for \( l = a_0 \) while the dashed curve omitted the factor \( c/c_l \) : \( |m_{0l}| \) for the \( a_0 \) wave shows a strong enhancement that peaks near \( ka = 46.3 \) where the value of \( \beta_{0l} = 3.17 \). It may be shown [1] that the value of \( \beta_{0l} \) at such a peak depends only weakly on the shell thickness and material parameters since the \( ka \) dependence is dominated by that of \( \beta_l \). The peak \( \beta_{0l} \) occurs close to where \( \beta_l = 1/\pi \).

Equation (1) is directly applicable to the calculation of the enhanced backscattering by tone bursts that are sufficiently short that the echoes associated with different values of the circumnavigation index \( m \) do not overlap. The incident burst must be sufficiently long that the effects of dispersion are weak. The time domain echoes for such tone bursts were calculated by a Fourier synthesis that used the exact partial wave series [1,2]. This was done for several bursts of different \( ka \). Figure 4 shows the resulting calculated amplitude for a 20 cycle burst with \( ka = 46 \). The dimensionless time units are \( T = ct/a \). The normalization is such that specular reflection by a fixed rigid sphere of radius \( a \) has a unit amplitude. The earliest contribution from the shell in Fig. 4 is a specular reflection of close to unit amplitude. This is followed by the \( m = 0 \) and \( m = 1 \) contributions of the \( l = a_0 \) wave as confirmed by comparison of their arrival times with ray theory. The amplitude of the \( m = 0 \), \( l = a_0 \) echo was similarly determined for several other values of \( ka \) and are plotted as the points in Fig. 3.

The ray model was also used to synthesize the exact form function \( f \) for steady-state backscattering in the region of the mid-frequency enhancement for the spherical shell considered in Fig. 2. The dashed curve in Fig. 5 shows the lift from the exact partial wave series and the solid curve shows \( |m_{0l}| \) where
The solid curve gives the ray model in Eq. (1) for the \( a_0 \)-guided wave echo amplitude in the region of the enhancement near \( ka = 46 \). The points are from the exact echo amplitude determined by a Fourier synthesis as illustrated in Fig. 4.

![Figure 4](image4.png)

Figure 4—Farfield scattered normalized pressure \( P(T) \) calculated by Fourier synthesis from the exact of partial-wave series for 20 cycle incident burst with a carrier \( ka \) of 46. The large second echo is enhanced as predicted in Fig. 3 for the empty 2.5% thick shell considered.

\[
f_{\text{ray}} = f_{\text{sp}} + f_{a_0} + f_{b_0} + f_{s_0}
\]

(2)

where \( f_{\text{sp}} \) is the specular contribution (here neglecting a small curvature correction) and \( f_{l} \) for \( l = a_0 \) and \( b_0 \) are leaky wave contributions computed as discussed previously [4]. In the region shown if \( l \) is negligible for \( l = a_0 \) but for \( l = a_0 \) it gives rise to the narrow resonance spikes. The enhancement is due to the \( l = a_0 \) term which is computed by summing ray terms of the form \( f_{ml} \) as in Eq. (1) including phase information. The result reduces to

\[
f_{l} = C_l \exp(\eta_l \sqrt{\pi \gamma_l}) \left( 1 + \exp(-2\pi \gamma_l) \right),
\]

(3)

where \( C_l = 8\pi \gamma_l (c/c_l) \exp(\phi_l) \), \( \gamma_l = xc/c_l - (\pi/2) \), \( x = ka \), and except for the expression for \( \phi_l \), the generalization to this case of a subsonic wave follows from the ray geometry in Fig. 1 and related considerations [1,3]. The approximation for \( \phi_l \) used in Eq. 5 was obtained by taking \( \phi_l \) to vary linearly with \( c/c_l \) as \( c_l \) approaches \( c \),

\[
\phi_l \quad (Rc < c_l < c) = F \left( \frac{3}{2} \right) - \left( \frac{3F}{5} \right)
\]

(4)

where \( F = [(c/c_l) \cdot R/(1 - R)] \), \( \phi_l \) becomes constant for \( c/c_l \) below a particular value of \( R < 1 \) so that \( \phi(c_l < Rc) = 3\pi/2 \). Note that \( F \) varies from 0 to 1. Results of Ho and Felsen [5] were used as a guide for formulating the limiting values of \( \phi_l \) although the value of \( F = 0.91 \) used for the synthesis in Fig. 6 was determined empirically. The comparisons in Figs. 3 and 6 confirm the general usefulness of the ray picture.

Figure 5—The dashed curve gives the exact \( f_{l} \) for steady-state backscattering by the same shell. The solid curve gives Eq. (2).

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BULK MODULUS OF POROUS MATERIALS FROM Biot’s THEORY AND CONTINUITY EQUATIONS

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LIST OF SYMBOLS

A: Biot’s elastic coefficient
b: viscous drag coefficient
C: Biot’s elastic coefficient
E: Transverse Young modulus of the solid frame
E': Longitudinal Young modulus of the porous solid
M: Biot’s elastic coefficient
N: Biot’s elastic coefficient
P: Frame pressure
Pf: Fluid pressure
Q: Biot’s elastic coefficient
R: Biot’s elastic coefficient
u: average displacement per unit volume in the solid
U: average displacement per unit volume in the fluid
\( \epsilon_{eff} \): effective strain
\( \phi \): porosity
\( \sigma \): stress tensor of the porous material
\( \sigma' \): stress tensor for the fluid in the porous solid
\( \nu \): Poisson coefficient of the solid frame

INTRODUCTION

Biot’s theory allows, in a general and rigorous way, a description of the sound propagation in porous materials. When this theory has been applied to experimental cases, one of the difficulties is to characterize mechanically the solid frame filled with a fluid. In some cases, the material can be considered as isotropic, and then the classical expressions for P, Q, and R (Biot’s elastic coefficients) can be used [1,2]. Under the experimental point of view, an important type of elastic material symmetry is the corresponding to the transverse isotropy. This kind of symmetry can be found in seabeeds, sedimentary rocks, sandstones, fabrics, fabric composites, etc... For these cases, there are eight elastic coefficients, but no analytical expressions for them.

To have analytical expressions for the Biot’s elastic coefficients as a function of the material and fluid properties is important in order to know and quantify the mechanical coupling between the phases. The influence of solid and fluid properties on the behaviour of the porous aggregate, and to be able to include viscoelastic effects. Now it will be applied the same theoretical procedure as in [2] to determine these elastic coefficients.

THEORETICAL PROCEDURE

As it was suggested by Depouillé [3], Biot’s theory provides the conditions of validity of the continuity equations traditionally used for predicting the acoustical properties of porous sound-absorbing materials. Now we will write the continuity equations for the material, allowing to obtain the requirements to be fulfilled by the Biot’s elastic coefficients.

The Biot’s equations are:

\[
\sigma_{ij,j} + \frac{\partial^2 \sigma_{ij}}{\partial t^2} \left( \rho_{11} u_{ij} \rho_{11} U_{ij} \right) - b \frac{\partial \sigma_{ij}}{\partial t} \left( U_{ij} - u_{ij} \right) \tag{1a}
\]

\[
\sigma_{ij} = \frac{\partial^2 \sigma_{ij}}{\partial t^2} \left( \rho_{12} u_{ij} \rho_{12} U_{ij} \right) - b \frac{\partial \sigma_{ij}}{\partial t} \left( U_{ij} - u_{ij} \right) \tag{1b}
\]

where \( u \) and \( U \) are the average displacements per unit volume in the solid and the fluid parts of the porous aggregate, \( \sigma \), the stress tensor and \( b \) the viscous loss coefficient. The dynamical coefficients \( \rho_{11}, \rho_{12}, \rho_{21}, \rho_{22} \) satisfy the following relations:

\[
\rho_{11}, \rho_{12} = (1-\Phi) \rho_s, \rho_{21}, \rho_{22} = \rho_f \tag{2}
\]

being \( \rho_s \) and \( \rho_f \) the density of the solid and liquid respectively and \( \Phi \) the porosity of the aggregate. Introducing in eq. (1a) and (1b) the stress-strain relations according to Biot for the transversely isotropic liquid-saturated porous solid case, the equation of motion can be derived [4,5].

Similar equations of motion can be obtained in a different way. This is the case in which, the two phases are considered separately. The procedure consists of developing the equation of motion from the continuity, the constitutive, and the force balance equations for the solid and the fluid. The obtained equations are [5]:

\[
(1-\Phi) \frac{\partial \sigma_{ij}}{\partial t} = (1-\Phi) \rho_{11} u_{ij} + \Phi \rho_f (k-1) \delta_{ij} \tag{3a}
\]

\[
(1-\Phi) \frac{\partial \sigma_{ij}}{\partial t} = (1-\Phi) \rho_{12} u_{ij} + \Phi \rho_f (k-1) \delta_{ij} \tag{3b}
\]

here \( k \) is the structure factor and \( \rho_s \), \( \rho_f \) the frame’s and fluid’s pressure respectively. These parameters have to be calculated for a transversely-isotropic media.

Taking into account the stress-strain relations for the solid [7], the constitutive equation, \( P_s = V_s \Delta \rho_s \), where

\[
\Delta \rho_s = \rho_s \left[ \delta_{ij} \frac{\partial u_{ij}}{\partial t} + \sum \frac{\partial u_{ij}}{\partial t} \right] - (1-\Phi) \left( \sum \frac{\partial u_{ij}}{\partial t} - \epsilon_{ij} \right) \tag{4}
\]

are the strain and \( V_s \) the volume of the frame, \( V_s \) the velocity of sound in the fluid and, considering a plane wave propagation along \( x \)-axis (no solid displacements in \( y \) and \( z \) directions), i.e.,

\[
\sigma_{xx} = -P_s \tag{5a}
\]

\[
\sigma_{yy} = -P_s \tag{5b}
\]

\[
\sigma_{yy} = -P_s + V_{E_x} \frac{\partial E_x}{\partial x} + \frac{\partial E_x}{\partial t} \tag{5b}
\]

\[
E_y = \frac{E_x}{V_{E_x}} \tag{5b}
\]

\[
E_y = \frac{E_x}{V_{E_x}} \tag{5b}
\]
\[ a_{xx} = -P_2 + P_1 \frac{E_z v_{zz} (v_{zz} + 1)}{E_z E_y v_{zz}} \quad (5c) \]

the expression for \( P_2 \) and \( P_1 \) may be obtained. In eq.(5) the minus sign on \( P_z \) and \( P_x \) consider the usual sign convention for stress (compression is negative). So we have

\[ P_2 = \phi - \frac{\frac{\partial u_x}{\partial x} - \frac{\partial u_y}{\partial y}}{W_1} - \frac{1}{\phi} \left( \frac{\partial u_z}{\partial z} - \frac{W_2}{W_1} \frac{\partial u_z}{\partial x} \right) \]

\[ \frac{v_{xx} + v_{yy}}{E_x} \frac{1}{W_1} \quad (6) \]

\[ P_1 = \phi - \frac{\frac{\partial u_x}{\partial x} - \frac{\partial u_y}{\partial y}}{W_1} - \frac{1}{\phi} \left( \frac{\partial u_z}{\partial z} - \frac{W_2}{W_1} \frac{\partial u_z}{\partial x} \right) \]

\[ \frac{v_{xx} + v_{yy}}{E_x} \frac{1}{W_1} \quad (7) \]

being:

\[ W_1 = \frac{1}{E_x} - \frac{D_{NN} T}{E_z E_x} + \frac{N_{NN} T}{E_z E_x} \quad (8) \]

where D and B are the coefficients of \( P_z \) in eq.(5b) and (5c) respectively.

\[ W_2 = \frac{-1 + 2D - v_{xx} (D - 1) - v_{xx} (2B + D - 1)}{E_x} + \frac{B}{E_z} \quad (9) \]

\[ W_3 = \frac{1}{E_x} - \frac{1 + v_{xx} (3v_{xx} - 1)}{E_z} \quad (10) \]

\[ \frac{v_{xx} T}{E_x} - \frac{1}{E_z} \quad (11) \]

The values of \( 2N + A, M \) and \( R \) are derived equating the left hand sides of equations (3) and (1):

\[ 2N + A = \phi - \frac{1}{W_1} \left( 1 + W_2 \phi \right) \frac{v_{xx} + v_{yy}}{E_x} \quad (12) \]

\[ M = (1 - \phi) W_1 \frac{v_{xx} + v_{yy}}{E_x} \quad (13) \]

\[ R = \phi W_1 \quad (14) \]

With the same procedure and for a propagation along z-axis the Biot’s elastic coefficients \( C \) and \( Q \) can be deduced.

The rest of the coefficients are calculated with the same procedure, but allowing lateral displacements, being for this proposal necessary to modify equations (5).

From these equations it can be seen that \( W_1 \) and \( 1/W_1 \) play the role of a modified fluid and solid bulk modulus respectively.

**NUMERICAL RESULTS**

For numerical evaluation a Nylon fabric with a porosity \( \phi = 0.6 \) has been considered. From the averaging method for composites \( \phi \) we obtain:

\[ E_1 = 3.0 \times 10^6 \text{ N/m} \]

\[ E_2 = 3.5 \times 10^6 \text{ N/m} \]

\[ v_{xx} = 0.47 \]

\[ v_{yy} = 0.3 \]

If the fluid was water (i.e. \( \rho \) = 1000 Kg/m\(^3\) and \( V_s = 1500 \text{ m/s} \), then:

\[ 2N + A = 1.216 \times 10^8 \text{ N/m} \]

\[ M = 7.88 \times 10^8 \text{ N/m} \]

\[ R = 1.14 \times 10^8 \text{ N/m} \]

\[ C = 7.84 \times 10^8 \text{ N/m} \]

\[ Q = 9.61 \times 10^8 \text{ N/m} \]

For the isotropic case \( 2N + A = C \) and \( M = Q \). In this case the anisotropic behaviour is more important in the solid frame \( 2N + A = C \) than in the mechanical coupling between the two phases \( M, Q \).

**CONCLUSIONS**

In this paper a theoretical procedure to obtain the Biot’s elastic coefficients of a transversely isotropic porous material has been developed. The procedure is based on a comparison between the equation of motion obtained by Biot and by the continuity, the constitutive and the balance equations. Analytical expressions are deduced for the elastic coefficients. Finally the bulk modulus for a porous aggregate as a Nylon fabric submerged into water has been evaluated.

**ACKNOWLEDGMENTS**

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**REFERENCES**


CHARACTERIZATION OF AN IMM ERED POROUS LAYER BY ULTRASONIC WAVES

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One of the frequency spectrum analysis of ultrasonic waves is made for non-destructive evaluation of porous-elastic materials. Characteristic behavior of reflecting power from a layer is an important clue to evaluate the fluid-filled porous materials. Referring to the former analysis[1], we obtain the closed form solution of the reflecting power intensity for normal incidence. Numerical examples, calculated by the solution, reveal the sharp spikes on the spectrum curves as a result of the characteristic nature of fluid-filled porous materials.

Many investigators in the field of geophysics have been studied on the elastic wave propagation in a fluid-filled porous media. Among them, it is interesting that the experimental examination by Poiria[2] of slow longitudinal waves of which slow phase velocity is special to the fluid-filled porous materials. His experimental result has verified the theoretical predictions of the slow wave by Biot[3]. Target for the present analysis within the Biot’s framework, is porous material characterization by the frequency spectrum of the reflection and transmission power.

BASIC RELATIONS

The equations of motion for fluid-filled porous materials are:

\[ N \nabla^2 \mathbf{u} + \text{grad} \left( (D + N) \nabla \cdot \mathbf{u} + Q \nabla \cdot \mathbf{U} \right) = \rho_1 \nabla^2 \mathbf{u} + \rho_2 \nabla^2 \mathbf{u} + \rho \nabla \cdot \mathbf{u} \quad (\hat{\mathbf{u}} - \mathbf{U}) \]  

\[ \text{grad} \left( Q \nabla \cdot \mathbf{u} - R \nabla \cdot \mathbf{U} \right) = \rho_2 \nabla^2 \mathbf{u} + \rho_2 \nabla \cdot \mathbf{U} \quad (\hat{\mathbf{u}} - \mathbf{U}) \]  

where \( \mathbf{u} \) and \( \mathbf{U} \) are respectively displacement vectors of frame solid and porous. \( D, N, Q, R \) are elastic constants for the material\( \alpha \). \( \rho \) is effective mass density and depends on the porosity and tortuosity. For example, one of the model[5] is

\[ \rho_1 = \rho + \rho' \delta (\alpha - 1) \]

\[ \rho_2 = \rho' \delta (\alpha - 1) \rho_1 \]

where \( \rho \) and \( \rho' \) stand for mass density of solid and fluid respectively. \( \rho \) is total average mass density. \( \delta \) is called tortuosity that represents the liquid flow resistance in a pore solid. The total stress and the fluid stress are

\[ \sigma_{11} = (D + Q) \varepsilon_1 + 2N \sigma_{11} \]

\[ \sigma = -\rho' \varepsilon + R \varepsilon \]

in which \( \varepsilon \) and \( \varepsilon' \) are porosity and fluid pressure. Strains in Eq.(4,5) will be given by

\[ 2 \sigma_{11} = U_{11} + U_{11} \quad \varepsilon = U_{11} \quad \varepsilon = U_{11} \]

COEFFICIENTS OF REFLECTION AND TRANSMISSION

Displacement potentials

For the analysis of wave motion in a porous solid, we use the additional scalar potential \( \Phi \). This displacement potential can represent the second longitudinal wave of which velocity is slower than the ordinary longitudinal wave velocity. Displacement vectors may be written by

\[ \mathbf{u} = \left[ \nabla \phi + \nabla \rho \cdot \nabla \times \mathbf{U} \right] \exp \left( i \omega t \right) \]  

\[ \mathbf{U} = \left[ \rho_1 \nabla \phi + \rho_2 \nabla \rho \cdot \nabla \times \mathbf{U} \right] \exp \left( i \omega t \right) \]  

The characteristic parameters \( a_2 \) and \( a_0 \) are determined so that the vectors may satisfy the governing equation (1)(2). When a particle motion is restricted in the xy-plane, equations (1)(2) are reduced into

\[ \nabla^2 \phi_1 = 0 \quad j = 1, 2, 3 \quad \text{and} \quad \phi_1 \]  

with replacements:

\[ \phi = \phi_j, \quad \rho = \rho_j, \quad \phi_0 = 0 \]

\[ \phi_0 = \rho_1 \phi_1 - \rho_2 \phi_2 \]

\[ \rho = \rho_1, \quad \rho_1 = \rho_2 \]

\[ \phi_0 = \rho_2 \phi_2 + \rho_2 \phi_2 \]

\[ \phi_0 = \rho_2 \phi_2 + \rho_2 \phi_2 \]

\[ \delta \] represents wave number for the \( j \)-th wave, and \( \rho \) is the average mass density of the porous material. \( A_j \) is the characteristic root[1] for the \( j \)-th wave. The first wave is corresponds to the fast longitudinal wave (say P1-wave) and the 2-nd wave is the slow longitudinal wave (say P2-wave). The 3-rd, the transverse wave. For normal incidence, no transverse wave is generated. Such a case, the k-th layer has the displacement potential expressed by

\[ \phi^{(k)} = \Sigma \delta_j \left( B_j^{(k)} E_j^{(k)} + C_j^{(k)} E_j^{(-k)} \right) \]

\[ E_j^{(k)} = \exp (\pm i \delta_j \left( z - a_j \right) \]  

with the unknown potential amplitudes, \( B_j^{(k)} \) and \( C_j^{(k)} \). Notation \( \Sigma \) stands for the summation \( j = 1, 2, 3 \). and lower surface of the k-th layer is assigned by \( z = d_k \). 0-th layer is corresponding incident side half space, and \( N+1 \)-layer is transmitted side half space. Then

\[ \phi^{(0)} = B^{(0)} E^{(0)} + C^{(0)} E^{(-0)} \]

\[ E_j^{(0)} = \exp (\pm i \delta_j \left( z - a_j \right) \]

\[ E_j^{(n+1)} = \exp (\pm i \delta_j \left( z - a_j \right) \]

Definition of the coefficients

- Coefficients of reflection and transmission, \( \Phi \) and \( \Phi_T \), are defined by the ratio of displacement amplitudes.

\[ \Phi = \frac{C^{(0)}}{B^{(0)}} \]

\[ \Phi_T = \frac{B^{(n+1)}}{B^{(0)}} \]

The energy conservation law holds for an attenuation free material. Expression of the law is

\[ \Phi \Phi_s + \Phi_T \Phi_s = 1 \]

This relation can be used to check the numerical results. The superscript \( * \) stands for the value of complex conjugate.

Closed form solutions

We, hereafter, consider the case that a porous single layer\( (n=1) \) is inserted between two elastic half spaces and sonified by normal incident waves. Some interfaces of fluid-filled media allow mass transmission. So, there are in general three kinds of continuity conditions depending on the combination of the media[6]. But both the elastic spaces do not allow the fluid flow from the layer to the elastic spaces. Then the continuity conditions, in our case, are

\[ \Phi_s = \Phi_s + s \]

\[ \Phi_s = \Phi_s \]

\[ \Phi_s = \Phi_s \]

\[ \Phi_s = \Phi_s \]

\[ \Phi_s = \Phi_s \]

\[ \Phi_s = \Phi_s \]

\[ \Phi_s = \Phi_s \]

Superscript bar represents for the value of homogeneous elastic space.

Quitting the detail manipulations, we summarize the expressions stem from the continuity conditions at both the interfaces, assigned by \( z = d_1 \) and \( 0 \), in the form of matrix:

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

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\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]

\[ \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} \Phi_s \\ \Phi_s \end{bmatrix} \]
where the following replacements are used with suppression of superscript (1):

\[ T_3 \cdot M_3 = (P^{(1)} + Q^{(1)}) \cdot \mu_3^{(1)} \left( Q^{(1)} + \bar{R}^{(1)} \right) \]

\[ F_{e1} = \sin(\theta_1 \xi_1), \quad F_{e3} = \cos(\theta_1 \xi_1) \]

It is possible to get the closed form solutions from Eq.(23) by making use of a mathematical processing computer software with some manipulations. The expressions of reflection and transmission coefficients have the following forms:

\[
\rho_T = \left[ \frac{2}{(1-F_{e1}F_{e2})} \right] \delta - \delta_2 (\mu_1 + 1) (\mu_2 + 1) / \delta F_{e2} \delta_2 (\mu_1 + 1) / \delta F_{e1} \delta_2 (\mu_1 + 1) / \delta \xi_1
\]

\[
\rho_T = \left[ \frac{2}{(1-F_{e1}F_{e2})} \right] \delta - \delta_2 (\mu_1 + 1) (\mu_2 + 1) / \delta F_{e2} \delta_2 (\mu_1 + 1) / \delta F_{e1} \delta_2 (\mu_1 + 1) / \delta \xi_1
\]

These results satisfy the energy conservation law, Eq.(21).

**NUMERICAL EXAMPLE AND THE SLOW WAVE VELOCITY**

The calculated results of reflection and transmission coefficients to the incident wave number are plotted in Fig.2. Sharp spikes are present periodically on the curve. Their locations are 0.95, 1.70, 2.55, 3.40, and 4.25 of normalized wave number. This sequence has period 0.85. Further, null reflections are also present by large period 3.14, that comes from the interference of ordinal (fast) longitudinal wave. Ratio of both the periodic coincides with the ratio of slow wave velocity to fast wave velocity. Then the slow velocity will be able to estimate from the fast velocity by this periodical behavior. The slow velocity carries additional informations for characterization of the porous materials. These calculations are easy because of the closed form solution, Eq.(23).

**Table 1** Input data for the numerical calculation

<table>
<thead>
<tr>
<th>Porous elastic layer</th>
<th>Elastic half space</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Elastic constants:</strong></td>
<td><strong>Elastic constants:</strong></td>
</tr>
<tr>
<td>( D = 4.45 \text{GPa} )</td>
<td>( D = 4.45 \text{GPa} )</td>
</tr>
<tr>
<td>( Q = 0.743 \text{GPa} )</td>
<td>( Q = 0.743 \text{GPa} )</td>
</tr>
<tr>
<td>( N = 2.76 \text{GPa} )</td>
<td>( N = 2.76 \text{GPa} )</td>
</tr>
<tr>
<td><strong>Density:</strong></td>
<td><strong>Density:</strong></td>
</tr>
<tr>
<td>( \rho_s = 2.65 \text{g/cm}^3 )</td>
<td>( \rho_s = 2.65 \text{g/cm}^3 )</td>
</tr>
<tr>
<td><strong>Porosity:</strong></td>
<td><strong>Porosity:</strong></td>
</tr>
<tr>
<td>( \beta = 35% )</td>
<td>( \beta = 35% )</td>
</tr>
<tr>
<td><strong>P-wave velocity:</strong></td>
<td><strong>P-wave velocity:</strong></td>
</tr>
<tr>
<td>( P_1 = 2.46 \text{km/sec} )</td>
<td>( P_2 = 2.46 \text{km/sec} )</td>
</tr>
<tr>
<td>( P_2 = 2.66 \text{km/sec} )</td>
<td>( P_2 = 2.66 \text{km/sec} )</td>
</tr>
</tbody>
</table>

* \( \rho = (1 + \beta^{-1})^{-1} \)

**Fig.1** Fluid-filled poroelastic layer between two semi-infinite elastic solid

**Fig.2** Coefficients of reflection and transmission

References
Sound Reflection from Porous Layers between Two Solids

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Introduction

For prospecting the reserve of oil and natural gas by seismic waves, it is necessary to understand the sound reflection characteristics from fluid-saturated porous medium bounded with solid rocks in both sides. So far quite intensive research has been conducted on sound reflection from an interface between fluid (or solid) and porous medium[4,5] and from porous layer embedded in liquid[6,7]. However, to the authors’ knowledge no reports are given on sound reflection from porous layer between two solids. The paper is intended to fill the gap. Some numerical results are presented to show the influence of the incident angle of sound wave and the fluid viscosity in the pores on the sound reflection coefficients.

Theoretical Model

Let us consider a fluid-saturated porous layer with thickness h between two isotropic semi-infinite solids, as shown in Fig.1. A plane longitudinal wave with displacement potential $\phi_0$ is assumed to become incident from the upper semispace onto the porous layer. In Fig.1, $\theta_1$ is the incident angle, $\phi_s$ and $\phi_t$ are the displacement potentials respectively for the reflected longitudinal and transverse waves, $\phi_x$ and $\phi_y$ are the displacement potentials respectively for the transmitted longitudinal and transverse waves.

The Biot’s theory predicts the existence of three kinds of bulk waves in a fluid-saturated porous medium[8], that is, the fast and slow longitudinal waves and transverse wave (the transverse wave travels faster than the slow longitudinal wave). In general, all the three kinds of waves may appear in a porous medium when a longitudinal or transverse wave in the upper solid semispace is incident onto the upper boundary. These waves are reflected again by the lower boundary. Finally three upward and downward waves are formed in the porous layer as shown in Fig.1. A transfer matrix of 6 X 6 order may be used to describe the relation between mechanical parameters at the opposite surfaces of the porous layer. The concerned mechanical parameters in our problem are as follows: the stress $\sigma_{xx}$ and $\sigma_{yy}$ and particle velocities $v_x$ and $v_y$ in the solid skeleton; the stress $\sigma_{xx}$ and particle velocity $v_z$ in the fluid.

If the porous layer is rigidly bounded with the solids, we have $v_x = v_x'$ at each interface of the layer with the solid semispace. After some algebraic operation, the following matrix can be derived:

$$
\begin{bmatrix}
v_x \\
v_y \\
\sigma_{xx}+\sigma_{yy}
\end{bmatrix} =
[C]
\begin{bmatrix}
v_x' \\
v_y' \\
\sigma_{xx}'+\sigma_{yy}'
\end{bmatrix}
$$

(1)

The elements of matrix [C] are not given here for saving space.

By using matrix method, the reflection coefficient of longitudinal wave may be written as

$$
R_{II} = \frac{D_{11}D_{11}-D_{12}D_{12}}{D_{11}D_{11}-D_{12}D_{12}}
$$

(2)

Similarly, when a transverse wave with displacement potential $\phi_y$ is incident from the upper semispace, the reflection coefficient of transverse wave is

$$
R_{II} = \frac{D_{22}D_{21}-D_{21}D_{22}}{D_{11}D_{11}-D_{12}D_{12}}
$$

(3)

In Eqs. (2) and (3) the matrix [D] is defined as:

$$
[D] = [A]^T[C][A]
$$

and the elements of matrix [A] are listed in the Appendix.
In obtaining the expressions for $R_1$ and $R_2$, the continuity conditions for stresses and particle velocities across each interface are applied.

**Numerical Calculation**

As an example we present numerically the reflection coefficients from a porous layer with porosity $\phi=0.283$ and thickness $h=0.5\text{cm}$ between isotropic solids. The related parameters in our calculation are taken as follows: for solid skeleton of the porous medium, the elastic moduli $K_s=0.407\text{Mpa}$, $\mu_s=0.297\text{Mpa}$ and the density $\rho_s=2.46\text{g/cm}^3$; for fluid in pores, $K_f=0.017\text{Mpa}$, $\mu_f=0.018\text{g/cm}^3$, the viscosity $\eta=100.0\text{c.p.}$ and the permeability $k=1.0\text{Darcy}$; for isotropic solid, the longitudinal and transverse wave speeds are respectively $C_l=6.0\text{Km/s}$ and $C_t=3.3\text{Km/s}$ and the density $\rho=2.53\text{g/cm}^3$; the frequency of sound wave is $f=0.5\text{MHz}$.

In Fig.2 and Fig. 3 is shown the dependence of reflection coefficients $|R_1|$ and $|R_2|$ on the incident angle $\theta_1$ and $\theta_2$, respectively. For transverse wave incidence, maximum and minimum values of $|R_1|$ appear at around $\theta_1=40^\circ$, indicating the possible excitation of Lamb wave modes in the porous plate at this angle. Two figures have also illustrate the difference between reflection coefficient with and without consideration of fluid viscosity (broken and solid lines respectively). Numerical calculation shows that at $f=0.5\text{MHz}$ the influence of fluid viscosity on reflection coefficient $|R_1|$ and $|R_2|$ is essential, which may be used to evaluate the oil quality in porous formation.

We have also computed the energy reflection and transmission coefficients for both longitudinal and transverse incidence. It has been demonstrated that when the fluid viscosity is not taken into account, the energy reflection and transmission coefficients of the different modes added together make one, which obeys the energy conservation law.

**Discussion**

The model proposed in this paper may be extended to more complicated layer structure where several different porous layers are located between two isotropic solids. In this case the transfer matrix of $6 \times 6$ order for each porous layer must be multiplied. The resulting matrix can be used to obtain the corresponding matrix $[C]$ in Eq. (1) and then to calculate the reflection coefficients by using Eq. (2) and (3).

**Appendix**

Matrix $[A]$ is defined as follows

$$A_{11} = A_{12} = A_{34} = A_{44} = \sin \theta_1 / c_1$$

$$A_{13} = A_{24} = \cos \theta_1 / c_1$$

$$A_{22} = A_{23} = -A_{44} = P \cos \theta_2$$

$$A_{23} = A_{34} = P \sin \theta_2$$

$$A_{41} = A_{42} = (c_1 / c_2)^2 \sin 2 \theta_1$$

**References**


![Fig.2](image2.png)

**Fig.2** Dependence of reflection coefficient $|R_1|$ on the incident angle

- Solid line—without viscosity in computation
- Dashed line—with viscosity in computation

![Fig.3](image3.png)

**Fig.3** Dependence of reflection coefficient $|R_1|$ on the incident angle

- Solid line—without viscosity in computation
- Dashed line—with viscosity in computation
TRANSMISSION OF SOUND WAVES IN TUBE BUNDLES

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ABSTRACT
Analytic studies have been undertaken in this paper to examine the transmission and reflection of sound waves in tube bundles. The tube bundles consist of parallel layers of tube rows submerged in water. The theory of wave scattering by an infinite grating has been applied to the tube rows, with modification to take into account the tube wall motions. The transmission and reflection coefficients for normal incidence case have been calculated for the tube row, and for tube bundles based on iterative equations. Sharp changes of the transmission and reflection coefficients around resonance frequencies have been observed. The results show that low transmission (high reflection) frequency regions may exist for tube bundles when the space between tube rows equals multiple of the wavelength.

1. INTRODUCTION
Tube bundles are used in heat exchangers and steam generators which are very important components in energy industry. It is often desirable to monitor and test tube bundles under operating conditions, and some of the surveillance techniques is based on sound waves transmission and scattering within the tube bundles. We consider sound transmission and dissipation in tube bundles which are constituted by a finite number of tube rows. Each tube row is an infinite grating consisting of identical tubes. The tube bundles are immersed in water. Some results of Twersky's on waves scattering by an infinite grating, with the modification to take into account tube wall motions are used to calculate the transmission and reflection coefficients for a tube row, and for tube bundles based on iterative equations.

2. GENERAL FORMULATION FOR A TUBE ROW
A tube bundle may be constructed by many layers of tube row, as shown in Fig. 1. We consider one of such tube rows as an infinite grating composed of parallel cylindrical shells extending along y axis with an equally separate distance d. All tubes are very long and parallel to z axis. The water is of sound speed c and density ρ respectively. A time-harmonic plane wave, denoted by \( p \), travels from left to right at an incident angle \( \theta_i \) in a polar coordinate system \((r, \phi)\) with its origin located at the centre of one of the tubes within the tube row, \( k = \omega/c \), being wave number. The time factor \( e^{+ \theta_i} \) will be dropped throughout our analysis. The incident wave will be scattered by the tube row and the total transmitted waves are given by \( p(r, \phi) \).

\[ p(r, \phi) = \sum \frac{2p_n}{kd} \sum A_n \sum e^{i(k_a - k)z} \cos \phi \]

\[ r \cos \phi > a, \]

and the reflected waves are

\[ p(r, \phi) = \sum \frac{2p_n}{kd} \sum A_n \sum e^{i(k_a + k)z} \cos \phi \]

\[ r \cos \phi < -a \]

which are essentially the sum of the incident and the diffracted plane waves at different angles \( \Phi_n \). In Eqs. (1a) and (1b),

\[ \sin \phi_n = \sin \phi_i + \nu l \]

and \( \cos \Phi_n = (1 - \sin^2 \phi_i)^{1/2} \), where \( \lambda \) is the wave length. The diffraction waves corresponding to complex \( \phi_n \) cannot propagate away from the tube array but they will damp out exponentially. The coefficients \( A_n \) are complex in Eqs. (1a) and (1b), which is a measure of amplitude of the diffraction waves, will be obtained by solving the system of algebraic equations

\[ A_n = \sum \frac{p_n}{kd} \sum \sum A_n \]

where \( F_{n,m} \) represents Schrödinger series. \( a_n \)'s are Fourier coefficients of scattering amplitude for an isolated tube, which contain information about acoustic behaviour for individual tube within the tube row. \( a_n \) can be obtained by

\[ a_n = \frac{D_n f_n(ka) - iI_n(ka)}{D_n H_n(ka) - iI_n(ka)} \]

where \( D_n \) is the tube wall impedance including the inner fluid loading, \( f_n(x) = \frac{a_n}{x} H_n(x) = \frac{a_n}{x} \).

3. RESULTS AND DISCUSSION
We consider the case in which \( \rho = 1000 \text{ kg/m}^3, \ c = 1460 \)

m/s and the tube material is steel. In our calculations the distance between tubes is \( d = 2.1 \text{ cm} \), and the frequency of the incident wave, denoted by the non-dimensional wave number \( ka = \omega/c \), is limited in a range of \( ka \leq 3(\lambda_{\text{m}} = 1.8 \text{ cm for } a = 0.85 \text{ cm}) \) to ensure that the wavelength is much large compared with the thickness of the tube wall. At low frequencies where \( \lambda > d \) and \( \nu \), only waves travelling with angles \( \phi_i \) and \( \phi_n + \pi \) are propagational modes while those at other directions are evanescent modes. The propagation waves in our calculation mainly pertain to mode \( \nu = 0 \). For this mode, the transmission coefficient of the tube row is

\[ T = p/p_{in} = 1 + 2 \sum A_n \]

and the reflection coefficient of the tube row is

\[ R = p/p_{in} = 2 \sum (-1)^n A_n \]

Fig. 2 shows numerical results of transmission coefficients for a tube row consisting of tubes with radius a = 0.85 cm and the wall thickness h = 0.1 cm. As it can be seen, the overall value of the transmission coefficient is high at frequencies below

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\[ ka = \omega a/c < 2, \text{ except at some very narrow frequency bands} \]

\[ |T| \text{ and } |R| \text{ drop sharply and jump up instantly with the values} \]

\[ \text{between zero and one. These frequencies are associated with} \]

\[ \text{the resonance modes of tubes within the row, whose mode} \]

\[ \text{numbers have been indicated by integers. The mode } n = 1 \text{ is} \]

\[ \text{induced by the inner fluid loading, and belongs to mode } n = 1 \]

\[ \text{but at a higher frequency}^{[3]} \]. When driving frequencies are in

\[ \text{the region over the critical value denoted by dashed line in Fig.} \]

\[ 2, \ \text{where the wavelengths are less than the space between the} \]

\[ \text{tubes, transmission coefficients drop to low levels since the} \]

\[ \text{scattering waves at other diffraction angles become} \]

\[ \text{propagational modes, and take some energy away from mode} \]

\[ v = 0. \]

A tube bundle may be modelled as many tube rows lined up in parallel with an equally spaced distance \( b \), see Fig. 1. The tube bundle of \( j \) tube rows tube can be regarded as a bundle of \( j-1 \) tube rows plus one more tube row. \( T_{j-1} \), \( R_{j-1} \), \( R_j \) are the known transmission and reflection coefficients for the single and the \( j-1 \) tube row respectively. The iterative equations for the transmission and reflection coefficients for \( j \)-row tube bundle \( |T_j| \) are given as follows\[^{[4]} \]

\[ T_j = \frac{T_j T_{j-1}}{1 - R_j R_{j-1} e^{ikb}}, \]

\[ R_j = R_{j-1} + \frac{T_j^2 R_{j-1} e^{ikb}}{1 - R_j R_{j-1} e^{ikb}}. \]

\[ |T_j| \text{ are plotted in Fig. 3 for the tube bundles which are of } j = 1, 3 \text{ and 6 rows with } b = d = 2.1 \text{ cm. The calculation has been} \]

\[ \text{conducted in a frequency range below the critical frequency} \]

\[ (ka = 2\pi a/d). \text{ It is interesting to note that multi-reflections} \]

\[ \text{within the tube bundle have drastically changed the} \]

\[ \text{transmission coefficients, and formed a low transmission zone} \]

\[ \text{where } |T_j| \text{ drops to a very low level. Being similar to that of} \]

\[ \text{multiplier interference filter in optics,}^{[6]} \text{ the low transmission} \]

\[ \text{zone is located approximately around the frequencies where} \]

\[ \delta = 2kb = m2\pi (m = 1, 2, \ldots), \text{which gives } ka = m\pi a/b = 1.27 \text{ for} \]

\[ m = 1, \text{ as shown in Fig. 3. Such frequency regions with low} \]

\[ \text{transmission for tube bundles have been observed in} \]

\[ \text{experiments.}^{[3]} \]

4. CONCLUSIONS

Some basic features of sound waves transmission and reflection in tube bundles have been considered and described in this paper. The results based on normal incidence \( (\phi = 0) \) show that the transmission and reflection coefficients of the tube row jump sharply around the resonance frequencies. When the tube rows are arranged in parallel to constitute the tube bundles, the transmission and reflection coefficients of such tube bundles can be calculated by an iterative method based on Eqs. (6) and (7). Multi-reflections between tube rows may cause low transmission (high reflection) in some frequency regions located around \( ka = m\pi a/b (m = 1, 2, \ldots) \). The existence of these low transmission zones, which have been observed in experiments, could be an important guideline of using surveillance techniques based on the sound transmitted through the tube bundles.

ACKNOWLEDGEMENT

The author highly appreciates valuable discussions with Dr. Maria Heckl. The support from the Science and Engineering Research Council of UK is acknowledged.

REFERENCES


Figure 1. Geometry of a tube bundle

Figure 2. Transmission coefficient of a tube row submerged in water. \( a = 0.85 \text{ cm}, h = 0.1 \text{ cm}, d = 2.1 \text{ cm}. \)

Figure 3. Transmission coefficient for tube bundles with row numbers \( j = 1, 3 \) and 6. \( a = 0.85 \text{ cm}, h = 0.1 \text{ cm}, b/a = d/a = 2.47. \)
RELAXATION TIME RELATED TO MOLECULAR STRUCTURE IN POLYURETHANES

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INTRODUCTION

Polyurethanes, like most high polymers, exhibit a broad, asymmetrical acoustic relaxation at their glass transition. Depending on the molecular structure of the polyurethane, this relaxation can be centered anywhere within a range of many decades of frequency. The purpose of this work is to establish the relation between the molecular structure of polyurethanes and the relaxation time for this process.

A series of polyurethanes with known variations in structure were synthesized and their acoustic properties measured with a resonant bar technique. The data is in the form of master curves of shear modulus and loss factor as a function of frequency. The master curves were fitted to the Havriliak-Negami analytical model to determine the average relaxation time, $\tau$, which determines the center frequency of the relaxation.

It is well known [1] that the location of the glass transition peak as a function of temperature is given by the glass transition temperature, $T_g$, measured in a differential scanning calorimeter. It is proposed here that there should be a correlation between the location in the temperature domain ($T_g$) and the location in the frequency domain ($\tau$). An Arrhenius plot of $\ln \tau$ vs $1/T$ verifies this assumption.

Since $T_g$ is an additive property [2], the correlation between $T_g$ and $\tau$ implies that $\tau$ is also an additive property. The principle of additive properties states that the overall relaxation time is made up of independent contributions from the bivalent components that make up the structure of the polymer. $\tau$ was found to be an additive property, and the relaxation time for each of the polyurethane components was determined. Having these component values, the average relaxation time for any polymer that can be synthesized from the known components can now be predicted.

EXPERIMENTAL

Polyurethanes were synthesized by first preparing a pre-polymer from the reaction between a polyglycol and a diisocyanate. This pre-polymer was then chain extended with a diol. The polyglycol used here were either polytetramethylene ether glycol or polypropylene glycol of molecular weight varying from 650 to 2900. The diisocyanate used was either 4,4'-diphenylmethane diisocyanate or 4,4'-dicyclohexylmethane diisocyanate. The pre-polymer had polyglycol/diisocyanate molar ratios varying from 1:3 to 1:6. The chain extender used was either 1,4-butanediol or propane diol with a variety of pendant groups. Stoichiometry of the chain extenders was adjusted so that there was a 5% excess of diisocyanate, ensuring a lightly crosslinked polymer. Details on the synthesis of these polymers have already been published [3,4,5].

Acoustic properties were determined using the resonance apparatus [6]. In this device, a sample in the shape of a bar, 10 to 15 cm long with square lateral dimensions of 0.635 cm, is excited into resonance at various harmonics. Typically, four or five resonant peaks can be determined in a nominal range of 1.5 decades centered at about 1 kHz. The real, $G'$, and imaginary, $G''$, parts of the complex shear modulus, $G'$, are calculated at each of these resonant frequencies. The measurements are repeated as a function of temperature from -60 to 70°C. Time-temperature superposition [7] is used to shift the measured data to form a master curve at a reference temperature of 25°C. Typical data [4] is shown as circles in Fig. 1.

![Figure 1. Typical complex modulus data for a polyurethane](image)

Glass transition temperatures were determined using a differential scanning calorimeter (DSC) module on a DuPont 9900 Thermal Analyzer at a scan rate of 10°C/min in an argon purge.

ANALYTICAL MODEL

The most successful analytical model to describe the frequency dependence of the glass transition is the Havriliak-Negami model [8]

$$\frac{(G''-G_0)}{(G''-G_\infty)} = [1+(i\omega \tau)^\alpha]^\beta$$

where $G_0$ is the limiting high frequency modulus, $G_\infty$ is the limiting low frequency modulus, $\omega$ is the angular frequency, $\tau$ is the average relaxation time, and $\alpha$ and $\beta$ are constants with values between zero and one.

Fitting of the data in Fig. 1 to eq. 1 was carried out using a non-linear least square routine. Fitted results are shown as solid lines in Fig. 1.
ADDITIVE PROPERTIES

The basic assumption of additive properties is that the properties of a polymer can be calculated from the properties of the component bivalent groups. This procedure has been demonstrated to be valid for numerous polymer physical properties and is documented in detail by Van Krevelen [9].

Two examples of additive properties are density and glass transition temperature. For density, which is the ratio of molar volume to molar mass, the equation is

$$\rho = \frac{\sum N_i V_i}{\sum N_i M_i}$$

where \(N_i\) is the number of component groups \(i\) in the repeat unit, \(V_i\) is the molar volume of component \(i\), and \(M_i\) is the molar mass of component \(i\). For \(T_g\), there are several possible forms of the additive property equation [2]. We will use

$$T_g^{-1} = \frac{\sum (N_i M_i / M)}{T_{g1}^{-1}}$$

where \(T_g\) is the polymer glass transition temperature and \(T_{g1}\) is the glass transition temperature of component \(i\).

The polymers considered here, constituent components are listed in Table I. For those components, \(V_i\) and \(T_{g1}\) can be determined using eqs. 2 and 3. The results are also listed in Table I. The polymer density calculated from the component values in Table I agrees with measured density within 0.5 percent while the calculated glass transition temperatures agree with measured values within 0.5 percent.

CORRELATION

Since \(\tau\) represents where along the frequency axis the transition occurs and \(T_g\) represents where the transition occurs along the temperature axis, the two properties should be related. This is verified in Fig. 2. From this Arrhenius plot, we find an activation energy of 123 KJ/mol.

Since \(\tau\) and \(T_g\) are correlated and since \(T_g\) is an additive property, \(\tau\) should be an additive property as well. Component values were determined and are listed in Table I. Average relaxation times calculated from these component values vary by 18 orders of magnitude and 4 orders of magnitude with measured results by an average of about 25 percent.

CONCLUSIONS

A series of polyurethanes of known, varying structure were synthesized and master curves for shear modulus and loss factor determined. Fitting this data to the Havriliak-Negami equation gives an average relaxation time. This \(\tau\) is correlated with the independently measured glass transition temperature, with an activation energy of 123 KJ/mol. In addition, this relation implies that \(\tau\) is an additive property. Component relaxation times vary from 10^{-16} s to 10^{-2} s. Using these component values, the average relaxation time for polyurethane polymers can be calculated to an accuracy of about 25 percent.

REFERENCES

3. J. V. Duffy, G. F. Lee, J. D. Lee, and B. Hartmann, in reference 1, Ch. 15.

Table I. Component Properties

<table>
<thead>
<tr>
<th>Component</th>
<th>M_i</th>
<th>(V_i)</th>
<th>(T_{g1})</th>
<th>(\tau_i)</th>
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<tr>
<td>-(CH_2)_n-O-</td>
<td>72</td>
<td>73</td>
<td>179</td>
<td>2.4E+16</td>
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<tr>
<td>-OC(OH)-MDI-HNCOO-</td>
<td>284</td>
<td>193</td>
<td>727</td>
<td>3.4E+11</td>
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<tr>
<td>-CH_2C(CH_3)_2CH_2-</td>
<td>70</td>
<td>98</td>
<td>315</td>
<td>9.0E+01</td>
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<tr>
<td>-CH_2C(CH_3)_2CH_2-</td>
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<td>146</td>
<td>259</td>
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<tr>
<td>-CH(CH_2CH_2CH_2-</td>
<td>91</td>
<td>122</td>
<td>222</td>
<td>3.0E-09</td>
</tr>
<tr>
<td>-CH(CH_2CH_2CH_2-</td>
<td>84</td>
<td>123</td>
<td>364</td>
<td>5.3E+02</td>
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<tr>
<td>-CH(CH_2CH_2CH_2-</td>
<td>126</td>
<td>162</td>
<td>348</td>
<td>7.9E+01</td>
</tr>
</tbody>
</table>

Units are: g/mol, cm^2/mol, K, and sec

![Figure 2. Arrhenius plot of relaxation time vs. reciprocal glass transition temperature](image-url)
ELASTIC WAVE DISPERSION IN A SYMMETRICAL TRILAYER

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INTRODUCTION

The research reported in this paper frames in a project on Lamb waves in multilayers. It aims at understanding more about the dispersion of these waves starting from the properties of the constituent layers. J. Laperre et al. [1] [2] [3] recently demonstrated for liquid an solid bilayers:

- that a Lambmode of the bilayer starts at normal incidence on a mode of one of the constituent layers
- that coupling between modes, switches the bilayer to the next higher or the next lower mode it was initially tuned to.

This second phenomenon can be explained by means of a simple mechanical model in which the bilayer is replaced by two masses connected by a spring [4].

The purpose of this article is to repeat this analysis for a symmetrical trilayer composed of two glass plates separated by a thin water layer.

MECHANICAL MODEL FOR A SYMMETRICAL TRILAYER

We first discuss the mechanical model so that we know what we have to look for when analysing the dispersion curves of a trilayer. The model consists of three masses connected by springs (see figure 1). The central mass $m_0$ is connected with two identical masses $m_1$ by means of springs with stiffness $k_0$. Two more springs with stiffness $k_1$ are needed to suspend the three masses so that they can vibrate.

\[ f_0 = \frac{1}{2\pi} \sqrt{\frac{k_0 m_1}{m_0}} \left(1 + \frac{k_0}{k_{0,1}}\right) \]  \hspace{1cm} (3)

The mechanical coupling coefficient $k_c$ is equal to:

\[ k_c = \frac{k_0}{2(k_1 + k_0)} \]  \hspace{1cm} (4)

According to equations 1 and 3, one of the resonance frequencies of the coupled oscillators coincides with the resonance frequency of the uncoupled mass $m_1$. As far as the Lamb modes of the glass/water/glass trilayer are concerned, this implies that we expect every mode of the isolated glass plate to be a mode of the trilayer.

We investigated the behaviour of the resonance frequencies as a function of $m_1$, using the following numerical data: $m_0 = 1$, $k_0 = 5$ and $k_1 = 100$. $k_c$ is equal to 0.024, which means that we are in the weak coupling limit ($k_c < 1$). The results are plotted in figure 2.

Figure 1: Mechanical analogue of the symmetrical trilayer

Figure 2: The behaviour of the resonance frequencies as a function of $m_1$ using $m_0 = 1$, $k_0 = 5$ and $k_1 = 100$.

By solving the system of equations governing these three coupled oscillators, we find that there are three resonance frequencies:

\[ f_\pm = \frac{f_0^2 + f_1^2}{2} \pm \frac{1}{2} \sqrt{(f_0^2 - f_1^2)^2 + 4k_1 f_0^2 f_1^2} \]  \hspace{1cm} (1)

and:

\[ f = f_1 \]  \hspace{1cm} (2)

The solid lines represent respectively $f_+$ and $f_-$ as indicated; the dash-line and the vertical dotted line represent $f_0$ and $f_0$ respectively. It turns out that in the weak coupling limit, $f_\pm$ are equal to either $f_0$ or $f_1$, except in the region where the dashed and the dotted line cross. There $f_+$ switches from $f_1$ to $f_0$, and $f_-$ switches from $f_0$ to $f_1$. This switching of the normal modes of the coupled oscillators between modes of the isolated oscillators is referred to as mode-coupling. It occurs when the resonance frequencies of the isolated oscillators approach each other. Far enough from the crossing point, $f_\pm$ almost coincide with $f_1$ which is also a resonance frequency of the coupled oscillators (see equation 2). The curve representing $f_1$ as a function of $m_1$ is so to speak, bifurcated.

In the next section we will demonstrate theoretically and experimentally that the Lamb modes of a symmetrical trilayer glass/water/glass shows a similar behaviour.
THE LAMB MODES OF A SYMMETRICAL TRILAYER

Before discussing the experimental results, we first present in figure 3 the numerical dispersion curves for a trilayer composed of two glass plates with a thickness of 1.73 mm, separated by a 0.82 mm thick waterlayer. They are obtained by solving the equation:

\[
\frac{\omega \cos(\theta_L)}{v_w} + 2\Omega = n\pi
\]

where \(\Omega\) represents half the phase change when a plane wave is reflected from the water/glass plate/air structure (see [5]). \(\theta_L\) represents the Lamb angle, \(\omega\) is the angular frequency, \(v_w\) is the sound velocity in water and \(d\) is the thickness of the waterlayer.

Each solid line corresponds to a particular Lamb mode and represents the frequency of this mode as a function of the Lamb angle. On the same figure we also plotted the numerical dispersion curves of the isolated glass plates (dashed-lines) with stress-free boundaries, and of the waterlayer (dotted lines) with rigidly held (displacement vector = 0) boundaries. Although this figure looks rather complicated at first sight, inspection of what happens at the crossing points of a dash- and a dotted line (of a mode of the glass plates and of the water layer) clarifies the picture a great deal, and confirms what we have found in the previous section. Namely:

- a Lamb mode of the glass plates is also a mode of the trilayer
- coupling between a mode of the glass and a mode of the water, switches a trilayer mode from a mode of the glass plate to one of the water layer, and vice versa.

Notice also that far from a crossing point between a dash- and a dotted line, the modes of the glass plates (dash-lines) are bifurcated as predicted by our mechanical model.

To study the dispersion of Lamb waves experimentally, we used a double-transmission measuring technique. It proceeds as follows: the trilayer is positioned in a watertank so that it can rotate in small steps around a horizontal axis.

A broadband ultrasonic transducer with a central frequency of 3.5 MHz insinuates the trilayer with ultrasonic pulses and picks up the double-transmitted echo's from a perpendicular plane mirror. For each angle of incidence, the echo's are gated and FFT analyzed. By plotting the frequency of the maxima in the FFT spectrum on the horizontal axis, and the angle of incidence on the vertical axis, we obtain the experimental dispersion curves of the trilayer. Figure 4 shows the experimental data (O's) for a trilayer composed of two glass plates with a thickness of 1.73 mm separated by a 0.35 mm thick water layer. On the same figure we also plotted the numerical dispersion curves of the trilayer (solid lines), of the isolated glass plates (dash-lines) and of the isolated water layer (dotted lines). We immediately see that not all the theoretically predicted Lamb modes are observed. We also ascertain that the switching of a trilayer mode between monolayer modes is clearly observed in only two cases indicated by circles on figure 4.

References

A8-6

ACOUSTIC RELAXATION OF FLUID SATURATED POROUS SAMPLE

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Introduction:
Acoustic relaxation in a sample of fluid saturated porous medium (abbreviated to FSPM layer) was observed by Spencer in natural sandstone; for the working frequency of his arrangement limited to below 400 Hz, most experimental points located below the peak frequency of relaxation curves. White and Dunn explained Spencer's experiments by Biot's theory, the physical mechanism is based on the flow of saturated fluid related to the solid frame, when the sample deformed statically or at a very low frequency, the fluid may move outside or inside the surface of the sample freely. So, the elastic modulus of the fluid has no contribution to the modulus of the sample. In this case, the elastic modulus of the sample is equal to the elastic modulus of the solid frame whatever the saturated fluid is. When the sample deformed at a high frequency, it doesn't enough time for the saturated fluid to flow outside or inside freely, so the elastic modulus of the fluid may contribute to the modulus of the sample and the modulus reaches a higher value near to that of the sealed boundary condition when the frequency high enough. In these two cases, the dissipation caused by viscous flow were small, for the low flow resistence in the first case and litter flow in the second case. The dissipation of viscous flow reaches a maximum at an intermediate frequency. The frequency range studied here, whatever "high" or "low", are the low frequency range in Biot's theory, for the viscous skin are much larger than the pore size in the frequency range studied here. The disagreement between theory and experiments are the experimental value of low frequency elastic modulus for same sample are different for different saturated fluid, in order to fit Spencer's experimental points, Dunn used selected elastic moduli for solid frame refer to different saturated fluid. White has listed some unknown mechanism accompanying the mechanism caused by the flow of saturated fluid. One may found the complicies in Spencer's experiments, for example, he get the relaxation curve for an unsaturated sample but has reserved in a humid environment. For further study of the acoustic relaxation of fluid saturated sample; Experimentally the artificial sample is used in the purpose for avoiding the complicity in natural samples; the viscosity (controlled by temperature) is selected as the variable for studying the change of attenuation and elastic modulus, the experimental points may cover the relaxation peak easier in this case than selecting the frequency as the variable reported before. Theoretically, based on Biot's theory, one dimensional approach on the pressure distribution along the thickness and the modulus of the FSPM plate at flexural vibration is expressed, numerical calculation by using measured parameters the theoretical curves agree with experimental points in two respects; the location of peak and the value of elastic modulus at low viscosity (corresponding to low frequency). The mechanism of acoustic relaxation, caused by the flow of saturated fluid proposed by White and Dunn, is examined to be valid in artificial samples with flexural vibration.

* Project supported by NSFC.

Complex Modulus of FSPM Plate at Flexural vibration:

1. Application of Biot's theory:
According to Biot's theory, the displacement of solid frame and fluid are designated by u and U respectively, the strains are \(\varepsilon_i = \left(u_i + \nu u_j \right)/2\) and \(\varepsilon = \nu u_i\) respectively, the stresses are \(\sigma_i\) and \(S = \beta \rho \varepsilon\) respectively, where \(\beta\) is the porosity, and \(\rho\) is the pressure, the total force \(t_i\) on unit surface of FSPM is

\[ t_i = \sigma_i - \beta \rho \varepsilon_i \]  

(1)

In elasto-statics of Biot's theory \(\sigma_i\) (or \(t_i\)) and \(S\) (or \(p\)) are linear functions of \(\varepsilon_i\) and \(\varepsilon\) respectively, four elastic constants are introduced to express these linear relations. There are two coupled equations in Biot's dynamical theory and three mass coefficients \(\rho_{11}, \rho_{22}, \rho_{23}\) are introduced, they are related to the frame, the fluid, and their coupling effect respectively. If the viscosity \(\eta\) of the fluid takes into account, we may replace \(\rho_{12}\) by a complex value \(\rho_{12}\) at Biot's low frequency range:

\[ \dot{\sigma}_{12} = \rho_{12} + \frac{\beta^2 \eta}{k\omega} \]  

(2)

where \(k\) is the permeability of frame, \(\omega\) is the circular frequency, \(i = \sqrt{-1}\). When the strain and stress are sinusoidal, we may eliminate the \(\Delta\) from the coupled dynamical equations, and get the following differential equations:

\[ \Delta \cdot \dot{\varepsilon} + \omega \omega \cdot \varepsilon = \omega \omega \cdot \dot{\varepsilon} + \omega \omega \cdot \varepsilon = 0 \]  

(3)

where the common factor \(\exp(\omega t)\) is neglected, and

\[ F = \frac{\omega^2 (\rho_{11} P - \rho_{12} Q)}{\rho R - Q^2} \]  

(4a)

\[ L = \frac{\rho^2 (\rho_{22} P - \rho_{23} Q)}{\rho R - Q^2} \]  

(4b)

where \(P = A + 2 i N, A, N, Q, R\) are the Biot's elastic constants.

It is worthy to point out that if in frequency near the relaxation peak of Spencer's experiments, the value of the second term in the right of equation(2) is always much larger than the first term, that is, in the experimental range the inertial term may neglect, this is the reason for the succession of Dunn's quasi-static assumption.

2. One dimensional flexural vibration of solid plate:
Fig.1 is the section of a bended plate, its crosssection is a rectangular, the coordinates \(x, y, z\) are parallel to the length, thickness and width respectively, and the \(y\)-axis coincide with the neutral surface(dash--dot line in Fig.1). If the radius of curvature \(r\) is much larger than the thickness of the plate \(d\), we have \(\sigma_{yy} = 0\). Assuming the width is large enough, we have \(c_{xx} = 0\). In this case the strain in the plate is \(\varepsilon_{xx} = E' c_{xx} / (1 - \nu^2)\), \(\nu\) is the Poisson ratio, for convenience we let \(E' = E / (1 - \nu^2)\). The torque \(M\) for unit width applied to the surface perpendicular to \(x\) may be expressed as:

\[ M = \int_{-d/2}^{d/2} E' c_{xx} dy = \int_{-d/2}^{d/2} \sigma_{yy} dy = \frac{2}{3r} E' d^3 \]

and

\[ E' = \frac{3r}{2d} \int_{-d/2}^{d/2} \frac{\sigma_{yy} dy}{2d} \]  

(5)

If the plate is made by FSPM, the \(\sigma_{yy}\) in equation (5) may be replaced by \(\dot{\varepsilon}_{xx}\) in (1), and the modulus may be complex.

3. The complex modulus of FSPM plate:
We solve equation(3) in the case of Fig.1, and get the expression for \(\dot{\varepsilon}_{xx}\) by Biot's elastic relations. Using \(\dot{\varepsilon}_{xx}\) instead
of \( \sigma_T \) in equation (5), we may get the expression of complex modulus \( E' \):

\[
E' = A_1 + A_2 (1 / kd - \cot \theta) \frac{kd}{kd}
\]

where

\[
A_1 = 4N (1 - \frac{N}{Q^2 + 2PQ + PR})
\]

\[
A_2 = -12N^2 (R + Q) \frac{P + Q}{Q^2 + 2PQ + PR} + \frac{Q}{PR - Q}
\]

The constants \( A_1, A_2 \) are real and expressed by Biot's elastic constants, but the \( K \) in (7) is complex, so the modulus is complex and may expressed by:

\[
E' = E_i' + iE_r'
\]

where \( E_i' \) and \( E_r' \) are the real part and the imaginary part of the modulus respectively, and the attenuation may be expressed by the inverse of the quantitative factor \( Q^{-1} \):

\[
Q^{-1} = \frac{E_i'}{E_r'}
\]

When the elastic constants are given, the modulus is a function of \( \omega \) or \( \eta \) through \( K \). It is clear from (7), the change of \( \omega \) and the change of \( \eta \) are equivalent. In our experiments, \( \omega \) is kept constant, and \( \eta \) is changed through temperature (the viscosity of the saturated fluid as a function of temperature has been measured). For comparison, the temperature is selected as the independent variable for the theoretical curve on the modulus and the attenuation \( Q^{-1} \). Fig. 3 is an example, the parameter used for calculation is measured previously.10

**Experiments:**

1. The sample of artificial FSPM:

   Our samples are sintered by bead of glass, the related constants are measured and expressed in C.G.S unit: \( \beta = 0.32; \rho_a = 1.62; \rho_m = 2.45; \kappa = 2.8 \times 10^{-4}; \gamma = 7.3 \times 10^{-4}; \kappa_a = 3.5 \times 10^{-1}; N = 4.0 \times 10^{-5}. \)

2. Saturated fluid:

   Three kinds of fluid are selected for saturation, their density \( \rho_v \) bulk modulus \( K_v \) and viscosity coefficient \( \eta \) measured as a function of temperature and expressed by the following empirical formulas:

\[
\rho_v = \rho_v + \rho \frac{A_1}{T}
\]

\[
K_v = K_v + B_1 e^B_2 T^2
\]

\[
\ln \eta = C_1 + C_2 \frac{C_3}{T^2}
\]

where \( t \) is temperature in centigrade degree and \( T = 273 + t \) is the Kelvin scale. The measured temperature coefficients in (12), (13), (14) for the fluid in following experiments are expressed in C.G.S unit: \( \rho_v = 0.916; \lambda = -6.34 \times 10^{-5}; \kappa_a = 2.14 \times 10^{11}; B_1 = -5.11 \times 10^5; B_2 = -4.5 \times 10^4; C_1 = -3.76; C_2 = -8.4 \times 10^2; C_3 = 1.2 \times 10^4. \)

3. Measurements of the relaxation of the sample of FSPM:

Fig. 2 is a sketch of the arrangement for measurements. The sample resonant at the basic frequency of flexural vibration, and is suspended at its node by two thin wires. Two piezoelectric plates \( P_1 \) and \( P_2 \) used as excitor and detector respectively, are glued to the sample near the node. The dimension of the sample is \( h \times h \times l = 0.365 \times 1.98 \times 15.2 \text{ cm}^3 \), where \( d \) is smaller than \( h \) and \( l \). Keep the exciting voltage on \( P_1 \) constant, we find the resonance frequency \( \omega_n \) at maximum voltage on the detector \( P_2 \), and the real part \( E_n' \) of the modulus is calculated by:

\[
\omega_n = \sqrt{\frac{E_n'}{\rho} (1.506 \alpha^2)}
\]

The imaginary part of the modulus \( E_n'' \) is measured by the logarithmic decrement \( \gamma \) at free vibration of the sample, and \( E_i \) may be calculated by:

\[
Q^{-1} = \frac{E_i}{E_r} = \frac{2\gamma}{\omega}
\]

Put the sample in a thermostat, the temperature may keep to \( +0.2 \text{ centigrade} \), and change it by 5 degree centigrade a step. Experimental results are expressed in Fig. 3, where the \( x \)-axis represent temperature and the \( y \)-axis represent \( E_i \) (scaled at the right) and the attenuation \( Q^{-1} \) (scaled at the left).

**Fig. 3** Attenuation \( Q^{-1} \) and elastic modulus \( E_i' \) as a function of temperature. "-" Theoretical curve, "-" Experimental points for dry sample, "-" Experimental points for the sample saturated by mineral oil.

Theoretical curves, calculated by using measured data are drawn in the figures for comparison. The theoretical curve and experimental points agree in two respects: the location of the relaxation peak and the elastic modulus at high temperature (corresponding to low viscosity or low frequency). There is no relaxation in unsaturated sample.

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**References:**

5. To be publish.
SOUND ATTENUATION IN WATER-SATURATED SAND

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INTRODUCTION

As suggested by some authors, the relation between sound attenuation α and frequency f in granular media is approximately linear within quite extensive frequency range [1] but to a power higher than the first when frequency is low enough [2,3]. Regarding particles approximately as smooth and spherical under condition $\nu R < 1$, where $\nu$ is compressional wave number, $R$ is average radius of particles, and taking account of both scattered and viscous waves' interactions as well as particle size distribution, the concentrated suspension theory established by one of the authors in 1980's [4-7] has reached a satisfactory agreement with some experimental results. However, when $\nu R$ is very large ($\nu R \approx 20-240$ in this experiment, where $\nu$ is viscous wave number [8]), the particles in nature should not still be regarded as spherical and smooth (see Figure 5) on the scale of the wavelength $\lambda$ of viscous wave ($\lambda \approx \nu R / \nu$) and it is expected that the theoretical results will obviously deviate from the experimental ones. Moreover, many previous experiments are not strict and systematic enough to show a trustworthy relationship between acoustic attenuation and frequency in wide frequency range, especially in high frequency range. Therefore strict experimental investigation as well as theoretical analysis was conducted on a sample of coarse riverbed sand in a high frequency range wide from 40 to 600 kHz.

EXPERIMENTAL METHOD

The experimental setup is illustrated in Figure 1. The tank was 60cm long, 40cm wide and 45cm high. After heating-cooling treatment to eliminate the disturbance of air bubbles in it [8], the sample, which was 35cm high and covered by a thin layer of water, was loaded in the tank. A pair of transmitting and receiving transducers, which were coaxial and fixed onto two stainless steel rods respectively, were placed at the half height of the sand. The rods were clamped onto a copper base that was also used as a vernier rule to read the distance $Z$ between the transducers. It was proved that sound signal and at each distance the measurements were repeated many times to obtain an average. Then the average was modified by a method to be mentioned below and a straight line was fit to the logarithms of the modified values vs. distances. Finally, $\alpha$ was calculated from the slope.

For the experiment in sand, $Z_0=3.2cm$ at maximum frequency 600kHz, and the minimum $Z$ was 15cm. Thus, the condition $Z/Z_0$ is always valid in the whole frequency range. Furthermore, the geometrical deviation of the quasi-spherical spread at practical far field point from the ideally spherical spread at infinite point was considered and a correction measurement was conducted in water at each frequency to get a geometrical correction factor. Provided the amplitude value measured in sand is modified by the factor, the error caused by the nonspherical spread is corrected. Since the far field condition $Z/Z_0$ was usually valid in our measurement, the geometrical deviations and the correspondent corrections were usually quite small (less than 5%).

Pulsed sinusoidal signals with a duration of 70 μs were used. It can be reasoned from the experimental conditions that the waves reflected from boundaries would never reach the receiver till more than 92 μs after the leading edge of the direct wave for the maximum $Z$, 30cm, when measuring $\alpha$ in sand. This delay time would be even longer for smaller $Z$. Eventually, a delay time less than 60 μs was chosen for reading the amplitudes of the direct waves so as to completely avoid the interferences of boundary reflections.

The amplitude variations of received signals at various distances from 150mm to 300mm were very large because of the strong attenuation in sand so that the linear dynamic range of the circuits must be wide enough to cover the maximum variation of the amplitudes at each frequency. The nonlinear distortions of the amplifiers were lower than 1% in this measurement.

A photograph of the riverbed sand is shown in Figure 2. Its density was 2.62 g/cm$^3$. The granular size distribution was measured by sieve method using standard cascade mesh sieves and fit with phi [3] normal distribution by weight. The standard deviation $\sigma_\phi = 1.27$ phi, or $\sigma = 0.41$ mm. The average diameter $\phi_\phi = 1.55$ phi, or $d = 0.34$ mm. The porosity was 0.45, or the volume concentration $\varepsilon = 0.25$. The sound speed in the water-saturated sand was measured to be 1.83 km/s.

EXPERIMENTAL RESULTS AND THEORETICAL CALCULATIONS

On the basis of the concentrated suspension theory [4,7], the dependence of $\alpha$ upon $f$ was calculated numerically. The experimental and theoretical results are illustrated in Figures 3, 4.

It seems that the curves of the concentrated
suspension theory are quite similar to the experimental ones in variation tendencies; nevertheless, there are distinct differences in numerical values and therefore it is reasonable to say that the rough surfaces and nonspherical shapes of the coarse sand particles at such high frequency will affect strongly on \( \alpha \).

It is shown in Figure 5 that not only the shapes of the sand particles are not spherical with many edges and corners, but also their surfaces are not smooth with many structures full of bumps and holes. This makes clear that there are no reasons to regard particles as spherical ones on the scale of the wavelength of viscous wave when \( \beta D > 1 \), still less as smooth. On the other hand, it can be seen from Figure 2 that the sizes of quite many particles are over an order so that \( HR \) is not far less than 1. When the attenuation resulting from scattering is calculated by the concentrated suspension theory, the contributions of higher-order scattered waves can not be neglected. A reasonable explanation can be given as following: on the one hand, the contact friction between nonspherical rough particles may be much stronger than that between spherical smooth particles (basically rolling friction); on the other hand, the surface areas of the former are greater than that of the latter when their volumes are the same, so do the viscous friction of the former. From above, we can expect that the sound attenuation given by the spherical particle model will be smaller than that from the experiment. It is worthy to note from Figures 4, 5 that the linear dependence of \( \alpha \) upon \( f \) is almost still maintained until \( f \) up to 400-500 kHz even though the mechanism is different from the situation under conditions of low frequency and small spherical particles, as just discussed in last paragraph.

CONCLUSIONS

When acoustic frequency is high enough to make \( DF \) much larger than 1 and \( HR \) not far less than 1, the theoretical \( \alpha \) values calculated on the basis of spherical smooth particle model will be much less than the experimental ones. It indicates that the nonspherical shape and rough surface factors can make acoustic attenuation increase greatly. Although the experiment still gave an approximately linear dependence of \( \alpha \) upon \( f \) at frequency 40-500 kHz, the physical explanation for nonspherical rough particles is quite different and therefore it would be necessary to take shape and surface factors etc. into account. There is some work [10] investigating the problem for ellipsoidal particle theoretically. However, it can not provide a correction as large as that expected in our situation. The hydrodynamical treatment for complex shapes and rough surfaces is very difficult, especially in case of large Reynolds number and when the particles rub against each other in the acoustic field. A semi-empirical method may be practicable for this problem.

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VALUES OF ACOUSTICS PARAMETERS IN N-ALKANE LIQUIDS FROM TONG'S EQUATION BASED ON SCHAFFS' THEORY

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INTRODUCTION

Up to now, study of the relation between microscopical characteristics and macro parameters in liquid material has been the feeblest but the most attractive subject in condensed-state physics. Since the interaction between molecules of liquid sharply depends on the concrete structure of the liquid, the behavior of molecular motion in liquid is still not quite clear. So, it has special significance to connect acoustic parameter with micro characteristics in liquid. In one aspect, it is an important method gaining micro information from macro characteristics. On the other hand, it is the only effective means to expound acoustic characteristics of liquid in the view of micro state. In 1986, the author of this paper based on Schaffs' theory derived out 3 expressions describing acoustic parameters of organic liquid at room temperature and under normal pressure. The expressions, including those of acoustic nonlinearity parameter $B/A$, temperature coefficient of ultrasonic speed and pressure coefficient of ultrasonic speed were successfully used in methanol, carbon tetra chloride, benzene and acetonitrile liquids. The estimated magnitude of $B/A$ value was in agreement with those of about 18 kinds of organic liquids. Especially, the characteristics of the coefficients in organic liquid were well explained by the expressions. However, because of the destitution of thermodynamic parameters, the usefulness of the expressions for other organic liquids was not tested. In this paper, we use the PVT data got by B.Hartmann to testify whether the expressions are suitable for n-alkane liquids.

VALUES OF THE ACOUSTIC PARAMETERS

Expressions based on Schaffs' theory

Schaaffs once got an equation between ultrasonic speed $C$ and molecule radius $r_m$ in organic liquid on the hypothesis that the molecules in the liquid were rigid balls, which is

$$ r_m = \sqrt[3]{\frac{3M}{16\pi\rho N}} \left[ 1 \frac{RT}{MC^2} \left( 1 \frac{MC^3}{3\pi RT} + 1 \right) - 1 \right] $$

where $M$ and $\rho$ are mole mass and density of the liquid respectively, $\gamma$ is the ratio of specific heat capacity under constant pressure $C_p$ with respect to that under constant volume $C_v$, $N$ is Avogadro constant and $R$ is gas constant. We introduced a non-dimensional parameter $X$, which is defined as $3V/16 \times \pi/2N$, reflecting structure condition of the liquid. With Ernst hypothesis, we got following expressions in the liquid:

$$ B/A = J(x) + J(x) $$

$$ (\Delta C/\Delta T) = \frac{1}{2} C \beta \left[ \frac{1}{T^2} - J(x) \right] $$

$$ (\Delta C/\Delta P) = \frac{1}{2} C \kappa J(x) $$

where

$$ J(x) = \frac{2(3-2x)^2 - 3(1-x)(3+2x)}{(3-x-1)(6+5x)} $$

and symbol $\beta = (\partial V/\partial T)/\rho V$ is isobaric expansibility, $\kappa = (\partial V/\partial P)/V$ is isothermal compressibility. Expressions (2)-(4) shows that $B/A$, $(\Delta C/\Delta T)\rho$, and $(\Delta C/\Delta P)\rho$ are closely related with the variable $x$, which is in nature the measure of the relative size of the average distance of molecule with respect to the average scale of the molecule radii. For most organic liquids, the thermodynamic parameters appearing in Eqs.(2)-(4) are not known. We use Hartmann's state equation to get the parameters in n-alkane liquids.

Thermodynamic parameters from Hartmann's PVT data

Hartmann once proposed a concised analytic state equation for n-alkane liquids:

$$ \rho V = T \frac{3}{2} - \alpha V $$

where $\rho = P/P_{cr}$, $T/T_{cr}$, $V = V/V_{cr}$, $P_{cr}$, $T_{cr}$ and $V_{cr}$ are constants for a given material. From above, we have

$$ \beta = \frac{T_0^3}{2 T_0 (1 + 5 \beta V^3)} $$

Because of $\gamma = \frac{\rho V}{T_0 (1 + 5 \beta V^3)}$, we can know

$$ \gamma = \frac{4 T_0 C_p V^3 (1 + 5 \beta V^3) - T_0 \rho V T^3}{4 T_0 C_p V^3 (1 + 5 \beta V^3)} $$

where $C_p$ is specific heat under constant pressure. Now, all thermodynamic parameters required for calculating Eqs.(3)-(4) can be computed.

Values of Eqs.(2)-(4) by Tong's expression

In order to avoid using molecule radius in computing $J(x)$, we rewrite $x$ with Eq.(1),

$$ x = 1 + \frac{1}{U + 2} $$

where $U = \sqrt{9 + \frac{3C_0T}{RT}}$

\[ \]
So, Eq.(6) can be written as
\[
J(x) = \frac{2 A^2}{3(\lambda - 3)} \quad J(x) = \frac{M \beta C^2}{\rho \alpha}
\]  
(11)

Eq.(11) implies \(J(x)\) and \(J(x)\) are determined by the value of \(C\). In the process of calculating Eq.(11), we use \(C_n\) which comes from Eq.(6):
\[
C_n = (1 + 5 \rho V^5 \rho V^5) \left(\frac{\rho V}{\rho V^5} - \frac{\rho V}{\rho V^5} \right)
\]  
(12)

to determine Eq.(11). Using above parameters, we computed the values of Eqs.(2)-(4) for pentane, heptane, octane, nonane, dodecan and hexadecane liquids. The results are shown in table 1 (where \(n\) is carbon number in a molecule of the liquid). Since Eqs.(2)-(4) only work at room temperature and under atmospheric pressure, we take T=293K, P=1atm in calculating the parameters. In table 2, some experimental data are listed. (Units of the parameters are as following: \(C\) in M/(S,AC/\alpha P)^2, in 10^-8M^2/sN, and (AC/\alpha T) in M/SK)

Table.1 Acoustic parameters by expressions (2)-(4)

<table>
<thead>
<tr>
<th>n</th>
<th>C_n</th>
<th>J(x)</th>
<th>J(x)</th>
<th>B/A</th>
<th>(AC/\alpha P)_T</th>
<th>(AC/\alpha T)_P</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>970</td>
<td>8.82</td>
<td>0.61</td>
<td>9.43</td>
<td>0.92</td>
<td>-4.75</td>
</tr>
<tr>
<td>7</td>
<td>1137</td>
<td>10.44</td>
<td>0.70</td>
<td>11.74</td>
<td>0.85</td>
<td>-5.29</td>
</tr>
<tr>
<td>8</td>
<td>1115</td>
<td>10.94</td>
<td>0.60</td>
<td>11.68</td>
<td>0.84</td>
<td>-4.71</td>
</tr>
<tr>
<td>9</td>
<td>1217</td>
<td>11.79</td>
<td>0.61</td>
<td>12.41</td>
<td>0.83</td>
<td>-4.85</td>
</tr>
<tr>
<td>12</td>
<td>1248</td>
<td>13.76</td>
<td>0.69</td>
<td>14.39</td>
<td>0.86</td>
<td>-5.51</td>
</tr>
<tr>
<td>16</td>
<td>1317</td>
<td>16.10</td>
<td>0.54</td>
<td>16.76</td>
<td>0.92</td>
<td>-6.50</td>
</tr>
</tbody>
</table>

Table.2 Experimental data from others^{(6)}

<table>
<thead>
<tr>
<th>n</th>
<th>C</th>
<th>B/A</th>
<th>(AC/\alpha P)_T</th>
<th>(AC/\alpha T)_P</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>1030</td>
<td>7.41*</td>
<td>0.72</td>
<td>-4.60</td>
</tr>
<tr>
<td>7</td>
<td>1158</td>
<td>9.19</td>
<td>0.69</td>
<td>-4.36</td>
</tr>
<tr>
<td>8</td>
<td>1185</td>
<td>9.37</td>
<td>0.65</td>
<td>-4.20</td>
</tr>
<tr>
<td>9</td>
<td>1227</td>
<td>9.53</td>
<td>0.62</td>
<td>-4.08</td>
</tr>
<tr>
<td>12</td>
<td>1298</td>
<td>9.55</td>
<td>0.56</td>
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</tr>
<tr>
<td>16</td>
<td>1367</td>
<td>9.40</td>
<td>0.51</td>
<td>-3.72</td>
</tr>
</tbody>
</table>

*In reference [6] this value is 9.82

DISCUSSION

From above tables, it can be seen that values of alkane liquids from Eqs.(2)-(4) are in agreement in magnitude with those from experiment. For example, estimated B/A values are about 10^8 in magnitude while experimental data are about 10 and estimated (AC/\alpha T)_P, (AC/\alpha P)_T values are closer to those of experiment than B/A values are. These results show that Schaffo's theory is correct in general to describe essential acoustic properties of n-alkane liquids. It's notable that Eqs.(2)-(4) lead to the results only with few hypotheses which did not consider concrete model of the liquids. This shows some general suitability of Schaffo's theory for organic liquids.

In tables 1 & 2, it can be seen that \(J(x)\) is so smaller than \(J(x)\) that B/A value is mainly determined by \(J(x)\). This result tests the predictions of reference^{(6)}. From the view of calculating the values quantitatively, Eqs.(2)-(4) are somewhat faulty for n-alkane liquids. The cause may perhaps appear in two aspects. One is that Schaffo's theory is too simple to describing all details of the organic liquids which are so complex in molecular structure that their properties are of differences among different kinds of liquids. The other cause may be that Hartmann's PVT data have some errors. For example, Hartmann pointed out that while predicting the coefficient \((AC/\alpha T)_P\) in the liquid, the errors is about 6%. It is necessary to point out that the tendencies of B/A value versus number of carbon atoms of a molecule of the liquid are different between Hartmann's prediction and Narayana's prediction. Hartmann's results show the B/A value increases at first and then almost remains constant while the number of carbon atoms in a molecule of the liquid increases. But Narayana's report shows the curve of B/A value versus number of carbon was a typical concave curve which has a minimal point. So, we think, after modified and compensated, Schaffo's theory has the possibility of estimating acoustic properties of n-alkane liquids accurately.

CONCLUSION

From all above, we conclude that Tong's expressions based on Schaffo's theory have some advantages in predicting acoustic characteristics of n-alkane liquids, especially in estimating the magnitude of the acoustic parameters. But in calculating the parameters quantitatively, there are still some disadvantages requiring modified.

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NEARFIELD EFFECT SETOFF METHOD FOR MEASUREMENT SOUND ABSORPTION COEFFICIENT

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1. INTRODUCTION

As the strong coherence of the sound wave in nearfield of a projector, the pressure field in the region is very complex. For the reason, many acoustic research and measurement works are limited in the plane wave field or in the other free field, such as to study sound scattering by a target, calibration of acoustic instruments and measurements of absorption coefficient in medium and so on. But in some situations, the work must be done in nearfield or in transition region. For example, to measure sound absorption coefficient in dense suspension fluid is the situation.

There is another thing that the hydrophones in practical applications usually have a certain size of their receptive face. Their maximum dimensions may be many times the size of the wave length, called "big" receiver. The receptive pressure by them is assembly pressure in local region and it will produce serious effect on measurement results.

According to coherence and incoherence of sound waves, the nearfield or the farfield is defined. In the optical diffraction theory [1], the nearfield is Fresnel diffraction region, the farfield is Fraunhofer region. There is not a clear limit between them. In practical application, if \( \lambda d \gg 1 \) ( \( \lambda \) - wave length, \( d \) - radius of the radiator, \( y \) is in ideal farfield. Oppositely, there is in nearfield or in transition region.

The numerical simulation method is used to study the nearfield and the smooth effect of pressure field by using a "big" receiver, to study the influence of measurement absorption coefficient. And nearfield effect setoff method (NBSM) is advanced.

2. RADIATION SOUND FIELD, RECEIPT PRESSURE AND ABSORPTION COEFFICIENT

There is not a accurate solution for the radiation field of a radiator. It is usually expressed by Green function integral. The projector of a piston has radius \( a \) and vibrates with simple harmonic motion normal to its face, the pressure at \( r \) in fig.1 can be written as [3]

\[
\rho_0(\mathbf{r}) = \frac{2\pi\rho_0}{\sqrt{n}} \left[ \int_{-\infty}^{\infty} \psi(\mathbf{r}') \frac{e^{i\mathbf{k} \cdot \mathbf{r}' - \omega t}}{\sqrt{\mathbf{r}' \cdot \mathbf{r}'}} \, d\mathbf{r}' \right] \tag{1}
\]

\( \rho_0 \) are medium density and velocity of sound, respectively. \( k \) is wave number, \( \omega \) is peak amplitude of transducer velocity. If the receiver is "big", the receptive pressure is integral of (1)

\[
\rho(\mathbf{r}) = \int_{-\infty}^{\infty} \psi(\mathbf{r}') \frac{e^{i\mathbf{k} \cdot \mathbf{r}' - \omega t}}{\sqrt{\mathbf{r}' \cdot \mathbf{r}'}} \, d\mathbf{r}' \tag{2}
\]

\( \rho(\mathbf{r}) \) is receiver radius. In the calculations, \( a^2 \gg 1 \) instead of \( b=0 \), called "point" receiver. The receptive pressure by it in fact is the radiation field of the projector.

If there is sound absorption medium, it is a complex number, \( k=\rho+i\sigma \). \( N \) is sound intensity attenuation coefficient in a wave length range. At \( 1, \omega \), the relationships are

\[
\rho(\mathbf{r}) = \frac{2\pi\rho_0}{\sqrt{n}} \left[ \int_{-\infty}^{\infty} \psi(\mathbf{r}') \frac{e^{i\mathbf{k} \cdot \mathbf{r}' - \omega t}}{\sqrt{\mathbf{r}' \cdot \mathbf{r}'}} \right] e^{-\sigma d / \rho} \tag{2}
\]

On the basis of definition of absorption coefficient [4],

\[
\sigma \approx 30\log_{10} N \quad \text{OR} \quad N = (3.583/\rho_0)^{1/2} \tag{3}
\]

Use (1),(2),(3), the radiation field of a projector and the nearfield effect on measurement absorption coefficient can be calculated.

3. CALCULATING RESULTS

In the calculations, all the length measurements were changed to dimensionless by multiplying \( \lambda / \rho_{02} \). The integrals were numerically evaluated on a VAX8650 computer and four points were sampled in the range of one wave length. The calculating range on the axis was limited in 0.02 < \( \tau < 4 \), step 0.02.

(1) Nearfield Pressure on Axis and Receptive Pressure by "big" Receiver

Figure 2, 3 are numerical results for \( a \lambda = 3, 9, b=0, 0.5 \lambda, 4 \lambda \), respectively.

![Fig.1 Coordinate system](image1)

![Fig.2 Magnitude of on-axis pressure, a/\lambda=3](image2)

![Fig.3 Magnitude of on-axis pressure, a/\lambda=9](image3)

The figures indicate that there is strong influence of the pressure on axis of \( \lambda < 1 \) along with various distance when used "point" receiver (b=0) because the coherence effect. In high frequency, the influence is strong. From the figures, it seems that there is the region of spherical regular attenuation of pressure in \( \lambda > 1 \). We shall see shortly that the region is not nearfield. It is clear that the pressure by "big" receiver is the longitudinal local superposition of pressure. For this reason, it greatly smoothes the distribution of nearfield pressure. Therefore, it has to be considered its effect when use a "big" receiver in nearfield.
2. Calculations Nearfield Effect for Measurement Absorption Coefficient

In the first, we suppose \( k = k_0 \sqrt{N/2} \) in (1), then calculate \( \alpha \) and \( N \) by (2)(3). It can be affirmed that if in free field, the counted \( N \) value equals \( N_0 \) as the all sound rays reaching on the receptive face at \( s = 2 \) have same propagation distance. But in the nearfield, it is not the case.

![Fig.4 Calculation value of N](image)

In the counting, \( N = 0.001, x = 1, 1.5, 2, 3 \) respective-ly, \( x = 0 < s \leq 4 \). Fig.4 show the results of calculations of \( N \) for \( d = 1 \) for different size receivers. It expresses that even if in the region of \( x = 1 \), no matter how "big" receiver is used, there is great difference between \( N \) and \( N_0 \). It is evident that the region of \( x = 1 \) is not sure farfield.

The reason can be expressed by the mathematical formula

\[
|P_k| = A \sum_{j=1}^{n} e^{-j2\pi \frac{Y_{kj}}{\lambda}}
\]

(4)

If \( k = 1 \) or 2 depends on at \( s = 2 \) or \( s = 1 \). If in the free field, \( N \)

\[
N = 2 \left( \frac{1}{\lambda} \sum_{j=1}^{n} \frac{e^{-j2\pi \frac{Y_{kj}}{\lambda}}}{Y_{kj}} \right) \left( \frac{1}{\lambda} \sum_{j=1}^{n} \frac{e^{j2\pi \frac{Y_{kj}}{\lambda}}}{Y_{kj}} \right)
\]

(5)

That is obvious that \( N \) value counted by (5) doesn't equal \( N_0 \) if there are big difference between \( Y_{k} \) and \( Y_{k,j} \).

4. NEARFIELD EFFECT SETOFF METHOD (NEISM)

The formal (6) expresses the nearfield effect on measurement absorption coefficient. The problem how to measure the coefficient accurately is how to find \( N_0 \) from \( p \) in (4). (4) can be written

\[
|P_k| = A \sum_{j=1}^{n} e^{-j2\pi \frac{Y_{kj}}{\lambda}}
\]

(6)

If \( N_0 \) is 0, it becomes

\[
|P_k| = A \sum_{j=1}^{n} \frac{e^{-j2\pi \frac{Y_{kj}}{\lambda}}}{Y_{kj}}
\]

(7)

In (8), because \( 0 < Y_{k,j} < \lambda \), \( N_0 \) has a small value, therefore, we have \( \exp(-N_0 Y_{k,j}) \approx 1 \). From (6)(9), the formula can be obtained

\[
|P_k| / |P_k| = e^{-j2\pi \frac{Y_{k}}{\lambda}}
\]

\[
|P_k| / |P_k| = \exp(-j2\pi \frac{1}{\lambda}) \approx N_0
\]

(8)

(8) is the mathematical expression of NEISM.

![Fig.5 NEISM Value of N, d=0.3](image)

![Fig.6 NEISM Value of N, d=0.9](image)

The figures (5)(6) show the numerical results of the method. The table 1 lists the relative error of \( N \) to

<table>
<thead>
<tr>
<th>( \frac{N}{N_0} )</th>
<th>( a )</th>
<th>( a )</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>1.0</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>1.5</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>2.0</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>2.5</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>3.0</td>
<td>0.56</td>
<td>0.6</td>
</tr>
<tr>
<td>3.5</td>
<td>0.56</td>
<td>0.6</td>
</tr>
</tbody>
</table>

\( * \), expresses the error < \( 1% \).

\( N_0 \) it indicates that the method not only widens the measurement range in nearfield but gives the accurate results. For the special situation, in high frequency and "big" receiver, it can do measurement at any region of the nearfield.

The work was supported by KAY STATE LABORATORY OF SOUND FIELD AND SOUND INFORMATION, NO. 19904.

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3. 梁华明, 张兰锋, 王曾:"压电换能器和换能器阵", (北京:北京工业, 1980)
SOUND ATTENUATION IN LINED CIRCULAR DUCTS WITHOUT FLOW

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** Faculty of Engineering, Kobe University, Rokkodai, Kobe, 657 JAPAN

1. INTRODUCTION

Lined circular and rectangular ducts have been studied for practical use because of the large amount of sound attenuation for their simple structure [1,2]. As far as we know, however, any appropriate design charts of sound attenuation with frequency based on rigorous calculation have been not reported. We have examined to give design charts as parameter with complex reflection coefficients at the surfaces of sound-absorbing materials.

In this report, numerical results of eigenvalues and sound attenuation are shown for different values of reflection coefficients.

2. THEORETICAL EQUATIONS

It is assumed that a circular duct without air flow is entirely lined with a single layer of sound-absorbing material of uniform density. The geometry of a lined circular duct and cylindrical coordinate $(r, \phi)$ are shown in Fig. 1. The relation between a reflection coefficient $R = |R| \exp(i\theta)$ and impedance $z_1 = \rho_0 c_k (\phi + i \chi)$ at $r = \alpha_1$ is expressed by the following equations:

\[
\phi = \frac{1 - |R|^2}{1 + |R|^2 - 2 |R| \cos \theta} \quad (1)
\]

\[
\chi = \frac{2 |R| \sin \theta}{1 + |R|^2 - 2 |R| \cos \theta} \quad (2)
\]

Introducing a variable $F$ instead of eigenvalues $\alpha_i$.

\[
F = \frac{\alpha_i^2}{2}, \quad (3)
\]

Rice [1] has given a nonlinear differential equation for eigenvalue:

\[
(F + \varepsilon^2) \frac{dF}{d\varepsilon} = F \quad (4)
\]

where $\eta = D/\lambda$ and $\varepsilon = (i \pi / 2) \eta (\rho_0 c / z_1)$.

He also has given a different type of equation

\[
(1 + \xi^2) \frac{dF}{d\xi} + F = 0 \quad (5)
\]

where $\xi = 1/\varepsilon$. Equation (4) gives $\alpha_i \varepsilon = \eta = 0$, 3.8317, 7.0156, ..., and eq.(5) $\alpha_i \varepsilon = 2.4048, 5.5201, \ldots$ for $j = 1, 2, 3, \ldots$, respectively.

The sound power at any cross-section in a lined circular duct is calculated by the equation

\[
E_{jk} = \frac{4}{\eta} \frac{\eta_0 c_k}{\eta} \exp \left[ - \frac{\omega}{c_k} \left| \alpha_j + \alpha_k \right| i \left( t_k - t_j \right) \right] \frac{(\alpha_j + i t_j)(\xi + \xi^*)}{Q_k \alpha_k (\alpha_k^2 - \alpha_j^2)} \quad (6)
\]

where $Q_k = (\xi^2 \alpha_k / \pi) \eta^2 - 1$. Both $\sigma$ and $\tau$ are function of $\alpha$ and $\eta$. The subscripts $j, k$ signify the pressure mode $j$ and the axial velocity mode $k$, respectively. The sound power at a cross section is obtained by summation of $E_{jk}$, which are not zero only for $j = k$. The sound attenuation between two points along the $z$ axis, separated by distance $L$, is estimated from respective sound powers.

3. NUMERICAL SOLUTION OF EIGENVALUE EQUATION

To get eigenvalues, eqs.(4) and (5) are solved for a large and small impedance, respectively. In series expansion method, the solution $F$ is expanded to the series.

\[
F = \sum_{n=0}^{\infty} a_n \cdot \varepsilon^n, \quad or \quad (7)
\]

\[
F = \sum_{n=0}^{\infty} b_n \cdot \xi^n \quad (8)
\]

The coefficients $a_n, b_n$ are determined from $\alpha_j \theta$ using recurrence formulas. The coefficient $a_0$ are found to be almost linear to $n$, which is shown in Fig.2. Thus we determined convergent radius from the slope of lines for each $j \text{ mode}$. Table 1 shows convergent radius $\gamma_0$ for eq.(7) and $\gamma_1$ for eq.(8).

4. COMPUTATIONAL RESULTS AND DISCUSSION

For given value of $L/\theta$, the sound attenuation was computed with parameters of amplitude $|R|$ and phase $\theta$ reflection coefficients. It was confirmed that the contributions from modes of higher order more than 5 are negligible.

Tables 2 and 3 show numerical results of eigenvalues calculated by both Runge-Kutta and series expansion method for $\eta = 1, L/\theta = 3$. Table 4 shows sound attenuation with $R$ and $\theta$. Sound attenuation is found to decrease with increasing $|R|$ for constant $\theta$. 

A8-10
5. CONCLUSION

The eigenvalues calculated by the series expansion method agree with those by the Runge-Kutta method. The sound attenuation decreased with magnitude of reflection coefficients.

REFERENCES


FIGURES AND TABLES

Fig.1. Geometry of a lined circular duct and cylindrical coordinates(r, θ).

Fig.2. Expansion coefficient a_n in equation (7).

<p>| Table 1. Convergent radius γ and γ for eqs. (1) and (3), respectively. |
|-----------------|-----------------|-----------------|-----------------|</p>
<table>
<thead>
<tr>
<th>j</th>
<th>γ</th>
<th>j</th>
<th>γ</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>0.92</td>
<td>6</td>
<td>0.18</td>
</tr>
<tr>
<td>2</td>
<td>0.92</td>
<td>7</td>
<td>0.18</td>
</tr>
<tr>
<td>3</td>
<td>0.46</td>
<td>8</td>
<td>0.18</td>
</tr>
<tr>
<td>4</td>
<td>0.31</td>
<td>9</td>
<td>0.18</td>
</tr>
<tr>
<td>5</td>
<td>0.24</td>
<td>10</td>
<td>0.18</td>
</tr>
</tbody>
</table>

| Table 2. Numerical results of eigenvalues for R=0.5<θ<90°, η=1 and L/D=3. |
|-----------------|-----------------|-----------------|
| Runge-Kutta     | Series-expansion |
| method          | method          |
| a               | 1.0719±1.4419   | 1.0719±1.4419   |
| b               | 0.6290±0.3044   | 0.6290±0.3044   |
| a               | 6.0305±0.1607   | 6.0305±0.1607   |
| a               | 10.1001±0.1101  | 10.1001±0.1101  |
| a               | 10.278±0.0839   | 10.278±0.0839   |

| Table 3. Numerical results of eigenvalues for R=0.9<θ<90°, η=1 and L/D=3. |
|-----------------|-----------------|-----------------|
| Runge-Kutta     | Series-expansion |
| method          | method          |
| a               | 0.2748±3.1580   | 0.2733±3.1874   |
| a               | 3.1992±0.0521   | 3.1995±0.0525   |
| a               | 6.6489±0.0364   | 6.6489±0.0366   |
| a               | 9.1969±0.0263   | 9.1979±0.0265   |
| a               | 13.1266±0.0205  | 13.1278±0.0206  |

| Table 4. Numerical results of sound attenuation (Att) for R=|R|<θ<90°, η=1 and L/D=3. |
|-----------------|-----------------|
| |R|<θ<90°| Att (dB) |
|-----------------|-----------------|
| 0.1<θ<10°       | 36.18 |
| 0.3<θ<20°       | 31.65 |
| 0.5<θ<30°       | 24.28 |
| 0.7<θ<40°       | 11.62 |
| 0.9<θ<90°       | 10.20 |
Cylindrical Liquid-Filled Acoustic Lens
Pu Lin Tai
Hangzhou Applied Acoustics Research Institute, China

I. INTRODUCTION

Recently, the cylindrical liquid-filled acoustic lens research and their application have drawn much attention. However, from the reference materials we have got, the lens work over several hundreds kilohertz, and the theory is based on the ray acoustics. Can the lenses mentioned above be used at lower frequencies? Which range of frequencies is suitable for? What is the relation between characteristic of lenses and properties of the liquid material within lenses? Until now, we have not seen any articles discussing such problems. By wave acoustics, the internal acoustic pressure field in lenses is derived and the possibility of the use of the cylindrical liquid-filled acoustic lens at lower frequencies discussed in the paper. The relation between characteristic of lenses and properties of liquid materials filled in the lens is also analyzed. The essential conclusions have already been confirmed by experiments.

II. THE MATHEMATICAL MODEL TO DESCRIBE THE ACOUSTIC PRESSURE FIELD

The cylindrical lens are made of two half-cylinders, as shown in Fig. 1, in which the thin-line presents the sound transmission window for the thin half-cylinder whose radius is a. It confines to the form of liquid and affects focusing property. Thick-line presents other half-cylinder of the lens, and the radius is not always a. The liquid filled in the cylinder is different from the one out of the cylinder. It is assumed that the density and sound velocity of liquid out of the cylinder are ρ0 and c0, respectively, and those in cylinder are ρ and c, respectively. If the cylinder is long enough and the shell thin enough, then we can neglect the influence of the cylinder vibration. As shown in Fig. 1, sound wave of perpendicular to axis z is incident along axis z.

If the incident wave is a plane wave, the incident wave at point r from the origin of the coordinate along the direction φ can be described as:

\[ P(x, y) = \sum_{n=0}^{\infty} \frac{1}{r} J_0(na) \frac{\partial J_0(na)}{\partial n} \cos \phi \]

where

- \[ e = \begin{cases} 1 & \text{when } n=0 \\ 2 & \text{when } n \neq 0 \end{cases} \]

\[ J_0, J_1 \] are Bessel functions.

\[ \frac{\partial J_0(na)}{\partial n} \] is length of sound wave in the medium out of the cylinder.

When the plane wave is incident on the cylinder, the scattered wave out of the cylinder is:

\[ P(x, y) = \sum_{n=0}^{\infty} \frac{1}{r} J_0(na) \frac{\partial J_0(na)}{\partial n} \cos \phi \]

where \( B_0 \) is underdetermined coefficient, \( n+1 \) is one kind of a step Bessel function.

The wave in the cylinder is:

\[ P(x, y) = \sum_{n=0}^{\infty} \frac{1}{r} J_0(na) \frac{\partial J_0(na)}{\partial n} \cos \phi \]

where \( B_0 \) is underdetermined coefficient.

In order to determine \( A_0 \) and \( B_0 \), we use boundary conditions of pressure and the velocity continuity. Consequently

\[ A_0 = \begin{bmatrix} 0 \end{bmatrix} \]

and

\[ J_0(na) \frac{\partial J_0(na)}{\partial n} \]

\[ J_0(na) \frac{\partial J_0(na)}{\partial n} \]

where \( n \geq 2 \pi / \lambda \) is the length of sound wave in the medium out of the cylinder.

\[ \frac{\partial J_0(na)}{\partial n} \]

is the length of sound wave in the medium out of the cylinder.

3. THE EXAMINATION OF RELIABILITY OF THE THEORY

The reliability of the theory is approved from the following experimental data. The radius of the cylinder "a" is 10 cm. The cylinder is filled with silicone oil.

Table 1 shows the experimental and theoretical results on focusing

<table>
<thead>
<tr>
<th>Frequency (kHz)</th>
<th>Theoretical Value</th>
<th>Experimental Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>31.5</td>
<td>7.2</td>
<td>7.4</td>
</tr>
<tr>
<td>50.0</td>
<td>9.7</td>
<td>9.7</td>
</tr>
<tr>
<td>100.0</td>
<td>10.4</td>
<td>10.4</td>
</tr>
</tbody>
</table>

2. Comparison of theoretical and experimental results on directivity

Fig. 2 shows the directivity of theoretical (the dotted line) and measured results (the solid line) of the cylindrical liquid-filled lens. The sound non-transmission window of the lens for the experiment adheres to absorptive rubber, thus, the back lobe of the measured directivity is small.

As stated above, the comparison of theoretical and experimental results is quite coincident. It can be concluded that the theoretical calculation is reliable.
focusing characteristic curves are different according to change of frequencies.

2. The higher frequency is, the better focusing is. There is obvious tendency to central energy at higher frequencies rather than at lower ones, which illustrate that lens' properties are not good at lower frequencies, and better at higher frequencies.

3. Generally speaking, focus changes with frequencies, and the higher frequency is, the greater the focus is. At last, the tendency is that the focus is determined according to the ray acoustics. However, quantity of the focus varies with liquid refraction index. The focus changes more greatly with frequency when the liquid refraction index is smaller (for example, silicone oil). The focus varies slowly with frequency when liquid refraction index is bigger (for example, Fluorocarbons).

For this reason, the appropriate frequency band of cylindrical lens filled with great refraction index liquid is broad. Table II shows the focus of the cylindrical liquid-filled lens (f = 10kHz).

<table>
<thead>
<tr>
<th>Frequency (kHz)</th>
<th>10.0</th>
<th>31.5</th>
<th>50.0</th>
<th>100.0</th>
<th>200.0</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Acoustic</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
</tr>
<tr>
<td>Tic</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
</tr>
<tr>
<td>Silicone oil</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
</tr>
<tr>
<td>Fluorocarbons</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
<td>15.0</td>
</tr>
</tbody>
</table>

4. The bigger refraction index of liquid filled in the cylinder is, the better focusing is, and the lower side lobe of lens directivity is. Consequently, when liquid refraction index is bigger, not only the appropriate frequency range is broad, but also the focusing is better.

5. When a point source hydrophone is put on symmetrical axis, the lens shows obvious directivity. However, the better focusing of the point with the hydrophone is, the better lens directivity is.

6. Comparing a planar array with the same effective aperture, the beam width of the cylindrical liquid-filled lens would be narrower and the side lobe is smaller. Table III shows the data of beam width under different conditions.

<table>
<thead>
<tr>
<th>Frequency (kHz)</th>
<th>10.0</th>
<th>31.5</th>
<th>50.0</th>
<th>100.0</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>The lens filled with silicone oil</td>
<td>24</td>
<td>10</td>
<td>7.2</td>
<td>3.2</td>
</tr>
<tr>
<td>The lens filled with Fluorocarbons</td>
<td>24</td>
<td>10</td>
<td>7.4</td>
<td>3.2</td>
</tr>
<tr>
<td>The lens filled with Fluorocarbons</td>
<td>24</td>
<td>10</td>
<td>7.4</td>
<td>3.2</td>
</tr>
<tr>
<td>The planar array whose width is 20cm</td>
<td>37.8</td>
<td>12.5</td>
<td>7.5</td>
<td>3.8</td>
</tr>
</tbody>
</table>

For acoustic lens filled with silicone oil, the conclusion mentioned above has been confirmed by experiments.

ACKNOWLEDGMENT

The experimental study of cylindrical liquid-filled acoustic lens is conducted by our research team. It is a pleasure to acknowledge Dr. Gong Chong-De, Tans Chong-Shou and Ms. Song Lan-Ting.

![Fig. 1 Cylindrical liquid-filled acoustic lens](image)

Fig. 1: Cylindrical liquid-filled acoustic lens

![Fig. 2 Comparison of theoretical (the dotted line) and experimental (the solid line) results on directivity (f=100kHz)](image)

Fig. 2: Comparison of theoretical (the dotted line) and experimental (the solid line) results on directivity (f=100kHz)

![Fig. 3 The focusing characteristic curves of acoustic lens (for two liquid materials: 1) silicone oil, 2) Fluorocarbons, f=100kHz)](image)

Fig. 3: The focusing characteristic curves of acoustic lens (for two liquid materials: 1) silicone oil, 2) Fluorocarbons, f=100kHz)

![Fig. 4 The directivity curves of acoustic lens (for two liquid materials: 1) silicone oil, 2) Fluorocarbons, f=100kHz)](image)

Fig. 4: The directivity curves of acoustic lens (for two liquid materials: 1) silicone oil, 2) Fluorocarbons, f=100kHz)

V CONCLUSION

Our study shows that although a point source hydrophone adopted in cylindrical liquid-filled lens is only one, the directivity of acoustic lens is better than that of the array with the same aperture. This acoustic lens can be used at the lower frequency range and also at the width band range. It will probably be used in engineering.
ELASTIC WAVE DIFFRACTION FIELD GENERATED BY A SOURCE WITH FINITE APERTURE ON SURFACE OF ANISOTROPIC MEDIUM

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INTRODUCTION

The angular spectrum theory, a current scalar theory, is used to analyze the diffraction field generated by a finite-aperture source on the surface of the anisotropic medium [1][2]. Such a theoretical analysis is insufficient and has many shortcomings. First, it adopts a scalar field model as in optical split diffraction, while elastic wave field in solid is a vector field. The diffraction patterns for various components of vector field are different generally. Secondly, for various wave modes, in addition to anisotropic difference of acoustic velocity that the angular spectrum theory estimates only, the diffraction patterns may be great different also. Thirdly, the nature of different sources are totally ignored in the angular spectrum theory. In fact, the 'exciting efficiencies' of different mechanical sources on surface have considerable distinction often. Last, the angular spectrum theory neglected actual physical field can only give a relative distribution of phenomenological field.

Recently, we developed generalized theory of surface excitation for complex piezoelectrics which have piezoelectricity and anisotropy [3][4]. According to the generalized theory, the diffracted elastic wave fields generated by a finite aperture source of the surface of anisotropic media were obtained adequately and precisely.

ONE-DIMENSION FINITE APERTURE SOURCE

Suppose a one-dimension source distribution on the surface, \( x_0 = 0 \), of anisotropic media

\[
T(x) = \{ T_{1} \} = \{ T_{1}, T_{2}, T_{3} \} = \{ T_{1}, T_{2}, T_{3} \} \quad (x_0 \leq x < a)
\]

which is infinitely extended along \( x \) and has a finite aperture of 2a. The elastic wave fields generated by this source distribution

\[
U(x_1, x_2) = \{ u \} = \{ u_1, u_2, u_3 \}
\]

are a dot product convolution of the Green's function matrix

\[
G(x_1, x_2) = \begin{bmatrix}
G_{11} & G_{12} & G_{13} \\
G_{21} & G_{22} & G_{23} \\
G_{31} & G_{32} & G_{33}
\end{bmatrix}
\]

and force source vector \( T(x) \) in real space

\[
u = \sum \nu^{(o)}
\]

\[
u^{(o)}(x_1, x_2) = \int G(x_1, x_2)T(x_1)\delta(x_2 - x_2)dx_1
\]

or an inverse Fourier transformation of the product of Green's function \( \hat{G}(\omega) = \{ G_{ij} \} \) and forces \( \{ T_\omega \} \) in \( k \)-space, which are Fourier transforms of \( G \) and \( T \) respectively.
\begin{equation}
T_\tau(x) = C \cdot \left( \frac{x_1}{2a} \right) e^{ix_1/2a}
\end{equation}

(18)

\begin{equation}
\mathbf{T}(a) = C \cdot \frac{\sin a a}{a a}
\end{equation}

(19)

in general cases, the nonuniform force source distribution for \( x_1 \), between \(-a \) and \( a \), can divide enough small element sources with width \( 2\Delta \), which have uniform amplitude and phase, to generate the element elastic wave fields, \( \delta u_m \). Later, the vector superposition of element fields can constitute all diffraction wave fields, i.e. \( u = \sum \delta u_m \). Furthermore, at the far field, \( \delta u_m \) can be obtained on using the stationary phase method as transforming to the polar coordinate system \((r, \theta)\):

For our exact theory,

\begin{equation}
\delta u_m(r, \theta) \sim \frac{\mu n}{\sqrt{2\pi}} \sum \frac{1}{a_n} \left( \frac{B_n(a_n)W_n(a_n)}{\det \Pi(a_n)} \right) \cdot \left( \frac{2\sin(a_n \Delta)}{a_n} \delta \tau \right) \cdot \exp\left[i(a_n \sin \theta + \gamma_n \cos \theta)\right]
\end{equation}

(20)

For angular spectrum theory,

\begin{equation}
\delta u_m(r, \theta) \sim \frac{\mu n}{\sqrt{2\pi}} \sum \frac{1}{a_n} \left( \frac{2\sin(a_n \Delta)}{a_n} \delta \tau \right) \cdot \exp\left[i(a_n \sin \theta + \gamma_n \cos \theta)\right]
\end{equation}

(21)

where \( a_n \) satisfy:

\begin{equation}
(\frac{d^2}{da^2})_{a = a_n} \gamma_n = -ig\theta
\end{equation}

(22)

and

\begin{equation}
(\frac{d^2}{da^2})_{a = a_n} \gamma_n = \cos \theta
\end{equation}

(23)

TWO-DIMENSION FINITE APERTURE SOURCE

The elastic wave fields, \( U \), generated by two-dimension surface distribution, \( T \), within finite region \( R \) are double dot–product convolution of the Green’s function matrix \( G \) and the force source distributions in real space or express as double inverse Fourier transformation of the product of Green’s function \( G \) and force \( \mathbf{T} \) in \( k \)-space, which are Fourier transform of \( G \) and \( T \) respectively. The theoretical expressions are omitted here for lack of space.

EXAMPLE FOR THE STRIP SOURCE ON SURFACE OF 6MM CRYSTAL

Assume the rotated symmetric axis of a 6mm crystal lies on substrate surface and is parallel to infinite extended direction, \( x_1 \) axis. On the surface is a strip source with width \( 2a \) extending infinitely in the direction of \( x_1 \). Such a source driving under uniform normal force, \( T_{10} \), can generate the longitudinal wave and a shear wave in crystal, while another shear wave mode is decoupled. The formulae of the diffraction fields for 6mm crystal are omitted here also. Only the numerical result are shown in Figures 1 and 2, which give the diffraction patterns at distance \( 25\Delta \), respectively for longitudinal and shear wave modes, generated by an uniform normal force source, \( T_{10} \), on surface of PZT–5H material with aperture \( 2a = 25\Delta \). Comparing the results of our theory corresponded to \( (b) \) in Figs 1 and 2 with those of the angular spectrum theory corresponded to \( (a) \) in Figs 1 and 2, we find that for \( u^{(1)} \) component of longitudinal wave mode they are similar and for \( u^{(2)} \) component have a significant difference, while for either component of shear wave mode, \( u^{(3)} \) or \( u^{(4)} \), they have essential distinction.

REFERENCES

The development of a surface acoustic wave-based method for nondestructive evaluation of materials has been described in previous works. In this paper, we consider the reflection of Rayleigh waves from a strip on an elastic half-space. The problem is solved in two dimensions, assuming uniform wave phase speeds and neglecting diffraction effects. The dispersion relations are given by the solutions of the following equations:

\[ \sqrt{K_0^2 - K_1^2} = \sqrt{K_0^2 - \xi_0^2} \]

where \( K_0 \) and \( K_1 \) are the wave numbers of the incident and reflected waves, respectively, and \( \xi_0 \) is the wave number of the surface wave. The root \( \sqrt{K_0^2 - K_1^2} \) corresponds to the scattered wave, and the root \( \sqrt{K_0^2 - \xi_0^2} \) corresponds to the transmitted wave. The solutions for the wave amplitudes are expressed in terms of the following functions:

- \( \Phi \) - the phase function
- \( \Psi \) - the amplitude function

The dispersion relations are obtained by substituting the above expressions into the general solution of the wave equation. The specific solutions are given by:

\[ P = \frac{4}{A} \left[ \begin{array}{c} \Phi \xi_0 \sin \theta - \xi_0 \Phi \cos \theta + \Phi \xi_0 \cos \theta - \Phi \xi_0 \sin \theta \\ \xi_0 \Psi \cos \theta - \Phi \Psi \sin \theta + \xi_0 \Psi \sin \theta - \xi_0 \Psi \cos \theta \end{array} \right] \]

where \( A \) is a constant determined by the boundary conditions. The displacement fields are expressed in terms of the wave amplitudes and phases, and the pressure field is obtained by taking the divergence of the displacement field. The general solution can be written as:

\[ \text{displacement field} = \Phi \xi_0 \sin \theta + \xi_0 \Phi \cos \theta + \Phi \xi_0 \cos \theta + \Phi \xi_0 \sin \theta \]

The stress field is obtained by taking the minus curl of the displacement field. The general solution for the stress field is given by:

\[ \text{stress field} = \nabla \times \left[ \begin{array}{c} \Phi \xi_0 \sin \theta + \xi_0 \Phi \cos \theta + \Phi \xi_0 \cos \theta + \Phi \xi_0 \sin \theta \\ \xi_0 \Psi \cos \theta - \Phi \Psi \sin \theta + \xi_0 \Psi \sin \theta - \xi_0 \Psi \cos \theta \end{array} \right] \]

The solutions for the displacement and stress fields are used to calculate the scattering and transmission coefficients. The displacement and stress fields are continuous across the boundary between the strip and the half-space.
Here $\varepsilon=\pi^2q^2\rho_A/\rho_s^2$ is the basic parameter of the problem that compares linear mass of the strip with the oscillating mass of the soundguide in the region of $\lambda_1, \lambda_2$ cross section ($\lambda_1, \lambda_2$ wavelength of bulk transversal wave, $D(q)$ is the Rayleigh determinant $D(q)=4q^{-2}\rho_s^2-(q^2+s^2)^2$, which is equal to zero in the pole points $q=\pm q_0$.

$I_{xx}$, $I_{yz}$ and $I_{yy}$ are the integrals which we calculate numerically.

\[
I_{xx} = \int_{-\infty}^{\infty} \frac{\Pi(q)}{D(q)} \left( 2q^2s^2 + 4psq^2 - (q^2+s^2)^2 \right) dq
\]

\[
I_{yy} = \int_{-\infty}^{\infty} \frac{\Pi(q)}{D(q)} \left( 2psq^2 + 4psq^2 - (q^2+s^2)^2 \right) dq
\]

\[
I_{yy} = \int_{-\infty}^{\infty} \frac{\Pi(q)}{D(q)} \left( 2psq^2 + 4psq^2 - (q^2+s^2)^2 \right) dq
\]

\[
\Pi(q) = \sin(qa)/(qa)
\]

For the heavy strip the maximum of reflection factor is independent of frequency and is determined by the Poisson factor of the medium & the angle of incidence. For $\theta=0.17$ (melted quartz) $R_{\max}=0.41$.

The theoretical data presented here show that rather great part of incident Rayleigh wave energy reflects from the heavy strip at the oblique incidence ( $R_{\max}=0.27$, at $\theta=0$, $\theta=0.27$). This fact must be taken into account in SAW technology.
ON THE DETERMINATION OF THE TRANSVERSAL WAVE VELOCITY IN PLATES

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INTRODUCTION

Ultrasonic velocity measurements can be made in either time or frequency domain. In the time domain approach, the time-of-flight of a short pulse between two faces of the specimen is measured [1]. In the frequency domain methods, the fast-Fourier-transform (FFT) analysis of the reflected signal provides minima from which the sound velocity can be calculated. The transmitted echo sequence can be used as well, but then the frequency of the maxima has to be used instead [2] [3].

Both approaches can be used to determine either the longitudinal wave velocity or the shear wave velocity. To measure the shear wave velocity, shear waves are generated either by means of a shear transducer mounted on the plate with a compliant or by insonifying it in a water tank at oblique incidence (typically a few degrees) [4]. In the latter case things are complicated by the fact that the observed shear resonances are generally weak and mixed up with the stronger longitudinal resonances.

In this paper we present a method for measuring the shear wave velocity in plates, which avoids this difficulty by working. The proposed procedure is tested on three samples: an aluminum, a glass and a polyethylene plate.

THEORETICAL APPROACH

We present a submersible technique to determine the shear wave velocity in plates. It takes advantage of the fact that above the first and below the second critical angle of the liquid/solid interface, only shear waves propagate in the whole thickness of the plate. Two lamellar evanescent [5] waves propagate along the plate axis, one is confined at the upper boundary and the other at the lower boundary. The potentials of these waves are given by (figure 1):

\[ \phi_{upper} = \exp(-kh_x x) \exp(ik_y x) \]  
\[ \phi_{lower} = \exp(kh_x x) \exp(ik_y x) \]  

with \( k_x \), equal to:

\[ k_x = \sqrt{\frac{k^2 - \frac{\omega^2}{v_x^2}}{v_y^2}} \]  

where \( \omega \) is the angular frequency and \( \theta \) is the angle of incidence, \( k_x \) is equal to \( \frac{\omega}{v_x} \) sin(\( \theta \)).

The decay constant (d.c.) of the lamellar evanescent waves is equal to \( k_x^2 \):

\[ \text{d.c.} = \frac{v_w}{\omega} \frac{v_i}{\sqrt{(v_i sin(\theta))^2 - v_x^2}} \]  

In the low frequency limit the decay constant is large compared to the thickness of the plate and the longitudinal wave propagating at the upper boundary of the plate reaches up to the lower boundary and vice versa. For the shear waves the phase of these boundaries therefore have a reflection coefficient, the phase of which is frequency dependent. In the high frequency limit on the contrary, the decay constant is small and the longitudinal waves do not reach up to the opposite boundary. The internal reflections of the shear wave sagaaging in the plate are then similar to reflections on a halfspace, i.e. with a reflection coefficient, the phase of which is now frequency independent. This allows us to calculate the shear velocity from the frequency difference \( f_{n+1} - f_n \) between two successive transversal modes of the plate. This is done starting from the dispersion relation of the N-th mode, which in the framework of a raymodel can be written as [6]:

\[ k_x d + \Omega = n \pi \]  

\( \Omega \) is the phase of the reflection coefficient of the solid/liquid boundary, \( n \) is the mode number and \( k_x \) is the x-component of the wave vector \( \vec{k} \) of the shear wave. It is given by:

\[ k_x = \frac{2\pi f_n v_i}{v_x \sqrt{1 - v_i^2 sin^2(\theta) / v_x^2}} \]  

With the angle of incidence \( \theta \) fixed between the first and the second critical angle, and in the high frequency limit, \( \Omega \) is frequency independent and can be eliminated. This is easily done by calculating from equation 5 and 6, the frequency difference \( f_{n+1} - f_n \) between two successive modes \( (n + 1) \) and \( n \):

\[ f_{n+1} - f_n = \frac{v_i}{2d \sqrt{1 - v_i^2 sin^2(\theta) / v_x^2}} \]  

This frequency difference can be measured on the FFT-spectrum of the reflected or transmitted signal, \( \theta \) and \( v_i \) are known so that \( v_x \) can be calculated.

EXPERIMENTAL VERIFICATION

The experimental procedure consists of two steps. In step one, the first critical angle is determined. In step two, the reflected or transmitted echo sequence is recorded and analyzed in the frequency domain for several angles of incidence above the first and below the second critical angle.

We investigated three samples: a 0.61 mm thick aluminum plate, a 2 mm thick glass plate and a 4 mm thick polyethylene plate. Two experimental methods were used: a bistatic Snell-DesCartes method and a monostatic double transmission method. The first one involves two broadband transducers placed according to the Snell-Descartes law. The second method used, is the double transmission technique described.
by Nagy et al [7]. It uses the same broadband ultrasonic transducer to generate the incident pulse and to pick up the double-transmitted echo from a perpendicular plane reflector. The minima in the spectrum of the reflected signal in the bistatic Snell-Descartes method and the maxima in the spectrum of the double-transmitted signal correspond to the frequency $f_a$ of equation (7).

Sample 1: Aluminum Plate of Thickness 0.61 mm.

The aluminum plate we investigated, is characterized by a longitudinal sound velocity of $6454 \pm 40$ m/s. This corresponds to a first critical angle of 13.4 degrees. The second critical angle which can not be measured precisely, is located around 28 degrees.

The aluminum plate is positioned in the farfield of a broadband transducer with a central frequency of 25 MHz and an operational frequency range from 4 MHz to 32 MHz. It isisonced with short pulses at ten different angles between 15 and 26 degrees (see table 1). The decay constant (equation 4) of the longitudinal waves generated in the plate, is at 4 MHz and 32 MHz equal to respectively 0.36 mm and 0.038 mm. For higher frequencies and larger angles of incidence it decreases further. This guarantees us that we are indeed in the high frequency limit for the whole measuring range because the decay constant is inferior to the thickness of the plate.

For each angle of incidence the double transmitted signal is recorded and analysed in the frequency domain. The frequency of the transmission maxima is listed in Table 1, with in the last column the corresponding shear velocity calculated by means of equation 7. Its average value is equal to 3100 $\pm$ 15 m/s. Using a 25 MHz shear transducer and the time-of-flight method, we obtained 3090 m/s $\pm$ 20 m/s. Both results are thus in good agreement.

<table>
<thead>
<tr>
<th>$\theta$</th>
<th>Position of the resonances in MHz</th>
<th>$\nu_r$</th>
</tr>
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<tbody>
<tr>
<td>15.5</td>
<td>5.2</td>
<td>7.9</td>
</tr>
<tr>
<td>15.4</td>
<td>5.9</td>
<td>8.6</td>
</tr>
<tr>
<td>17.5</td>
<td>6.4</td>
<td>9.8</td>
</tr>
<tr>
<td>18.7</td>
<td>7.0</td>
<td>10.4</td>
</tr>
<tr>
<td>19.9</td>
<td>7.4</td>
<td>11.0</td>
</tr>
<tr>
<td>21.0</td>
<td>7.8</td>
<td>11.8</td>
</tr>
<tr>
<td>22.2</td>
<td>8.2</td>
<td>13.4</td>
</tr>
<tr>
<td>23.4</td>
<td>8.8</td>
<td>13.8</td>
</tr>
<tr>
<td>24.6</td>
<td>9.6</td>
<td>14.5</td>
</tr>
<tr>
<td>25.7</td>
<td>10.6</td>
<td>16.4</td>
</tr>
</tbody>
</table>

TABEL 1

Sample 2: Glass Plate of Thickness 2 mm.

The glass plate is characterized by a longitudinal sound velocity of 5480 $\pm$ 26 m/s. This corresponds to a first critical angle of 15.6 degrees. Using a 10 MHz transducer and the same experimental procedure as for the aluminum plate we found a shear-wave velocity of 3281.6 m/s. The details are given in Tabel 2. Using a shear wave transducer we obtained a value of 3242 m/s, which is in close agreement with the average value obtained in tabel 2.

Sample 3: Polyethylene Plate of Thickness 4 mm

The polyethylene plate is 4 mm thick and is characterized by a longitudinal velocity of 2347 m/s. This corresponds to a first critical angle of 39 degrees. The transversal velocity is expected to be smaller than the sound velocity in water, so that there is no second critical angle. Because the plate is relatively thick, we could use a 1 MHz transducer and still be

<table>
<thead>
<tr>
<th>$\theta$</th>
<th>Position of the resonances in MHz</th>
<th>$\nu_r$</th>
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</thead>
<tbody>
<tr>
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<td>7.1</td>
</tr>
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<td>17.8</td>
<td>6.8</td>
<td>10.1</td>
</tr>
<tr>
<td>18.9</td>
<td>7.4</td>
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</tr>
<tr>
<td>20.1</td>
<td>8.4</td>
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<td>21.3</td>
<td>9.2</td>
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</tr>
<tr>
<td>22.5</td>
<td>9.9</td>
<td>14.5</td>
</tr>
<tr>
<td>23.6</td>
<td>10.6</td>
<td>16.1</td>
</tr>
</tbody>
</table>

TABEL 2

Table 2 gives the results of the transmission experiment. The average shear velocity is now equal to 939.2 $\pm$ 7.1 m/s. These values are again in good agreement with the shear wave velocity 921.6 $\pm$10 m/s, obtained with a shear transducer and the time-of-flight method.

<table>
<thead>
<tr>
<th>$\theta$</th>
<th>Position of the resonances in MHz</th>
<th>$\nu_r$</th>
</tr>
</thead>
<tbody>
<tr>
<td>48.0</td>
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<td>0.54</td>
</tr>
<tr>
<td>51.0</td>
<td>0.42</td>
<td>0.552</td>
</tr>
<tr>
<td>54.0</td>
<td>0.43</td>
<td>0.600</td>
</tr>
<tr>
<td>57.0</td>
<td>0.48</td>
<td>0.549</td>
</tr>
<tr>
<td>60.0</td>
<td>0.50</td>
<td>0.572</td>
</tr>
</tbody>
</table>

TABEL 3

References
NEW RESULTS BY THE PHASE GRADIENT METHOD IN THE STUDY OF LAMB MODES

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INTRODUCTION

The phase of the reflection coefficient of a plate immersed in water has been very few exploited [1]. However, when a Lamb mode is excited, the phase sharply changes from \(-\pi/2\) to \(\pi/2\). Moreover, as it has been shown for cylindrical shells [2], the derivative of the phase with respect to the frequency \(N\) exhibits peaks at the resonances frequencies, whose amplitude have the very interesting property to be inversely proportional to the half-width of the resonance. In addition, this derivative can be simply written with regard to the derivatives of the phase versus the longitudinal celerity \(C_L\), the shear celerity \(C_T\) and the sound celerity in water \(C\).

THEORETICAL BASIS

The expression of the phase is provided by the exact expression of the reflection coefficient which follows:

\[ R = \frac{C_C C_T - \tau^2}{(C_A + j\tau)(C_S + j\tau)} \]  

(3) and can be written:

\[ \phi = \arctan \frac{\tau}{C_S} - \arctan \frac{\tau}{C_A} \]

\( \tau \) is the ratio of the acoustical impedances and \( C_S \) and \( C_A \) respectively give the symmetric and antisymmetric Lamb modes of the free plate.

The expression for the frequency derivative can then be written as:

\[ \frac{\partial \phi}{\partial N} = \frac{\tau \frac{\partial C_S}{\partial N} - \frac{\partial C_A}{\partial N} \tau \tau}{C_A + \tau^2 \frac{\partial C_A}{\partial N} + C_S + \tau^2 \frac{\partial C_S}{\partial N}} \]

\( \tau \) being independent of the frequency \( N \).

At the vicinity of a resonance at the frequency \( N_R \), we can compare the exact expression of the frequency derivative to its expression derived from the Resonance Scattering Theory (RST) formalism [3]:

\[ \frac{\partial \phi}{\partial N} = G \frac{G(N - N_R)^2 + \Gamma^2}{2d} \]

(2) where \( G = \frac{1}{C_I} \) (\( d \): thickness of the plate) and \( \Gamma/2 \) is the half-width of the resonance at the frequency \( N_R \). When \( N = N_R \), \( \partial \phi/\partial N = C/2/\Gamma \). So, in addition of the determination of the resonances frequencies, the frequency derivative of the phase allows to obtain very simply their half-width.

RESULTS

We have studied the frequency derivative for a 2 mm-thick aluminum plate whose celerities are \( C_L = 6380 \) m/s, \( C_T = 3100 \) m/s and density \( \rho \) is 2800 kg/m³. The sound speed in water is \( C = 1470 \) m/s. Fig. 1 shows the frequency derivative versus frequency, at the incidence angle \( \theta = 5^\circ \). It exhibits a series of peaks located at the excitation frequencies of the modes A1, S1, S2, A2, S3 and A3. We note that the peaks related to shear waves (A1, A2, S3) have a larger amplitude than the peaks related to longitudinal waves. It appears that their half-width is smaller, as well as their reemission, these two parameters being linked [4]. This is experimentally confirmed: the reemission due to shear waves is less easy to detect at small incidence angles than the ones due to longitudinal waves. Beyond the first critical angle \( \theta = 13.32^\circ \), the longitudinal waves become evanescent. We observe on figure 2, plotted at \( \theta = 20^\circ \), that all the peaks have the same amplitude and they are equidistant. Everything happens as if there was a interference between shear waves. It is possible to calculate by means of a ray model [5] the frequency interval between two peaks:

\[ \frac{2d}{\sqrt{C_I^2 - \sin^2 \theta}} \]

It allows to determine experimentally the shear celerity of a material.

We have studied the influence of the parameters \( C_L \) and \( C_T \) on the resonances whose frequency are obtained by the phase derivative method. This type of study has already been performed with cylindrical shells [6]. At two given angles \( \theta \) and \( \zeta \), we have plotted the resonance frequencies when \( C_L \) ranges from 5300 m/s to 9300 m/s (fig. 3 and 5) and when \( C_T \) ranges from 1500 m/s to 4500 m/s (fig. 4 and 6), all the other parameters remaining constant in the two cases. The curves of resonance frequencies have been plotted for the modes A1, S1, S2, A2, S3 and A3. At \( \theta = 2^\circ \), when \( CL \) varies, we observe that the distribution of the frequencies for each mode is not random, but is composed of steps and slopes. We also note that the "shear" modes, like A1, S1, A2, S2, exhibit steps in a large range of \( CL \); their unsensibility to variations of \( CL \) seems to confirm the shear character of these modes. On the other hand, the "longitudinal" modes, like S2 and in a least way A3, exhibit slopes in a quite large domain of CL. For two close modes like (S1, S2) or (S3, A3), when the evolution of the frequencies of a mode (for instance S3) passes from a slope to a step, the evolution of the frequencies of the other one (A3) exhibits the inverse transformation. The changing of state appears approximately for the actual value of CL. Assuming that the steps can be extended to zero, we can consider that the plate is a fluid one. From the ordinate of the steps, we calculate a celerity in the plate of about 3000 m/s.
At the same incidence angle, we observe when \(\theta\) varies (fig. 4), a similar phenomenon: the curve \(N_R\) versus \(\theta\) for a mode exhibits steps and/or slopes. The curves associated to shear modes predominantly show slopes, while those associated to longitudinal modes are rather constant, but in a little \(\theta\) range. The latter modes seems to be nevertheless quite sensible to the shear celerity. We note that the modes \(S_1, S_2\) and \(A_2\) have a common step. If we consider that the steps correspond to the case of a fluid plate, we can obtain a celerity in the plate of about 6500 m/s.

At \(\theta=11^\circ\), we note for a variation of either \(CL\) or \(CT\) (fig. 5 and 6), quasi similar phenomena, except that the steps and the slopes are less pronounced. Moreover, we do not always find for each mode similar shapes. For instance, the mode \(A_2\), whose curve at \(\theta=11^\circ\) is a step, exhibits slopes at \(\theta=11^\circ\). At that time, this type of study only gives an indication about the prevailing character of a Lamb mode, either longitudinal or shear, particularly at small incidence angle.

CONCLUSION

We have shown that the method consisting in deriving the phase versus the frequency of the reflection coefficient of a plate embedded in a fluid allows to completely and simply characterize the resonances related to the Lamb modes by means of their frequency and their half-width. It avoids a "heavy" search of the poles of the reflection coefficient in the complex plane. The study of the frequential derivative when the longitudinal or the shear celerity varies, gives first results about the determination of the main character of the Lamb waves. This is to be completed by the study of the derivatives of the phase with respect to the celerities which is in progress.

BIBLIOGRAPHY

THICKNESS VIBRATIONS OF PIEZOELECTRIC PLATE EXCITED BY PARALLEL FIELD

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The thickness vibrations of piezoelectric crystal plate with the exciting field perpendicular to the major surfaces of the plate, and also to the plane wavefront, have been discussed in detail [1,2]. And this kind of exciting field is called perpendicular field. In this paper, parallel field excitation of thickness vibration of coupling modes in homogeneous-anisotropic piezoelectric crystal plate with electrodes on lateral surfaces will be investigated. Here, the exciting field is parallel to major surfaces of the plate and also to the wavefront, which is called the parallel field. That is similar to the vibrations of the piezoelectric bar loaded by an electric impedance in [3,4].

ANALYSIS OF THICKNESS MODES BY PARALLEL FIELD EXCITATION

The piezoelectric vibration system connected to the electric exciting source with the constant voltage $V_0$ is as shown in Fig.1.

![Fig.1 Piezoelectric vibration system excited by a parallel field](image)

where the differentiation to a space coordinate is meant by the comma before the index, and the differentiation to time is meant by the dot over the variable. The summation convention for repeated tensor indices is employed. The mechanical boundary conditions and electrical terminal conditions are:

\[ T_{ij}(x_i) = 0, \quad x_i = \pm \frac{t}{2}; \]

\[ D = 0, \quad x_i = \pm \frac{t}{2}; \]

\[ V_3 = \int \frac{W}{2} E_3 dx_3; \]

\[ I_3 = \int \frac{t}{2} \frac{t}{2} B_{3} dx_1 dx_2; \]

\[ V_0 = V_3 + I_3 Z; \]

Deriving procedure is similar to that in [2]. the eigen equation is obtained:

\[ \left| \begin{array}{c}
T_{ij} - \frac{1}{2} C_{ij} \end{array} \right| = 0, \quad (1)
\]

where:

\[ C_{ijkl} = C_{ijkl}^F + (\epsilon_{ijk} \epsilon_{ikl}/\epsilon_{ii}^F). \]

It yields three real positive roots, i.e. the eigenvalue \( \lambda(n=1,2,3) \). They are the effective elastic constants corresponding to three vibration modes of the plate. The effective elastic constant which depends on piezoelectric constants is called piezoelectrically stiffened one, while that which does not depend on any piezoelectric constants is called unstiffened one. And, the relevant vibration modes are also called stiffened or unstiffened correspondingly [5]. The velocity and wave number are obtained as:

\[ V_n = (\lambda_n / \rho) \lambda_n; \]

\[ \kappa_n = \lambda_n / V_n; \]

\[ \omega_n = 2\pi f_n. \quad n = 1,2,3. \]

The frequency equations satisfying the mechanical boundary conditions and electrical conditions are obtained:

\[ B_{kn} \left[ C_{ijkl} \frac{\lambda_n^k \lambda_n^l}{Z} \right. \left. \cos \kappa_n x_3 \right] = \frac{\epsilon_{ijk} \epsilon_{ikl}}{Z} \frac{\lambda_n^k \lambda_n^l}{Z} \sin \kappa_n x_3 \left. \right| = 0. \quad (2) \]

where:

\[ B_{kn}: \text{Amplitude Ratio}; \]

\[ Z = \sqrt{\omega_n C_{33} \frac{Z}{Z}}; \]

\[ C_{33} = \frac{t}{w} \epsilon_{33}^F; \]

\[ \epsilon_{ijk} = \epsilon_{ijk}^{(k)} \text{ for } (1 - \frac{t}{w} \epsilon_{ij}^{(k)}); \]

\[ \epsilon_{33} = \epsilon_{33}^F \text{ for } (1 - \frac{t}{w} \epsilon_{33}^F). \]

It is very important to know that the roots of eq. (2) are the function of impedance $Z$. It is shown
that, for this system, the resonant frequencies of either the fundamental tone or overtones of the three vibration modes of plate can be changed by changing the impedance $Z$. This is the principle of piezoelectrically adjustable frequencies [3.4]. Therefore, both of sound velocity and wave number are also the functions of impedance $Z$.

When $Z=0$, i.e. the electric terminal of the plate is short-circuited, the resonant frequency equation is obtained from eq.(2):

$$\cos^2 \frac{kz}{2} = 0.$$  

(3)

It is shown that the resonant frequency of a stiffness mode excited by parallel field is equal to the antiresonant frequency of that mode of the same plate with perpendicular field excitation (see eq.(9.71) in [1]).

When $Z$ goes to infinity, the electrical terminal of the plate is open-circuited. Then, the antiresonant frequency equation is obtained as:

$$\text{Bkn} \left( \frac{C_n}{2} \cos \frac{kz}{2} + \frac{\sigma_{jk} \epsilon_{jk}}{\epsilon_\mu} \sin \frac{kz}{2} \right) = 0.$$  

(4)

Eq.(4) looks, in form, like eq.(9.69) in [1], which is the resonant frequency equation under perpendicular field excitation. But, the signs in front of the sine terms in them are opposite to each other. It is shown that the frequency behaviors of vibration of the plate are different under two kinds of exciting field.

From eqs. (1), (2), (3) and (4) one reaches the following conclusions:

a). The impedance $Z$ will affect the vibration characteristics of a piezoelectric crystal plate. (Such as resonant frequencies, effective coupling coefficient, stress and displacement distribution and so on.)

b). Using a parallel field, three thickness vibration modes exist in a piezoelectric plate, and are coupled with each other in general as in class C1(1).

c). The ratio of the resonant frequency of any overtones to fundamental tone is an odd number, i.e. they are in harmonic relation. It is convenient to use this harmonic relation for measuring electroelastic constants of crystals by dynamic method.

d). For the class with higher symmetry, some piezoelectric or elastic constants are zero. Consequently some modes relevant to those zero constants may not couple with the other or even they would become unshifted. Thus, the equations will be simplified. The unshifted modes can’t be excited by parallel field. But, fortunately, they may be excited by the parallel field in some cases. For example, it has been applied to piezoelectric measurement of piezoelectric crystals in classes C2v(3m), C6(6), C3(3), D3(32) and C2v(2mm) [6.7.8.9.10].

RESULTS IN THE MOST GENERAL CASE

In order to obtain the solution for the most general case, let’s consider the vibration system like that in Fig.1 and let $\gamma_{ij}$, $C_m$. The set of equations for the thickness vibrations of the piezoelectric plate under a parallel field excitation is as follows:

$$\epsilon_{ij} = C_{ij} \epsilon_{ij}$$  

(5)

where:

$$\epsilon_{ij} = \epsilon_0 + \sigma_{ij} \epsilon_{jk} \epsilon_{jk} \epsilon_{ij} / \epsilon_\mu,$$

$$\text{Bkn} \left( \frac{C_n}{2} \cos \frac{kz}{2} + \frac{\sigma_{jk} \epsilon_{jk}}{\epsilon_\mu} \sin \frac{kz}{2} \right) = 0.$$  

(6)

where:

$$\text{Bkn: Amplitude Ratio:}$$

$$\frac{\sigma_{jk} \epsilon_{jk}}{\epsilon_\mu} = \frac{1}{Z} \frac{\epsilon_\mu}{\epsilon_\mu}.$$  

$$\epsilon_{jk} \epsilon_{jk} = (1 - \frac{\epsilon_\mu}{\epsilon_\mu}) \epsilon_{jk} \epsilon_{jk}.$$  

$$\epsilon_\mu = \frac{1}{Z} \frac{\epsilon_\mu}{\epsilon_\mu}.$$  

$$\text{CON} \cos \frac{kz}{2} = 0.$$  

(7)

$$\text{CON} \cos \frac{kz}{2} = 0.$$  

When $Z \to \infty$.

$$\text{Bkn} \left( \frac{C_n}{2} \cos \frac{kz}{2} + \frac{\sigma_{jk} \epsilon_{jk}}{\epsilon_\mu} \sin \frac{kz}{2} \right) = 0.$$  

These forms of equations are convenient to apply to the studying on the rotated plate with any orientation.

The performance of the piezoelectric plate with frequencies piezoelectrically adjustable by the parallel field excitation would be discussed in the other paper.

REFERENCES

GEL RIGIDITY MEASUREMENT USING SURFACE WAVES

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Gels have attracted considerable interest. One of the physical properties which characterize gels is its small rigidity. The rigidity of typical gel is \( \sim 10^4 \) Pa, whereas that of solid is \( \sim 10^10 \) Pa. There exist some methods for measuring gel rigidity, however few methods have been reported which can monitor the variation of rigidity as gel network formation proceeds. Sol-gel transition in gelatin was observed using the present rigidity measurements with surface waves. The rigidity was also measured in agarose gel.

§ 1. THEORY

On a sol, surface tension waves propagate since sol is a liquid state. When the sol becomes gel, small rigidity emerges. It is expected, therefore, that the surface waves have the nature of both surface tension waves and elastic waves (Rayleigh waves). The dispersion relation of the surface waves on a gel will be then

\[
\omega^2 = \frac{\gamma}{\rho} k^3 + \frac{0.91}{\rho} G k^2
\]  

(1)

where \( G \) is the rigidity of gel. The first term of the right hand side of eq.(1) is due to the surface tension and the second term is due to the elasticity. Equation (1) could be correct if the elasticity is dominant for the return force of surface waves on a gel. The phase velocity \( V \) of the surface waves is derived from eq.(1) as

\[
V^2 = V_t^2 + \left( \frac{0.91}{\rho} \right) G
\]  

(2)

where \( V_t \) is the velocity of surface tension waves. Equation (2) suggests that the rigidity of gel can be obtained from the measurement of surface wave velocity.

§ 2. APPARATUS

The block diagram of the experimental system is shown in Fig.1. Surface waves are excited using a PZT bimorph, the edge of which touches on the sample surface. A burst wave of four cycles is applied to the bimorph with a function generator (HP 8116A). The pulsed surface waves propagating on the sample is detected by using an optical deflection method. A light from He-Ne laser with the output power of 5 mW is introduced to the sample surface almost vertically. The displacement of the surface waves modulates the reflected light which is detected with a position sensitive photodiode. The detected signal is proportional to the gradient of the surface

displacement. The sample is contained in a laboratory dish which is enclosed by a water jacket for controlling the temperature within 0.5 °C. The water jacket and the bimorph are set on a motor-driven slide stage, and so the detection position can be varied. The time position of a specific peak of the detected signal on a digital oscilloscope is recorded at several propagation distances, providing the phase velocity of the surface waves.

§ 3. MEASUREMENTS OF GEL RIGIDITY

3.1. GELATIN

Sol-gel transition was studied in gelatin with the measurement of surface wave velocity at gelatin concentrations of 4, 5 and 7 wt%. Ageing effect and frequency dependence of the surface wave velocity were also observed at concentration of 7%. The gelatin powder (Sigma Chemical, G2500) was dissolved in distilled water and the solution was heated to 60 °C for 1 hr. The sample was cooled at a rate of 10 °C/hr from 60 to 20 °C.

The surface wave velocity measured vs. temperature is shown in Fig.2. The measured frequency was 200 Hz. The velocity in the sol state (at higher temperatures) is constant and increases with decreasing temperature in gel state. The sol-

Fig.1. Block diagram of the apparatus.

Fig.2. Surface wave velocity vs. temperature in gelatin. The velocity increase represents sol-gel transition.
gel transition temperature $T_g$ is about 36 °C for 7% concentration. This increase is possibly due to the emergence of elasticity which reflects the gelation. Figure 3 shows $V^2 - V_t^2$ (this value is proportional to the rigidity) plotted against reduced temperature $(T_g - T)/K$ for 7% gelatin. The value $V_t$ was substituted for that in the sol state, because Kikuchi et al. reported that the surface tension remained almost constant when gelatin sol became gel. Figure 3 indicates the value $V^2 - V_t^2$ is proportional to $(T_g - T)^{2.2}$. The exponent 2.2 agrees with the elasticity exponent 1.98 based on the percolation transition theory by de Gennes.

The ageing effect of frequency dependence of 7% gelatin gel are represented in Fig.4. The surface wave velocity in the frequency range 200 - 1600 Hz was measured 1 hr., 2 hrs., 3 hrs., 4 hrs and 1 day after the sample was cooled down to 20 °C, at constant temperature of 20 °C. The velocity linearly increases with frequency, and increases with time. The frequency dependence may be interpreted by a relaxation phenomenon.

4.2 AGAROSE

The agarose solution was heated to 90 °C for 2 hrs., and stored at 5 °C for 20 hrs to obtain gel. The surface wave velocity was measured at concentrations of 0.5, 1 and 1.5 wt %, at temperatures of 10, 20, 30 and 40 °C. The frequency dependence of the velocity is plotted in Fig.5 for 1% agarose. The velocity is constant in the frequency range 400 - 2400 Hz. The velocity is considerably higher than that of 0.5 w/s in the sol state, suggesting the surface waves on agarose gel are originated in elasticity. The rigidity of the gel at 40 °C is calculated to be $1 \times 10^{12}$ Pa from eq.(2), which is in agreement with the value obtained from other method.

Fig.4. Ageing effect of frequency dependence of the surface wave velocity in 7% gelatin gel at 20 °C. The open circles indicate the values immediately after cooled down to 20 °C from 60 °C. The other symbols denote the values after 1 hr, 2hrs, 3hrs, 4hrs and 1 day ageing.

Fig.5. Frequency dependence of the surface wave velocity in 1% agarose gel at 10 - 50 °C.

CONVERSION OF THE NATURAL MODES OF A LIQUID/SOLID INTERFACE BY A DIHEDRAL.

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Introduction.

In this paper, we study the waves which can propagate on an infinite interface separating a liquid and a perfect elastic solid. Originally, this research was started by Stoneley who established the equation verified by the interface waves. Later on, Scholte found a particular solution corresponding to an undamped interface wave. We have written the problem into a linear, homogeneous system, numerically resolved. Among the solutions of this system are the Scholte and the Rayleigh interface waves. Briefly, we recall the main characteristics of the Scholte-Stoneley wave and of the leaky Rayleigh wave. The experimental study of the Scholte wave diffraction by a dihedral edge shows two kinds of phenomena. The first one is a simple phenomenon of diffraction and the second one is a phenomenon of mode conversion. Reciprocally, if a bulk wave reaches a dihedral at an angle close to the Rayleigh one, the conversion into Scholte waves occurs.

1. Interface waves.

The interface waves are described using evanescent waves with complex wave vectors. Each component of vibration in the liquid is described by:

\[ u_0 = U_0 \exp(j(\kappa_0 + j\Omega_0)t) \]

It was proved that vectors \( \kappa_0 \) and \( \kappa_0' \) are perpendicular when the mediums are not absorbent \([1]\). \( \kappa_0 \) represents the propagation direction of the wave and \( \kappa_0' \) the decrease direction of the amplitude. In the solid, the descriptions are similar, but require a summation of the contributions of a lamellar (L) and a torsional (T) wave.

The Scholte wave is a liquid/elastic solid interface wave that theoretically propagates without attenuation because the component \( \kappa_0' \), of this wave is zero. Its celerity is always lower than the sound celerity \( c_s \) in water and than the celerities \( c_L \) and \( c_T \) of the longitudinal and transverse waves in the solid substrate. In the fluid medium, the amplitude of this evanescent wave decreases as \( \exp(-\alpha k x) \). \( \alpha = 2\pi F/c_0 \) is the wave vector in the OX1 direction, \( F \) is the frequency and \( \alpha = k_0''/k_0' \), see figure 1.

The coefficient \( \alpha \) depends a lot on the solid medium and is related directly to the celerities by the relation:

\[ \alpha^2 = 1 - (c_L/c_0)^2 \]

with \( c_L \) the Scholte wave celerity and \( c_0 \) the longitudinal wave celerity in the fluid. The wave vibrations in the liquid and solid mediums are in the sagittal plane and of elliptic polarization. The great axis of vibration is perpendicular to the interface in the substrate and parallel to it in water. In this case, the components \( u_L \) and \( u_T \) are in quadrature and \( u_L = iu_T \). The wave energy is mostly in the liquid medium (more than 99%).

The Rayleigh wave has elliptic vibrations in the liquid and solid mediums and these vibrations are in the sagittal plane. In the fluid medium, the amplitude increases exponentially when we move off the interface, see figure 2. So, the generalized Rayleigh wave radiates energy in the fluid in the \( \beta \) direction (with respect to the perpendicular to the surface) and its amplitude decreases during its propagation at the interface.

2. The Scholte wave diffraction by an elastic dihedral.

We study the diffraction phenomena which appear in the interaction of a Scholte wave and a dihedral edge. We have limited the study to the case of a duraluminum solid and a Scholte wave propagating in a direction perpendicular to the dihedral edge.

The Scholte wave is excited by an interdigital transducer (I.D.T.) realized on an amorphous substrate. The I.D.T. is stucked at the plate extremity in the same plane than the plate [2].

When the Scholte wave meets the dihedral edge, two phenomena appear. The first one is the diffraction which is connected essentially to the liquid contribution of the incident Scholte wave. The second one is a mode conversion phenomenon which is a function of the solid components of the wave and, therefore depends on the solid substrate and on the dihedral angle [3].

On figure 3 we present the directions of the waves radiated by the plate extremity due to the incident Scholte wave.
Figure 3: The Scholte wave diffraction by a dihedral. Representative scheme.

\[ V \] is the main direction of the radiated bulk wave. \( R_r \) and \( R_t \) are the directions of reemission of the leaky Rayleigh waves, \( \beta \) the Rayleigh angle and \( \theta \) the direction of observation. On figure 4, we have plotted, versus angle of observation \( \theta \), the maximal values of the signal received by a classical immersed transducer rotating around a dihedral edge of angle \( \gamma = 45^\circ \).

Figure 4: The Scholte wave diffraction by a dihedral of angle \( \gamma = 45^\circ \). Signal measured in the far field of the edge as a function of \( \theta \).

The leaky Rayleigh waves are observed at \( \theta \) equal to \(-75^\circ\) and \(120^\circ\). The Scholte wave is transmitted on the second face of the dihedral and diffracted by the supplementary dihedral (angle \( \gamma' \)) at an angle of observation close to \(-135^\circ\). This phenomenon is reversible [4]. The ultrasonic beam emitted by a transducer has a structure close to the generalized Rayleigh wave structure. When such a beam is incident on a dihedral, in the Rayleigh direction, it gives rise to a Scholte wave. In order to verify the existence of this Scholte wave, we measure the acoustic field at the end of the horizontal plate with a micro-transducer of 0.6 mm diameter. We observe, at the proximity of the plate, an exponential decay in the vertical direction characterizing the Scholte wave, see figure 5.

Figure 5: The exponential decrease (between points A and B) of the Scholte wave in the fluid.

Conclusion.

We have experimentally shown that the interaction of a Scholte wave with a dihedral gives rise to Rayleigh waves on both faces of the dihedral. Using the verified reciprocity of this phenomenon, we have found out a new technique to generate a Scholte wave on liquid/solid interfaces.

References.

ON THE EXISTENCE OF RAYLEIGH-TYPE SAW IN PERIODIC STRUCTURES WITH RESONATING ELEMENTS

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Abstract: Propagation of Rayleigh-type acoustic waves in semi-infinite elastic substrates supporting an array of periodically-spaced scatterers is investigated. The scatterers are modeled by vibrating resonators with eigen-frequency \( \omega_0 \). The analysis is based on the concept of phased line-source-arrays and the periodic Green's functions. For frequencies \( \omega \) with \( \omega \approx \omega_0 \) and \( \omega < \omega_0 \) (Bragg frequency), the existence of slow Rayleigh waves is predicted. Under these conditions the wave-energy is dominantly localized in resonators rather than in the supporting substrate. For \( \omega \approx \omega_0 \) with \( \omega < \omega_0 \) we find Bragg-reflection and strong damping of slow waves. When \( \omega_0 \) approaches \( \omega \) a complicated interaction mechanism takes place, which leads to a stopband-widening in Brillouin-diagram. In the case \( \omega_0 > \omega \), SAW is strongly damped due to efficient generation of bulk acoustic waves at frequencies \( \omega \approx \omega_0 \).

I. INTRODUCTION

Propagation properties of Rayleigh waves in periodic gratings have been the subject of extensive theoretical and experimental research [1], [2]. This is due to the fact that periodic reflecting structures play a fundamental role in a number of surface acoustic wave (SAW) devices; i.e. resonators, reflective-array compressors, narrow band-filters. Recent works [3],[4] show that by choosing specific electrode-profiles and/or large-amplitude gratings, the reflection efficiency can be essentially improved. In 80s, considerable attention was paid to reducing the insertion loss of SAW devices. One attractive technique for realization of low loss devices is the design of strongly reflecting elements [5],[7]. Considering surface transverse waves interaction with an isolated, large-amplitude protuberance, localized shear horizontal surface shape resonances can arise [5]. Recently the existence of sagittal acoustic surface shape resonances has been theoretically proven [8]. In this paper the main features of acoustic-wave-grating interaction are carried out by modelling the electrodes as resonators. The resonators vibrate at their characteristic eigen-frequencies and are periodically-spaced attached to the substrate surface. A two-dimensional analysis (\( \partial^2 \xi / \partial x^2 = 0 \)) is performed. It is useful to characterize the scatterers by eigen-frequency \( \omega_0 \) and the parameter \( \varepsilon = M/\rho a^2 \) rather than their mass and elasticity. The quantities \( M \) and \( \rho \), respectively, denote the mass of the scatterers and the mass-density of the substrate material. The \( \lambda \) denotes a characteristic wavelength. One essential result in this paper is that we show the existence of slow Rayleigh-type waves (SRTWs) in structures described above. It is shown that SRTWs exist, if scatterer eigen-frequency \( \omega_0 \) is less than the Bragg-frequency \( \omega_B \) of the periodic system. Further, it is shown that the wave energy is dominantly localized in the resonators and not in the supporting substrate which couples the resonators [9].

II. THEORETICAL BACKGROUND

II.1 Statement of the Problem: Let us consider an isotropic halfspace occupying the region \( y \leq 0 \) and supporting a periodic array of line resonators which are parallelised to \( x \)-axis. The resonators are periodically-spaced at distance \( p \) and are as narrow strips (\( 2a << p \)) attached to the surface \( (y = 0) \). Further, the resonators are assumed to move either vertically (parallel to \( y \)-axis) or horizontally (parallel to \( x \)-axis). The influence of the resonators on the elastic semi-space can then be modelled by stresses acting vertically or horizontally on the surface. The physical quantities including the acting forces have then the form

\[ \psi(x, y) = \sum_{m=-\infty}^{\infty} \phi_m e^{im\pi x/a} = \sum_{m=-\infty}^{\infty} \psi_m e^{im\pi x/a}, \]  

with \( \phi_m = \phi_0 + 2m \pi p / \lambda \). Here, \( \phi_0 \) is the wavelength of the eigen-modes and is allowed to be real- or complex-valued.

II.2 Concept of Phased Line-Sources: We now seek for the simplest force-distribution which excites a field pattern according to Eq.(1). In a non-periodic case, this problem leads to the concept of an isolated line-force excitation with resulting field-distribution known as Green's function. In a natural way the introduction of an array of phased line-sources suggests itself

\[ F(x) = \int \Phi(x,y) \sum_{m=-\infty}^{\infty} \delta(x - mp) F_m. \]

The resulting field-distribution will be called periodic dyadic Green's function.

II.3 Periodic Dyadic Green's Functions: We begin with the excitation of an elastic semi-space by an harmonic force-density distribution \( F \) applied to the surface \( F(x,0) = F_0 e^{im\pi x/a} \) and \( F(x,\pi/a) = F_1 e^{im\pi x/a} \) with \( F_0(\pi/a) = F_1(\pi/a) = F_2 \exp(i\pi x/a) \) and \( F_1(\pi/a) = F_2 \exp(i\pi x/a) \). The resulting dyadic Green's function is [9]

\[ G = (1/\mu D) \left( \begin{array}{cc} k^2 \lambda & -i\mu \phi(a^2 + \pi^2 - 2na) \\ i\mu \phi(a^2 + \pi^2 - 2na) & k^2 \lambda \end{array} \right) \]

(3)

Here, \( \mu \) denotes the Lame constant. The \( k_1, k_2, \mu, \kappa \) and \( D \) are defined as follows: \( k_1 \): wave-number of shear bulk wave, \( k_2 \): wave-number of longitudinal bulk wave, \( \kappa = \sqrt{\mu^2 - k_1^2} \), \( \mu = 4 \eta^2 c_0 - (\kappa^2 + \pi^2) \) is the Rayleigh determinant. In the latter, \( \eta \) is considered as a complex quantity.

Let us return to the aforementioned periodic system of phased line-sources. We write \( \phi(x) = \phi_0 e^{im\pi x/a} \sum_{m=-\infty}^{\infty} (F_m - \nu_m \Delta) \) where \( \sigma \) stands for one of the stress-tensor-components \( \sigma_{xx} \) or \( \sigma_{yy} \), \( \nu_0 \) is a normalization factor and \( f(x) \) is the stress distribution in the fundamental period, \( (f(x) = \nu_0 \text{zero on the interval } [a] \text{ only}) \). Using the representation \( \sum_{m=-\infty}^{\infty} (F_m - \nu_m \Delta) = \sum_{m=-\infty}^{\infty} V_m \exp(i\pi x/a) \), we obtain \( \sigma = \nu_0 \exp(i\pi x/a) \sum_{m=-\infty}^{\infty} V_m \exp(i\pi x/a) \), \( \nu_0 \) is Fourier transform of \( f(x) \). The wave-number \( Q \) is defined as \( Q = 2\pi / p \). Thus in contrast to harmonic generation, we now deal with the wave-generation in the medium by a collection of harmonic \( \sigma = \nu_0 \exp(i\pi x/a) \), where \( \nu_0 = q + iM \). We define \( \nu_0, \kappa, \lambda \) and \( D \) evaluated at \( \nu_0 = q + iM Q \). Assume uniform stress-distribution within the interval \( [a] \text{ near } a, \text{ and denote the total-force per unit length in } z\text{-direction by } F_0 \). For a Dirac delta stress-distribution, \( \text{i.e., } a, 0 \text{ obtain } \nu_0 V_0 \rightarrow F_0 / p \). The resulting field-distribution is called periodic dyadic Green's function \( \Phi \) in real space

\[ \Phi(x, y) = \sum_{m=-\infty}^{\infty} (1/\mu D_m) \left( \begin{array}{cc} k^2 \lambda_m & -i\mu \phi(a^2 + \pi^2 - 2na_m) \\ i\mu \phi(a^2 + \pi^2 - 2na_m) & k^2 \lambda_m \end{array} \right) \]

(4)

III. DISPERSION EQUATIONS

III.1 Semi-space Substrate Loaded by an Array of Periodic Resonators: We first define the following characteristic quantities \( \nu_0 \) and \( F_0 \): \( \nu_0 \) stands for substrate displacement at the point, where the resonator is attached to the surface and \( F_0 \) denotes the force acting at this point. Further \( K \) and \( M \) denote the rigidity and the mass of the resonators, respectively: \( \omega_0 = \sqrt{K/M} \) is the resonator eigen-frequency. Restricting the resonator-motion to vertical direction, and the action of the force \( F_0 \) to the narrow but finite strip \( |x| < a \), we obtain \( \omega_0 = \sqrt{K/M} \) as the solution. From the condition for a non-trivial solution \( (\nu_0 \neq 0) \) we obtain the dispersion equation. For the discussion below it is useful to define \( \varepsilon = K/2\mu \). A system of horizontally vibrating resonators can be analogously tackled.
III.2 Heavy Strips: We now concentrate on wave-propagation in periodic systems supporting heavy, narrow strips. The strip located in the fundamental period occupies the range $|z| < a << p$. The force acting on the substrate in the strip region has two components $F_{x} = m \omega_{0}^{2}u_{x}$. The mass of the strip per unit length. We find

$$u_{x} = G(u_{x}, F_{x}) \frac{m \omega_{0}^{2}u_{x}}{2a} + G(u_{x}, F_{y}) \frac{m \omega_{0}^{2}u_{y}}{2a}, \quad (5a)$$

$$u_{y} = G(u_{y}, F_{x}) \frac{m \omega_{0}^{2}u_{x}}{2a} + G(u_{y}, F_{y}) \frac{m \omega_{0}^{2}u_{y}}{2a}. \quad (5b)$$

For non-trivial solutions the determinant of (5) has to vanish.

IV. NUMERICAL RESULTS

It is convenient to use normalized frequency $\eta = \omega / \sigma_0 = k_l / (Q/2)$ and normalized wave-number $\xi = \eta (Q/2) = q / \sigma_0$. In the numerical calculations carried out here, the substrate material was assumed to be fused quartz glass with a Poisson ratio $\sigma = 0.17$. In Fig.1, solutions of dispersion-equation are shown for eigen-frequency $\eta_0 = 0.2$ and $\epsilon = 0.3$. The Rayleigh wave essentially slows down for frequencies slightly lower than $\eta_0$. The main part of the energy is carried by the resonator while the substrate acts as a coupling between the resonators. If the frequency increases slightly but still remains below $\eta_0$, we find a new Bragg stopband for the slow Rayleigh waves discussed above (Fig.1). Near $\eta_0$, the waves are strongly attenuated [9]. Increasing the frequency up to ultimate vicinity of $\eta_0$, no solution is obtained. Increasing the frequency further, for frequencies higher than $\eta_0$, we again observe propagating waves. The dispersion curve starts to exist at a point on the line $\eta = \eta_0$. The following interpretation suggests itself: The propagating waves are quasi surface-skimming transverse-waves with polarization in sagittal plane, accompanied by evanescent compressional modes. The system of resonators attached to the surface makes possible the appearance of these solutions which have no counterpart in elastic semi-space with free boundary. At higher frequencies these waves become more and more localized near to the surface and behave like Rayleigh waves. Since, in this case, the interaction is not very strong ($\epsilon = 0.3$) and $\eta_0 < \eta_B$, Bragg stopband is disconnected from the radiation zone. Fig.2 shows the dispersion-curve for a more massive resonator $\epsilon = 1$. Here, the Bragg stopband and the scattering region are not separated. When $\eta_0$ approaches $\eta_B$, the behaviour of dispersion curves becomes more complicated, Fig.3. Two aspects may be emphasized. (i) The waves with frequencies lower than $\eta_0$ are strongly damped [9]. (ii) For the same parameter $\epsilon$ as in Fig.1, we observe stopbands widened. The overall behaviour of the dispersion curves for horizontally vibrating resonators is quite similar to the curves given in Fig.1. SAW interacts more weakly with horizontal resonators than with vertical ones [9].

VI. CONCLUSION

Propagation of Rayleigh-type waves in elastic halfspace supporting a periodic array of resonators is investigated. At frequencies close to the resonator eigen-frequency $\omega_0$ ($\omega < \omega_0 < \omega_B$) the existence of slow Rayleigh waves is predicted.

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INFLUENCE OF THE EXTREMITIES OF A SEMI INFINITE PLATE ON LAMB WAVES PROPAGATION.

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INTRODUCTION

A lot of studies deal with the propagation of Lamb waves on either a free elastic plate (in vacuum) or an embedded one. In most cases, the plate is considered to be infinite; these investigations are interested here in the effect of the limitation of the plate length. At the plate extremity, different kinds of phenomenon can occur: reflection of the incident Lamb wave (see 1 for the free plate), radiation of a bulk wave in water (2,3), and mode conversion into other Lamb waves at the same frequency but with a different propagation direction (see 4, 5 for the free plate). This paper presents a study of the first two phenomena which are not greatly perturbed by the third one.

THE EXPERIMENTAL SET-UP

The three broadband transducers A, B, C of Fig. 1 are of center-frequency 2.25 MHz and all lie in a same rigid plane. The end sections of the plates used have right angles and the plates thickness is either 4 or 5 mm. The transducer acts as an emitter in all experiments. The ultrasonic beam axis intersects the plate at a given angle less than the first critical angle; the transducer aperture angle is about 15° at 2.25 MHz. The experiments are conducted with broadband signals. For all frequencies such that the projection of the incident wave vector along the interface is equal to a particular Lamb wave vector, that Lamb wave is excited. As it propagates in the x direction, it reemits energy in water under an angle c = α, until it reaches the plate extremity, where refraction and scattering in water occur. We present three experiments, labeled "Exp1", "Exp2" and "Exp3". In all of them, the received time signal is recorded for different positions of the receiver, and its amplitude spectrum calculated (see e.g. [6]). Peaks are observed on the spectra, at frequencies corresponding to the excited Lamb modes. The amplitude of each Lamb mode peak is plotted as a function of the position of the receiver.

"Exp1": The receiving transducer C rotates around the plate extremity (−90° ≤ c ≤ +90°) and catches the radiated bulk wave in water.

"Exp2": The receiving transducer B is set at different positions, such that the x2 distance of Fig.1 is increased by steps of 1 cm. The measured signal is related to the reemission of the excited Lamb waves.

"Exp3": The receiving transducer is also A. The recorded signal corresponds to the reemission of the excited Lamb waves after their retroreflection at the extremity of the plate. The A transducer is set at different positions such that the x1 distance of Fig.1 increases by steps of 1 cm.

ANALYSIS OF THE SCATTERED BULK WAVES

We analyse the "Exp1" results. For each given value of c, we observe significant radiation patterns at the plate extremity at different frequencies. Each pattern is associated with a Lamb mode by pointing out the values of θ and 2θ on the dispersion curve. The radiation pattern of the A1 Lamb mode is then explained with the help of a model based on the Huygens principle. We consider the radiation to be mainly due to the u(x, z) component perpendicular to the plate end section. We use its expression for a free plate [1]. The measured pressure at an observation point in the far field of the plate and at angle α is then compared to the modulus of g(α):

\[ g(\alpha) = \int u(0, z) \ e^{-i \omega z} \, dz \]

k is the wave vector associated with the bulk waves in water. Fig. 2 presents the curves of the modulus of g(α) versus α for the A2 and S3 modes. The experimental diagrams (amplitude in dB) obtained for the same two modes are given on Fig.28-38. The comparison between these curves and the theoretical ones shows that the model takes into account the symmetry and the positions of the extremas of the radiation patterns.

ESTIMATION OF THE RETROREFLECTION COEFFICIENT OF LAMB MODES

This study requires the achievement of both the "Exp2" and the "Exp3" experiments.

* "Exp2" experiment: the measurements are done for non-zero values of x2 extending from about 5 to 15 cm, in order not to be affected by the specularly reflected echo. For each Lamb mode, we plot the amplitude A (in dB) of its peak in the amplitude spectrum versus the x2 distance. The so obtained points align on a straight line, the equation of which is obtained by a least square method. An example is shown on Fig.4, where the least mean square straight line is drawn from x2 = 5 cm to 15 cm. If we continue this straight line for x2 values extending from 0 to 5 cm (dashed line), we measure the A0 amplitude. The relation between A and x2 may then be written as:

\[ A = 20 \, \log\, (1 + x^2) \]

* "Exp3" experiment: the peak amplitude B of each mode is plotted in dB units versus distance x1; a least mean square straight line is drawn and its equation calculated. This equation is then identified:

\[ B = 20 \, \log\, (1 + x^2 / 2) \]

where R is the retroreflection coefficient of the Lamb wave considered (the two transducers A and B are supposed to have the same frequency response). The calculated slope of the least mean square straight line provides a value of the retroreflection coefficient c which is different from the one obtained by the "Exp2" experiment. The
calculated ordinate of a point at abscissa \( x_1 = d \) provides a way to estimate the retroreflection coefficient \( R \), at \( 2\overline{d} \), with \( \overline{d} \) equal to 2d. Table 1 shows the results obtained for eight Lamb modes excited at angle \( 8 = 8^\circ \) on a 5 mm thick aluminum plate.

At the water/metal interface of the plate extremity, the stress components in the \( x \) direction must be continuous. In the solid part of the interface, these components are summation of those associated with the incident Lamb wave, those associated with the retroreflected Lamb mode, and those due to mode conversions (which are neglected). In the fluid part of the interface, the \( T_{xx} \) component is zero, while the \( T_{xx} \) component is \(-p\) (\( p \) = pressure due to the radiated bulk wave). In the limit case of a purely transverse incident Lamb wave for which \( T_{xx} = 0 \), there is no way to ensure the continuity of \( T_{xx} \) if one considers a radiated bulk wave. A Lamb wave should have the maximum of its integral reflection coefficient. It seems then reasonable to expect the retroreflection coefficient to have maxima for modes with a \( T_{xx} \) stress component much larger than the \( T_{xx} \) one. Following Viktorov, we choose as a characteristic of the end stress (as a whole) for any mode the integral square of the stress, and we calculate the ratio \( m \):

\[
m = \frac{\int T_{xx}^2 \, dz}{\int T_{xx}^2 \, dz}
\]

The values of the \( m \) ratio are given in the last column of table 1. Actually, the values of the retroreflection coefficient do increase with \( m \), as expected.

CONCLUSION

We have presented experiments on two phenomena occurring when an incident Lamb wave reaches the plate extremity: the scattering of bulk waves in water and the retroreflection of the incident Lamb wave. A simple model allows the prediction of the radiation patterns obtained in the first experiment. The value of a theoretically defined ratio of integral squares of the stress components is shown to be indicative of those modes which have a greater retroreflection coefficient.

REFERENCES


<table>
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Table 1
THE APPLICATION OF INVARIANT ELASTIC CONSTANTS IN STUDYING THE ACOUSTIC AXES OF TETRAGONAL AND TRIGONAL CRYSTALS

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INTRODUCTION

It is convenient to describe the elastic properties of solids by constants which are invariant under coordinate transformation. These constants are known as invariant elastic constants (IC). It has the advantages of not dependent on the coordinates being used with respect to a particular set of crystallographic axes.

The existence of the acoustic symmetry axis have been studied experimentally in a series of scheelite crystals in great detail (1). Several fluoride compounds have also been reported (2). It has been shown that referring to the acoustic axes considerably simplifies the interpretation of the elastic behavior of Ti and TII, RI and RII group crystals.

In the following we are going to show that in the IC formulation the acoustic symmetry comes out naturally because of their unique transformation properties being used.

ACOUSTIC AXES IN TI, TII AND RI, RII GROUP CRYSTALS

In an elastic solid with an arbitrary Cartesian frame of reference, we have

$$\sigma_{ij} = C_{ikjm} \varepsilon_{km} \quad (i,j,k,m=1,2,3)$$

where $\sigma_{ij}$ are the components of stress tensor and $\varepsilon_{km}$ are the strain tensor and $C_{ikjm}$ are second order elastic constants. From the symmetry relationships of the crystals in TI, TII groups the numbers of the elastic constants are reduced to 6 and 7 respectively. They can be represented by the 6x6 matrices as follows:

$$\begin{pmatrix}
C_{11} & C_{12} & C_{13} & 0 & 0 & C_{16} \\
C_{12} & C_{11} & C_{13} & 0 & 0 & -C_{16} \\
C_{13} & C_{13} & C_{33} & 0 & 0 & 0 \\
0 & 0 & 0 & C_{44} & 0 & 0 \\
0 & 0 & 0 & 0 & C_{44} & 0 \\
C_{16} & -C_{16} & 0 & 0 & 0 & C_{66}
\end{pmatrix}$$

(in TI groups, $C_{16} = 0$)

Similarly for RI, RII groups:

$$\begin{pmatrix}
C_{11} & C_{12} & C_{13} & C_{14} & C_{15} & 0 \\
C_{12} & C_{11} & C_{13} & -C_{14} & -C_{15} & 0 \\
C_{13} & C_{13} & C_{33} & 0 & 0 & 0 \\
C_{14} & -C_{14} & 0 & C_{44} & 0 & -C_{15} \\
C_{15} & -C_{15} & 0 & 0 & C_{44} & C_{14} \\
0 & 0 & 0 & -C_{15} & C_{14} & C_{66}
\end{pmatrix}$$

(in RI groups, $C_{15} = 0$)

In an earlier paper by Khatkevich (3) he described the classification of elastic constants of TI, TII (RI, RII) group crystals in terms of the acoustic axes. He found that the only difference between TI and TII groups (RI, RII) are by an rotation of $\phi$ about the c-axis. For TI, TII crystals:

$$\tan 4 \phi = \frac{C_{16}}{C}$$

where $C = C_{11} - C_{12} - 2C_{66}$

in the case $\phi = 0$, $C_{16} = 0$:

For RI, RII crystals:

$$\tan 3\phi = \frac{C_{15}}{C_{14}}$$

in the case $\phi = 0$, $C_{15} = 0$;

In other words, when $\phi = 0$, TI and TII are identical (and $\phi = 0$, RI and RII are identical). In so far as their acoustic properties are concerned.

As mentioned earlier, the concept of the acoustic axis has been demonstrated experimentally (1,2). In the following we are going to show that these expressions can be obtained by using IC and the results are in full agreement with that of (3).
THE APPLICATION OF THE INVARIANT CONSTANTS

IC are the set of elastic constants which does not dependent on the coordinate system chosen (4). The numbers of IC for each crystal classes are the same as the regular elastic constants, but the values are different. After some simple calculations, the relationships between the two sets can be obtained. For TI, TII groups:

\[ C_{11} = C_{22} = \lambda_1 + 2 \lambda_2 + \lambda_3 \]
\[ C_{33} = \lambda_1 + 2 \lambda_2 + \lambda_4 + 2 \lambda_5 + 4 \lambda_6 \]
\[ C_{23} = C_{31} = \lambda_1 + \lambda_5 \]
\[ C_{55} = C_{44} = \lambda_2 + \lambda_6 \]
\[ C_{16} = \lambda_7, \quad C_{66} = \lambda_2, \quad C_{12} = \lambda_1 \]

(for TI groups, \( C_{16} = \lambda_7 = 0 \))

In tetragonal crystals, by using these equations and IC we can write e.g., \( C_{11} \) as a function of \( \phi \), where \( \phi \) is the angle of rotation about c-axis as follows:

\[ C_{11} = \lambda_1 + 2 \lambda_2 + 3/4 \lambda_3 \]
\[ + (1/4 \lambda_3 \cos 4 \phi + \lambda_7 \sin 4 \phi) \]

which turns out to be the same expression obtained by Khatkevich with

\[ \tan \phi = -C_{16}/(C_{11} - C_{12} - 2C_{66}) \]

For RI, and RII group crystals, we have:

\[ C_{11} = \lambda_1 + 2 \lambda_2, \quad C_{13} = \lambda_1 + \lambda_4 \]
\[ C_{33} = \lambda_1 + 2 \lambda_2 + \lambda_3 + 2\lambda_4 + 4 \lambda_5 \]
\[ C_{44} = \lambda_2 + \lambda_5, \quad C_{12} = \lambda_1 \]
\[ C_{66} = \lambda_2, \quad C_{15} = \lambda_7 \]

(for RI group crystals: \( C_{15} = \lambda_7 = 0 \)). Again by using the transformation equations of IC, we obtained the same expressions for \( \phi \), i.e., \( \tan 3 \phi = C_{15}/C_{14} \) as that of (3).

Thus we have demonstrated that by using IC on the rotation \( \phi \) about the acoustic axis, we get the identical results as that of Khatkevich. This is because of the fact that the IC are not dependent on any particular coordinate systems being used. A simple rotation of \( \phi \) can relate the two crystal groups directly.

CONCLUSION:

We have shown that using the IC, it is easy to obtain the expression for \( \phi \), which in turn gives the relationships between TI and TII group crystals. Similarly for RI and RII group crystals. It comes out naturally as the characteristics of IC. The results are in full agreement with that of Khatkevich (3). We are currently investigating other properties of IC with the hope that further physical insight may be obtained in relating the elastic property to the other physical properties of solids.

ACKNOWLEDGEMENT

The valuable help and discussions with Professor Y. Li is highly appreciated.

REFERENCES


ULTRASONIC BEHAVIOR IN SUPERIONIC GLASS (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub>

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Over past ten years, the superionic glasses, having a great technological interest, had been intensively studied, and many efforts had been made to obtain the information about structure role on the ionic dynamics and the possible coupling mechanisms. The purpose of the following: 1) Ultrasonic attenuation and velocity behavior in this system, 2) What role does the AgI play in the ion conduction and influence on glass microscopic structure? 3) To examine the universality of IDR theory for superionic glass materials.

EXPERIMENTAL PROCEDURES

Series superionic glass (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub> (x = 0.50, 0.55, 0.60, 0.65) samples for this study were prepared with <br>8×13×5mm<sup>3</sup> by liquid nitrogen rapid quenching technique. All of the samples were examined by X-ray diffraction to confirm their amorphism, and polished optically in correct shape for ultrasonic measurements. Both ultrasonic attenuation and velocity were measured by MATEC 7700 ultrasonic modulator and receiver system with pulse echo technique in the temperature range 100-300K, at frequencies 5, 10 and 14MHz.

EXPERIMENTAL RESULT AND DISCUSSION

Sound Velocity

The AgI content dependence of the sound velocity for longitudinal wave at room temperature shows in Fig.1.

Fig.1 The sound velocity in (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub> as a function of AgI content x at room temperature

Fig.2 The temperature dependence of the sound velocity for longitudinal and shear waves in (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub>

Fig.2 shows that sound velocities are almost same for two different polarized waves in the measured temperature range. This phenomenon implies that there is a same dispersion mechanism for two kinds of ultrasonic waves in (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub> superionic glass.

Ultrasonic Attenuation

Ultrasonic attenuation as a function of temperature in (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub> (x = 0.50, 0.55, 0.60, 0.65) are shown in Fig.3.

Fig.3 The temperature dependence of ultrasonic attenuation in (AgI)(Ag<sub>3</sub>PO<sub>4</sub>)<sub>1-x</sub>
It is obvious that there is a broad peak in measured temperature range. The height of the peak is getting bigger and bigger with increasing AgI content, and the temperature of the peak is getting lower and lower. To examine the character of these peaks, the temperature–dynamic experiment of ultrasonic attenuation at different frequency 5, 10 and 14MHz in (AgI)\textsubscript{x}Ag\textsubscript{3}PO\textsubscript{4}Ag\textsubscript{3} has been carried out, and the result is shown in Fig.4.

![Graph showing ultrasonic attenuation dependence on temperature](image)

**Fig.4** The ultrasonic attenuation dependence on temperature at different frequency in (AgI)\textsubscript{x}Ag\textsubscript{3}PO\textsubscript{4}Ag\textsubscript{3}.

The fact that the temperature of the peak shifts to higher and the height of it gets bigger with increasing frequency verifies that the broad peak is a typical relaxation one. The behavior of ultrasonic attenuation in Fig.3 indicates that with AgI content progressive increasing the total number of the mobile silver ions Ag\textsuperscript{+} which will join in the thermal activated relaxation processes is getting more and more, so that ultrasonic attenuation gets an increase. Furthermore, it should be pointed from the lower right corner in Fig.3 that as AgI content x > 0.60, \( a_{\max} \) the ultrasonic attenuation of peak, increases faster. This is because of a strong interaction between PO\textsubscript{4} tetrahedral chain and AgI sublattice. It is this strong interaction that AgI sublattice may become a path and part of silver ions in Ag\textsubscript{3}PO\textsubscript{4} will move in the path and absorb more ultrasonic energy.

**Theory and Fitting Data**

By the IDR theory, dissipation and relaxation phenomena in condensed matter associated with the existence of some ubiquitous very–low–energy excitation in the system. These excitations exhibit an infrared divergence–like response to transition of ultrasonic behavior induced by a time–varying stress field.\(^3\) The following ultrasonic attenuation formula was deduced:\(^2\)

\[
a = A\beta e^{-\gamma/(\omega,\tau_{\text{rel}})}
\]

(1)

where A is constant, \( \beta = 1/(kT) \), \( \theta_{\text{rel}}(e) \) is Levy function, \( \omega \) is angular frequency of ultrasonic wave, n is infrared divergence index and \( \tau \) is relaxation time which obeys Arrhenius Equation:

\[
\tau = \tau_{\text{rel}} \exp \left( \frac{E^*}{kT} \right)
\]

(2)

where \( E^* \) is apparent activated energy.

For a fixed \( \omega \),

\[
(\Delta T)_{\text{max}} = \frac{Z\theta_{\text{rel}}(e)}{[Z\theta_{\text{rel}}(e)]_{\text{max}}}
\]

(3)

Choosing n and \( E^* \) in equation (3), we can fit the theoretical curve of equation (3) to ultrasonic experimental data. The values of the parameters obtained in ultrasonic field are given in the Table 1, and the fitting curves are shown in Fig.5.

**Table 1.** Values of the most probable apparent activated energy \( E^* \) and infrared divergence index n in ultrasonic field

<table>
<thead>
<tr>
<th>x</th>
<th>n</th>
<th>( E^*(eV) )</th>
<th>( r_{\text{rel}}^{-1} )</th>
<th>( E^*(eV) )</th>
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<tr>
<td>0.55</td>
<td>0.648</td>
<td>0.385</td>
<td>5.87 \times 10^{-4}</td>
<td>0.380</td>
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<tr>
<td>0.60</td>
<td>0.635</td>
<td>0.364</td>
<td>5.03 \times 10^{-4}</td>
<td>0.362</td>
</tr>
<tr>
<td>0.65</td>
<td>0.625</td>
<td>0.349</td>
<td>4.14 \times 10^{-4}</td>
<td>0.350</td>
</tr>
</tbody>
</table>

Where \( E^* \) is the apparent activated energy in electric field.

![Graph showing fit of IDR theory](image)

**Fig.5** The fit of IDR theory to ultrasonic experimental data

We can see from Table 1 that the apparent activated energy \( E^* \) is close to each other, and they are both on the decrease with increasing AgI content. It states that there are similar microscopic mechanisms for the motion of the mobile silver ions Ag\textsuperscript{+} in electric field and stress field. Fig.5 indicates that experimental data agree with IDR theory at above 200K. Perhaps, there is other kind of mechanism below 200K for this system material.

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References
THE ELECTRO–ELASTIC CONSTANTS FOR THE PbB₉O₁₇ PIEZOELECTRIC CRYSTAL

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INTRODUCTION

The Crystals of the borate series are new nonlinear optical crystals. They belong to orthorhombic system, C₃ᵥ point group. There are nine independent elastic, five independent piezoelectric, and three independent dielectric constants in this crystal. Calculation formulas were derived from elementary theory of thickness mode vibrations for piezoelectric thin plates. The electro–elastic constants and electromechanical coupling coefficients were obtained from experimental measurements.

THEORY

Thickness mode vibrations of thin plates

For thickness vibrations of piezoelectric plates with thickness t, normal vector n, the nonzero solution satisfied the differential equations and the boundary conditions exist if the following secular equation hold: ⁰
\[ |C''_{xy} \cdot C''_{yx} - C''_{xx} | = 0 \]  (1)
where \( C''_{xx} \) is a piezoelectrically stiffened elastic constant, \( C'' \) is an effective elastic constant.

The transcendental frequency equation which satisfies boundary conditions is given by
\[ |\beta_{x} | [C''_{xy} \cdot C''_{yx} \cdot \cos k_{y} \cdot \cos k_{z} \cdot \frac{e_{nx}}{e_{nz}} \cdot \sin k_{z} \cdot \frac{t}{2} ] = 0 \]  (2)
For high overtones, Eq.(2) is closely approximated by
\[ \bar{C}_{x} = 4\pi (f_{nm} / m)^{3} \quad n = 1, 2, 3 \]  (3)
where \( f_{nm} \) is the frequency of the nth overtone, \( \rho \) is density.

The plates which are different orientations must be selected. For a selected orientational plate three high overtone resonant frequencies \( f_{nm} \) may be measured, thus all of the electro–elastic constants will be obtained by Eq.(1) and (3).

Analysis of vibration for \( C''_{xx} \) thin plate

From Eq.(1), we can get following secular determinants for six different vibration modes.

\[ \begin{vmatrix} C''_{xx} - \bar{C} & 0 & 0 \\ 0 & C''_{yy} - \bar{C} & 0 \\ 0 & 0 & C''_{zz} + e_{15}^{2} / \varepsilon_{11} - \bar{C} \end{vmatrix} = 0 \]

It is a cubic equation in \( \bar{C} \) and yields three real positive roots:
\[ \bar{C}_{x} = \bar{C}_{y}, \quad \bar{C}_{y} = \bar{C}_{z}, \quad \bar{C}_{z} = \bar{C}_{x} + e_{15}^{2} / \varepsilon_{11} \]  (4)
It shows that a electric field direction parallel to the major surface can excite one unstiffened pure extensional mode, a field direction perpendicular to the major surface can excite one stiffened pure shear mode.

Similarly, three roots of the Eq.(1) for a ⁰ X–cut are
\[ \bar{C}_{x} = \bar{C}_{xz}, \quad \bar{C}_{y} = \bar{C}_{yz}, \quad \bar{C}_{z} = \bar{C}_{xz} + e_{33}^{2} / \varepsilon_{33} \]  (5)
on one unstiffened pure extensional mode, one stiffened pure shear mode, respectively.

For a ⁰ Z–cut, three roots of the Eq.(1) are
\[ \bar{C}_{x} = \bar{C}_{xz}, \quad \bar{C}_{y} = \bar{C}_{yz}, \quad \bar{C}_{z} = \bar{C}_{xz} + e_{33}^{2} / \varepsilon_{33} \]  (6)
two unstiffened pure shear modes, one stiffened pure extensional mode, respectively.

For 45* rotated X–cut, one root is \( \bar{C}_{x, y} \), the sum and product of other two roots \( \bar{C}_{x, 1} \) and \( \bar{C}_{x, 2} \) are written as follows:
\[ \Sigma_{x} = \bar{C}_{x, 1} + \bar{C}_{x, 2} = \frac{C_{11} + C_{12}}{2} + \frac{C_{11}^{2} + C_{12}^{2}}{2} \]  (7)
\[ \Pi_{x} = \bar{C}_{x, 1} \cdot \bar{C}_{x, 2} = \left( \frac{C_{11} + C_{12}}{2} \right)^{2} + \frac{C_{11}^{2} + C_{12}^{2}}{2} \]  (8)
\[ \bar{C}_{x, y} = \frac{C_{11} + C_{12}}{2} + \frac{C_{11}^{2} + C_{12}^{2}}{2} \]  (9)
It expresses that parallel–field can excite two unstiffened extensional shear coupling modes, perpendicular–field can excite one stiffened shear mode.

Similarly, for a 45* rotated Y–cut and 45* Z–cut, these roots are:
\[ \bar{C}_{y, 1}, \quad \Sigma_{y} = \bar{C}_{y, 1} + \bar{C}_{y, 2}, \quad \Pi_{y} = \bar{C}_{y, 1} \cdot \bar{C}_{y, 2} \]
\[ \bar{C}_{z, 1}, \quad \Sigma_{z} = \bar{C}_{z, 1} + \bar{C}_{z, 2}, \quad \Pi_{z} = \bar{C}_{z, 1} \cdot \bar{C}_{z, 2} \]

Calculation of electro–elastic constants

Formulas found electro–elastic constants for \( C_{xx} \) from several cut–plates are as follows: ⁰ X–cut

Elastic constants
\[ C_{11}'' = \bar{C}_{x} = \frac{f_{nm}^{2} \cdot \varepsilon_{11}}{m^{2}} \]  (10)
\[ C_{33}'' = \bar{C}_{x} (1 - K_{13}^{2}) \]  (11)
Piezoelectric constant
\[ e_{15} = K_{13} (\bar{C}_{x} \cdot e_{11}^{2})^{1/2} \]  (12)
Effective elastic constant

\[ \bar{C}_{33} = 4p f_{1m}^2 \cdot t/m^3 \]

0° Y-cut

\[ C_{xx} = \bar{C}_{33} = 4p f_{1m}^2 \cdot t/m^3 \]  \hspace{1cm} (13)

\[ C_{yy} = \bar{C}_{33}(1 - K_{ab}) \]  \hspace{1cm} (14)

\[ \epsilon_{33} = K_{ab}(\bar{C}_{33} + \epsilon_{33}')^{1/2} \]  \hspace{1cm} (15)

\[ \bar{C}_{33} = 4p f_{1m}^2 \cdot t/m^3 \]

0° Z-cut

\[ C_{zz} = \bar{C}_{33}(1 - K_{ab}) \]  \hspace{1cm} (16)

\[ \epsilon_{33} = K_{ab}(\bar{C}_{33} + \epsilon_{33}')^{1/2} \]  \hspace{1cm} (17)

\[ \bar{C}_{33} = 4p f_{1m}^2 \cdot t/m^3 \]

45° rotated X-cut

The effective elastic constants

\[ C_{33} = 4p f_{1m}^2 \cdot t/m^3 \], \[ C_{xx} = 4p f_{1m}^2 \cdot t/m^3 \]

The elastic constant \( C_{33} \) can be obtained from Eq.(7)

\[ C_{33} = \Sigma_{C} - (C_{11}^e + C_{33}^e) / 2 \]  \hspace{1cm} (18)

\[ C_{11} \] can be found from Eq.(19) which is derived from Eq.(8).

\[ C_{11} + [2\Sigma_{C} - (C_{11}^e + C_{33}^e)]C_{11}^e + [4\Pi_{C} - (C_{11}^e + C_{33}^e)\Sigma_{C} + C_{11}^e C_{33}^e] = 0 \]  \hspace{1cm} (19)

45° rotated Y-cut

To measure \( f_{33} \) and then calculate \( C_{yy}, C_{yy}, \epsilon_{33} \) and \( C_{33}^e \) are respectively obtained by quadratic equations connected with \( \Sigma_{C} \) and \( \Pi_{C} \).

For a 45° rotated Z-cut plate, by the same as above procedure \( C_{11} \) and \( C_{11}^e \) can be found out.

EXPERIMENT AND RESULTS

The single crystal of Pb\(_3\)B\(_2\)O\(_5\) has been grown by ourself. It is transparent, the density is 5.87g/cm\(^3\). The experiments for above seven kinds of different cut plates have been carried out with the transmission line method \( \Sigma_{C} \) \( (1) \)

All of the elastic constant \( C_{33}^e \) and piezoelectric stress constant \( \epsilon_{33} \) can be determined by Eq.(10) to (21).

The electromechanical coupling factor can be obtained from the ratios of measured fundamental and overtone resonant frequencies, \( k_{33} \) is yielded by Eq.(21).

Dielectric constants can be obtained from capacitance measurement with low and high frequencies for the 0° X-cut, 0° Y-cut and 0° Z-cut plates, respectively.

The results are summarized in Table 1.

<table>
<thead>
<tr>
<th>TABLE 1.</th>
<th>Constants of Pb(_3)B(_2)O(_5)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Elastic stiffness constants (×10(^{11})N/m(^2))</td>
<td></td>
</tr>
<tr>
<td>( C_{11}^e )</td>
<td>( C_{12}^e )</td>
</tr>
<tr>
<td>3.28</td>
<td>0.50</td>
</tr>
<tr>
<td>Compliance Constants (×10(^{12})m(^3)/N)</td>
<td></td>
</tr>
<tr>
<td>( S_{11}^e )</td>
<td>( S_{12}^e )</td>
</tr>
<tr>
<td>2.96</td>
<td>-4.44</td>
</tr>
<tr>
<td>Piezoelectric Stress Constants (C/m(^2))</td>
<td></td>
</tr>
<tr>
<td>( e_{11}^e )</td>
<td>( e_{22}^e )</td>
</tr>
<tr>
<td>1.56</td>
<td>0.89</td>
</tr>
<tr>
<td>Piezoelectric Strain Constants (×10(^{-11})C/N)</td>
<td></td>
</tr>
<tr>
<td>( d_{13} )</td>
<td>( d_{24} )</td>
</tr>
<tr>
<td>1.53</td>
<td>0.70</td>
</tr>
<tr>
<td>Dielectric Constants</td>
<td></td>
</tr>
<tr>
<td>( e_{11} )</td>
<td>( e_{22} )</td>
</tr>
<tr>
<td>24.2</td>
<td>13.3</td>
</tr>
<tr>
<td>Electromechanical Coupling Factors</td>
<td></td>
</tr>
<tr>
<td>( k_{15} )</td>
<td>( k_{24} )</td>
</tr>
<tr>
<td>0.33</td>
<td>0.23</td>
</tr>
</tbody>
</table>

The experimental accuracy of the above values depends on measuring accuracy of a series of high order resonance frequencies, on a positive identification of any one mode sequence and on the dimension, produced flatness and parallelism of samples. In addition the high order overtone frequencies should be selected as high as possible. As finding out above quadratic equations, the sign of the relative constants should be chosen as specified by the IRE standard on piezoelectric ceramic crystals, and according to the consistency between the data.

The measuring results show that there are one stiffened pure shear mode for 0° X-cut plate, coupling factor \( k_{33} \) = 0.33. It implies that this cut would make a shear wave transducer, and exist as stiffened pure mode of vibration for other several cut plates. It appears to be useful for transducer application.
ACOUSTICAL INVESTIGATION OF MOLTEN ALKALI METALS

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INTRODUCTION

Metals are elements comprising over three fourths of the periodic table. Based on distinctions given in the Metals Handbook (12), 85 of the 118 elements are metals. The fundamental ultrasonic measurements in any material are those of the propagation velocity and the amplitude absorption coefficient (12). Three reviews have appeared in the literature in recent years devoted to the propagation of sound in liquid metals (12-15).

The acoustical investigations of molten alkali metals are of interest in both technological and theoretical problems. Their use as liquid metal coolants involves problems strongly connected with their thermodynamic properties. A systematic investigation of molten Li, Na, K, Rb and Cs metals has been attempted in this article, using experimental surface tension and density data, taken from literature (15). The obtained results have been used to explain the behaviour of these metals in the temperature range under investigation.

THEORY

Using experimental surface tension and density data taken from literature (15), several acoustical and thermodynamical parameters for molten alkali metals (e.g. Li, Na, K, Rb and Cs) have been evaluated using standard relations.

According to the Auerbach relation (16) the velocity of sound (u), obeys the equation:

\[ u = \sqrt{\alpha/\beta} \]  

(1)

Where \( \alpha \) and \( \beta \) are surface tension and density of the system under study.

Using the Auerbach relation, ultrasonic velocities are evaluated. Taking these ultrasonic velocities and density data, acoustical parameters such as adiabatic compressibility (\( \kappa_a \)), molar sound velocity (\( \kappa \)), molar adiabatic compressibility (\( \kappa_h \)) and acoustic impedance (\( Z \)) have been evaluated using standard relations reported by Blatt et al. (16). The available value (V\( _m \), free volume (V\( _f \)), internal pressure (P\( _i \)), van der Waal's constants (a & b) and intermolecular free length (L\( _{\text{free}} \)) are also determined using standard formulae reported elsewhere (16).

The empirical Eykman equation (18) is used to determine Eykman constant (\( \kappa \)). The ratio of molar sound velocity (\( \kappa \)) and geometrical volume (\( \beta \)) is also determined using expressions reported earlier (16).

The alkali metals studied here are investigated in the following temperature range.

<table>
<thead>
<tr>
<th>Metal</th>
<th>Temperature range</th>
</tr>
</thead>
<tbody>
<tr>
<td>Lithium</td>
<td>Li</td>
</tr>
<tr>
<td>Sodium</td>
<td>Na</td>
</tr>
<tr>
<td>Potassium</td>
<td>K</td>
</tr>
<tr>
<td>Rubidium</td>
<td>Rb</td>
</tr>
<tr>
<td>Cesium</td>
<td>Cs</td>
</tr>
</tbody>
</table>

Some of the representative results for these metals for the extreme ends of temperature ranges (denoted as M-1 and N-1, where M is the corresponding metal) are reported in Table 1 to Table 3. All the parameters listed in these tables are in SI units except reported otherwise.

<table>
<thead>
<tr>
<th>Metal</th>
<th>Temperature range</th>
</tr>
</thead>
<tbody>
<tr>
<td>Li</td>
<td>1955 K - 1500 K</td>
</tr>
<tr>
<td>Na</td>
<td>70 K - 1120 K</td>
</tr>
<tr>
<td>K</td>
<td>728 K - 975 K</td>
</tr>
<tr>
<td>Rb</td>
<td>758 K - 945 K</td>
</tr>
<tr>
<td>Cs</td>
<td>767 K - 915 K</td>
</tr>
</tbody>
</table>

<p>| Table 1. Some Acoustical Parameters of molten alkali metals |
|-------------|------------------|------------------|------------------|------------------|</p>
<table>
<thead>
<tr>
<th>Metal</th>
<th>( u )</th>
<th>( \kappa_a )</th>
<th>( \kappa_h )</th>
<th>( \kappa )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Li</td>
<td>1.42</td>
<td>1.28</td>
<td>0.33</td>
<td>21.42</td>
</tr>
<tr>
<td>Na</td>
<td>2.48</td>
<td>1.25</td>
<td>0.44</td>
<td>18.05</td>
</tr>
<tr>
<td>K</td>
<td>6.08</td>
<td>1.25</td>
<td>0.50</td>
<td>14.18</td>
</tr>
<tr>
<td>Rb</td>
<td>6.72</td>
<td>1.25</td>
<td>0.41</td>
<td>9.95</td>
</tr>
<tr>
<td>Cs</td>
<td>8.79</td>
<td>1.28</td>
<td>0.40</td>
<td>8.52</td>
</tr>
<tr>
<td>Cs</td>
<td>9.82</td>
<td>1.27</td>
<td>0.46</td>
<td>7.94</td>
</tr>
<tr>
<td>Cs</td>
<td>12.18</td>
<td>1.30</td>
<td>0.51</td>
<td>6.72</td>
</tr>
</tbody>
</table>

| Table 2. Some thermodynamical properties of molten alkali metals |
|-------------|------------------|------------------|------------------|------------------|
| Metal | V\( _m \) | \( P_i \) | a | b | \( L_{\text{free}} \) |
|----------|------------------|------------------|------------------|------------------|
| Li       | 8.74             | 4.50             | 5.22             | 10.85            | 18.14            |
| Na       | 8.06             | 5.79             | 6.10             | 13.37            | 29.78            |
| K        | 9.19             | 6.80             | 7.50             | 17.24            | 39.15            |
| Rb       | 6.41             | 6.53             | 6.84             | 16.74            | 29.75            |
| Cs       | 4.97             | 6.53             | 5.84             | 16.74            | 29.75            |

<p>| Table 3. Some thermodynamical parameters of molten alkali metals |
|-------------|------------------|------------------|------------------|------------------|</p>
<table>
<thead>
<tr>
<th>Metal</th>
<th>V( _m )</th>
<th>( E_{\text{cal}} )</th>
<th>( \kappa )</th>
<th>E</th>
</tr>
</thead>
<tbody>
<tr>
<td>Li</td>
<td>1.62</td>
<td>13.56</td>
<td>1.28</td>
<td>0.33</td>
</tr>
<tr>
<td>Na</td>
<td>2.48</td>
<td>11.25</td>
<td>0.44</td>
<td>18.05</td>
</tr>
<tr>
<td>K</td>
<td>6.08</td>
<td>11.25</td>
<td>0.50</td>
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<td>6.72</td>
<td>11.25</td>
<td>0.41</td>
<td>9.95</td>
</tr>
<tr>
<td>Cs</td>
<td>8.79</td>
<td>11.25</td>
<td>0.40</td>
<td>8.52</td>
</tr>
</tbody>
</table>

RESULTS AND DISCUSSION

The calculated values of some acoustical parameters of molten alkali metals at extreme and temperatures of the temperature range under investigation for each alkali metal are reported in Table 1. A perusal of the results reveals that the values of ultrasonic velocity decreases as we go down in the group of alkali metals i.e. from Li to Cs. The values of \( \kappa_a \) for each metal increase with rise in temperature and these values increases with rise in atomic number of the alkali metal as well. The high values of ultrasonic velocities in these molten metals can be understood in the light of the fact that supercooling vapour pressure on these systems varies between 29 and 480 Torr. The systematic increase in \( \kappa_a \) values for these molten alkali metals is as expected and is in agreement with those reported for other molten metals (12). A strong temperature dependence of \( \kappa_h \) and \( \kappa \) is observed for each metal. The decreasing trend of \( \varepsilon \) and \( \kappa \) with rise in temperature is observed. However, \( R \) and \( W \) values rise as we go from one alkali metal to another in this group of metals in the periodic Table. This may be attributed to the increase in volume of...
the metal due to increase in its atomic number thus increase in size of the metal atoms is well reflected in the large magnitudes of R and M. This is as expected. The temperature dependent behavior of R and M point to the associated nature of the metal under study. This is in line with results reported for similar systems by Bhatti et al. (12). The effect of temperature rise is to cause thermal expansion of the system which is reflected in the decreasing trend of acoustic impedance (Z) in each case. The observed trend of Z is in accordance with the results reported in literature (13) for other metals.

The two traditional concepts of the liquid state regard a liquid either a very dense gas or a solid on which the long range order between molecules has been broken down. In considering the propagation of sound in liquids and solids, the gas model is found particularly suitable as it enables to predict the sound velocity qualitatively. The free volume (Vf) and internal pressure (Pi) values determined for the system under study are extension of the concepts "valid for solids" and "in the liquids" (14). The free volume values obtained for these metals are of the order 10^10 m^-3 in the temperature range under study. The high value of Pi observed for these metals are in the order as found in the case of other highly associated liquids (14), and molten metals (12). The free volume values obtained from the relation V = V0 + V show an increase similar to that observed for Li for these metals. The van der Waal's constant a and b show a rapid increase in their values with rise in atomic number of the metals under study. The rise in intermolecular free length (L) with rise in temperature can be understood in terms of weakening of quasi-crystalline or remnant crystal structures in these systems due to thermal agitation. The rise in atomic number also leads to an increase in the intermolecular free lengths, as is evident from Table 2. A linear temperature dependence with positive temperature coefficient is observed for the parameters Vf and Lf, whereas Pi and a values do not change much with temperature rise in each case. The trends of variation of a and b, with rise in temperature, as observed in previous studies are in conformity with trends observed for other metals (12). The increase in available volume (Va) of the alkali metals with rise in temperature (and also with rise in atomic number) has supports and supplements the conclusions arrived at above. The observed changes in Va, Vf and Lf are caused by the weakening of covalent interactions in these liquids with rise in temperature of the system and also with increase in atomic number. For unassociated liquids KfKf is expected to be nearly unity, while its value should be very small for associated liquids (15) and melts (16). The small value of ratio KfKf for these molten metals point to the presence of strong interaction in these. In case of molten metals a high degree of covalent interactions would be expected in the near and next near neighbors. The increase in the ratio KfKf (Table 2) with rise in temperature (and also with rise in atomic number) show the disruption of ionic order due to thermal agitation and size effects. The value of the excess energy of association per mole (over and above the van der Waal's contribution) E decreases with rise in temperature. It has been shown that for liquids with strongly interacting molecules such as hydrogen bonded systems (17) and melts (18) the E values are known to be very high and tend to fall with rise in temperature. The large E values (Table 3) obtained in present studies confirm the presence of a highly associated molten state of Li, as compared to Cs, due to the strong metallic bonding among its atoms. This can be understood in terms of the smallest size of the Li atoms as compared to other metallic atoms of this group of metals. Schaeff (17) has shown that in sound transmission, the molecular sound volume (B) plays a more important role than atomic sound velocity (V). It has been shown that the ratio of molecular sound volume (B) and actual volume of the molecules per mole (B) is around 1.5 for unassociated liquids and is lower for polar liquids. In the present investigation, B/B increases from 1.78 to 1.38, pointing to weakening of interatomic interactions in these systems with rise in temperature or with increase in atomic number. The Schaeff parameter (Ec) however, remains almost constant with temperature change for a given system though it reports a rapid fall with rise in atomic number of alkali metals.

REFERENCES

DIFFRACTION D'UNE ONDE PLANE PAR UN RESEAU D'OBSTACLES CYLINDRIQUES PLONGES DANS UNE COUCHE VISCOELASTIQUE

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G.E.R.D.S.N., 83140 Six Fours, France.

1- INTRODUCTION

On étudie un modèle théorique d'un matériau acoustique constitué par un réseau de diffuseurs incorporés dans une lame viscoélastique à faces parallèles. Ce type de matériau absorbant est utilisé pour le traitement anéchoïque des surfaces sous-marines. Le problème a été étudié, dans l'hypothèse d'ondes scalaires par Vidoret [1]. L'ambiance technique d'éléments finis par Hladky-Hennion [3]. La méthode proposée s'appuie sur une théorie du réseau de diffraction exposée dans [4].

2- THÉORIE

Les données du problème sont résumées dans la fig. 1. L'onde incidente est une onde plane monochromatique (facteur temporel en exp(-jωt)).

Fig. 1. Géométrie de la lame.

On distingue 5 régions : la région 1, délimitée par la face soumise à l'onde incidente, est un fluide (milieu marin). Les régions 2 et 3, séparées de manière fictive par le plan du réseau, sont constituées par un même milieu élastique (ou viscoélastique caractérisé par un module d'Young complexe). Les obstacles, région 5, sont également de nature élastique. Enfin, le domaine inférieur 4 est un fluide.

2.1- ÉTUDE DE LA LAME A FACES PARALLÈLES

On peut assimiler l'ensemble constitué par la lame et son environnement à un milieu stratifié comportant trois interfaces. Considérons une onde incidente sur une des interfaces, l'onde envisagée peut être soit plane homogène, soit périodique de même période que le milieu. Les ondes transmises et réfléchies seront également périodiques, on peut les développer en modes élastiques :

\[ \Phi_n^\pm = \{ \exp \{ \pm jk x \sin \alpha + y \cos \alpha \} \} \]

\[ k = \cos \alpha = k \cos \alpha + \sin \alpha \]

\[ \alpha = 1 - \cos \alpha \]

Dans les domaines occupés par un fluide, le champ scalaire est défini à partir de la pression \( p \). En projetant cette fonction sur la base modale on obtient l'expression :

\[ p = \sum_{n=-\infty}^{\infty} (\Phi_n^+ + \Phi_n^-) \]

qui dépend de 2 vecteurs \( X = [X_1] \) et \( X = [X_2] \).

Dans les régions élastiques, le champ acoustique est exprimé à partir des potentiels \( \Phi \) et \( \Psi \) des ondes longitudinales \( P \) et des ondes transversales \( S \) (tels que le déplacement \( u = \Phi + \Psi \wedge k \), \( k \) vecteur unitaire de oz). Après avoir remplacé \( (k) \) par \( k_p \) ou par \( k_s \) dans la formule (1), on obtient des décompositions de \( \Phi \) et \( \Psi \) analogues à (2) qui dépendent de 4 vecteurs : \( X_1 \), \( X_2 \), \( X_2 \) et \( X_2 \). Afin d'obtenir des représentations du champ formellement identiques dans les divers domaines, nous conviendrons de rassembler en un seul vecteur : \( X \) et \( X \) d'une part, \( X \) et \( X \) d'autre part.

Le champ monore ainsi représenté subit des discontinuités à la traversée des interfaces que l'on peut traduire par des coefficients de réflexion et de transmission. Les faces externes de la lame sont des dioptries dont la théorie est bien connue. L'interface du réseau sera étudiée au § 2.2.

Après avoir numéroté les interfaces de 1 à 3, nous introduisons des notations résumées dans la figure suivante:

\[
\begin{align*}
X_1 & = Y_1 \\
X_2 & = Y_2 \\
X_2 & = Y_3 \\
X_3 & = Y_4
\end{align*}
\]

La traversée de la i-ème interface se traduit par les équations :

\[ X_2^i = X_1^i + T_{12}^i Y_1^i, \quad Y_3^i = T_{13}^i X_1^i + R_{13}^i Y_1^i \]

Les vecteurs d'"entrée" des faces 2 et 3 résultent de la propagation des vecteurs de "sortie" des faces 1 et 2. En introduisant un "propagateur" \( t(e) \) fonction des quantités exp \( \{j(k)\} \) et \( \{i(k)\} \), on peut écrire les équations suivantes :

\[ X_2 = t_1(e_1) Y_1, \quad Y_3 = t_2(e_2) Y_1 \]

Les ondes réfléchies et transmises par la lame sont définies par les vecteurs \( X_1 \) et \( Y_3 \). Ces inconnues peuvent être calculées en résolvant le système linéaire constitué par les équations (2) et (3), et les conditions aux limites : \( X_1 \) donné, \( X_3 = [5] \) pour une onde plane, et \( Y_3 = 0 \).

2.2 THEORIE DU RESEAU

Plaçons nous dans une section droite de la lame, soit \( (x_1,y) \), un système d'axes tels que \( x \) soit le centre d'un obstacle, \( y \) coïncidant avec l'axe du réseau. Au voisinage de cet obstacle les potentiels des ondes \( P \) et des ondes \( S \) peuvent être développés en série d'ondes cylindriques, régulières à l'origine jusqu'à s'agir d'ondes incidentes, singulières pour les ondes diffusées. Les coefficients des ondes constituent des vecteurs désignés par \( B \), \( B \), pour les ondes incidentes et par \( A \), \( A \), pour les ondes diffusées. On définit des T-matrices de diffraction telles que :

\[ A = T BB + T SB, \quad A = T BB + T SB \]

La lame étant soumise à une onde plane incidente, les ondes diffusées par les différentes inclinations.
ne diffèrent que par un facteur de phase constant d'une inclusion à sa voisine. L'ensemble des ondes diffractées est donc entièrement déterminé par la donnée de 2 vecteurs colonne $\mathbf{A}^T$ et $\mathbf{A}^T$. Ces vecteurs peuvent être calculés en écrivant le problème de diffraction relatif à un élément arbitraire du réseau.

Un élément particulier est soumis à l'onde incidente et à une onde d'interaction égale à la somme des ondes diffractées par les autres éléments du réseau. L'onde d'interaction s'exprime en fonction des deux vecteurs $\mathbf{A}^T$ et $\mathbf{A}^T$ suivant une relation $\mathbf{S}^T\mathbf{A}^T$ pour les ondes $P$, et $\mathbf{S}^T\mathbf{A}^T$ pour les ondes $S$. $\mathbf{S}^T$ et $\mathbf{S}^T$ sont des matrices dont les coefficients s'expriment au moyen de série de Schrödinger [4].

La figure 2 est un exemple de calcul des coefficients de réflexion et de transmission spécifiques (modes 0) d'une onde plane tochant normalement sur un réseau d'événements cylindriques dans de la résine époxide, (par raison de symétrie il n'y a pas de couplage entre les modes $P$ et les modes $S$ d'ordre 0). Les données sont : $\alpha = 0$, $\mu_1/a = 1.5$, $\rho = 1.160$ kg/m$^3$, $c_p = 2770$ m/s, $c_s = 1363$ m/s.

2.3 - CALCUL DES T-MATRICES DES DIFFUSEURS

Nous devons calculer les matrices de transition définies en (5). Lorsque la section droite du diffuseur est un cercle ou une ellipse le problème peut être traité analytiquement. Afin d'apporter une solution générale nous proposons une méthode d'éléments finis de frontière dont le principe est le suivant :

Les potentiels $\Phi$ et $Z$ des ondes $P$ et $S$ sont exprimés au moyen de potentiels de couche mixte du type:

$$ f_c \mu \left( \frac{\Omega}{\Omega_0} + ik G \right) d \Omega \text{ (6 fonction de Green),}$$

(9)

 dont les densités inconnues $\mu_0$ et $\mu_0$ sont définies sur la frontière $C$ du domaine de l'obstacle. L'utilisation de couches mixtes assure l'unicité des représentations de $\Phi$ et $Z$, et l'écriture des conditions aux limites conduit à un système d'équations intégrales dont la solution est unique. Les équations intégrales sont discrétisées et résolues par collocation.

Les conditions aux limites font intervenir les dérivées partielles des potentiels de couche jusqu'à l'ordre deux. En dérivant sous le signe somme on obtient des équations divergentes qui doivent être calculées au sens des parties finies. Ce calcul est facilité en considérant des éléments de frontière rectilignes, ce qui revient à approcher la section de l'obstacle par un polygone. Les intégrales se réduisent à des intégrandes convergentes après intégration par partie. Les formules obtenues font intervenir des termes en figure le dérivées premières des densités de couche. Pour obtenir une expression correcte de ces termes nous avons choisi d'approcher les densités par des fonctions linéaires par intervalle.

3. CONCLUSION

Nous avons décrit un modèle mathématique dont l'objet est la prédiction des caractéristiques de réflexion et de transmission d'une onde plane par un réseau de diffuseurs inclus dans une lame viscoélastique. Cette approche est une alternative aux techniques purement numériques. Par rapport à celles-ci, elle a l'avantage de conduire à un code de calcul peu volumineux, qui peut être installé sur un micro-ordinateur.

BIBLIOGRAPHIE


USING MULTI-NODE METHOD TO DEAL WITH DISCONTINUOUS POINTS IN SOUND FIELD CALCULATION

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ABSTRACT

The surface Helmholtz formulation has been used for obtaining approximate solutions of the exterior steady-state acoustic radiation problem for an arbitrary surface. This paper presents a method for bodies with non-smooth surface, by putting up more than one node on the discontinuous point. Numerical results for radiation from bodies sitting on a rigid, infinite plane are given to verify this method.

INTRODUCTION

The boundary integral equation (BIE) method has been used extensively as a numerical technique to solve acoustic problems.\(^1\)\(^-\)\(^4\) To deal with bodies with corners and edges, a simple integral form for the evaluation of the coefficient C(P) in the Helmholtz integral equation was derived in Ref. 8. For other approaches to take care of corners and edges, we cite the earlier work by Leis,\(^11\) for the particular case of scattering from a cube, and the more recent developments by Costable et al.\(^12\)\(^-\)\(^14\) and Hebecker\(^15\) and Seybert.\(^16\)

In this article, we present a method for bodies with corners and edges. This method called the Multi–Node Method (MNM) is different from traditional approaches to take care of discontinuous points which modified the coefficient C(P) depended on the Green’s function as well as the local geometry. We introduce the concept of using more than one node, generally two or three nodes, on the point. Several test examples were run for the radiation from bodies on a rigid, infinite plane. Numerical results show good agreement with analytical solutions or numerical solutions of equivalent full-space problems.

I. BOUNDARY INTEGRAL FORMULATION

The starting point is the Helmholtz integral equation, which can be written in the form

\[
C\Phi + \int_{s} \frac{\Phi}{n_{s}} dS = \int_{s} \frac{u}{n_{s}} ds
\]  

(1)

This equation is valid in an acoustic domain B' exterior to a finite body B with boundary surface S and inward normal n. In Eq. (1), \(\Phi\) is the velocity potential satisfying the Helmholtz equation \(\nabla^2 \Phi + k^2 \Phi = 0\) in B' and the Sommerfeld radiation conditions in the farfield, \(u\) is the Green’s function. For a full-space problem, the Green’s function is

\[
u = e^{-i\kappa r} / (4\pi r)
\]

(2)

where \(k\) is the wavenumber and \(r\) is the distance between any two points P and Q. The coefficient \(C_{in}\) in Eq.(1) is 1 for P in the acoustic domain B' and zero for P in the body B. If P is on the boundary surface S, and there is a unique tangent plane at P on S, then \(C_{is}\) is 1/2.

Furthermore, if there exists an infinite plane S, as shown in Fig.1, a half-space Green’s function should be used in place of \(u\) in Eq. (1) to remove the contribution of S to the boundary integral equation.\(^10\) The half-space Green’s function, denoted by \(u\), takes the form

\[
\begin{align*}
\Phi_{in} &= e^{-i\kappa r} / (4\pi r) + R_{in} e^{-i\kappa r} / (4\pi r) \\
\Phi_{is} &= e^{-i\kappa r} / (4\pi r) + R_{is} e^{-i\kappa r} / (4\pi r) 
\end{align*}
\]

(3)

where \(R_{in}\) is the reflection coefficient of the infinite plane and \(r\) is the distance between Q and the image point P, with respect to S. The reflection coefficient \(R_{is}\) is equal to 1 for a rigid, infinite plane or -1 for a soft, infinite plane. The boundary integral equation for a half-space problem thus becomes

\[
C_{is}\Phi + \int_{s} \frac{\Phi}{n_{s}} dS = \int_{s} \frac{u}{n_{s}} ds
\]

(4)

II. THE MULTI-NODE METHOD

Figure 2 shows the geometry of the problem in three dimensions. Node N is on the corner of the body surface (show in Fig.2(a)). Number 1, 2, 3 refers to the different element. We put up 3 nodes \(N_1, N_2\) and \(N_3\) to replace node \(N\). These nodes have the same coordinate values but belong to different elements, have different boundary conditions. As Fig. 2 (b), node \(N_1\), on element 1 has the velocity potential value \(\Phi_1, \Phi_{n1}\) and normal \(n_1\), node \(N_2\) on element 2 has the same \(\Phi\) but different normal \(n_2\) and \(\Phi_{n2}\), node \(N_3\) on element 3 also has the same \(\Phi\), different \(n_3\) and \(\Phi_{n3}\). We can establish three boundary integral equations about these three nodes. This method to deal with discontinuous points is called MNM.
III. NUMERICAL RESULTS

A number of test cases were run to demonstrate the application of MNM for bodies in contact with an infinite plane. The first test case is the comparison of MNM with the analytical solution of the radiation from a pulsating hemisphere of radius \(a = 1\) sitting on a rigid, infinite plane. The solution for radiation with the surface velocity \(v(a) = 1\) prescribed on the hemisphere at frequency \(f = 100\) Hz is shown in Fig. 3. Fig. 3 also shows the analytical solution. The second case gives the comparison of MNM with numerical solution in full space. Because the problem is equivalent to a pulsating sphere in full space, due to symmetry. The third case is the comparison of MNM with the method provided by Ref. 16, which calculated the coefficient \(C(P)\) when the node was on the corner or on the edge.

![Fig. 2 Nomenclature for the Multi–Node Method](image)

![Fig. 3. Plot of \(|\Phi|\) at a distance \(x\) from the center of the hemisphere.](image)

![Fig. 4. Plot of \(|\Phi/\Phi_{\text{in}}|\) at a distance \(x\) from the center of the hemisphere (\(\Phi_{\text{in}}\) is the analytical solution).](image)

![Fig. 5. Plot of \(|\Phi/\Phi_{\text{in}}|\) at a distance \(x\) from the center of the hemisphere.](image)

IV. CONCLUSIONS

The Multi–Node Method has been given to solve the acoustic problems when the surface exist discontinuous points. More than one node are set up on the corner or on the edge. Numerical tests for the radiation from bodies sitting on a rigid, infinite plane have been carried out. By comparing the numerical results with either an analytical solution or the numerical solution of an equivalent full space problem, we have confirmed the validity of this method.

REFERENCE

EFFECT OF SPATIAL SAMPLING IN THE CALCULATION OF ULTRASONIC FIELDS GENERATED BY PLANAR SOURCES

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INTRODUCTION

Spatial sampling has been widely used for the calculation of ultrasonic fields generated by planar sources. Compared with other methods [1-3], this method is simple, can cope with sources with complex geometry and has no computational burden with the advent of fast computing machines. This method was illustrated for a rectangular piston by Ochelbre [4].

Quantitative analysis, however, on the aliasing errors which occur along with the spatial sampling has not been given. This analysis is very important in that we can guarantee the validity of the calculated sound field with the analysis. Therefore, we first analyze the aliasing errors and then perform numerical analysis for the case of rectangular piston radiators. The result can be utilized as a barometer for the calculation of ultrasonic fields with spatial sampling methods.

I. BASIC THEORY

Sound field generated by a planar source vibrating in an infinite baffle can be obtained by the well-known Rayleigh integral,

\[ p(x, y, z) = -\text{ip} \int \frac{v(x', y') e^{i k r}}{2 \pi r} dx' dy' \]

\[ = -\text{ip} \int v(x', y') \nabla g(x, y, z) \]

\[ R = \left( x-x' \right)^2 + \left( y-y' \right)^2 + \left( z-z' \right)^2 \]

where \( v(x', y') \) is the velocity distribution in the source plane, the symbol \( * \) denotes 2-dimensional convolution and \( g \) is the Green function

\[ g(x, y, z) = \frac{e^{-i k r}}{2 \pi i x^2 + y^2 + z^2} \]

If we sample the source distribution \( v(x', y') \) to \( \tilde{v}(x', y') \):

\[ \tilde{v}(x', y') = \sum_{n, m} v(x', y') \delta(x' - m \Delta x) \delta(y' - n \Delta y) dx' dy' \]

where \( \Delta x' \) and \( \Delta y' \) denote sampling interval in \( x' \) and \( y' \) axis in the source coordinates respectively. And the acoustic pressure \( p \) is modified to \( \tilde{p} \):

\[ \tilde{p} = -\text{ip} F^{*} \left[ \sum_{n, m} \tilde{V}(k_x - \frac{2 \pi m}{\Delta x}, k_y - \frac{2 \pi n}{\Delta y}) G \right] \]

\[ = p - \text{ip} F^{*} \left[ \sum_{k_x, k_y} \tilde{V}(k_x, k_y) \right] G \]

where \( \tilde{V}(k_x, k_y) \) and \( G(k_x, k_y) \) denote 2-dimensional Fourier transforms of \( v(x', y') \) and \( g(x, y, z) \), respectively. And \( G(k_x, k_y) \) is

\[ G(k_x, k_y) = \frac{e^{ik_x x}}{k_x} \left[ \frac{1}{k_x^2} - \frac{k_x^2}{z^2} \right] \left( k_x^2 + k_y^2 \leq k^2 \right) \]

Since \( \tilde{G} \) is not bandlimited, we can't obtain exact sound field by spatial sampling. But most of the sources which are practically useful retain beam patterns such that energy be concentrated in its main beam and thus, we can calculate sound field almost exactly with spatial sampling method. So we illustrate the validity of spatial sampling for the case of rectangular piston radiators.

II. SAMPLING IN THE RECTANGULAR PISTON RADIATOR

For a rectangular piston with dimension \( a \times b \) in an infinite baffle, the source distribution \( V(k_x, k_y) \) is

\[ V(k_x, k_y) = \text{absinc} \left( \frac{a k_x}{2 \pi} \right) \text{sinc} \left( \frac{b k_y}{2 \pi} \right) \]

From Eq.(4), we see that the source function of rectangular piston is modified to

\[ \tilde{V}(k_x, k_y) = \sum_{n, m} \text{asinc} \left( \frac{a}{2 \pi} k_x - \frac{2 \pi m}{\Delta x} \right) \text{sinc} \left( \frac{b}{2 \pi} k_y - \frac{2 \pi n}{\Delta y} \right) \]

For simplicity, if we consider the first part on the right side of Eq.(7) and let \( k_x = 2 \pi / \Delta x \) \& \( a = k / \Delta x \) and \( h = a k / 2 \), then for \( \alpha \) integer and after some algebra

\[ \sum_{n, m} \text{asinc} \left( \frac{a}{2 \pi} k_x - \frac{2 \pi m}{\Delta x} \right) \text{sinc} \left( \frac{b}{2 \pi} k_y - \frac{2 \pi n}{\Delta y} \right) \]
For efficient sampling i.e. \( h = \pi \), and from Eq. (9) we get

\[
|v|^2 = \sum_{n} \left[ \frac{1}{2} a(k_n - mk_f) \right]^2 \sum_{n} \left[ \frac{1}{2} b(k_n - mk_f) \right]^2
\]

\[
= \sum_{n} |v_n|^2
\]

(10)

where \( V_n \) is the source function of the \( m,m \) -th replicated source and \( m=n=0 \) yields the original source. Eq. (10) shows that at efficient sampling there occurs no cross term between each source and therefore, from Eq. (9) we see

\[
\Pi = (\text{power radiated by the original source}) + (\text{power radiated by replicas themselves})
\]

\[
= \Pi_{0,0} + \sum_{n \neq 0} \Pi_{n,n}
\]

(11)

And the aliasing error \( E \) can be defined as

\[
E = 10 \log_{10} \frac{\sum_{n \neq 0} \Pi_{n,n}}{\Pi_{0,0}}
\]

(12)

Eq. (12) was calculated and illustrated in Fig. 2 for rectangular piston with sampling frequency \( f=1/\Delta x=1/\Delta y=2-16 \) \([1/\lambda]\) assuming \( a=b=L \), for simplicity, and the integration was carried out by Gauss-Kronrod \((GK)\) pair [8]. The error becomes smaller as sampling frequency increases and at constant sampling frequency, better results are obtained with larger sources.

Fig. 3 illustrates the trends of aliasing errors for sampling frequencies 2, 4, 8 and 16 \([1/\lambda]\) as source dimension changes. At the constant sampling frequency the aliasing error decreases as source becomes larger because more ideal simple sources are generated across the source aperture via spatial sampling.

![Fig. 1. Modification of source function with sampling.](image)

![Fig. 2. Aliasing error for sources of \( L=1\lambda, 2\lambda, 5\lambda, 10\lambda, 50\lambda \text{ and } 100\lambda \).](image)

![Fig. 3. Trend of aliasing error as source dimension changes for sampling frequencies of 2, 4, 8, and 16 \([1/\lambda]\).](image)

III. CONCLUSION

Sampling effect of planar piston radiators has been analyzed and quantitative analysis has been performed on rectangular pistons. The 'efficient sampling' was found to be such that the source size multiplied by sampling frequency should be odd integers. As source dimension becomes larger the aliasing error decreases at constant sampling. And larger sampling frequency yields more accurate data as it should do from the physical point of view.

Since the radiated power doesn't depend on the distance from the source, the error in the calculated sound pressure becomes negligible as we move away from the source plane. Moreover, for near field pressure we could obtain fairly accurate sound field for radiative solution. For reactive solution in the near field, however, the modification factor blows off in the vicinity of \( \sin(\alpha \pi) = 0 \) and we can't guarantee the accurateness of the calculated sound field until the reactive solution becomes negligible when compared to the radiative solution. However, for efficient sampling \(( h = \pi \)), the modification factor converges to 1 as \( \sin(\alpha \pi) \to 0 \) and therefore, we can count on the calculated sound field even in the near field.

Reference

Diffraction on Focused Field of Short Ultrasonic Pulses

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1. Introduction

As well-known, the short pulse ultrasound has been widely employed now-a-day in the ultrasonic medical imaging technique in order to improve the longitudinal resolution. However, the diffraction patterns of short ultrasonic pulses must be considerably different from that of continuous waves because a short pulse includes wide frequency components. Although, studies have been contributed to the subject of short ultrasonic pulses[1], only a few analyses were involved in their focused field[2]. In the present work, we will give a systematic analysis for the short pulse focused field excited by a Gaussian transducer and a piston transducer based on the linearized KZK equation.

2. Theory

It is known that the KZK equation under the parabolic approximation takes the diffraction, absorption and nonlinearity of the sound wave into account consistently. Since the present work is limited only to treatment of the linear field, the linearized KZK equation can be expressed as follows[3],

\( \left( \frac{i}{\omega} \frac{\partial}{\partial \xi} - k_0 \frac{\partial}{\partial \tau} - k_0 \frac{\partial}{\partial \tau} - k_0 \frac{\partial}{\partial \tau} \right) \bar{E} = 0 \)  

(1)

Here \( p = \rho c \), \( \rho_0 \) and \( c \) represent the sound pressure, static density and sound velocity respectively. \( r = (\tau - t)/c \) stands for a dimensionless retarded time. \( \omega_0 \) is the angular frequency. \( \omega_0 \) represents the dimensionless axial distance. \( t_0 \) stands for the absorption coefficient of medium corresponding to \( \omega_0 \). In the cylindrical coordinate,

\( \frac{\partial^2}{\partial \xi^2} - \frac{1}{\rho^2} \frac{\partial^2}{\partial \tau^2} - k_0 \frac{\partial^2}{\partial \tau^2} \)

in which \( \rho \) is a radial coordinate and \( \omega_0 \) represents a reference radius. By letting \( x/a = \omega_0 \) the Fourier transform of \( p \) may be expressed as,

\( \bar{p}^i(\xi, \omega_0, \tau) = \frac{1}{\sqrt{2 \pi}} \int \bar{p}(\xi, \omega_0, \tau) e^{-i \omega_0 \tau} d\tau \)  

(2-a)

and its inverse transform is

\( \bar{p}(\xi, \omega_0, \tau) = \frac{1}{\sqrt{2 \pi}} \int \bar{p}^i(\xi, \omega_0, \tau) e^{i \omega_0 \tau} d\tau \)  

(2-b)

Putting (2-b) into (1), then introducing a new variable \( \omega_0 x/2 \), and taking \( \xi = \xi' \), we may have the following equation,

\( (4j \frac{\partial}{\partial \xi'} - \frac{\partial^2}{\partial \omega_0^2} + \frac{3j}{2}) \bar{P}^i(\xi, \omega, x) = 0 \)  

(3)

the solution of Eq.(3) is solved as

\( \bar{P}^i(\xi, \omega, x) = \frac{2e^{-ix}}{j \omega} \int \exp \left( j \frac{\xi'^2}{2 \omega_0} \right) J_0 \left( \frac{2 \xi'}{\omega_0} \right) \bar{q}(\xi', x) d\xi' \)  

(4)

Here \( J_0(y) \) is the zero-order Bessel function. \( \bar{q}(\xi', x) \) is the Fourier transform of a distribution function \( q(\xi', x) \) at the source \( x = 0 \). If the distribution function is taken as \( q(\xi', x) = \delta(\xi')f(x) \), the Fourier transform of the corresponding acoustic field can be obtained as

\( \bar{P}^i(\xi, \omega, x) = P^i(\xi, \omega, x) f'(x) \)  

(5)

in which

\( \bar{P}^i(\xi, \omega) = \frac{2e^{-ix}}{j \omega} \int \exp \left( j \frac{\xi'^2}{2 \omega_0} \right) J_0 \left( \frac{2 \xi'}{\omega_0} \right) \bar{q}(\xi', x) d\xi' \)  

(6)

and \( f'(x) \) is the Fourier transform of a time function \( f(x) \) of source excitation. By putting (6) into (2-b) the ultrasonic field can be finally computed.

For the plane transducer with an acoustic concave lens, under the assumption of long focal length, the distribution function of excitation at source \( q(\xi') \) may be transformed as expression [4],

\( \bar{q}(\xi') = q_0(\xi') e^{-\rho_0 \xi'^2/2} \)  

(7)

where \( q_0(\xi') \) represents the radial distribution function of excitation at the plane source, \( \rho_0 \) is the geometrical focal length of the acoustic lens and \( r_0 = \omega_0^2(1-c^2)/2c \) is the modified Rayleigh distance. \( c_0 \) stands for the sound velocity of the lens. Obviously, when \( \rho_0 = 0 \) or \( q_0(\xi') = \delta(\xi') \) the focusing effect of the acoustic lens disappears.

(1) Gaussian sound field

Suppose that the radial distribution of sound source is expressed as a Gaussian function i.e., \( q_0(\xi') = \exp(-B \xi'^2) \), in which \( B \) is the Gaussian coefficient of the sound source. Putting the expression and (7) into (6) gives

\( \bar{P}^i(\xi, \omega, x) = \frac{A_\rho}{B} e^{-\frac{\xi'^2}{2 B} - \frac{\rho_0 \xi'^2}{2 B} - \frac{\rho_0 \xi'^2}{2 B}} \)  

(8)

where

\[ A_\rho = B^2(2\omega_0^2 + (1 - \omega_0^2)/\omega_0)^{-1} \]

and

\[ B^2 = \frac{(1 - \omega_0^2)/\omega_0}{(1 - \omega_0^2)/\omega_0} \]

(9)

(2) Piston sound field

For piston transducer, as we know, (7) cannot be
integrated analytically, except for the sound fields on the axis as well as at the geometric focal plane. However, it was shown by Wen et al.[5] that any sound field can be expressed as the superposition of a series of Gaussian fields with different parameters. Thus we may let

$$q(x) = \sum_{m} (A_{m} e^{j m \phi}) e^{-[(m \phi) / f_{0}]}$$

for the piston source, where $A_{m}$, $B_{m}$, $f_{m}$, and $s_{m}$ etc. have been determined in the article[5]. If the focused field is taken into account, $q(x)$ may be expressed further as

$$q(x) = \sum_{m} (A_{m} + j B_{m}) e^{-[(m \phi) / f_{0}]} e^{-j m \phi}$$

in which $1/\sigma_{m} B_{m} = \sigma_{m} A_{m}$ and $\sigma_{m}$ may be regarded as an equivalent dimensionless geometric focal length.

Putting (11) into (8), we can obtain

$$D(x) = \sum_{m} (A_{m} + j B_{m}) \sqrt{\frac{m_{0}}{\sigma_{m}}} e^{-\frac{x^2}{\sigma_{m}}}$$

in which $A_{m}$ and $B_{m}$ can be written in the similar expression to $A_{m}$ and $B_{m}$ in (10) respectively.

3. Results and Discussion

Here we choose $f(x)e^{-\frac{x^2}{\sigma_{m} - s_{m}}}$ as a wavefront function of source excitation, which is extensively employed in the ultrasonic medical imaging technique, where $Q$ is the quality factor. Since the bandwidth of a practical transducer is finite and generally is limited to within the frequency range of two octaves, in the present calculation the integral are taken from $x=0.5$ to $x=2$. Fig. 1 shows that the axial distribution of relative intensity for piston transducer with ignoring $s_{m}$, $\sigma_{m}$ is taken 0.2. Curves in the figure are plotted for $Q=1$, $2$, and $4$, respectively. From those, it can be observed that the maximum acoustic intensity or the physical focus for short ultrasonic pulses always occurs at the location ahead of the geometric focus, which is similar to that for the continuous wave and caused by the diffraction [6]. However, the tendency of increase and the shift of the peak toward the geometric focus becomes observable with the decrease of Q value for the short pulse excitation.

The curves of normalized radial distribution of acoustic intensity are plotted in Fig. 2, with the same parameters as those used in Fig. 1. The figure shows that the radial distribution of acoustic intensity becomes narrow with the decrease of $Q$ value, which is in correspondence with the increase of peak on the axis observed in the figure 1. Figures for Gaussian transducer are omitted here owing to the limitation of space, however the similar tendency of results is predicted. The similar analysis was performed for short tone-burst excitation. The similar results are observed. By and large, the results show that the shorter of pulse, the better of focusing effect. Therefore, it seems to be recognized that the short ultrasonic pulses, which are employed extensively at present in ultrasonic medical imaging technique to raise the longitudinal resolution, are capable of raising the lateral resolution too.

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GENERATION OF STRESS PERTURBATION MONOPOLE WAVE IN METAL BY ELECTROMAGNETIC METHOD

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First, the electromagnetics method for generation of stress perturbation waves is reviewed. Then a concrete system is constructed and the generated waves are observed by using a high speed optical displacement measuring device. The results are in good agreement with the theoretical predictions.

The effects of superposition of proper two waves for the improvement of monopoleness is also shown. Finally, a system for estimation of elastic limit inside metal by using the generator is constructed and observed results are discussed.

1. Electromagnetic Transducer and Its Response

The force in metal surface generated by a current in the coil which is placed just on a half infinite metal plane as is shown in Fig.1 is said to be proportion to the square of the current $I$.

Fig.1 Shape of the transducer with coil

Hence the stress waves which are generated by a disk shaped coil can be estimated as has been shown by L.F.Bresse. The response for a link shaped transducer can also be estimated as the difference of responses for two disk transducers with different radiuses.

2. Time Domain Monopole Pump Waves Generation

2.1 Single Coil System

The schematic construction of a single coil transducer pump wave generation system is shown in Fig.2. The stress wave in the metal is essentially proportional to the derivative of the surface force. But the control of the shape of current to the desired form is difficult, hence the generation of completely monopole stress waves by the system can not be expected.

2.2 Double Coil System

To get monopole stress perturbation wave a new double coils system is introduced. The schematic construction of the double coils system driven by two current sources is as is shown in Fig.3.

Fig.2 Construction of the system of stress perturbation wave generation and the corresponding optical measurement system

Fig.3 Construction of double coil stress perturbation wave generation system

The basic principle of monopolization of stress wave is the cancellation of the large dip following to the main peak due to the first coil by the use of the first peak of the second coil.

3. Comparison of Experimental Results with Results of Computer Simulation

3.1 Single Coil System

First, the stress perturbation waves are observed for a single coil system of disk shape $a=10\text{mm}$ of Fig.1, at the distance of $x=20\text{mm}$ of Fig.2 in aluminium. The shape of square of the used electric current is as shown in Fig.4.

Fig.4 Waveform of square of electric current used for stress perturbation generation (surface force on the metal)

An example of observed results and that of computer simulation are shown in Fig.5. They show quite good agreements, although a deep dip is appeared following to the main peak.

3.2 Double Coil System

Now the effect of double coils system is examined by computer
simulation. Here $a_1=10\text{mm}$, $a_2=7.5\text{mm}$, $a_3=5\text{mm}$ of Fig.3 are used and the current of the second coil is adjusted so that the cancellation of the first dip due to the first coil is realized in an optimum way in the measure of least mean square error.

The results are as shown in Fig.6. We can see clearly the effect of the second coil for the increase of monopoleness of the equivalently generated stress waves.

Fig.6 Stress perturbation waves generated by double coil system

**4. Application for Estimation of Elastic Limit of Aluminum**

The stress perturbation wave generation system is applied for the estimation of elastic limit inside of metal. It is combined with the probe wave system of 50MHz which detects the change of ultrasound velocity generated by the pump wave as is shown in Fig.7.

It uses the nonlinear dependence of ultrasound velocity on the stress and the fact that the sign of the velocity change for stress perturbation is changed around the elastic limit of the metal.

The detected phase change is shown in Fig.8. We can see the start of decrease of the phase shift value around the elastic limit which is estimated from the stress strain relation.

**Reference**


UNCONVENTIONAL SUPERCONDUCTIVITY IN UPt$_3$

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Ultrasound methods have been used to study the superconductor UPt$_3$, which is believed to be unconventional. Evidence of this is seen from the low temperature behavior of the ultrasonic attenuation and the existence of at least three superconducting phases. A complete phase diagram for UPt$_3$ has been obtained from high resolution velocity measurements.

1. INTRODUCTION

Superconductivity in "conventional superconductors" can be well explained by the microscopic theory, the BCS theory, based on (a weak coupling) electron-phonon interaction for the pairing mechanism. The superconducting order parameter is a (complex) scalar function, and is usually taken to be isotropic in k-space; this reflects the so called L=0 pairing of the BCS state. If we neglect the Meissner phase for $H < H_c$, there is only one superconducting phase in the H-T plane described by a single order parameter.

Ultrasound attenuation has been an important technique for the study of superconductivity. The superconducting state is characterized by the existence of a gap $\Delta(T)$ in the energy spectrum. For a isotropic energy gap in k-space, quasiparticles thermally excited at finite temperatures obey the Fermi distribution and its interaction with sound wave leads to an exponential temperature dependence for the attenuation $\alpha$ in the superconducting state, with $\alpha/\alpha_0 = 2[\exp(\Delta(T)/k_BT) + 1]^{-1}$.

In the mixed state of a type II superconductor, the normal core of the vortex line contribute to the sound attenuation. The attenuation increases with increasing field exceeding $H_c$ and becomes equal to the normal state value when the field reaches $H_c$. In the region that the applied field H is close to $H_c$, the attenuation depends on $H$ ($H > H_c$ - H) as $\alpha/\alpha_0 = 1 - H/H_c$ in the clean limit (t > $\xi$) and is linearly dependent on the field, $\alpha/\alpha_0 = 1 - H/H_c$ in the dirty limit ($t < \xi$).

The heavy fermion (HF) systems provide a system having exceptional superconducting properties. No adequate theory is yet available to explain the superconductivity in these materials. These are intermetallic compounds having electronic heat capacities $10^2 - 10^3$ times larger than that observed in most metals. The Pauli susceptibility at low-temperatures is also ~100 times larger. Both these indicate a very large effective mass for the conduction electrons.

The resistivity and the sound attenuation both continue to change rapidly with temperature down to very low temperatures, rather than staying constant as in most metals. At low temperatures, the resistivity is of the form $\rho = \rho_0 + AT^2$, where $A$ is large, indicating that electron-electron (spin-spin) scattering is important in the compounds. The large heat capacity jump at the superconducting transition and the large critical field slopes $dH_c/dT$ at $T_c$ in the HF superconductors indicate that it is the heavy electrons that are responsible for the superconductivity.

Neutron diffraction and muon spin resonance experiments on UPt$_3$ [1] show an antiferromagnetic (AFM) ordering at $3^\circ K$ with a very weak ordered moment of 0.07$\mu_B$ which lies in the basal plane. The AFM ordering is thought to play an important role in the subsequent superconducting ordering [2,3].

Ultrasonic measurements have played an important role in the investigation of superconducting properties in UPt$_3$. They have shown that the energy gap has a higher-order symmetry (possibly d-wave) and that there are at least three distinct superconducting phases in UPt$_3$. These observations, along with others, show that UPt$_3$ is a very strong candidate for unconventional superconductivity which can not be adequately accounted for by the simple phonon mediated BCS coupling mechanism.

2. SOUND MEASUREMENTS IN UPt$_3$

In the superconducting state of UPt$_3$, the attenuation does not drop exponentially, but follows a power law behavior, $\alpha(T)$, where $n$ may be 1, 2 or 3, depending on orientation, polarization and sample. This is seen in Fig. 1 (data from Qian et. al. [4]), where a $T^2$ fit has been made to the attenuation. These temperature dependences of the ultrasonic attenuation [4,5,6] in the superconducting state have been interpreted as evidence for a gap with line or point nodes, and hence lead to speculations of unconventional superconductivity.

For polarization in the basal plane (i.e. along the c-axis) the transverse sound attenuation [7] is linear with temperature; for polarization along the a-axis, the attenuation follows a $T^3$ behavior. These temperature dependences suggest that the gap has point nodes along the c-axis, and line nodes in the basal plane, suggesting a d-wave symmetry of the order parameter.

In high quality samples, there is a very sharp velocity change [8] at $T_c$, accompanied by a sharp attenuation peak. The transition width is ~14 mK. Shear wave measurements indicate only a slight change in slope at the transition. A typical temperature sweep of the velocity for two different fields is shown in Fig. 2. At $T_c$, there is a sharp signature that coincides with the susceptibility transition (dashed line) to within 10mK. There is also a change in the slope of the velocity, before and after $T_c$. In addition, at about 60mK below $T_c$, there is a small anomaly (marked by arrows in Fig. 2). A very slow temperature sweep through this anomaly (see inset) reveals a small velocity jump of about 3 ppm. This lower signature (denoted by $T_s'$) is at the position of the second (lower) heat capacity jump, seen earlier [9,10].

The change in the velocity at the superconducting transition...
temperature is related to the heat capacity jump, and the depression of $T_c$ with pressure. The changes in the $T_c$ of $UPt_3$ with hydrostatic pressure and uniaxial stress have been measured; one obtains 24 mK/kbar for a hydrostatic pressure and 26 mK/kbar ($\sim$ 0 mK/kbar) for a uniaxial stress along the c-axis (in the basal plane). Using previously measured values of $\Delta C$ [9], $c_0$, and $c_0$ [11] we obtain $\Delta V/V_0 = 18$ ppm, and $\Delta V/V_0 = 55$ ppm.

The field dependence of the longitudinal sound attenuation, $a$, at a frequency of 300 MHz and at $T = 50$ mK and along the c-axis is plotted in Fig. 3. At about 0.6 $H_c2$ there is an unusual peak in the attenuation which has not been observed in any conventional superconductor. This peak has been referred to as the $H_n$ peak (FL - flux lattice) by Qian et al. [4]; they saw no frequency dependence of this peak between 75 and 300 MHz. The position of the peak has been mapped in the field-temperature plane and was found to move to lower fields with increasing temperature. The shape of the resulting curve led to speculation that the peak may be a signature of a transition between two different superconducting states. The peak is seen only for longitudinal sound measurements, implying that the phase transition couples only to density oscillations.

**Fig. 4.** Velocity in magnetic field sweeps.

Anomalies in the velocity are also observed by sweeping the field. As the field is increased, the velocity decreases monotonically. Field sweeps at $T = 65$ mK and 330 mK are shown in Fig. 4. At the position of the $H_n$ attenuation peak, there is a dip in the velocity of about 3 ppm; at $H_c2$ there is a distinct signature in the velocity as the sample becomes normal. No heat capacity anomaly (latent heat or heat capacity jump) has been seen at the position of the $H_n$ signature. Field sweeps at higher temperatures show a shift of the $H_n$ to lower fields. The $H_n$ signature in the velocity was also observed in temperature sweeps at higher temperatures where $dH_n/dT$ is steepest. None of the anomalies were observed in transverse sound experiments.

The velocity measurements revealed all the known signatures that had previously been identified as possible phase transitions, the $H_n$.

$T_c^n$ (at the position of the lower heat capacity jump [10]) and $H_c6$.

Figure 5 shows the phase diagram obtained by Adenwalla et al. [8] from velocity measurements. The velocity signatures remain sharp and unambiguous to within the width of the transition. This phase diagram (the first to be obtained by a single measuring technique on the same sample) indicated that $UPt_3$ has three distinct superconducting phases. These three phases and the normal phase intersect at a tetracritical point on the $H_n$ curve, for fields both parallel and perpendicular to the c-axis. There is a kink in $H_c2$ at the tetracritical point. Comparable results have been obtained by other groups [12].

**Fig. 5.** Phase diagram of $UPt_3$.

### 3. CONCLUSION

There are many other experiments that indicate $UPt_3$ has several superconducting phases in the H-T plane. The existence of multiple superconducting phases is a clear indication that $UPt_3$ is an unconventional superconductor. However, the exact nature of these phases has not yet been established, but it is believed that the superconducting order is of a higher order (d-wave or maybe even p-wave). Each superconducting phase is described by a separate (nonvanishing) order parameter. The effect of the AFM ordering (at 5K) is to introduce a symmetry breaking field that lifts the degeneracy of the superconducting states, and this leads to the two low field phases II and III. Solutions of the appropriate Ginzburg-Landau equations in a magnetic field give interesting results (3) for the structure of the vortices; it is possible to have doubly quantized vortices and the "H_c6" line may correspond to a transition between two different vortex states.

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### References

MODE CONVERSION TO FOUR BULK WAVES IN PIEZOELECTRIC SOLIDS

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INTRODUCTION

The quasi-static approximation, applied to the reflection of plane acoustic waves in piezoelectric solids, yields a set of four coupled equations for the particle displacements and electrostatic potentials [1]. For any incident wave there are four partial waves in the reflected field whose amplitudes, together with that of an evanescent electric wave on the vacuum side of the boundary, are determined by the mechanical and electrical boundary conditions. It is generally supposed that one or more of these reflected modes is evanescent. In this paper we show that for certain materials and reflection geometries, all four may be evanescent waves. This requires that the longitudinal sheet of the acoustic slowness surface possess negatively curved regions (possible because of piezoelectric stiffening of the elastic constants), thus permitting a reentrant longitudinal mode. We report calculations on Rochelle salt which show that for an incident longitudinal wave, the reflection coefficient for the fourth bulk wave can be as large as unity.

ANALYSIS

We consider the reflection of a plane electroacoustic wave of slowness $S$ at the free surface of a piezoelectric half-space. Reference axes are taken with $x_2$ aligned along the inward directed normal to the surface, $x_1$ along the intersection of the surface with the plane of incidence and $x_3$ normal to the plane of incidence. Phase matching of the incident and reflected waves and evanescent electric wave on the vacuum side of the interface requires that all these waves have $S_3 = 0$ and the same value of $S_1$ [1]. Outgoing waves within the medium are solutions of the secular equation which correspond either to bulk waves ($S_2$ real) with inwardly directed energy flux vectors, or evanescent waves ($S_2$ complex) which fall off exponentially with distance into the medium.

At moderate angles of incidence ($k_3 > v/c$, where $v$ is the velocity of sound in the medium and $c$ is the velocity of light) the trace velocity $1/S$ lies in the sonic regime, and the quasi-static approximation [1,2] implies that the electric field component of the wave is derivable from a scalar potential $\phi$. The secular equation for the equations coupling the particle displacements in the medium and $\phi$ takes the form [3]

$$\Delta = \begin{vmatrix} \Gamma_{11} - \rho & \Gamma_{12} & \Gamma_{13} & \Gamma_{14} \\ \Gamma_{21} & \Gamma_{22} - \rho & \Gamma_{23} & \Gamma_{24} \\ \Gamma_{31} & \Gamma_{32} & \Gamma_{33} - \rho & \Gamma_{34} \\ \Gamma_{41} & \Gamma_{42} & \Gamma_{43} & \Gamma_{44} \end{vmatrix} = 0,$$

where $\Gamma_{ij} = C_{ij}^r S_j S_m \gamma_{ij} = e_{im} S_j S_m \varepsilon = \tau_{jm} S_j S_m$, and $j$ and $m$ take on the values 1 and 2 in the sums. The materials constants involved in the coupling are the elastic modulus tensor at constant electric field $C_{ij}^{r_0}$, the piezoelectric stress constants $e_{1p}$, and the electric permittivity tensor at constant strain $\tau_{jm}$. Equation (1) has eight solutions, four of which pertain to the reflected field. They correspond to electroacoustic waves having, in general, electric field as well as particle displacement components.

For small $|S_1|$, Eq. (1) factorizes approximately. One of these waves is evanescent and predominantly electrical in character, with $|S_2|$ being given by

$$\epsilon = \Gamma_{11} S_1^2 + 2 \Gamma_{12} S_1 S_2 + \Gamma_{13} S_2^2 = 0.$$

The other three solutions are piezoelectrically stiffened bulk acoustic waves. With increasing $|S_1|$, the coupling effect of the $\gamma$ terms on the electrical wave grows larger and it acquires more of a mechanical character. For $|S_1|$ comparable to $1/v$, Eq. (1) implies that all four waves are, in general, of a mixed electroacoustic character.

On the vacuum side of the boundary the same analysis yields an exponentially decaying evanescent wave satisfying $S_1^2 + S_2^2 = 0$.

Taking the amplitude of the incident wave $A^i$ as unity, the amplitudes $A^j$ of the outgoing waves are obtained from the boundary condition equations (vanishing of the traction forces $e_{21}$ and continuity of $E_1$ and $D_1$). The energy reflection coefficient $a_{ij}$ for bulk mode $\nu$ is given by

$$a_{ij} = -P_2^i/P_2^j,$$

where

$$P_2 = 1/2 \omega A^i S \frac{\partial A^j}{\partial n_1} U_i U_j,$$

is the normal component of the energy flux vector [4], $\omega$ is the angular frequency, $n = S/S$ is the wave normal, $U$ is the unit polarization vector for the mode, $A$ is the displacement amplitude and

$$A_{ij} = C_{ij}^{r_0} \frac{e_{1p} e_{1m}}{p^m} n_p n_q,$$

is the piezoelectrically stiffened Christoffel tensor.

The ray of a wave is directed along the outward normal to the slowness surface, and its projection in the plane of incidence is therefore normal to the slowness curve in that plane. Reflected modes are represented by points where the normal to the slowness curve points into the medium [5], and for some of these $S$ may point towards the surface.

In non piezoelectric solids the equation for the slowness surface is of degree 6 in the components of $S$. The outer two transverse (T) sheets of this surface may possess regions where one or both of the principal curvatures is negative but the inner longitudinal (L) sheet is entirely convex [6]. A critical values of $S_1$ is where the ray is parallel to the reflecting surface, and with an infinitesimal change in $S_1$ this limiting wave becomes evanescent. Where the $S_1$ axis passes through a concave region of the slowness section the critical value is reentrant in the sense that with increasing $S_1$ an evanescent wave changes back into a bulk wave [7,8].

In the case of piezoelectric solids the equation of the slowness surface is of degree 8, and when the electromechanical coupling is strong enough, as it is for Rochelle salt, the L sheet may also be negatively curved. Figure 1 shows schematically the real and complex solutions of the slowness equation in a situation that admits reentrant L modes. Real solution are denoted by solid lines, while the dashed lines represent decaying complex solutions. The dotted straight line (1) of
Figure 1. Schematic section of the slowness surface of a piezoelectric solid allowing an L reentrant mode. The slope 1 corresponds to $s_1^2 + s_2^2 = 0$. For small $s_1$ the evanescent electric wave solution given by Eq. (2), is denoted by the dotted straight line (2'). With increasing $s_1$, this solution changes in character from almost purely electrical to electroacoustic. The magnitude of $s_2$ falls below (2), with its imaginary part decreasing and reaching zero at the reentrant point a. For $s_1$ between a and b, there are four bulk wave solutions to the slowness equation with rays pointing into the medium, two of them corresponding to piezoelectrically stiffened longitudinal waves. At point b one of these L waves becomes evanescent and at point c the second one becomes evanescent.

APPLICATION TO ROCHELLE SALT

Figure 2 shows a section of the slowness surface of Rochelle salt for which the orientation of the $S_1$-axis with respect to the crystallographic axes is $(\theta, \phi) = (45^\circ, 60^\circ)$, where $\theta$ and $\phi$ are the spherical polar and zenithal angles respectively, and that of the $S_2$-axis is $(90^\circ, -30^\circ)$. The $S_1$-axis, as can be seen, passes through a shallow concave region of the L sheet giving rise to an L reentrant mode.

In Fig. 3 we have plotted the variation of the energy reflection coefficients $\alpha_{\parallel L}$ ($\nu = L, L', FT, ST$) of the bulk waves reflected from the solid--vacuum interface as a function of the angle of incidence $\theta_0$ or its negative complement $\theta_0 - 180^\circ$, of an incident L wave. Energy balance requires that the reflection coefficients sum to unity, which we have used as a check on the calculations. For some angles of incidence there is only one reflected bulk longitudinal wave while for others there are two.

For small angles of incidence there are only three reflected bulk waves, one L and two T. At point b, which corresponds to a reentrant critical value of $s_1$, the second bulk longitudinal wave $L'$ emerges. Even though the amplitude of this wave is initially finite, the reflection coefficient starts from zero, since $P_2$ is zero for this limiting wave. For similar reasons, the reflection coefficients of the limiting waves at points c, g, and h also vanish.

Figure 2. Incident plane section of the slowness surface of Rochelle salt.

At point c the reflection coefficient of $L'$ reaches unity while the coefficients of the other three reflected waves drop to zero. The incident wave and $L'$ at this point are indistinguishable and the boundary conditions are satisfied by combining these limiting waves in antiphase. A similar phenomenon occurs at points d and f. At point e, which corresponds to an exit critical value of $s_1$, the one L wave becomes evanescent and there are only three bulk waves between e and g. At g the second L wave emerges, and it grows in intensity, while the first L wave decreases in intensity, reaching zero at point h where it is limiting. Beyond h there are only three reflected bulk waves.

Figure 3. Energy reflection coefficients $\alpha_{\parallel L}$ for the free surface of Rochelle salt.

References
PHOTOACOUSTIC STUDY OF RESONANCE ABSORPTION ON CORRUGATED SURFACES

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The resonance absorptions due to excitations of surface plasmons on corrugated silver surfaces or guide-waves on corrugated GaAs surface have been studied by photoacoustic (PA) method. The effect of the groove depth of the corrugated surface on the shape of the absorption curves has been obtained. On the other hand, the differential method based on rigorous electromagnetic theory is used to calculate the parameters of the grooves and the optical absorptivity of the sample. The theoretical results are in good agreement with those of experiment.

INTRODUCTION

Recently, the studies of resonance absorption of incident light on corrugated surfaces, such as periodic metallic and dielectric gratings, have been attracted considerable interest[1-2] due to the discovery of linear and nonlinear optical effects of grating surfaces[3]. The resonance absorptions on corrugated surfaces include two different modes: the resonant excitation of surface plasmons and guide-wave resonances[4]. Both of these have the applications in nonlinear optics and optical bistable systems. Inagaki et al. studied plasma resonances on corrugated silver surfaces by PA method[5]. Since the PA technique directly measures the optical absorption, it is easily to be used to study the resonance absorption in some conditions, such as conical diffractions, in which the conventional optical techniques are difficult, even not possible, to be used.

In this paper, we present the PA investigation of two kinds of corrugated surface resonances. First, the resonance absorption of the incident light due to surface plasmons on corrugated silver surfaces has been observed and the effect of the groove depth on the shapes of the resonance peaks has been discussed. Second, the resonance absorption on corrugated GaAs surface due to the guide-wave resonance is also found. In addition, in order to explain the experimental results, the differential method based on rigorous electromagnetic theory is applied to calculate the parameters of the grooves and the optical absorptivity of the sample. The relationship between the optical absorptivity and the received PA signal has also been obtained. Moreover, the numerical calculations in accordance with the experimental condition have been carried out and theoretical results are in good agreement with those of experiment.

EXPERIMENT

In order to study the resonance absorption of an incident light on the diffraction gratings, several grating samples composed of metals and semiconductors are prepared. All the samples have the same period $d=2000nm$ but different groove depths $2h$. The grating period $d$ is determined by measuring the angle position of the 1-order diffraction light under normal incidence, and the groove depth $2h$ can be evaluated from the 1-order diffraction intensity.

A modulated p-polarized laser beam illuminates the grating sample in a PA cell. The incident light geometry is shown in Fig.1. As the sample stage is rotated, the PA signal from the diffraction grating is a function of the incident angle. Alternatively, the PA signal can also be detected by PZT transducer, which is bonded on the back side of the sample.

At first, as an example of the plasma resonance absorption, we investigate a series of silver film samples with periodic corrugations, which are composed of a series ofphotomask gratings covered by 200nm thick silver films. A 632.8nm He-Ne laser beam with modulation frequency 329Hz impinges upon the sample. In this case, the PA cell with gas-microphone is used to detect the PA signal of the sample. The absorption curves are shown in Fig.2. It can be seen that the resonance absorptions due to excitation of surface plasmons occur at three discrete incident angles 4.5°, 23.5°, and 45.5°. Meanwhile the shapes and the heights of the absorption peaks change dramatically with the groove depths of the corrugated surfaces, but the positions of the peaks are essentially independent on the depths of the grooves. Therefore, the shape of the absorption peak depends on the groove depth and an optimal groove depth exists for which the peak gets itself maximum.

Second, for the guide-wave resonance, GaAs samples with corrugated surface produced by ion-etching are studied. A 514.5nm argon ion laser beam with middle power and modulation frequency 114KHz is used to illuminate the GaAs sample, and the PA signal is detected by a PZT transducer. For different azimuth angles (see Fig.1), the absorption curves as the function of incident angle are measured as shown in Fig.3. For 0°, due to guide-wave resonance, the absorption curve exhibits three peaks at 14.5°, 29.5° and 48.5°. In Fig.3 it can also be seen that the resonance angles shift to higher incident angles as the azimuth angle varied from 0° to 90°, but no resonance absorption peak is found as 90° because the electric-field component of the incident light is parallel to the groove grating.

THEORY

When a monochromatic light beam modulated periodically illuminates a semiconductor wafer, three kinds of wave are excited in the sample. They are: an electron-hole wave, a thermal wave and an acoustic wave. Compared with the thermal diffusion length of the sample, the groove depth of the corrugated surface is much smaller. Therefore, it is reasonable to assume
Fig. 2: Surface plasmon resonance absorption in corrugated silver surface for p-polarized light as a function of incident angle. The groove depths are (1) \( h=10 \) nm, (2) \( h=24 \) nm, (3) \( h=65 \) nm. Dot line is theoretical result.

Fig. 3: Guide-wave resonance absorption on corrugated GaAs surface for p-polarized light. The azimuth angles are (1) 0°, (2) 30°, (3) 45°, (4) 90°.

A thermal smooth surface of the sample. For simplicity, we also assume that the sample is opaque and the optical absorption only occurs at the surface of the sample. With an one-dimensional theoretical model, the temperature \( T(x) \) and the photogenerated carrier density \( N(x) \) as the functions of distance \( x \) beneath the sample surface are given by [6]:

\[
T(x) = \frac{Q_0}{Q} \exp(-qx)
\]

\[
N(x) = \frac{P_0}{D} \exp(-px)
\]

where the thermal wave vector \( q \) and the electron-hole wave vector \( p \) are defined as follows:

\[
q^2 = \frac{-j\omega\varepsilon}{\kappa}
\]

\[
p^2 = \frac{(1+j\omega\tau)}{D}
\]

\( Q \) is the heat source and \( P_0 \) is the electron-hole source term. \( D \) is the ambipolar diffusion coefficient.

As pointed out in section II, the important grating parameters are the grating period \( d \) and the groove depth \( 2h \). The grating period can be determined by measuring the angle position of the 1-order diffraction light under normal incidence. On the other hand, Pockrand et al. [8] have given a formula to determine groove depth \( 2h \) from the intensity of the 1-order diffraction light, but it is only valid for sinusoidally corrugated surfaces of perfect conductor. In order to evaluate the groove depth from the 1-order diffraction light intensity, an analysis based on rigorous electromagnetic theory is developed for the groove shapes with sinusoidal, rectangular or any other periodic profile, as well as the grating may be composed of metal, dielectric and semiconductor.

The calculated absorptivity as a function the incident angle for a sinusoidally corrugated silver surface with a period 2000 nm is also shown in Fig. 2. From Fig. 2 we can see that the theoretical result (dot line) is in good agreement with that of experiment, including the correct prediction of the small peak at 4.6°.

Furthermore, the calculations have also been done for rectangularly corrugated GaAs surface. The calculated angles of the resonance peaks also agree with the experimental results. But the shapes of the absorption curves appear some discrepancies, it is due to the lack of the sufficient information about the dielectric function of doped GaAs wafer.

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REFERENCES

PHASE FOURIER TRANSFORM INFRARED PHOTOACOUSTIC SPECTROSCOPY

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INTRODUCTION

Fourier Transform Infrared Photoacoustic Spectroscopy (FTIR-PAS) is now widely used to analyse any kind of materials, even those for which conventional FTIR is not able to obtain spectra.

The depth profiling capability of PAS is also possible with FTIR-PAS since the frequency modulation may be changed by using different mirror speeds in such a way that the sample thermal diffusion length is changed. Applications of FTIR-PAS to opaque or diffuse materials and to powdered or liquid samples make this technique very useful and permit to obtain new information on these materials and samples.

But PAS presents some inconvenient such as signal saturation and band distortion due to sample dilation, optical diffusion or various thermal or mechanical heterogeneities. The physical nature of one material (powdered or bulk) may cause artefacts like the change of the relative band intensities or a different background signal between bands, thus making quantitative PAS analysis difficult. Moreover, with the limited frequency domain available in FTIR, the sample depth profiling is limited to a few μm in most solids. Near surface analysis (a few 100 Å) is thus impossible.

Using our Phase Fourier Transform Infrared Photoacoustic Spectroscopy (PHAS/FTIR-PAS) we overcome all the photoacoustic inconvenient and artefacts mentioned before and improve greatly the depth profiling of the sample.

PHASE FOURIER TRANSFORM PHOTOACOUSTIC SPECTROSCOPY

Photoacoustic signal is a complex one and the real part is in phase with the modulation whereas the imaginary part is quadrature. By using a complex Fourier transform we are able to obtain, from the interferogram, the two components of the photoacoustic signal. Using first a thermally thick carbon sample as a reference to correct the interferogram from the phase apparatus and to normalise the spectrum, we apply the complex transform to the sample interferogram with an iterative procedure to obtain the actual apparatus phase for the analysed sample.

The two components of the photoacoustic signal are now used to obtain:
- the δμ product, δ being the absorption coefficient and μ the thermal diffusion length;
- the absorption coming from a surface layer less than a tenth of the thermal diffusion length and
- the absorption coming from the remaining thermal diffusion length or the remaining sample thickness if the sample is thermally thin.

With PHAS/FTIR-PAS the δμ product is free of the effects due to the sample dilation, the resolution is improved and the band intensities are more accurate since photoacoustic saturation occurs at higher values of absorbance. Moreover our technique improves greatly the depth profiling of the sample since we probe the thermal diffusion length itself in addition to modifying it by frequency modulation variations.

EXAMPLES

We have chosen three examples to illustrate the advantages of the PHAS/FTIR-PAS.

Figure 1 shows the in-phase and in-quadrature components of the photoacoustic spectrum of polyethylene between 1000 and 4000 cm⁻¹ and figure 2 shows the photoacoustic (amplitude) and δμ spectra, the latter being obtained from the data of Fig. 1. Clearly the δμ spectrum shows less photoacoustic saturation and a better resolution than those of the amplitude spectrum without the characteristic photoacoustic background signal. Moreover, in Fig. 1 a large absorption band between 3000 and 3300 cm⁻¹ appears only in the in-quadrature spectrum. This indicates that this absorption occurs on the sample surface only. This absorption is due to the water adsorbed on the sample surface corresponding, for this non-polar polymer, to a thickness of about a few hundred angstroms.

Figure 1: In-phase and in-quadrature photoacoustic spectra of polyethylene.

Figure 2: Amplitude and δμ spectra of polyethylene.
Figure 3 is the photoacoustic and $\beta\mu$ spectra of asbestos, a very heterogeneous sample. Sample dilation is very important since interstitial gas will expand more than the solid. The photoacoustic spectrum demonstrates a large background signal and important photoacoustic saturation. Our $\beta\mu$ spectrum does not possess these photoacoustic artifacts. Figure 4 shows the spectra obtained from two asbestos samples, one being more compacted than the other. $\beta\mu$ spectra are quite the same, whereas the two photoacoustic spectra have different background signals and relative band intensities. This clearly demonstrates the effect of sample texture on the photoacoustic signal and also the necessity to have a method that corrects the photoacoustic signal for the photoacoustic artifacts.

Figure 4: Amplitude and $\beta\mu$ spectra of asbestos samples; COMP: compacted sample.

The last example is a thermally thin sample: a thin layer (less than 1 μm) of plasma-polymerized hexamethyldisiloxane, a polar polymer, coated on a CaF$_2$ substrate. In this case, sorption results in a monolayer of water (1.5 Å thick) and a penetration of water into the bulk of the polymer (a volume signal). Studying the water band absorbance by PHAS-FTIR-PAS we are able to measure surface and volume water absorption. Figure 5 shows this absorption as a function of the relative humidity. Our results are in very good agreement with already published results about water sorption on polymers. It is important to note that, in this case, the sample is thermally thin and that the surface signal arises from a thin layer only a few angstroms thick.

Figure 5: Surface and volume water absorption for a thin layer of PP-IMDSO on CaF$_2$.

CONCLUSION

Phase Fourier Transform Infrared Photoacoustic Spectroscopy is a very powerful technique to analyze complex materials and to obtain precise informations on surface analysis. The common photoacoustic artifacts are suppressed and the photoacoustic saturation occurs at higher absorption level, resulting in better quantitative measurements. The PHAS/FTIR-PAS technique improves greatly the depth profiling capability of the photoacoustic method, probing the thermal diffusion length itself, even with thermally thin samples.
OPERATION OF A PHOTOACOUSTIC CELL WITH MICROPHONE AND PREAMPLIFIER AT VERY LOW PRESSURES AND CRYO TEMPERATURES

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INTRODUCTION

In order to measure collisional deactivation rates in binary mixtures of methane, laser-pumped at 3.3 μm, we have constructed a photoacoustic cell which contains a commercial condenser microphone and preamplifier which are exposed to pressures as low as 0.001 Torr and cryogenic temperatures. This cell has produced measurements of vibrational relaxation lifetimes[1] which agree well with other published values. However, we have observed unusual effects while operating at the extremes of low pressure or low temperature. We wish to discuss preliminary explanations for these unusual microphone signals.

The ability of a microphone to function at very low gas pressure is important in that it extends the range where photoacoustic detection is valid. Photoacoustics [2], itself, is the acoustic detection of radiation absorption. Photoacoustic detectors have higher sensitivity than the more commonly used photomultipliers etc. when the total absorption is low. A low absorption case of considerable current importance is the detection of air pollutants since they are invariably in such low concentrations that the total absorption is low.

EXPERIMENTAL

The photoacoustic cell used was a 2.5 cm diameter cylinder, 20 cm long, constructed from copper with Brewster’s angle windows of ZnSe seated on indium seals for good thermal contact. A stainless steel 1/2 inch microphone (Bruel & Kjaer 4134) suspended flush at the cell’s inner surface was surrounded by a cooling jacket containing the microphone/preamp arm. The entire cell assembly was contained in a 9 inch glass cross, evacuated to prevent window condensation and provide thermal and acoustic isolation.

A pulsed Nd/YAG (Spectra Physics) laser fed a dye laser whose output was heterodyned with a portion of the original beam in a LiNbO3 crystal to produce a 3.3 μm Gaussian-profile beam. The output beam was wavelength tunable and brought to coincidence with the photoacoustic cell axis by a 4 mirror system aligned by a HeNe laser.

DISCUSSION

Low Pressure Operation

The microphone response when the cell contained 2% methane in argon is shown in fig.1. The solid curve shows the microphone signal as a function of time when the total gas pressure is 2.5 Torr. A typical cylindrical pressure wave is evident (the part of the signal due to vibrational relaxation is small and can be neglected in fig.1).

The dotted curve in fig.1 is for a pressure of 0.016 Torr. (For both curves, 5000 microphone responses were averaged.) At this low pressure, the acoustic wave approximation may no longer be a valid description [3]. A qualitative change is now evident in the cell's microphone response.

Since the laser excitation pulse is very short, both curves in fig.1 represent the impulse response of the system made up by the photoacoustic cell and microphone. (Only 2.5 ms of the response is shown, which actually extends over approx. 15 ms. The laser pulse is repeated every 100 ms and most of our recordings were for a duration of 10 ms.) The unusual low pressure effect is that the impulse response changes markedly as the pressure is lowered.

At low pressures, the gas flow between the microphone membrane and the backplate is reduced causing a resonance in the high frequency response[4]. At times we observed highly underdamped oscillations, presumably as a consequence of this. However, the impulse response at 0.016 Torr in fig.1 shows an anomalous response, i.e., one which is slower. This anomaly is currently under investigation.

Besides the change in response as the pressure varies, we also observed an, as yet, unexplained repeating pulse train response after the initial 15 ms. This pulse train response consisted, for example at 2.5 Torr, of a repeat of the response shown by the dark curve in fig.1 but with less distinct features, i.e., the peaks became blurred. The total pulse train repeated itself three times over 45 ms. It does not correspond to the acoustic transit time of any part of the cell or connecting tubeulation.

Low Temperature Operation

At temperatures near 77 K we observed uncontrolled oscillations. We have described these in detail to the microphone manufacturer who was unable to explain them.
Conclusions

It appears that a condenser microphone yields useful acoustic information even at 0.02 Torr. At that pressure, the main signal peak is still proportional to the optical energy absorbed by the gas in the photoacoustic cell when the laser beam passes through it. However, many features remain unexplained when operating a microphone at very low pressures or temperatures. Chief among these is an anomalous change in impulse response as the total pressure is lowered.

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Fig. 1 Microphone Signal When Laser Pulse Absorbed In CH₄
POLY (VINYLIDENE FLUORIDE) TRANSUDERS FOR THE DETECTION OF LASER INDUCED TRANSIENTS

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INTRODUCTION

Radiation induced thermal and acoustic transients are used extensively for the nondestructive evaluation of samples. Photothermal (PT) and photoacoustic (PA) testing, as these methods are called when employing light and lasers in particular for the excitation of the transients, have a large number of advantages compared to conventional techniques. PT analysis has the advantage of noncontact generation of a well defined instantaneous heat source at the surface or in the volume of the sample. This facilitates the solution of the heat diffusion equation and hence the derivation of the thermal characteristics of the sample from experimental data. In a similar fashion laser induced ultrasound allows the convenient noncontact generation of high frequency ultrasonic waves in a sample. In contrast to conventional ultrasonic transducers PT generation of ultrasound is always noncontact and free of transducer ringing. Coupling of the sound pulse from the transducer into the sample, subject to acoustic impedance matching, etc., when employing a transducer is no problem with laser generated ultrasound since the sound is generated in the sample! These features of PT and PA analysis might be only a convenience when studying bulk samples. Noncontact analysis of samples in a vacuum system or hostile environment or the continuous monitoring of a manufacturing process are just two examples. For thin films samples, however, the photothermal generation of the probe pulse is frequently the only viable alternative since the wavelength of the probing pulse should be substantially less than the sample thickness and no other transducers of sufficient bandwidth are available. With a sufficiently short and powerful laser pulse, absorbed in a thin layer, ultrahigh frequencies are easily generated making PT and PA analysis of thin films relatively straightforward. Ultrashort Pulse Excimer lasers are ideal for this type of application. Due to the short pulse duration of these lasers and the short absorption length of Excimer laser radiation in most materials source broadening of the heat source or the ultrasonic source is negligible. For thermal analysis an almost ideal instantaneous point, line or planar source of heat can be realized greatly simplifying the analysis of data and allowing a spatial resolution of tens of nanometers. For ultrasonic analysis frequencies up to tens of GHz can be readily generated allowing analysis of samples with a resolution of the order of hundreds of nanometers.

One application that takes advantages of ultrashort thermal pulses generated by the absorption of light from a picosecond laser pulse is the mapping of charge or polarization profiles in electrets with extremely high depth resolution.

Charge and polarization profiles of thin dielectric polymer films are of considerable scientific and technological interest. Photoconductors in copiers and laser printers are one area of application, electret materials in radiation dosimeters, microphones and loudspeakers another area and temperature or pressure sensors still another. In all these applications knowledge of the charge distribution or the profile of the polarization is crucial. Therefore, a large number of methods has been developed to determine the charge, polarization and electric field distributions in dielectric films. Laser induced thermal and pressure pulses are commonly utilized for high resolution depth profiling. At the same detection bandwidth thermal pulses have, due to their shorter wavelength, the advantage of higher spatial resolution as compared to acoustic pulses.

Pyro- and piezoelectric sensors play an important role in the detection of radiation induced thermal and acoustic phenomena. Sensitivity and rise time of pyroelectric calorimeters and piezoelectric ultrasound transducers are, for example, determined by the charge or polarization profile of the electret material. For utmost sensitivity, for example, a large bulk polarization would be desirable. For fast rise time of the transducer the polarization at the surface of the electret is the critical parameter. For a sensor with high bandwidth and good fidelity it is important that the polarization is homogeneous throughout the poled material. In addition to the polarization other material properties are important for a sensor: For an ultrasonic transducer acoustic characteristics such as acoustic impedance and ultrasonic damping or electrical characteristics such as dielectric permittivity or electrical impedance are of concern. For a pyroelectric temperature sensor the corresponding thermal properties such as specific heat of the sensor material or its thermal diffusivity become the material parameters of interest.

Poly (Vinylidene Fluoride) (PVDF), a ferroelectric polymer, lends itself to the design and implementation of ultrafast pyro- and piezoelectric transducers for the detection of laser induced thermal and acoustic transients due to excellent thermal and acoustic properties. For high performance transducers, however, poling of these films has been a major problem. A depolarized layer with a thickness of the order of 100 nm has been limiting the rise time of PVDF transducers to 100 ns in thermal detectors and 10 ns for acoustic sensors. To take full advantage of polymer electrets for high bandwidth transducers improved poling efficiency of the electrets is crucial. Recent progress in the poling of PVDF films allowed improvements by several orders of magnitude.
EXPERIMENT

Samples were prepared using high purity PVDF granules from Pennwalt as starting material. The PVDF granulate is dissolved at slightly elevated temperature in dimethyl formamide (DMF). The solvent has to be distilled to obtain a conductivity of less than 0.3 \( \mu \)Mhos. Films are then spin coated at approximately 200 rpm on pyrex substrates. The coated substrates are subsequently then vacuum dried at 80°C overnight. Free standing films with thickness of the order of 10 nm are then obtained by peeling the dried film from the substrate. These films can then be processed directly into sensor elements or undergo corona charging. For the purpose of this study only one silver electrode with a thickness of 120 nm was deposited onto the film according to a procedure described elsewhere in more detail. The electrode serves as light absorber for the laser pulse and to collect the induced electrical charges. Since PVDF is practically transparent for excimer laser light the heat source was generated directly at the metal-PVDF interface by exciting the laser through the PVDF. A customized Lambdaphysik PSL4000 laser system that provides a pulse of approximately 20 ps duration and up to 20 mJ energy at a wavelength of 248 nm was utilized to generate a short thermal pulse via the absorption of the light in the highly reflecting metal electrode. To ensure that a one dimensional model of the signal generation process describes the experiment a 5 mm diameter spot in the center of the foil was excited by the laser. The induced charge was amplified by a custom high impedance preamplifier with a bandwidth of 600 MHz and one single transient recorded via a Tektronix 7104 Oscilloscope with a digital camera system. The 1 GHz bandwidth of the T209 amplifier plug-in together with the 7B10/15 time bases and the resolution of the camera system result in a time resolution of 10 ps and a rise time of 350 ps for the data acquisition system. The overall bandwidth is therefore currently limited by the preamplifier to 600 MHz.

Pressure pulses for the analysis of the polarization profile were generated by absorbing the laser pulse in a tungsten rod. The foil under study was coupled to the rod by a thin water film. The piezoelectric signal was recorded by the same electronics that were used for pyroelectric analysis of the polarization profile. The rise time of the observed photothermal and photoacoustic pulses were limited by the electronics to 350 ps. Using a slightly modified set up surface acoustic waves were generated and detected with a rise time of the order of 5 ns, mainly limited by the geometry of the experimental set-up.

A pyroelectric calorimeter\(^4\) is characterized by a polarization that is homogeneous throughout the electret film, the pyroelectric thermometer is unpoled and exhibits only a surface charge that resides at the interface. As mentioned above the performance of electrets as thermal detectors is directly related to their performance as ultrasonic detectors. Using the above information on the polarization profile of the homogeneously poled electret one would predict a rise time of the order of 1 ps in the ultrasonic mode; for the surface charge transducer a 1 ps pulse width is expected. Due to the limited detection bandwidth experimental verification of these data is, however, not yet feasible. The foil with the surface charge has another attractive feature for ultrasonic applications: since the signal generating layer and the remainder of the film are identical in all mechanical respects the interface does not reflect any acoustic energy due to an ideal match of acoustic impedances.

Applications of transducers using PVDF of other ferroelectric polymers as the ferroelectric sensor are discussed with examples in calorimetry thermometry and ultrasonic NDE of materials.

REFERENCES


APPLICATION OF THE INVERSE BORN APPROXIMATION TO EXPERIMENTAL THERMAL WAVE IMAGES

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INTRODUCTION

Thermal wave images have sometimes been criticized because their resolution is often less than that of high-quality ultrasonic images. This reduction of the resolution arises from the lateral diffusion of heat in the thermal waves scattered from subsurface features of the sample. In this paper we illustrate a technique for removing the resultant blurring of pulsed thermal wave images of planar defects through the reconstruction of the shape of the scatterer. The algorithm developed produces a convolution of a "heat spread" function with a function which describes the shape of the scatterer. The Fourier transform of the surface temperature contrast can then be expressed as a simple product of the Fourier transforms of the heat spread function and the shape function of the scatterer. Use of the algorithm consists of performing a two-dimensional spatial Fast Fourier Transform (FFT) on the experimental image of the scatterer, dividing the resulting transform by the transform of the known heat spread function, and finally doing an inverse FFT to obtain the shape of the scatterer. A practical difficulty is caused by image noise which arises from the finite accuracy of the imaging process, which is typically digitized at an 8-bit level, and from high-frequency detector noise. This difficulty can be overcome by applying a variety of low-pass filters to the Fourier transform of the experimental image.

THE INVERSION ALGORITHM

Our inversion algorithm is based on a Green's function model in which the defect is described as a planar scatterer at some given depth below the surface of the sample. The heat source is assumed to be uniformly distributed over the sample's surface, and to occur as a short pulse at time, \( t = 0 \). Our model describes the undisturbed propagation of the pulse from the surface of the sample down to the defect, and the first-order reflection of the pulse at the surface of the defect. It is thus equivalent to the first-order Born approximation. With this model it is possible to express the theoretical temperature contrast, which is defined as the surface temperature resulting from the scattering with the time-dependent background following the incident pulse having been removed, as

\[
\Delta T(x,y,t) = -\frac{A}{2\pi} \left( \frac{1}{4\pi \alpha t} \right)^{1/2} \frac{\partial}{\partial t} \int \frac{dx'dy'}{\sqrt{|x-x'|^2 + (y-y')^2 + k^2/4}}
\]

\[
\times \left\{ \frac{(x-x')^2 + (y-y')^2 + k^2/4}{4\alpha t} \right\}^{1/2} f(x', y')
\]

Here \( \alpha \) is the thermal diffusivity of the material, which is defined as the ratio of the thermal conductivity to the heat capacity per unit volume of the material, and \( f(x,y) \) is a function which describes the planar shape of the defect. The constant \( A \) is determined by the power per unit area in the original heat pulse. This formula takes account of both the dispersion of the incident pulse in its propagation down to the defect, and its scattering and subsequent propagation back to the surface. A detailed derivation of the equation, as well as similar equations for different boundary conditions at the defect will be published elsewhere. The inversion algorithm depends on the fact that this equation is a convolution of two functions of the form

\[
\Delta T(x,y,t) = \int dx'dy' g(x - x', y - y', t) f(x', y')
\]

(2)

The algorithm for inverting the scattering is very simple. The experimental thermal wave image first is converted to a differential image by subtracting the time-dependent background leaving only the experimental equivalent of \( \Delta T(x,y,t) \), i.e., the thermal wave "echo" image. This is followed by the performance of a two-dimensional FFT on both the experimental echo image and the function \( g(x,y,t) \), with the value of the time \( t \) chosen to correspond to the delay time at which the experimental image was taken. The Fourier transform of the \( f(x,y) \) is then given in principle by the well-known convolution theorem as the ratio of the transform of \( \Delta T(x,y,t) \) to the transform of \( g(x,y,t) \). All that remains to acquire an image of the shape function \( f(x,y) \) is the performance of an inverse FFT on the ratio of the transforms. However, although this procedure functions flawlessly for simulated images, it fails when applied to experimental images. This failure is due to the fact that the experimental image contains a considerable amount of high frequency spatial noise resulting both from detector noise, and from the digitization (usually 8-bit) of the image signal. The function \( g(x,y,t) \), on the other hand, is quite smooth and contains almost no high frequencies. When the transform of the experimental image is divided by the transform of \( g(x,y,t) \), the high frequency noise it contains is exaggerated to the point where it completely obscures the image of the defect. Fortunately this problem is easily overcome. All that is required is that the transform of the experimental image be filtered with a low-pass filter (Gaussian, Wiener, etc.), prior to or during the division by the transform of \( g(x,y,t) \). The application of the inverse FFT then produces a reconstructed image of the defect. The only adverse effect of the filtering is a slight rounding of the corners of the defect in the image. The complete inversion calculation takes about one minute on a Macintosh IIix.

RESULTS

In this section we will display images of fabricated defects, both before and after reconstruction. The images displayed in this section were obtained with the use of our Box-car Video Thermal-Wave Imager at Wayne State University. These are pulse-echo images, i.e., those for which the thermal waves are launched by pulsing a set of high-power flash lamps, and in which the returning thermal wave image is recorded by an infra-red video camera. A background level averaged over a

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Fig. 1 A pulse-echo thermal wave image of two 12.5 mm holes milled into the back surface of a piece of opaque plastic. The wall between the two holes is 1 mm thick.
Fig. 2 The result of the application of the inversion algorithm to the image shown in Fig. 1. Note the resolution of the 1 mm wall between the two holes.

The featureless region of each image has been subtracted from the raw data to yield an image corresponding to the differential contrast temperature given by Eq. 1. The first image, shown in Fig. 1, is of two 12.5 mm holes milled into the rear surface of a thick piece of opaque plastic. The centers of the holes are 13 mm apart, so the wall between them is 1 mm thick. The bottoms of the holes are 1.7 mm from the front surface of the sample. This image shows the pulse-echo thermal wave image of the bottoms of the two holes as viewed from the front (blank) side of the sample. The blurring of the image due to the diffusion of heat parallel to the plane of the sample has rendered the thin wall between the two holes rather indistinct. In Fig. 2 we show the result of the inversion of the image in Fig. 1. It should be noted that the 1 mm separation of the two holes is now clearly resolved.

Figure 3 shows a pulse-echo thermal wave image of a sample designed to provide a simulation of corrosion in an aircraft lap joint. This sample consists of two 2.4 mm thick plates of aircraft aluminum alloy with two rectangular air-gap areas between the two plates. The air gaps were fabricated by machining away the metal within the rectangles on the surface of one of the plates before it was joined to the other plate. The air-gap thicknesses were chosen to simulate 50%, and 25% corrosion of the metal. Note that the deeper the air-gap is beneath the surface, i.e. the less the "corrosion", the more indistinct its thermal wave image becomes. Nonetheless, when the inversion algorithm is applied to the image, both "defects" appear with quite sharp edges and corners, and both with the same (correct) size. This inverted image is shown in Fig. 4. In this image we have used a display technique which causes the defects to appear to be raised from the surface to emphasize the similarity of their sizes and shapes.

Fig. 4 The result of applying the inversion algorithm to the thermal wave image shown in Fig. 3. Note the clearly defined edges and corners of the defects, and the fact that they both appear to be the same (correct) size.

CONCLUSIONS

We have demonstrated the possibility of performing inverse scattering calculations on experimental thermal wave images to improve their resolution and to obtain quantitative shape and size information about the scatterer. The technique is simple and fast. The method is currently being extended to include the effects of multiple scattering and of any possible anisotropy of the thermal properties of the material. The latter capability will permit its use for sizing of defects in highly anisotropic composites like graphite-epoxy layups.

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PHOTOACOUSTIC FREQUENCY-DOMAIN DEPTH PROFILING OF CONTINUOUSLY INHOMOGENEOUS CONDENSED PHASES. APPLICATIONS OF THE INVERSE PROBLEM.

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I. Theoretical

The depth-profiling thermal-wave problem for continuously thermally varying condensed phases (photo thermal PT) detection has been addressed. Unlike earlier treatments, [1-3] well-posed, direct, simple and convenient expressions for the frequency dependence of the PT signal are obtained based on the classical mechanical concept of the Hamilton-Jacobi thermal harmonic oscillator (THO) [4]. It can be shown that the solution to the thermal-wave inhomogeneous problem

$$\frac{d}{dx} \left[ k(x) \frac{d}{dx} T(x) \right] - i \alpha(x) c(x) \frac{d}{dx} T(x) = -\frac{1}{2} \eta \beta I_0 e^{Rt} \times S(0)$$

with the well-known quantities $T$ (ac temperature); $\eta$ (nonradiative efficiency); $\beta$ (absorption coefficient); $I_0$ (incident irradiance) in a semi-infinite condensed phase medium with continuously variable thermal conductivity, $k(x)$; density, $\rho(x)$; and specific heat $c(x)$, gives rise to a THO field. In the limit of photo thermally saturated media and assuming samples with known surface properties, and a specific convenient thermal diffusivity profile

$$\alpha(x) = \alpha_0 \left[ 1 + \Delta e^{-2\theta} \right]^{2}; \Delta = \frac{\alpha_0}{\alpha_n} - 1,$$  

the normalized PT signal amplitude, $|M(\omega)|$, and phase, $\Delta \phi(\omega)$, becomes

$$|M(\omega)| e^{i \Delta \phi(\omega)} = 1 + \frac{1}{4} R^{1/2}(\omega) \exp \left[ \frac{(1+i)\sqrt{\alpha_0}}{2 \sqrt{2} q} \ln \left( \frac{\alpha_n}{\alpha_0} \right) \right]$$

where as a reference we assumed a homogeneous material with constant thermal diffusivity $\alpha_0$. In the bulk of the inhomogeneous sample the thermal diffusivity $\alpha_0(x)$ assumes an arbitrary value $\alpha(x) = \alpha_0$. $R(\omega)$ is the ratio of thermal diffusivities:

$$R(\omega) = \frac{\alpha_0(\omega)}{\alpha_0}$$

The essence and versatility of the technique lie in its treatment of differential subsurface layers corresponding to a measurement at modulation frequency $f_1 = \omega_0/2\pi$ with local $\alpha_0(x)$ values determined by redefining (upgrading) the values for $q$ and $R(\omega)$ found at a neighboring frequency, $f_j$, differing from $f_1$ by $\Delta f = f_1 - f_j << f_1$. Thus, local variations of $\alpha_0(x)$ are accounted for, as manifested by normalized signal differences between neighboring frequencies, and the original single exponential decay for $\alpha_0(x)$, Eq. (2), becomes irrelevant in the determination of the global depth profile: this is determined as an inverse problem from successive values of the entire set of the experimental data dependences on $\alpha_0$ and can be any function of $x$, numerical or in closed form. A test $\alpha_0(x)$ profile input to Eq. (3) according to Eq. (2) was used. Excellent reconstruction of the same profile from the simulated amplitude ratio and phase difference data with a homogeneous reference sample with $\alpha_n = \alpha_0$, was thus obtained.

In the case where the surface thermophysical properties of the sample are not known, or have been altered through some process (e.g. thin-film deposition, laser processing, surface damage), the bulk value of the thermal diffusivity, $\alpha_n$, becomes the reference parameter of interest. Under these conditions, the normalized PT signal becomes

$$|M(\omega)| e^{i \Delta \phi(\omega)} = \frac{1}{R(\omega)} \left[ 1 - \frac{1}{4} R^{1/2}(\omega) \right] \times \exp \left[ \frac{(1+i)\sqrt{\alpha_0}}{2 \sqrt{2} q} \ln \left( \frac{\alpha_n}{\alpha_0} \right) \right]$$

where a monotonically increasing thermal diffusivity depth profile ($\alpha_n > \alpha_0$) of the form

$$\alpha_n(x) = \alpha_0 \left[ 1 + \Delta e^{-2\theta} \right]^{2}; \Delta = 1 - \left( \frac{\alpha_n}{\alpha_0} \right)^2$$

is assumed, corresponding to the presence of a surface layer with poorer thermal transport properties than the bulk. The reference sample is one with thermophysical properties similar to the bulk of the sample under investigation. Although Eq. (6) is valid for a monotonically increasing thermal diffusivity depth profile, arbitrary $\alpha_n(x)$ profiles can be handled by the method outlined above.

II. Experimental

Photoacoustic gas-cell experiments were conducted covering both cases of condensed phase samples with a) known surface diffusivity $\alpha_n$ and b) known bulk diffusivity, $\alpha_n$. In the former case (a), the method was used to obtain quantitative depth profiles of thermal diffusivity decreases extending 20-30 $\mu m$ below the surface of the liquid crystal 8CB [5].

The experimental apparatus was described elsewhere [5]. An 8CB sample without an applied transverse magnetic field was the reference of the same sample in the presence of a magnetic field B. A 3.39 $\mu m$ He-Ne laser was chosen, a radiation strongly absorbed by 8CB, so that the sample was approximately photoacoustically saturated. Figure 1 shows the reconstructed thermal diffusivity depth profile $\alpha_n(x)$, with [6] $\alpha_n(0) = 1 \times 10^{-7} m^2/s$.

Fig. 1  Reconstructed thermal diffusivity profile of a nematic 8CB liquid crystal sample with an applied transverse magnetic field B, using $\Delta f = 1Hz$; $B = 1.65kG$ at $T = 37.5^\circC$.

The results indicate that the application of the transverse magnetic field across the 8CB nematic liquid crystal generates a region in the bulk with lower thermal transport properties than near the surface, the "bulk" behavior commencing at and below 26-28 $\mu m$ from the surface. Although this is the first time that such a variation is identified, yet it is consistent with the lowering of the nematic values of $k(T)$ in the presence of a transverse magnetic field in 8CB, when the conductivity is measured photoacoustically [7] at low $f$. The depth profile of Fig. 1 is con-
considered to be clear evidence of an interface-induced alignment of the long molecular axes perpendicular to the surface as a long-range phenomenon (tens of microns). This novel effect was shown to be completely removed at 43°C (isotropic phase). The present results may be significant for future studies of the long-range positional order of liquid crystal molecules due to interface-molecule interactions and of the effect of surfaces in phase transition behavior.

In case (b), several samples of low carbon steel (0.042% C) and stainless steel (Type 301) were laser-processed at various beam spot diameters and scanning speeds. The processing was performed with a CO₂ laser with 1 kW CW output at speeds varying from 1 to 2 inches/sec, and at beam spot diameters from 0.3 to 0.9 mm. In order to achieve wide enough laser processed area, some overlapping was allowed between laser scans.

Samples were machined to the shape of the sample holder inside the cell, from the original laser-processed specimens. Reference samples were machined from unprocessed original material. For each sample surface, the PA signal amplitude and phase were recorded in the modulation frequency range between 10 Hz and 500 Hz. The frequency increment was 2 Hz. The lower frequency limit was dictated by the requirement that the air column above the sample be thermally thick. This avoids any signal behavior complications due to propagation of heat through the walls of the cell [8]. The upper limit was determined by the signal-to-noise ratio and complications due to cavity resonances in the cell-microphone assembly. For each surface, data was averaged over three to five experimental runs to reduce random noise.

In order to compare PA measurement results with the microstructure of the surface layers of laser-processed samples, specimen cross-sections were characterized by optical microscopy. After the PA experiments the samples were cut normal to the laser processed direction for metallographic examinations. Hardened zones were determined by normal etching techniques and by microhardness testing, using a Vickers indenter with 100g load. The hardened zones obtained with all tests were of lenticular shape that is found when processed with a defocused low order mode laser beam. From the metallographic photographs (optical) and microhardness profile, the depths of hardened zones were measured.

![Graph](image)

**Fig. 2** Reconstructed thermal diffusivity profile of laser processed (a) low carbon steel sample and (b) stainless steel sample.

Processing parameters for the carbon steel sample are 1 inch/sec scan speed and 0.9 mm spot size, and for the stainless steel sample laser scan speed was 1 inch/sec and the spot size 0.3 mm. The laser processed surface has been compared with the unprocessed reference, of the same material. The data clearly show the difference in the signal due to laser processing.

The algorithm described was then used to obtain thermal diffusivity profiles from the low carbon steel data. The reconstructed profile for the laser processed low carbon steel sample is shown in Fig.2(b). This profile shows that the thermal diffusivity starts decreasing from a depth of about 130μm which is in agreement with the independently determined hardness profile. Bulk thermal diffusivity αₐ (=17.3 x 10⁻⁶ m²/s) for this sample was calculated from the thermal conductivity (κ), specific heat (c) and density(p) values for this type of carbon steel [9]. Diffusivity has monotonically dropped to 9.7 x 10⁻⁶ m²/s at a depth of 50 μm (corresponding to the highest frequency) from the surface starting from the bulk value at about 130 μm. A similar trend in thermal diffusivity decreasing with increase in hardness in carbon steel has been reported earlier [10].

Figure 2(b) shows the profile for the laser processed stainless steel sample. Bulk diffusivity (α₀) calculated from κ, c and p values obtained from Ref.[9] for stainless steel is 4.1 x 10⁻⁶ m²/s. Thermal diffusivity starts to decrease around 200 μm from the surface which is, again, in good agreement with the hardness profile. The rate of decrease in diffusivity is found to have changed (become steeper) at depths shallower than 55 μm which may correspond to the equiaxed crystal layer seen in the optical metallograph of this specimen.

Figure 3 shows the thermal diffusivity profile of a stainless steel sample laser processed with a lower power density where the scan speed was the same (1 inch/sec) but the processing laser beam spot size has increased to 0.7 mm. This reconstruction shows a thinner damaged layer than the profile in Fig.2(b), as expected.

In summary, in case (b) we have established a reliable technique to obtain thermal diffusivity depth profiles of a layer extending from the surface into the bulk from the frequency domain surface data. We have also observed that the thermal wave technique as represented by photoacoustic detection is very sensitive to surface hardening due to laser processing in steels. Previous work [10] on steel surface hardening (achieved by increasing the carbon concentration at high temperatures) has shown that hardening reduces the thermal diffusivity of steel. The profiles obtained are in agreement with the earlier work and with observations made by destructive methods such as the optical microscope cross-sectional imaging and the microhardness test, and have better resolution compared to microhardness test.

![Graph](image)

**Fig. 3** Reconstructed thermal diffusivity profile of stainless steel sample laser processed at a lower power density compared to the sample in Fig. 2(b).

**REFERENCES**

PULSE-ECHO THERMAL WAVE TOMOGRAPHY

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INTRODUCTION

Pulse-echo ultrasonic techniques are widely used for the purpose of nondestructive evaluation (NDE) of materials. Knowledge of the elastic wave velocities in a material permits time-of-flight measurements to be converted into measurements of the subsurface depths of various subsurface defects which are strong elastic wave scatterers. In recent years, the authors and others have utilized an analogous pulse-echo technique to study thermal wave scattering from subsurface defects [1-4]. In this technique, the waves are launched as the result of energy deposited at the sample surface from a suitably pulsed heat source, and the resulting time evolution of the surface temperature is monitored by means of an infrared (IR) camera. Descriptions of the detailed methods for processing the data stream from the camera in real time, so as to effect box-car or lock-in image processing, have been presented elsewhere [1-4].

Vavilov, together with some of the authors, [5] extended the box-car method, further exploiting the time-dependent information by making tomographic thermal wave images. The tomographic processing method operated on a sequence of thermal wave images which were recorded and stored on disk memory, following the flash-heating pulse. It involved a numerical search for the temporal peak of the thermal wave echo, pixel-by-pixel. Since this process involves extensive post-processing and a search through a large number of stored images, it is memory-intensive and slow. As will be demonstrated below, we have since developed a real-time thermal wave tomographic method which accomplishes the same result, but does so with real-time techniques which avoid the storage of a large number of images.

TECHNIQUE

The one-dimensional time dependent solution to the heat equation in the time domain following flash heating of the surface at time \( t = 0 \) has the form,

\[
T(x) = \frac{C}{\sqrt{4\pi \alpha t}} \cdot e^{-\frac{x^2}{4\alpha t}},
\]

where \( C \) is a constant which is related to the amount of heat deposited on the sample surface (assumed to be located at \( x = 0 \)). At \( t = 0 \), this pulse is infinitely narrow and located at the surface \( x = 0 \). As the pulse propagates into the sample, because of the high frequency dependence of the velocity of its thermal wave Fourier components, the pulse undergoes extensive broadening. Because of the strong attenuation suffered by thermal waves, the amplitude of the pulse also decreases very rapidly as it propagates and broadens. When such a thermal pulse encounters a subsurface scattering region, e.g. a boundary across which the thermal impedance changes, a reflected pulse (the thermal wave "echo") is returned to the surface. The surface temperature decrease in the presence of such a subsurface scatterer is slower than that predicted by Eq. (1). A typical cooling curve corresponding to some point on the surface can be thought of as being comprised of two components. One component is simply that which would occur if the material underneath that region contained no thermal wave scatterers. The other component results from the thermal wave reflected from any such scatterer. The experimental set up is shown in Fig. 1. This system is controlled by a computer (Sun workstation, Macintosh II, or 486 clone) which is used to synchronize the components of the system, download the software for the real-time processor.

![Fig. 1. Block diagram of the Box-car Imaging System.](image)

and to display and post-process the resulting thermal wave images. The heat source consists of a battery of up to eight 6.4 kJ, 2 ms pulse duration, xenon flash lamps, which are fired simultaneously by a signal from the computer. The surface temperature of the sample is monitored as a function of time by an infrared video camera (Inframetrics Model 600), operating in the 8μm to 12μm region of the IR. The stream of analog data coming from the camera is digitized in the real-time processor (DataCube, Percecepts), which is programmed to carry out the box-car gating in real time. In this mode of operation of the system, an image, or series of images, is acquired at fixed times after the occurrence of the flash. This type of processing is analogous to that of a boxcar averager, such as is used for pulse-echo ultrasonics, but has been extended to the full video image, in real-time. The gate is set so as to capture the returning thermal wave echoes from any subsurface defects when their contrast with the background is near its maximum value.

As an example of thermal pulse propagation and reflection, in Fig. 2 we show experimentally obtained heating and cooling curves corresponding to several subsurface flat bottom holes at different depths in a polymer test specimen. It can be seen from Fig. 2 that there are varying cooling rates over different regions, as expected on the basis of the description above. In Fig. 2, the curve with the fastest cooling rate corresponds to a surface region beneath which there are no subsurface scatterers. In our tomographic method we make use of this curve as a reference and subtract this reference from the time dependence of every pixel of the image as it is generated. Examples of such subtracted curves (expanded in scale) are shown in Fig. 3 for the six selected regions corresponding to the cooling curves of Fig. 2. Because these curves have had the background curve subtracted, they correspond to plots of the reflected thermal waves from the subsurface holes. Although the peaks are broadened because of the diffusive nature of the heat propagation, the peaks of the subtracted curves (see Fig. 3) are seen to vary both in height and in temporal position, with the peaks corresponding to the deeper holes being lower and later. As can be seen in Fig. 4, this relationship is quadratic in depth, as expected for thermal diffusion.

![Fig. 2. Experimental heating and cooling curves for a plastic sample with six flat-bottomed holes milled into the back surface. The lowest curve is a reference taken over a region with no subsurface defects. The other six curves are from areas directly over the six holes.](image)
Fig. 3. Plots of the returning thermal wave echoes from the six holes in the sample of Fig. 2. These are obtained by subtracting the reference curve in Fig. 2 from the remaining six curves.

Fig. 4. Plot of the peak times of the thermal wave echoes from Fig. 3 as a function of the squares of the depths from which the echoes originate.

The curves shown in Fig. 3 are representative of six thermal wave echoes, the peak of each returning to the surface at a different time and with a different amplitude. The essence of our real-time thermal wave tomographic method, is to carry out a pixel-by-pixel determination of the arrival times and amplitudes of the echoes. This is accomplished instrumentally by first smoothing the curves to avoid spurious peaks caused by noise in the signal, then applying a numerical algorithm to determine the peak times (and amplitudes). The resulting 512 x 480 collection of peak times is stored as a pseudo-image in a buffer for viewing. In addition, a second buffer contains a "peak contrast" image. This second image displays the amplitude of each individual pixel of the echo as it appeared at the time when it arrived at the surface. Currently, we have the ability to produce both tomographic and peak contrast images simultaneously with a single heating flash, and with post-heating sampling rates up to 10 Hz.

RESULTS

It is instructive to compare the results of thermal wave pulse-echo imaging, with conventional ultrasonic pulse-echo imaging. Such a comparison is shown in Figs. 4 and 5. Figure 4 shows a thermal wave tomogram of a graphite-epoxy impact damage sample. This particular impact caused some damage directly under the impact site, with concentric rings of damage around that damage. The first ring of damage is shallower than the damage directly under the impact site, with the remaining rings getting progressively deeper in a more-or-less conical pattern. The deepest ring is decidedly asymmetrical. This thermal wave tomogram can be compared to the ultrasonic pulse-echo image shown in Fig. 5. This image was obtained by scanning a printed color image of the ultrasonic image, followed by a modification of the resulting gray-scale to obtain the closest correspondence to the depth scale used in the thermal wave tomogram of Fig. 4. This particular ultrasonic image happens to have lower resolution than the thermal wave image, but the general features of the two are clearly the same.

Fig. 5. Ultrasonic pulse-echo image of the same graphite-epoxy impact damage sample whose thermal wave tomogram is shown in Fig. 4.

ACKNOWLEDGEMENTS

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ANALYSIS OF SURFACE ACOUSTIC WAVE MEASUREMENT ACCURACY USING INTERFEROMETRIC DETECTION

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I INTRODUCTION

Laser generation and detection of ultrasound has developed rapidly in the last few years since it offers a non-contacting technique for non-destructive testing. One application that has not been fully exploited is the accurate measurement of surface acoustic wave (SAW) velocity. Accurate measurement of SAW velocity has important applications such as determination, of thin film thickness, variation of adhesion between film and substrate and measurement of residual stress.

This technique does not use the coupling fluid which is necessary in scanning acoustic microscopy, this should, at least in principle, make it possible to achieve higher measurement accuracies of velocity and material attenuation than is possible with the acoustic microscope which has achieved 1 part in $10^4$ [1]. Clearly also the techniques also has the merit of being non-contacting which makes it suitable for real time in process monitoring.

II ANALYSIS OF SAW WAVE MEASUREMENT ACCURACY

Continuous wave (CW) excitation has the disadvantage that the acoustic amplitude is lower than with pulse excitation, but, nevertheless, it more readily lends itself to single frequency measurements. We confine our discussion to CW generation. In particular we consider detection systems using dual-beam interferometers. The basic principles of single-beam and dual-beam heterodyne interferometry are described in [2][3][6][7].

1. Direct dual-beam heterodyne interferometry

The basic principle of this system has been described in [6], we will make our discussion more definite by determining the signal expected at the detector. (Fig1)

\[
i_s = 2\pi A^2 k_b \sin\left[2\left(\omega_1 - \omega_2 \right) + \omega_2 x_1 \right] + k_b x_1 x_2
\]

Signal to noise ratio is given by [4]:

\[
i_s = 2\pi A^2 \sqrt{\frac{\eta}{2}} \sin\left[2\left(\omega_1 - \omega_2 \right) + \omega_2 x_2 \right]
\]

\[
\frac{\eta}{2} \cos k_b x_2 + \cos k_b x_2 (x_1 + \Delta x)
\]

We can thus see that the signal is proportional to the difference in displacements between the two probe positions i.e. the total variation in path length. Expanding $i_s$, we obtain the amplitude $i_s = k_b \Delta x$. $A = A_1 \sqrt{2-2\cos k_b x_2}$.

\[
A = A_1 \sqrt{2-2\cos k_b x_2} = 2A_1 \left| \sin \frac{k_b x_2}{2} \right|, \text{ where } A_1 = k_b \Delta x
\]

giving $i_s = k_b \Delta x \left| \sin \frac{k_b x_2}{2} \right| \sqrt{4 \pi f_0 \eta + \frac{2}{B_h \gamma}}$

If we take $\eta = 0.8$, $\delta f_0 = 0.1 \text{Hz}$, $f_0 = 5 \text{MHz}$, $B_h = 100 \text{kHz}$, $\lambda = 6280 \AA$, $\beta = 0.3$, we have: $i_s = 1.2 \times 10^3 \text{A}$.

Under the condition that the signal to noise ratio (SNR) is large, the relative measurement error caused for a shot noise limited system is:

\[
\Delta = \frac{\Delta x}{\Delta x} \left| \frac{d_i}{d_i} \right| \Delta x \left| \frac{d_i}{d_i} \right| \Delta x \left| \frac{d_i}{d_i} \right| \Delta x
\]

This equation implies that that the error is related to the noise and the gradient of the rate of change of signal with displacement at that particular point. If points A (peak), B (bottom) and C (half amplitude point) on figure 2 are examined we can see that around position A, $d_i/d_i$ is very large and this is necessary to consider the second derivative to obtain an estimate of the measurement error. The region of maximum signal is thus the worst region for accurate measurement. At point B, $d_i/d_i$ is a maximum even though the signal to noise ratio is poor at that point.

This gives $\Lambda = \frac{\Delta x}{\Delta x} \left| \frac{d_i}{d_i} \right| \Delta x$

In the region between B and C, we assume $\left( \frac{\Delta x}{\Delta x} \right) = 1$. $\Delta = \frac{\Delta x}{\Delta x} \left| \frac{d_i}{d_i} \right| \Delta x$

\[
\Delta \sim \frac{2B_h \gamma}{\eta \eta_t} \frac{1}{k_b \delta f_0^2 \eta_t (x_1 + \Delta x)}
\]

$n = 1.2$ for attenuation and non-attenuation respectively.

\[
\text{f}_1(\Delta x) = \sqrt{1 + \exp(-2\mu \Delta x) - 2 \exp(-\mu \Delta x) \cos k_b \Delta x}
\]

\[
\text{f}_2(\Delta x) = \exp(-\mu \Delta x) \cos k_b \Delta x + \exp(-\mu \Delta x) k_b \Delta x
\]

\[
\frac{\eta}{2} \cos \frac{k_b \Delta x}{2} \text{ without attenuation}
\]

The computed results of $f_1(\Delta x)$ show that at $\Delta x = 1.1$, $f_1(\Delta x)$ has a maximum, $f_2(\Delta x)$ has no such a point since it increases monotonically with $\Delta x$.

The computed results also show that for $\mu = 0.0035/\lambda_0$, the attenuation of $\text{Si}_N\text{H}_4$ at $1000 \text{MHz}$, $f_1(\Delta x) = 660$, at $\Delta x = 3000 \text{nm}$, where $f_2(\Delta x) = 1800$. This means that attenuation becomes $f_1(\Delta x)$ smaller at $\Delta x = 100 \text{nm}$. $f_1(\Delta x)$ and $f_2(\Delta x)$ are nearly the same, so that attenuation will have little effect on the measurement accuracy. Another effect of attenuation is shown in figures 3 and 4.

\[
\text{FIGURE 3. Plot of } f_1(\Delta x)
\]

\[
\text{FIGURE 4. Plot of } f_2(\Delta x)
\]
Figure 3 is enlarged plot of $f_1(\Delta x)$ for $\mu = 0.0035$. Figure 4 is enlarged plot of $g_2(\Delta x)$ for $\mu = 0$ (no attenuation). In fig. 3, we obtain two peaks instead of just one in fig. 4 for every wavelength of beam displacement. This means attenuation has changed the gradient of the change of signal $dv/d(\Delta x)$ at point B (fig. 2) from maximum into minimum. As a result, the best measurement point is shifted from B to towards C (fig.2).

The relative error in SAW velocity measurement is:

$$\frac{dv}{v} = \frac{\Delta x}{\Delta x} = \frac{2\Delta v}{\Delta x} = \frac{\Delta \omega}{\omega} = \frac{\Delta f}{f}$$

Take $B=100$Hz, $\lambda = 6680$A, $\eta=0.8$, $\beta=0.3$, $I_0=5$cmW, we get:

$$\frac{dv}{v} = \frac{\Delta x}{\Delta x} = \frac{2\Delta v}{\Delta x} = \frac{\Delta \omega}{\omega} = \frac{\Delta f}{f}$$

If $\Delta x = 0.1A$, $\mu = 0.0035$, $\Delta x_{min} = 5000$A, $dv = 4.0 \times 10^{-6}$

This implies it is possible to measure SAW velocity at the accuracy of better than $10^{-5}$ for an acoustic displacement of 0.1A.

2. Indirect dual-beam heterodyne interferometry

This is shown in fig. 5, where each probe beam interferes with the reference beam which is reflected from mirror. We thus have two heterodyne interferometers in parallel. The system detects the phase of the surface wave.

**FIGURE 5. Configuration of Indirect dual-beam heterodyne interferometry**

The received signals are:

$$I_{s1} = 2V_0 \delta^2 \Delta \phi \sin(2\omega t + \alpha + \theta + \phi_1)$$
$$I_{s2} = 2V_0 \delta^2 \Delta \phi \sin(2\omega t + \alpha + \theta + \phi_2)$$

We can then compare their phase to obtain the phase difference.

The signal to noise ratio of each interferometer is:

$$\frac{I_{s1}}{I_{n1}} = k_1 \frac{2V_0 \delta^2 \Delta \phi}{\sqrt{2} \Delta \phi}$$
$$\frac{I_{s2}}{I_{n2}} = k_2 \frac{2V_0 \delta^2 \Delta \phi}{\sqrt{2} \Delta \phi}$$

Once again, we take $\Delta x = 6680$A, $\delta = 0.1A$, $\alpha = 30\lambda_{max} = 3000$A, $\eta = 0.8$, $I_0 = 5$cmW, we obtain $\Delta x_{min} = 900$, so $\Delta x_{min}^2 = 1$.

The measurement phase error (fig.6):

$$tg \phi_1 = \frac{I_{s1}}{I_{n1}} \cdot tg \phi_2 = \frac{I_{s2}}{I_{n2}}$$

The total error $\Delta \phi = \sqrt{\frac{(2 \Delta \phi_1)^2}{(2 \Delta \phi_2)^2}}$.

Increasing $\Delta x$. If there is attenuation, we have $I_0 = I_0 \exp(-\mu \Delta x)$.

Under the condition of large SNR:

$$\Delta x = -\frac{\ln(1 + \exp(-2\mu \Delta x))}{\sqrt{2} \Delta x}$$

We define a function $f_2(\Delta x)$:

$$f_2(\Delta x) = \frac{\Delta x}{\sqrt{2} \Delta x}$$

The relative error is:

$$\Delta x = f_2(\Delta x) \exp(-2\mu \Delta x)$$

Its plot is shown in fig. 7.

**FIGURE 6.**

**FIGURE 7.** Plot of $f_2(\Delta x)$

$\Delta x$ has a maximum at $\mu = 1$. For example, if $\mu = 0.0035$, $\Delta x_{min}^2 = 5000$, and the SNR here is $900 \exp(-1) = 300$, so the result is still valid. We also note that the minimum error decreases with attenuation as expected.

The relative error for SAW velocity measurement is:

$$\frac{dv}{v} = \frac{\Delta x}{\Delta x} = \frac{2\Delta v}{\Delta x} = \frac{\Delta \omega}{\omega} = \frac{\Delta f}{f}$$

Assuming $\mu = 0.0035$, $\Delta x = 0.1A$, $\Delta x_{min}^2 = 5000$, gives $\Delta x = 1.7 \times 10^{-6}$.

**III CONCLUSION**

We have considered SAW detection using dual-beam interferometry. The system sensitivity is close to $10^4$ A/W/Hz. In the presence of attenuation there exists a particular beam separation at which best accuracy can be obtained. The relative error for velocity measurement is:

$$\frac{dv}{v} = \frac{f_1(\Delta x)}{f_1(\Delta x)}$$

for direct interferometry, and

$$\frac{dv}{v} = \frac{f_2(\Delta x)}{2\Delta x}$$

for indirect interferometry.

This implies for both direct and indirect interferometry, relative accuracy $10^{-5}$ can be achieved for 0.1 A SAW amplitude. The indirect system has the added advantage that it allows independent measurements of attenuation and the probe beams can be kept fixed, which should be generally favorable for SAW measurements.

**REFERENCES**


THERMOELASTIC GENERATION OF ULTRASOUND BY PULSE LASER IN SOLIDS

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Ultrasonic pulses are generated when a pulse laser irradiates a condensed matter sample. Many authors have studied this general problem, both theoretically and experimentally, since 1963 (1-3). But the arguments of these papers based on the "Slow" nature of thermal conduction. Recently McDonald (4) studied this problem with hyperbolic thermal conduction equation. But he only discussed the temperature field in an infinite medium and assessed the normal displacement with an assumption $C_T=C_1$ which is not so accurate according to $C_T^{\max}=C_1^{\max}/3$. Therefore we use the hyperbolic thermal diffusion equation and Navier–Stokes equation to find the temperature field and normal displacement field generated by a laser pulse in a semi–infinite medium. The theoretical results of normal displacements on the epicenter and on the surface in the sample and experimental results of the ultrasonic group and phase velocity in several samples are presented.

1. TEMPERATURE FIELD

When a $\delta(t)$ laser pulse is incident on a sample surface, a pulse heat source is formed and a rapid thermal expansion can be produced to generate detectable ultrasonic waves in the sample. The thermal power density of the heat source is given by

$$ g(\vec{r},t) = \delta(\vec{r})\delta(t)/2\pi r $$

(1)

Instead of the traditional parabolic thermal diffusion equation [1,3], the hyperbolic thermal diffusion equation is used to find the temperature distribution in a semi–infinite homogeneous and isotropic medium in order to avoid getting the impossible physical result that the thermal wave velocity would be larger than sound velocity in the medium when the laser pulse width is short enough. Then the temperature increment $T(\vec{r},t)$ is determined by

$$ T(\vec{r},t) = \delta(\vec{r})/\kappa \left[\frac{1}{C_T^3}(\nabla T/\alpha T) + \frac{1}{\alpha}(\alpha T/\alpha t)\right] \nabla^2 T $$

(2)

and

$$ T(\vec{r},0) = \delta(\vec{r})/\kappa \left[\frac{1}{C_T^3}(\nabla T/\alpha T) + \frac{1}{\alpha}(\alpha T/\alpha t)\right] \nabla^2 T = 0 $$

(2a)

where $C_T$ is thermal wave velocity, $\alpha$ and $\kappa$ are thermal diffusivity and conductivity of the medium, respectively.

Using Laplace–Henkel transform, the transform solution $\tilde{T}(s,p,z)$ is obtained from Eq.(2):

$$ \tilde{T}(s,p,z) = \exp(-K_{p}z) \exp(-K_{s}r) $$

(3)

where

$$ K_p = \frac{s}{C_T}, K_s = \frac{p}{\alpha}, K_r = \frac{p}{4\pi \kappa} $$

(3a)

$$ T(x,t) = \int_{0}^{\infty} \tilde{T}(s,p) \exp(-st) \frac{dp}{p}, \quad \tilde{T}(s,p) = \int_{0}^{\infty} T(r,t) \exp(-pr) \frac{dr}{r}, \quad \tilde{T}(s,p) = \int_{0}^{\infty} T(r,t) \exp(-pr) \frac{dr}{r} $$

(3b)

Therefore the temperature field is obtained by inverse transform:

$$ T(r,t) = [H(t-z/C_T) \exp(-C_T^{2}t/2\alpha) / 4\pi k] \ast \int_{0}^{\infty} \tilde{T}(s,p) \exp(-st) \frac{dp}{p}, \quad \tilde{T}(s,p) = \int_{0}^{\infty} T(r,t) \exp(-pr) \frac{dr}{r} $$

(4)

where $H(t)$ is Heaviside step function.

2. NORMAL DISPLACEMENT FIELD

The Navier–Stokes equation for the displacement $\vec{U}$ (Ur,Uz) at point $(r,\theta, z)$, when the temperature field is restricted on the sample surface, is

$$ \mu \vec{V} \cdot \vec{U} + (\lambda + \mu) \vec{V} \cdot \vec{U} - \rho \frac{d^2 \vec{U}}{dt^2} = 0 $$

(5)

where $\lambda$ and $\mu$ are Lamé constants, and $\rho$ is density.

Introducing the scalar and vector potential $\phi$ and $\vec{A}$, such that

$$ \vec{U} = \vec{\nabla} \phi + \vec{A}, \quad \vec{V} = \vec{\nabla} \times \vec{A}, \quad \vec{V} \cdot \vec{A} = 0 $$

(6)

Let $A_x = -\phi/\alpha t$, we get

$$ \frac{\partial \phi}{(1/C_T^2 \phi)} / \alpha t = 0, \quad \frac{\partial \phi}{\alpha t} / \alpha t = 0 $$

(7)

and the initial conditions and boundary conditions are

$$ \phi(r,0) = \phi(r,0) / \alpha t = 0, \quad \phi(r,0) = \phi(r,0) / \alpha t = 0 $$

(8)

$$ \tau_{rr}(r,0) = \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} + \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} + \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} = 0, \quad (1/C_T^2 \phi) / \alpha t = 0 $$

(9)

$$ \tau_{rr}(r,0) = \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} + \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} + \frac{\partial r}{\partial r} \frac{\partial \phi}{\partial r} = 0 $$

(10)

where $C_T^2 = \lambda + 2\mu / \rho, \quad C_T^2 = \mu / \rho$, and $\alpha_T$ is the thermal expansion coefficient.

Applying Laplace–Henkel transform to Eqs. (7)-(10), we obtain

$$ U_z = -[(\lambda + 2\mu) / 4\pi \kappa \mu] \exp(-s/C_T) \exp(-p/C_T) $$

(11)

where $k_1 = p^2 - (s/C_T)^2$, $i = 1$ or 2.

Using Cagniard method and $K_x = \int_{0}^{\infty} T(s)/C_T^2 s^2 + p^2 dp$, the epicenter displacement $Uz(z=0)$ is obtained:

$$ U_z = B' \cdot (C_T^2 t^2 + \delta_T) H(t-z/C_T) * \left[ \frac{2t^2 + \delta_T}{2t^2 + \delta_T} \right]^{1/2} $$

(12)

where $\delta_T = (C_T^2 t^2 + \delta_T) / (2t^2 + \delta_T)$.
\[ 0, \quad (0 \leq t < C_2 / C_1), \]
\[ \left( 2B \int_0^{y_2} \frac{1}{[(C_2 / C_1)^2 - y_1^2]} \left( C_2 y_1^2 - (C_2 / C_1)^2 y_1^2 \left( C_2 / C_1 \right)^2 \right) dy \right) \]
\[ \left( (1 - 2y_1^2) + 16y_1^2 (1 - y_1^2) y_1^2 \right) \]
\[ \left( (1 - 2y_1^2)^2 + 4y_1^2 (1 - y_1^2) y_1^2 \left( C_2 / C_1 \right)^2 (y_1^2 - 1) \right) \]
\[ \left( (1 - 2y_1^2)^2 + 16y_1^2 (1 - y_1^2) y_1^2 \left( C_2 / C_1 \right)^2 \right) \]
\[ \left( (C_2 / C_1)^2 - y_1^2 \right), \]
\[ \left( C_2 / C_1 < t < 1 \right), \]
\[ U_2(t) = \left( C_2 y_1^2 - (C_2 / C_1)^2 y_1^2 \right) \]
\[ \left( (1 - 2y_1^2)^2 + 16y_1^2 (1 - y_1^2) y_1^2 \left( C_2 / C_1 \right)^2 \right) \]
\[ \left( (C_2 / C_1)^2 - y_1^2 \right), \]
\[ \left( C_2 / C_1 < t \right). \]

The normal displacement waveforms \( U_2(\tau = 0) \) and \( U_2(\tau = 0) \) are shown in Fig.1(a) and (b), respectively. They are in good agreement with the experimental results Fig.1(c) (3) and (d) (6).

3. MEASUREMENT OF ULTRASONIC VELOCITY

A laser-ultrasound experimental system is shown in Fig.2. The ultrasound pulse generated by a N\(_2\) pulse laser (337 nm, 8 ns and 6-8 mJ) and the train of echo pulses in the sample are detected by a PVDF film or PZT coupled with the sample. The ultrasonic pulses are recorded by a digital oscilloscope Gould 4072 and sent to a microcomputer to process. The correlation method and spectrum analysis technique are used to obtain the group and phase velocity of several samples. The experimental values of group velocity of AI and Steel are

\[ C_{\text{Al}} = 6417.74 \text{ km/s}, \quad C_{\text{steel}} = 5.941 \text{ km/s}, \]

respectively. The phase velocity of AI and Steel are shown in Table 1. They are in good agreement with the theoretical values:

\[ C_{\text{Al}} = 6417.74 \text{ km/s} \quad \text{and} \quad C_{\text{steel}} = 5.941 \text{ km/s}. \]

**Fig.2 Laser-ultrasound experimental system.**

<table>
<thead>
<tr>
<th>Table 1: Experimental results of phase velocity</th>
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</thead>
<tbody>
<tr>
<td>Sample</td>
</tr>
<tr>
<td>AI</td>
</tr>
<tr>
<td>Sample</td>
</tr>
<tr>
<td>Steel</td>
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</table>

Therefore the laser-ultrasound method is a very powerful new technique for nondestructive testing.

- This work is Supported by National Natural Science foundation of China.

**Reference**


**Fig.1** Theoretical results (a) \( U_2(\tau = 0) \), (b)\( U_2(\tau = 0) \) and experimental results (c) \( U_2(\tau = 0) \), (d) \( U_2(\tau = 0) \).
PHOTOACOUSTIC NONDESTRUCTIVE EVALUATION OF STRESS DISTRIBUTION IN SAMPLE

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Nondestructive evaluation of stress distribution in materials is a very important in both characterization of the quality of mechanical parts and constructional components, and their inspection in operation. As we know the optoelastic, acoustoelastic[1], Raman Spectroscopy[2], X-ray[3] and neutron diffraction method[4] can be used to measure the distribution of stress. Recently the SEAM images about vickers indentation[5] show that the thermoelastic method is also a very powerful to evaluate the stress field in sample. In this paper, one-dimension theory of thermoelastic detection is discussed and some experimental results are presented.

1. THERMOELASTIC COUPLING HEAT CONDUCTION EQUATION

The entire field of continuum mechanics is governed by three basic physical conservative laws, namely, the laws of conservation of mass, momenton and energy. In solid mechanics, an additional constitutive law relating the stress to strain and temperature is required to fully describe the system. For a Hookean material, these laws may be written as

\[ \int \rho v dv = \int \rho_0 v dv, \]

where \( \rho \) and \( \rho_0 \), \( v \) and \( v_0 \) are stressed and unstressed density and volume of material, respectively.

\[ \sigma_{ij} + \rho_0 \dot{v}_i = \rho_0 \ddot{v}_i, \]

where \( \sigma_{ij} \) is the stress tensor, \( v_i \) the body force per unit mass, and \( u_i \) is the displacement vector.

\[ (\dot{U} - R) = \sigma_{ij} \dot{e}_j - q_{ij}, \]

where \( \dot{U} \) is the rate of change of internal energy per unit mass, \( R \) the rate of heat produced per unit mass by internal heat source, \( \dot{e}_j \) the strain rate, and \( q_{ij} \) is the heat flux through the surface of the body whose outward direction normal is \( \eta_j \).

\[ \sigma_{ij} = 2\mu\varepsilon_{ij} - \lambda\varepsilon - \beta T_{ij} \delta_{ij}, \]

where \( \varepsilon \) and \( \mu \) are Lame constants, expansion, \( \varepsilon = a_\varepsilon \) is the first strain invariant, \( \beta = (3\lambda + 2\mu)/2\mu a_\varepsilon \) is the coefficient of linear thermal expansion, \( T_1 = T - T_0, T_0 \) is the room temperature, and \( \delta_{ij} \) is the Kronecker delta.

Now we introduce the Helmholtz free energy \( \Phi \):

\[ \Phi = U - TS, \]

where \( S \) is the specific entropy. Using Eq(3) and the second law of thermodynamics

\[ q_{ij} = \rho_0 (R - T\dot{S}), \]

it may be concluded that

\[ \sigma_{ij} = \rho_0 \Phi / \varepsilon_{ij}, \]

\( S = -\Phi / \varepsilon T. \)

because \( c_\varepsilon \) and \( T \) are independent variables for a elastic material.

Substituting Eq.(7) into Eq.(5) we obtain

\[ q_{ij} = T(\sigma_{ij} / \varepsilon T) \dot{e}_j - \rho_0 C_{ij} \dot{T} + \rho_0 R. \]

where \( C_{ij} = -T^2 \Phi / \varepsilon T^2 \) is the specific heat at constant deformation.

Using Eq.(4) and following expression

\[ e_{ij} = [\sigma_{ij} - \lambda \delta_{ij} / (2\mu + 3\lambda)] / 2\mu + \varepsilon T, \]

where \( s = s_0 \) is the first stress invariant, and omitting higher order terms, we get

\[ \sigma_{ij} = T\left[ \frac{1}{2\mu} [\sigma_{ij} + \frac{2\lambda}{\mu + 3\lambda} \delta_{ij} + \frac{\lambda^2}{(2\mu + 3\lambda)^2} \delta_{ij}] \right] \]

\[ + \frac{2\mu}{sT} \frac{s s}{(2\mu + 3\lambda)^2} \]

\[ - \rho_0 C_{ij} \dot{T} + \rho_0 R. \]

For the special case of uniaxial loading, we have \( \sigma_{11} = \sigma_{22} = 0, \sigma_{33} = s_0 + s_1, \) and \( \sigma_{ij} = 0 \) (in\( \varepsilon \)). Then the Eq.(10) can be reduced as

\[ q_{ij} = T_s \left[ \frac{1}{2\mu} \frac{s E}{s T} \delta_{11} - \rho_0 C_{11} \dot{T} + \rho_0 R. \right] \]

According to the Fourier law, the thermoelastic coupling equation of heat conduction is obtained

\[ -\frac{s E}{s T} \left[ 1 - (k_{ij} s X_i) T \right] = T_s \left[ \frac{1}{2\mu} \frac{s E}{s T} \delta_{11} - \rho_0 C_{11} \dot{T} + \rho_0 R. \right] \]

where \( k_{ij} \) is thermal conductivity tensor.

2. TEMPERATURE FIELD

A disk sample of thickness 2\( l \) and radius \( b \) is simply supported at \( r = R \) and is prestressed along the \( z \) axis with stress \( s_0 \). A chopped laser beam is incident on the surface \( z = -l \) of the sample and a heat source \( H(z,t) \) is formed in it. In one-dimension model it can be expressed as

\[ H(z,t) = \gamma \frac{1}{E} \exp \{ \alpha (z - \gamma (z + l)) \} \]

where \( \gamma \) is the optical absorption coefficient, \( \lambda_0 \) the amplitude of laser intensity, \( \omega = 2\pi f \), and \( f \) is the chopping frequency.

Then Eq.(12) can be written as

\[ -\frac{s E}{s T} \frac{s T}{s T} = T_s \left[ \frac{1}{2\mu} \frac{s E}{s T} \delta_{11} - \rho_0 C_{11} \dot{T} + H(z,t), \right] \]
where $k$ is thermal conductivity along $z$ direction and $E$ is Young's modulus.

Owing to the chopping frequency $f$ is low, the Beltrami–Michell equation [6]

$$
s^2 z / a^2 = [-2Ea / (1 - \nu)]s^2 T / a^2
$$

(15)

can be used to find the relation between $T_1$ and $s_t$. Using initial conditions $s = s_0$ and $T = T_0$ at $t = 0$, the solution of Eq.(13) is found as

$$
s_t = [-2Ea / (1 - \nu)]s_1,
$$

(16)

Therefore Eq.(14) can be expressed as

$$
s_T \frac{s^2}{a^2} + \frac{1}{s_t} \frac{s}{a} = \frac{-\gamma s}{k} \exp[-\gamma(z + l) + i\omega t],
$$

(17)

where

$$
1 / s_t = 1 / s + 1 / a, \quad s = k / \rho_s C_s, \quad \text{and} \quad (1 / a) = [2EaT_2 / k(1 - \nu)]^{1/2} \frac{sE}{\rho_s C_s} - a.
$$

(18)

According to the boundary conditions

$$
s_T / a = 0 \quad (z = 0), \quad s_T = 0 \quad (z = l),
$$

the temperature field is obtained:

$$
T_1(z, t) = Q \cdot \left[ A \cdot \exp[s(z + l)] + B \cdot \exp[-s(z + l)] \right] \exp(i\omega t)
$$

(19)

where

$$
Q = -\gamma s / (k(\gamma^2 - s^2)), \quad s = \omega / \sigma, \quad \gamma = \eta / \sigma, \quad A = 1 / \left[ \exp(-2\omega l) - \exp(-2\nu l) \right], \quad B = \exp(2\omega l) + \exp(-2\nu l)
$$

(19a)

3. DISPLACEMENT FIELD

According to the Navier–Stokes equations, constitutive equation (4), and $\sigma_m = 0$, the displacement $U$, $U_2$) in a thin sample, whose surfaces are considered as stress free, is determined by following set of equations:

$$
s^2 U / a^2 + (1 / r)[sU / a] - (U / r^2) = 0,
$$

$$
sU / a = [(l + v)(3v - 1) / (1 - v)]sT_1 - [v / (1 - v)][sU / a + U / r],
$$

$$
U / a = -sU / a.
$$

(20)

and boundary conditions:

$$
\int_{r} \sigma_n dz = 0 \text{ and} \int_{r} \sigma_n dz = 0 \text{ at} \ r = b, U_2(R, z, t) = 0.
$$

Then the displacement field $U$ is solved and the component $U_1$ is

$$
U_1(r, z, t) = [\sigma(1 + v) / 2(1 - \nu)](N + 3M, z / i^2) \gamma,
$$

(21)

where

$$
M t = \int_{-l}^{l} sT_1 dz, \quad \text{and} \quad N t = \int_{-l}^{l} T_1 dz.
$$

(22a)

4. OUTPUT SIGNAL OF PZT

A thin piezoelectric transducer, polarized along its thickness direction, is used as a detector to receive the PA signal and its surfaces are considered as stress free. Then from piezoelectric equations, displacements continuous at the coupling surface $z=1$ between the sample and PZT, and Eq.(21), the output signal of PZT, when its output terminals are open, can be found to be (17):

$$
V = \epsilon_{3p} \epsilon_{33}^* / \epsilon_3^* \epsilon_{33} (1 - \nu) \frac{s^2}{a^2} (N + 3M, z / l)
$$

(23)

If the sample is strong absorptive and thermally thick, the output signal of PZT can be simplified as

$$
V = \frac{-2\epsilon_{3p} \epsilon_{33}^*}{\epsilon_3^*} (1 + \nu) \frac{s^2}{a^2} (N + 3M, z / l)
$$

(24)

or

$$
V = \frac{V_0}{V} = \frac{K}{\gamma}
$$

(24a)

where $V_0$ is the output signal of PZT when prestress $s_0 = 0$, $\epsilon_3^*$, and $\epsilon_3^*$ are coefficients of the PZT related to its piezoelectric, dielectric and compliance constants, $K = \frac{m}{(m - 1)}, \quad m = \frac{2Ea T_0}{\rho_s C_s (1 - v)}$, and $n = \frac{1}{\epsilon_3^*} (\epsilon_3^* / \epsilon_3)$. Therefore the stress distribution in the sample can be evaluated from Eq.(24).

![Fig.1 Experimental results about stress distribution of Vickers indentation on A1.](image)

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REFERENCE
Application of the Method of Kichhoff Integral Formula to Acoustics Scattering from Arbitrary Shape Body.
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§ 1 Kichhoff integral formula used in scattering problem

According to the free-space Green's function, in a closed field V and its boundary S, we have

$$\int_{S} (\Phi_{\omega} - \psi_{\omega}) ds = \int_{S} (\psi \Delta \psi - \psi \Delta \Phi) dv$$

(1)

In Eqn(1), $\Phi$ and $\psi$ are successive in space V and on surface S respectively, and they all have their first and second partial derivatives. In acoustics, $\Phi$ and $\psi$ are simple harmonic functions, therefore, Helmholtz function could be written as

$$\int_{S} (\Phi_{\omega} - \psi_{\omega}) ds = 0$$

(2)

In Eqn(1) and (2), $\partial \Phi / \partial n$ represents the normal derivative with respect to the observational point.

We can lead into the following auxiliary function.

$$\psi = \frac{e^{ir}}{r}$$

To handling the problem of scattering, a volume V formed by four closed surfaces could be assumed as

$S_{1}$: enclosed surface area of scatterer.
$S_{2}$: a micro global surface area enclosing receiving point $M(\vec{R})$, its radius $r = 0$.
$S_{3}$: enclosed surface surrounding the source point.
$S_{4}$: global surface whose radius $R = \infty$, on the surface of $S_{2}$.
$\Phi = 0$, $\partial \Phi / \partial n = 0$.

In accordance with the statements above, we have

$$\int_{S} (\Phi_{\omega} - \psi_{\omega}) ds = \int_{S_{2}} (\Phi_{\omega} - \psi_{\omega}) ds + \int_{S_{3}} \Phi_{\omega} \phi_{\omega} ds + \int_{S_{1}} \phi_{\omega} \psi_{\omega} ds = 0$$

(3)

Consider $\psi = e^{ir} / r$, $e \to 0$, and $\Phi = \Phi_{\omega} + \Phi_{\psi}$ Eqs(3) can be written as:

$$\Phi_{\omega} (\vec{R}) = \frac{-1}{4\pi} \int_{S} (\Phi \frac{\partial}{\partial n} \frac{\omega n}{r} - \phi \frac{\partial}{\partial n} \frac{\nu n}{r}) ds$$

(4)

where $r$ is the distance between the receiving point $M(\vec{R})$ and the microunit ds.

The Eqns(4) is the Kichhoff Equation applicable for solving the problem in scattering field. On the left hand side, $\phi_{\omega}$ is the distribution in the scattering field, on the right hand side, $\phi_{\omega}$ is the distribution on the surface of scatterer, and $\partial \psi / \partial n$ is the distribution of normal derivative of $\psi$ on the surface of the scatterer.

Coordinate the two-dimension kichhoff equation could be derived as follows:

$$\Phi_{\omega} (M') = \frac{i}{4\pi} \int \frac{\partial \Phi}{\partial n} \cdot H_{0}^{(1)} (kr) - \Phi_{\omega} \frac{\partial}{\partial n} H_{0}^{(1)} (kr) dl$$

(5)

§ 2 Solving acoustical potential function by means of Kichhoff formula

Theoretically, if we know the distribution of sound pressure and vibration velocity on the surface of scatterer, we can find out the potential function of arbitrary point in the scattering field by using the Eqns(5) and (4).

By moving the observing point M to the surface of scatterer (the corresponding point is M'), we can obtain $\phi_{\nu} (M')$ by solving integral equation as follows:

$$\Phi_{\omega} (M') = \frac{1}{2\pi} \int_{r} \frac{r}{r} \frac{e^{i(r'\omega)}}{r} \cdot \frac{e^{i(r'\omega)}}{r} dl$$

where $r(M',s)$ is the distance between the M' and micro unit ds.

In numerical calculation, to avoid the complicated process of solving the integral equation, we can simplify it by a method of solving algebraic equations.

The method can be expressed as follows:

The surface of scatterer can be divided into finite micro elements as ds1, ds2, ..., dsn, whose centers are M1, M2, ..., Mn respectively.

If the vibration velocity of each unit is known as $\phi_{\omega} (S_{i}) / \omega n = \nu_{\omega} (S_{i})$

The integral equation could be written as:

$$\Phi(M) + \frac{1}{2\pi} \sum_{i=1}^{n} \Phi(M') \frac{\partial}{\partial j} (e^{i(r'\omega)}/r) dsj$$

$$= \frac{1}{2\pi} \sum_{i=1}^{n} \nu_{\omega}(S_{i}) e^{i(r'\omega)/r} dsj$$

(7)

Having solved the liner equations, we can obtain dispered results of the surface potential functions.

$\phi_{\omega}(M_{1}), \phi_{\omega}(M_{2}), ..., \phi_{\omega}(M_{n})$

The formula (4) can be written in the calculated form.

$$\Phi(M) = \frac{1}{4\pi} \sum_{i=1}^{n} [V_{\omega} \frac{e^{i(r'\omega)/r}}{r} - \Phi(M') \frac{\partial}{\partial j} (e^{i(r'\omega)}/r)] dsj$$

(8)

For vj and $\Phi$ of the finite micro area are available, therefore $\Phi(R)$ can be calculated easily.

§ 3 The approximate numerical method.

I) Integral formula
If the dimension of scatterer and radius of curvature of scatterer are much larger than the wave length, it can be considered that the relationship between sound pressure and vibration velocity is as same as that in the circumstance of plane wave, we have

\[ P = \rho_s C \cdot V_s \]

then

\[ \Phi_s = \frac{1}{ik} \cdot \frac{\partial \Phi_s}{\partial n} \]

Substitute into equation (4), we have

\[ \Phi_s(\hat{R}) = -\frac{1}{4\pi \hat{R}^2} \cdot \frac{\partial \Phi_s}{\partial n} \cdot \frac{1}{r} \cdot \frac{\partial \theta}{\partial n} \cdot \frac{1}{r} \cdot \frac{\partial \Phi_s}{\partial n} \cdot \frac{1}{r} \cdot \frac{\partial \theta}{\partial n} \cdot ds \] (9)

Assume that, on the surface of scatterer, the angle between the incident direction \( \hat{n} \) and the normal line direction of \( ds \) is \( \theta \), and the angle between the direction of \( \hat{R} \) which is from \( ds \) to the receiving point \( M(\hat{R}) \) is \( \theta_r \). Also, the far-field condition is considered in which \( k > > 1/r \), then

\[ \Phi_s(\hat{R}) = \frac{1}{4\pi} \int_s V_s \cdot \frac{\partial \Phi_s}{\partial r} \cdot \frac{1}{1 + \cos \theta_r} \cdot ds \] (10)

Coordinately, the formula in the two dimensions space can be deduced as follows:

\[ \Phi_s(\hat{R}) = \frac{1}{\sqrt{8\pi R}} \cdot \int_s V_s \cdot cos\theta_i \cdot \frac{1}{\sqrt{r}} \cdot \left[ e^{i\theta_r} - \frac{1}{r} e^{i\theta_r} \right] \]

\[ + \frac{e^{i\theta_i}}{\sqrt{r}} \cdot \cos \theta_r \cdot dl \] (11)

II) Calculating formula

By use of formula (10) and (11), we can figure out any sound field formed by incident wave scattering on scatterer. To apply these formulas, we have studied the scattering field formed by incident plane wave scattering on a solid arbitrarily shape body.

Consider the circumstance shown in the following graphics.

The essential parameters are as follows,

- \( ds_j \): a certain micro unit whose central coordinate are \( x_j, y_j, z_j \)
- \( \hat{n} \) : the direction of incident wave
- \( n \) : external normal direction of \( ds_j \)
- \( \theta_{\hat{n}} \) : the angle between the direction of receiving point and \( \hat{n} \)
- \( \theta_{\hat{n}} \) : the angle between the incidence and

We can change the Eqs (10) into the following form

\[ \Phi_s(M) = C \cdot \sum_{i=1}^n \cdot a_i \cdot \cos(\omega t - \varphi_i) \]

where

\[ C = P_s \cdot \triangle s / 4\pi \cdot \rho_s \cdot c \]

\[ a_i = \cos \theta_i (1 + \cos \theta_r) \]

\[ \varphi_i = k(x_i + r_i) \]

Let \( A_i = c \cdot a_i \)

Then

\[ \text{As}(M) \cdot \cos(\omega t - \varphi_i) = \sum_{i=1}^n A_i \cdot \cos(\omega t - \varphi_i) \]

By means of addition of cosine function or by alteration of time periodical function, we can obtain \( \varphi_i(M) \) and \( \text{As}(M) \)

The direct sound field arrived on \( M \) is

\[ \varphi(M) = \varphi_i \times e^{-i(\omega t - \varphi_i)} \]

The whole sound field is

\[ \varphi(M) = \varphi_i(M) + \varphi_i(M) \]

III) Example of calculation

To apply the deduced formula into practice, we have studied the sound field formed by the plane wave incidenting an unlimited cylinder and an arbitrary shape body.

1. Plane wave incidenting into cylinder. The essential parameters are as follows \( a = 2 \), \( ka = 5 \)

The distance between observed point and central axis is 40

By calculating result on computer and studying the graphics drawn by graphic machine, we discover that the graphics concerned is almost as same as that calculated by means of classic method.

2. Plane wave incidenting into a regular triangular prism

According to the approximate condition, contribution points are those that can be "irradiated" by the light emitted from point \( M \)

§ 4 Conclusion

The aim of the paper is to solve the problem of stiffness body scattering. By means of Kichhoff formula method, the integral function has been deduced, and the approximate method with numerical calculation has been proposed. In order to demonstrate the formula's validation, we lay much emphasis on studying the incidence of plans wave incidenting on cylinder and compare with the result calculated by classic method, so that the validity of calculating the scattering field by using kichhoff function has been proved. In the end, the velocity graphics of scattering field of triangular prism formed by plane wave has been given out.
ULTRASONIC SOUND PRESSURE REFLECTED FROM SEVERAL REGULAR REFLECTORS

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1. INTRODUCTION

It is different from the other monographs used higher and special mathematics, in this article, the demonstrating of formular of sound pressure reflected from several regular reflectors have been completed by using elementary mathematics on the basis of some basic physical concepts and laws of geometric ultrasonic optics which are known well.

2. BASIC CONCEPTS OF PHYSICS AND GEOMETRIC ULTRASONIC OPTICS FORMULAS

For plane waves, the acoustic pressure doesn't change in his propagation. For spheric waves, if the acoustic pressure at distance 1 from source is P₁, then the acoustic pressure at distance a from source is Pₐ = P₁/√a. For cylindric wave, if acoustic pressure at distance 1 from source is P₁, then the acoustic pressure at distance a from source is Pₐ = P₁/√a. The geometric optics law of reflection of spheric wave on a spheric or cylindric surface (Fig. 1):

\[
\frac{1}{SA'} + \frac{1}{SA} = \frac{2}{SC} = \frac{1}{\frac{f}{a}} = \frac{1}{f} \tag{1}
\]

The acoustic pressure at the distance a from piston oscillator is P = PoF₀/λ/a. Po is initial sound pressure, λ is wavelength, F is the area of crystal with dia. D.

3. REFLECTION OF PLANE WAVE ON SPHERIC AND CYLINDRIC SURFACE (Fig. 2)

The sound pressure is P₁ at distance 1 from A. According to reversibility of sound, we can consider the pressure Pₛ at point S and the pressure Pₓ at point B are set up by point source A. So, Pₛ = P₁/|F|, |F| = f/2. And,

\[
P_{\infty} = P_{\text{plane}} \cdot \frac{|F|}{|X - F|} \tag{2}
\]

Now consider cylindric reflector. At S, pressure Pₛ = P₁/|F|. According to reversibility of sound, we can consider acoustic pressure Px at point B is set up by cylindric source at A. Hence,

\[
P_{\infty} = P_{\text{plane}} \cdot \frac{|F|}{\sqrt{|X - F|}} \tag{3}
\]

4. REFLECTION OF SPHERIC WAVE ON SPHERIC AND CYLINDRIC REFLECTOR (Fig. 3)

From (1), b = a|F|/|a - F|. If the pressure at distance 1 from O is P₁, then Pₛ = P₁/|F|. If the pressure at distance 1 from A is P₁', then Pₛ = P₁/|F| = P₁'/|F|/|a - F|/|F|/a. So, Pₛ = P₁|F|/|a - F|. According to reversibility of sound, we can consider the pressure Px at point B is set up by point source A. So,

\[
P_{\infty} = \frac{P₁}{a} \cdot \frac{|F|}{x - F(l + \frac{x_a}{a})} \tag{4}
\]

The cylindric surface can be regarded as double curvatuer surface. One radius of curvature is infinite great, the other is radius of cylindric surface. Consider the reflection of spheric wave on the cylindric surface can be regarded as the geometric average of both reflection of two surface. The reflection on plane can be obtained by setting f in (4) to be infinite great. So,

\[
P_{\infty} = \frac{P₁}{a} \cdot \frac{f}{l + \frac{x_a}{a}} \tag{5}
\]

Hence, acoustic pressure Px of spheric wave after reflecting on cylindric surface is,

\[
P_{\infty} = \frac{P₁}{a} \cdot \frac{|F|}{\sqrt{(l + \frac{x_a}{a})[x - F(l + \frac{x_a}{a})]}} \tag{6}
\]
5. ACoustic Pressure RECEIVED by PUSh—catch Probe AFTER HAVING BEEN REflected BY seVERAL reguLAR REFLECTORS

Cylindrical Reflector
Let \( x = a, f = r/2(\ell_r) \), cylindrical radius in (6).

\[
P_{\text{cylinder}} = \frac{P_0 F_0}{2 \alpha} \frac{\sqrt{1 + F^2}}{a - f} \quad (7)
\]

For Fig. 4(a), \( a = R_1 - R_2, r = -R_2 \), so,

\[
(P_{\text{cylinder}}) = \frac{P_0 F_0}{2 \alpha (R_1 - R_2)} \sqrt{\frac{R_2}{R_1}} \frac{R_2}{R_1} \frac{R_2}{R_1} (R_1 = R_2) \quad (8)
\]

For Fig. 4(b), \( a = R_1 - R_2, r = -R_1 \), so,

\[
(P_{\text{cylinder}}) = \frac{P_0 F_0}{2 \alpha (R_1 - R_2)} \sqrt{\frac{R_1}{R_2}} \frac{R_1}{R_2} \frac{R_1}{R_2} (R_1 = R_2) \quad (9)
\]

Large Plane Bottom
Let \( a = 1 \), so \( P_1 = \text{PoF}_0/\lambda \), substitute \( P_1 \) and \( x = a \) into formula (5), so,

\[
P_x = \frac{P_0 F_0}{2 \alpha} \quad (10)
\]

Plane Bottom Hole With Dia. \( \varphi \) or Small Square With Side Length \( L \)
Consider these small reflectors as new oscillator, by double times using \( P = P_0 F_0 / \lambda / a \),

\[
P_{\text{plane bottom hole}} = \frac{P_0 F_0 F_f}{\lambda^2 a^2} \quad (11)
\]
in above formula, \( F_f = (\varphi / 2)^2 \) or \( F_f = L^2 \).

Small Sphere With Dia. \( \varphi \)
In this case, \( f = -\varphi / 4, x = a, P_1 = \text{PoF}_0 / \lambda / x = \text{PoF}_0 / \lambda / a \), substitute these into (4), so,

\[
P_0 F_0 \varphi \quad (\varphi \ll a) \quad (12)
\]

Transcendent Drilled Hole With Dia. \( \varphi \)
When hole length is above the width of sound beam, \( x = a, f = -\varphi / 4, P_1 = \text{PoF}_0 / \lambda \), substitute these into (7),

\[
P_{\text{long drilled hole}} \approx \frac{P_0 F_0}{2 \alpha} \frac{\sqrt{\varphi}}{2 a} (a \gg \varphi) \quad (13)
\]

When hole length is smaller then width of beam, suppose the length \( L \) and diameter of hole is the same order of magnitude, that is \( = L, \alpha \), so, \( f = -\varphi / 4, P_1 = \text{PoF}_0 / \lambda \). Consider reflection of sphere wave on short transcendent hole is equal to geometric average of the both reflection on small sphere with dia. and small square with side length \( L \). So,

\[
P_{\text{short drilled hole}} \approx \frac{P_0 F_0 L}{2 \alpha} \frac{\varphi}{\lambda^3} \quad (14)
\]
ANALYSIS AND CONTROL OF MUZZLE BLAST

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When a bullet is launched from barrel of a gun at relatively high initial velocity, the muzzle blast flowfield and gaint explosive sound are generated at the muzzle. These physical phenomena are analysed sketichly as follows.

A large amount of the gunpowder gas at high pressure and high temperature fluxes from the muzzle, after a bullet is shot from a gun. The gunpowder gas is released suddenly, it spreads freely. The fluid flow velocity rapidly accelerated approximately 3 times. Thus the gunpowder gas has been surpassing a bullet, they get together to form the cloud of gases. Furthermore the gunpowder gas spreads toward both sides at large angle. This supersonic jet of the gunpowder gas compresses surrounding air with great violence, it creates air pressure and density to jump. It follows that strong the muzzle blast wave is formed. The pressure of wavefront is very high. The jet boundary is in compressed air space. The structural model of the muzzle blast flowfield has been shown in Figure 1.

Figure 1 shows the results:
1. The cloud of gases and blast wave which have shape of cap. They are formed by the gunpowder gas in front of a bullet.
2. Muzzle blast wave. The expanding gunpowder gas compresses surrounding air, thus the airblast is generated. After the airblast separates from the gunpowder gas, it travels independently.
3. The gunpowder gas jet. The centre of jet is unsteady supersonic jet core. After a bullet is discharged from a gun about 70 μs, shock wave and Mach disc emerge in jet area. Size of Blique Shock changes with the time, because process of flux of the powder gas is an exhaustive process from fixed vessel. The length of Shock Bottle is 10-20 times one caliber. The gunpowder gas in core of a gun has been exhausted after 400μs. The Shock Bottle and Mach Disc contracts gradually until it is removed. By this time a bullet has past the cloud of gunpowder gas.

The validity of a model of the muzzle blast flowfield was proved by experiments.

It is evident that the airblast formed by the gunpowder gas jet is main cause of generating the muzzle noise. This supersonic gunpowder gas jet is source of energy which manufactures the airblast, it pushes the airblast by virtue of gunpowder gas-air contact surface outside jet boundary, and continually supplies energy to the airblast. Shock Bottle which envelopes jet core has sound baffle, thus the jet core is an emissive source of the muzzle noise in essence.

The process of propagation of the muzzle blast may divide into three sections: near field, middle field, and far field. The intensity of blast is high enough to destroy human ear organ in near field. The blast is in state of strong pressure wave, its superpressure is higher than 1kg/cm². The blast enters middle field with increase of propagation distance. The superpressure of wavefront is reduced into GAF<1kg/cm². The blast is in state of weak pressure wave. The weak blast continuously travels forward, finally it is attenuated into sound wave. This moment superpressure P is tending toward zero. The wave speed remains constant i.e. sound speed.

Some submachine gun. The initial velocity of a bullet is 710m/s, the pressure and temperature of the gunpowder gas at the muzzle are 500kg/cm² and 1700K respectively. Measure confirmed that data of the muzzle noise may be summarized in Table. The points of measure can be presented in Figure 2.

<table>
<thead>
<tr>
<th>Position (Position in deg)</th>
<th>30°</th>
<th>90°</th>
<th>120°</th>
</tr>
</thead>
<tbody>
<tr>
<td>peak sound pressure level</td>
<td>162.6</td>
<td>161.0</td>
<td>146.9</td>
</tr>
<tr>
<td>effective sound pressure level</td>
<td>142.0</td>
<td>136.4</td>
<td>126.3</td>
</tr>
<tr>
<td>sound level (dB)</td>
<td>146.5</td>
<td>147.7</td>
<td>145.0</td>
</tr>
</tbody>
</table>
The peak frequency of the muzzle noise is 500 Hz and the pulse length is 430 μs.

Such strong noise endangers directly health of gunners.

The muzzle silencer is a device which weakens the muzzle noise. Therefore it is crucial to weaken the muzzle noise to reduce pressure, velocity, and temperature of gurnpowder gas. The muzzle silencer can solve better than problems. Figure 3 is a muzzle silencer of the latest design. The outside diameter of silencer is 26 mm and length is 104 mm. The silencer is a extended tube which is attached to barrel of gun. There are three decompression chamber (1, 3, 4) in the tube. The high pressure gurnpowder gas from a gun barrel expands in decompression chamber. There is a section of spiral flake in middle of the tube (A-A'B'). It is possible that the gurnpowder gas revolves round spiral flake. The stroke of gas flow is increased. Thereby the pressure of central place is reduced and the escaping velocity of gas is decreased. Spiral structure can eliminate yet the regenerated noise. The eighty-four small holes are arranged on outside wall of the second decompression chamber 3 (BB-CC). Diameter of hole is 1 mm and distance between holes is 5 mm. To attain the goal of silence noise, the principle of frequency conversion is used with the help of small holes. In order that shock wave surface is destroyed and whistler-type noise is removed, we have cut a gap which has the shape of the letter V on surface of the third decompression chamber 4 (CC-DD).

The pressure of gurnpowder gas and Mach number of gas flow in the silencer may calculated by fundamental formula

\[
\frac{p_0}{p} = (1 + \frac{k-1}{2} \frac{\rho}{\rho_0})^{\frac{k-1}{k}} \quad (1)
\]

\[
\frac{T_0}{T} = 1 + \frac{k-1}{2} \frac{\rho}{\rho_0} \quad (2) \quad C = \sqrt{\frac{R \kappa_0}{\rho_0}} \quad (3)
\]

\[
\frac{X}{d} = \left[1 + \frac{k}{k+1} (\frac{\rho}{\rho_0})^{\frac{k-1}{k}} \right]^{\frac{1}{k}} \quad (4)
\]

where \( X \) is Mach number, \( p_0 \) the pressure, \( T_0 \) the temperature, \( k \) the heat insulation index, \( C \) the sound speed, \( g \) the acceleration of gravity, \( R \) the constant of gurnpowder gas, \( d \) the caliber, and \( X \) the distance between the muzzle and calculating point.

The silencing value of spiral flake may be calculated by equation (5)

\[
\Delta L = 20 \log \left(1 + \frac{\rho_0}{\rho} \sin \frac{\pi X}{L} \right) + 10 \log \left(1 + \frac{\rho_0}{\rho} \right) \quad (5)
\]

where \( \Delta L \) is the silencing value, \( \rho \) the dilation ratio, \( L \) the length of decompression chamber, \( \rho_0 \) the equivalent length of spiral flake, \( n \) the number of spiral flake, \( \pi = 2 \pi l/c \).

Sound level from the silencer is estimated by experimental equation (6)

\[
L = 93.5 + 20 \log \frac{D}{D_0} + 20 \log \frac{\rho_0}{\rho} + 10 \log \left(1 + \frac{\rho_0}{\rho} \right) + 10 \log \left(1 + \frac{\rho_0}{\rho} \right) \quad (6)
\]

where \( D \) is diameter of silencer mouth, \( p \) the pressure of gurnpowder gas (BB-CC), \( p_0 \) the environment pressure, \( \rho_0 = 1 \text{kg/cm}^2 \), \( h \) the Mach number, \( \rho = 1 \), \( n \) the number of holes, \( X = 0.165D/l_0 \).

The calculating results show that we attach a silencer to submachine gun, the predicted value of muzzle noise is 140.3 dB. The predicted value of silencing value is 6.26 dB, the actual measuring value is 8.0 dB.

This silencer controls effectively the muzzle blast and decreases the muzzle noise, it provides scientific basis for design of new silencer.

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Reflection of Sound Wave from a Vibrating Plane

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Introduction

It is known that Doppler effect can be determined through a simple coordinate transformation when sound is reflected by an object with uniform motion. But the problem becomes very complicated when the object moves arbitrarily. Reflection of normally incident electromagnetic wave at a vibrating plane of conductor has been investigated by Censor, Cooper, Van Bladel and De Zutter [1]-[3]. With direct extension of the method to acoustics, Censor has studied several cases that harmonic plane wave are scattered by vibrating boundaries of simple geometry [4]-[6]. His results show that reflecting wave contains frequency components of summation and subtraction of frequencies of incident wave and boundary vibration, as well as other components of higher orders. Picquette and Van Buren have a different opinion. They believed that appearance of components of any new frequency in reflecting wave must be a result from consideration of nonlinear effect of the medium [7]. Therefore, there is a need to get a quantitative concept that in what degree the minor vibration of reflecting object could affect the frequency composition of reflecting wave.

The reflection of normally incident plane wave at a vibrating rigid plane is analyzed on in this paper.

1. Wave Propagation and Observation in Motion

In order to find the acoustic variable induced at the vibrating surface by incident wave, let us first consider the situation that acoustic variable is detected by an observer with arbitrary movement. To make analysis case, it is supposed that plain sound wave propagates along x-axis in its positive direction, the sound source locates at x=0, and observer's motion is also along x-axis and has no effect on sound wave.

Let one acoustic variable (for instance, the particle velocity of sound) at x=0 be \( f(t) \). Assume that \( f(t) \) is non-zero only at a time interval of \([0,T]\), shown as in fig. 1a. Then the wave function corresponding to \( f(t) \) is \( f(t-x/c) \), which defines a spatial distribution of the acoustic variable at a fixed moment \( t \geq T \), that is,

\[
 f(t-x/c) = \begin{cases} \frac{f(tc_0-x)}{c} & \text{if } c_0-T \leq x \leq ct, \\ 0 & \text{otherwise} \end{cases} \quad (1)
\]

Then the acoustic variable detected as a function of time by a stationary observer at \( x_0 \) (\( c_0 \geq x_0 \)) is

\[
 f_2(t) = \begin{cases} \frac{f(tc_0-x_0)}{c} & \text{if } c_0-T \leq x_0 \leq ct, \\ 0 & \text{otherwise} \end{cases} \quad (2)
\]

shown as in fig. 1b, with the shade part denoting the received sound signal.

The transmission of sound signal from a stationary source to a motionless observer can be regarded as the following procedure. Due to the effect of propagation of the wave, the acoustic variable as a function of time at the location of the source is incessantly transformed into a definite spatial distribution. The observer gets the corresponding acoustic variable by recording the spatial distribution that continuously arrives at him at the speed of c.

Because the spatial distribution is fully determined by the source and the sound speed in the medium, and the relativity exits between the propagation of spatial distribution and the movement of observer, we can keep the spatial distribution fixed and let the observer move in the opposite direction of the wave propagation to obtain the same result.

A shift of \( ct_0 \) of the spatial distribution of equation (1) in the minus direction of x-axis gives

\[
 f_4(x) = \begin{cases} f(-x/c) & \text{if } -ct_0 \leq x \leq 0 \\ 0 & \text{otherwise} \end{cases} \quad (3)
\]

shown as in figure 1c. If the function \( f(t) \) is defined on interval of \([0, \infty)\), then we can extend the above spatial distribution to \((\infty, 0)\). Thus, for any given function \( f(t) \) of acoustic variable of a source, a definite spatial distribution can be determined. By the relativity, the acoustic quantity detected by a stationary observer at \( x_0 \) is equivalent to that of a moving observer of the fixed spatial distribution. Let the initial position of the moving observer be at \( x_k \), then the acoustic quantity he observed is gained by substituting his position into equation (3). The nonzero part of it is

\[
 f_4(t-t_0) = \int_{t_0}^{t} \frac{f(tc_0-x)}{c} \, dt 
\]

where \( t_0 = x_0/c \) is the time that the observer initially detects the distribution.

It is easy to extend the above result to enclose the situation when the observer moves with arbitrary velocity \( v(t) \) along the x-axis. With substitution of the total velocity into equation (4), the observed acoustic variable is given as

\[
 f_4(t-t_0) = \int_{t_0}^{t} \frac{f(t-x/c)}{c} \, dt 
\]

where \( t_0 \) is the time that signal is first detected. Because the total velocity \( c-v(t) \) varies with time, \( t_0 \) usually is not equal to \( x_0/c \). Extension is also not difficult to include the case when the direction of wave propagation differs from the observer's.
2. SOUND REFLECTION FROM A VIBRATING PLANE

Let us now consider the problem of reflection of normally incident plane wave from a rigid plane with vibration. At the surface of vibration, the boundary condition

\[ u_t(t_0) + u_y(t_0) = 0 \]  

should be met, where \( u_t(t) \) is the particle velocity induced by incident wave and \( u_y(t) \) the particle velocity causing reflecting wave. Suppose the particle velocity of the incident wave is \( u_t(t-x/c) \) and the vibrating velocity of the boundary \( v(t) \), from equation (5) and (6) we have

\[ u_y(t) = u_t(t-x/c) \int_{c}^{t} v(T)/c \, dT \]  

in which \( t_0 \) is the time the wave initially arrives at the boundary. If the summation of velocities of \( u(t) \) and \( u_y(t) \) still meets the request for emission of a sound wave of small amplitude, then we can determine the reflecting wave and the emitting wave caused by boundary vibration itself by the boundary condition of particle vibration given at a stationary position (for example, the balance position of vibration of the reflecting plane). Thus, we can write down the forms of these directly from \( v(t) + u_y(t) \); they are \( u(t + (x-x_0)/c) \) and \( v(t + (x-x_0)/c) \) respectively, where \( x_0 \) is the balance position of vibration of the plane.

The effect on frequency composition of reflecting wave by vibrating boundary is usually concerned. To carry out more detail calculation, we suppose the incident wave and the vibration of reflecting boundary are both harmonic in the following discussion.

Let the particle velocity of incident wave be

\[ u_t(x, t) = u_0 \cos \omega (t - x/c) \]  

and the displacement of the vibrating plane

\[ x(t) = x_0 \sin \omega t \]  

Substitution of equation (8) and (9) into (7) gives

\[ u_y(t) = u_0 \cos(\omega t - \alpha \sin \omega t + \varphi) \]  

where

\[ \alpha = 2 \pi f_s / \lambda \omega \]  

\[ \varphi = \omega t_0 - \alpha \sin \omega t \]  

with \( \lambda \omega \) being the wavelength of incident wave and \( t_0 \) the time the wave initially arrives at the vibrating plane.

We further suppose \( \alpha < 1 \). An expansion of equation (10) gives

\[ u_y(t) = u_0 (\cos(\omega t - \varphi) - \frac{\alpha^2}{2} \sin^2 \omega t + \frac{\alpha^3}{6} \sin^3 \omega t - ...) + \sin(\omega t - \varphi)(\alpha \sin \omega t - \frac{\alpha^3}{3} \sin^3 \omega t - ...)) \]  

Obviously, it contains components with frequencies of \( \omega \), \( 2 \omega \), \( 3 \omega \), \( 4 \omega \), etc. The first three terms are

\[ -u_0 (1 - \frac{\alpha^2}{2} \sin^2 \omega t + \frac{\alpha^3}{6} \sin^3 \omega t - ...) \cos(\omega t - \varphi) \]

\[ u_0 \left( \frac{\alpha^2}{2} \sin^2 \omega t - \frac{\alpha^3}{6} \sin^3 \omega t - ... \right) \cos((\omega + \omega) t - \varphi) \cos((\omega - \omega) t - \varphi) \]

\[ u_0 \left( \frac{\alpha^2}{2} \sin^2 \omega t - \frac{\alpha^3}{6} \sin^3 \omega t - ... \right) \cos((\omega + 2 \omega) t - \varphi) \cos((\omega - 2 \omega) t - \varphi) \]

It seems that the magnitudes of new frequency components of \( \omega \), \( 2 \omega \), \( 3 \omega \), etc. as proportion to the displacement of vibrating plane to the wavelength of incident wave. The magnitudes of the first three terms versus \( f_s / \lambda \omega \) are listed in table 1 to help us to get a quantitative concept.

<table>
<thead>
<tr>
<th>( f_s / \lambda \omega )</th>
<th>( \alpha )</th>
<th>( \omega )</th>
<th>( \omega + \omega )</th>
<th>( \omega - \omega )</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.001</td>
<td>0.0063</td>
<td>0.04D</td>
<td>-50.0D</td>
<td>-106.09D</td>
</tr>
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<td>0.0628</td>
<td>0.04D</td>
<td>-30.0D</td>
<td>-66.09D</td>
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<td>-0.9D</td>
<td>-10.5D</td>
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</tr>
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<td>0.15</td>
<td>0.942</td>
<td>-2.18D</td>
<td>-7.56D</td>
<td>-17.99D</td>
</tr>
</tbody>
</table>

Table 1: The magnitudes of frequency components of first three order versus \( f_s / \lambda \omega \) where the magnitude of incident wave is taken as 1 (corresponding to 0dB)

It seems from table 1 that if \( f_s / \lambda \omega \) is less than 0.001, the magnitudes of frequency components of summation and substraction are more than 50dB lower than that of the incident wave so that they have no practical significance. But if this ratio is larger, they may not be ignored.

4. REFERENCES


UNDERWATER ACOUSTICS

B1  The Application of Neural Networks to Underwater Acoustics
B2  Under Water Acoustic Signal Processing
B3  Fisheries Acoustics
B4  Propagation
B5  Propagation, Reverberation
B6  Noise, Attenuation
B7  Scattering
B8  Scattering, Reflection
B9  Signal Processing
B10 Signal Processing
BP  Poster Papers
APPLICATION OF NEURAL NETWORKS TO SEISMO-OCEANIC ACOUSTIC TOMOGRAPHY

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INTRODUCTION

Recently there has been an increase in the activities of offshore oil exploration around the world. To save cost and to reduce the risk of offshore oil exploration, it would be useful to have a crosssection image of the earth's structure below the ocean. We call this technique seismo-oceanic acoustic holography because it involves the transmission of sound from the ocean and then the seabed. Seismo-oceanic tomography has to combine ocean acoustic tomography with seismic tomography. Ocean acoustic tomography is performed by measuring the travel time of acoustic pulses transmitted in the sound channel over several nearly horizontal paths through the portion of the ocean being investigated. Then for seismo-oceanic tomography, the travel time of the sound pulse through the seabed has to be measured and together with the travel time through the ocean, the total travel time has to be obtained.

FORWARD PROBLEM

Ocean acoustic tomography differs from medical ultrasound tomography and seismic tomography in that sound travels long distances of thousands of kilometres during the travel-time measurements stage and sound will be refracted and diffracted in three dimensions and this will affect the measurements. The problem of knowing how to depoly sensors and how to interpret their measurements is known as the "forward problem". To solve the "forward problem", one needs to use various ocean models. Solving sound propagation in ocean, an inhomogeneous and random medium is a nonlinear problem. There are various approaches used such as ray theory, normal mode theory, parabolic equation (PE) approximation and statistical approach etc. In this paper the PE model will be used because it can correctly treat the physics of diffraction, caustics, leakage of energy into and out of surface ducts. The transmission loss (TL) predictions of the PE model in range and depth is useful for tomographic inversion and in the planning of sampling strategies. The transmission loss within the seabed will need to be calculated. The sum of the two transmission losses is needed. First, for the oceanic part the transmission loss in the time domain is examined in the presence of ocean currents. The transmission loss functions in the time domain are computed from the PE obtained by using the Split-Step Fourier (SSF) algorithm. The PE method is used to model the effects of ocean currents on sound propagation. The CW pressure due to a source at point $A$, range $x_A$ and depth $z_A$ is written as

$$P_{A_{\text{OA}}} = \frac{B_0}{r} \left( \frac{r}{c_s} \right)^{1/2} e^{ik_b r}$$

where $r=x-x_A$, $k_b = w/c_s$, $c_s$ is the reference sound speed, $B_0$ is the magnitude of PE, $A_0$, $2\pi Z \rightarrow 0$, $k_b = \text{dimensionless \ envelope}$, $v(x,x)Z$ is horizontal component of the vector current and

$$U(x,x_z) = \frac{1}{2} \left[ 1 - \frac{2}{c_s^2} (x_o, z) \right] - \left( L \angle x, \angle z \right)/k_o$$

where $x(x_o, z)$ = sound speed in the vertical plane between source and receiver and $L(x, z)$ = volume loss. The time-domain transmission loss, including current effects, may be computed from

$$TL_{\text{AB}}(t) = -10 \log \left[ \frac{P_A(t)}{P_A(t)} \right]^2$$

(3)

where $P_A(t)$ = pulse response function

$$P_A(t) = \left( \frac{R_0}{b^2} \right)^{1/2} \int_0^{\infty} \int_0^{\infty} \phi_a(w, b, b') P_{A_{\text{OA}}} \exp \left( i w T_A - k_b T_A \right) \frac{dw}{d(u)}$$

(4)

and $P_{A_{\text{OA}}} = \text{peak pressure at reference range } R_0$ if $\phi_a(w, b, b')$ the solution of the parabolic equation is computed.

To compute the transmission loss through the seabed, we use the layered earth model (Goupillaud) model. The two-way transmission loss (down and back up through the interface) is given by

$$z^2 = -1$$

(5)

where $z = \text{reflection Coefficient$. For K layers, the transmission loss will be

$$TL_K = (1-c_1^2) (1-c_2^2) \ldots (1-c_K^2)$$

(6)

Hence the total transmission loss will be

$$TL = TL_{\text{AB}} + TL_K$$

(7)

INVERSE PROBLEM: RECONSTRUCTION OR TOMOGRAPHY USING NEURAL NETWORKS

Neural networks are naturally suited to perform the inversion of tomography data. The backpropagation neural network is used to implement the filtered backprojection algorithm used in the reconstruction of tomogram. This is done by extending linear parallel distributed processing (FDP) network for implementing the filtered backprojection algorithm of the x-ray case to the nonlinear ultrasound case which includes diffraction effect. To implement the linear, FDP network, the object function has to be discretized. The sampled modified projection serves as the input to the network. The connection weights contain the filter information shifted to perform the convolution. The network multiplies the input by the value of the connection weight and the weighted signal is then passed on to the summation function. This process is repeated in the next layer. Therefore the network realizes the object function at a point (x,y) for any object.

Starting with this linear network at a high frequency a smooth transition to ultrasound and diffraction case can be made by gradually lowering the frequency and introducing a nonlinearity after the network's summation functions. This non-linearity, a characteristic property of neural network can "learn". In this case, the non-linearity is a sigmoid function. After slightly reducing the frequency, the network can be trained again. This gentle process can
be repeated until a network has been trained that compensates for diffraction and avoids local minima in weight space. Instead of entering the value of the connection weights, the neural network can learn the correct weights when repeatedly presented with example input-output pairs.

In the backpropagation network (Fig 1) used in the experiment, the network was constructed with three fully connected layers and a bias. The connections between the processing elements hold weights, values multiplied by the signal passing through the connection. When the weighted signals reach the processing element, they are added together internally before passing the resultant signal through nonlinearity (Fig 2).

The process can be represented by

$$x_j^s = f\left( \sum_{i=0}^n w_{ji}^s x_i^{(s-1)} \right) = f[\theta_j]$$  \hspace{1cm} (8)

and

$$f[\theta_j] = \frac{1}{1 + \exp(-\theta_j)}$$  \hspace{1cm} (9)

for sigmoid.

The layers are indexed by $s$, $s-1$ etc. The weights are usually initialized to a random value. Training the network involves altering the weights to a compact value that gives the desired output.

Training begins by presenting an input $i$ to each of the elements in the input layer. These values propagate through the network to the output layer according to eqns (8) and (9). At the output layer, the local error $e$ is calculated by comparing the network output $\theta$ with the desired output $d$.

$$e_k = O_k - \theta_k$$  \hspace{1cm} (10)

where $k$ is the index of the output processing elements. The connection weights $w$ are then altered according to

$$\Delta w_{ji}^{s} (t) = c_1 e_j^s x_i^{s-1} + c_2 \Delta w_{ji}^{s} (t-1)$$  \hspace{1cm} (11)

where $\Delta w_{ji}^{s} (t-1)$ is the previous weight change and $c_1$ and $c_2$ represent learning coefficients. The local error and the delta weight are calculated by propagating back through each layer. Then the weights are adjusted by adding the corresponding weight change. This process controls learning. The input data for our case is the transmission loss.

CONCLUSION

Neural network can "learn" characteristics about a scanned object from its projection data. Neural networks trained from completely randomized weights require many training cycles. This method does promise to provide greater accuracy because the neural network does not assume weak scattering. The neural network can also be fine-tuned, according to instrument-dependent variables.

REFERENCES


COMPARISON OF SONAR DISCRIMINATION BY AN ECHOLOCATING DOLPHIN AND A COUNTERPROPAGATION NEURAL NETWORK

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The capability of an Atlantic bottlenose dolphin to discriminate wall thickness differences of hollow cylinders by echolocation was studied by Au and Pawloski (1992). A standard cylinder of 6.35 cm wall thickness was compared with cylinders having wall thicknesses that differed from the standard by ±0.2, ±0.3, ±0.4, and ±0.8 mm. All cylinders had an O.D. of 37.85 mm, and a length of 12.7 cm. The standard and a comparison target, separated by an angle of ±1°, were presented simultaneously at a range of 8 m and the dolphin was required to indicate the location of the standard target. The standard target was paired with a different comparison target for ten consecutive trials apiece. The experiment was conducted in the operating field and in the presence of broadband masking noise.

In this study, a counterpropagation artificial neural network was used to examine the broadband echo features from the same cylinders used in the dolphin experiment. Features of the echoes were determined by passing them through a bank of constant-Q filters. Omitting-of-Q filtering was chosen because the dolphin's auditory system can be modeled as a bank of constant-Q filters (Johnson, 1969). The objectives were (1) to determine if a counterpropagation network could discriminate the target echoes using features from the constant-Q filter bank, (2) compare the performance of the counterpropagation network with the that of the dolphin, and (3) determine Q-values needed by the network for comparable performance as the dolphin.

Roitblat et al. (1989) used a counterpropagation network to emulate a dolphin performing a sonar discrimination task. The network performance was perfect with echoes collected in a test pool, and was 97% correct when selective “natural” echoes resulting from the dolphin’s sonar emissions were used. However, the dolphin’s task was not difficult; the animal’s average performance was 94.5% correct. The discrimination task in this study was considerably more difficult with the dolphin’s accuracy varying from 98% to 52% correct. Moore et al. (1991) used a backpropagation network and consecutive “natural echoes” to discriminate the same targets used by Roitblat et al. (1989). The backpropagation network achieved performances between 90 and 93% correct.

I. PROCEDURE

Target echoes were collected with a planar transducer that projected and received the echoes. The transducer was mounted on the dolphin’s pen so that the target measurements would be made in the same environment and under similar conditions as for the dolphin. A simulated dolphin signal with a peak frequency of 117 kHz was projected and a 16-bit analog-to-digital converter at a rate of 1 kHz was used to digitize the echoes. Each echo consisting of 1024 points was stored on computer disk. Five disk files, with 10 consecutive echoes per file or 50 echoes were collected for each target.

Target features were determined by passing each echo through a bank of N contiguous constant-Q filters. Each echo consisted of the energy from each filter normalized to the output of the filter with the maximum energy. From the definition of Q, the frequency boundaries of the ith constant-Q filter can be expressed as

$$f_i = \frac{20 \cdot 10^{(Q-1) f_{0,i}}}{10^{(Q-1)}}$$ (1)

where $f_i$ is the upper frequency and $f_{0,i}$ is the lower frequency limit of the filter. Let $f_0$ be the upper frequency of the $N$th filter and $f_1$ the lower frequency of the last filter of a bank of constant-Q filters, then from Eq. 1 the relationship between $f_0$ and $f_1$ can be expressed as

$$f_0 = \frac{20^{(Q-1)}}{10^{(Q-1)}} f_1$$ (2)

For a specific Q, three parameters can be varied, $f_0$, $f_1$, and N. A frequency of 150 kHz was used for $f_0$ to coincide with the bottlenose dolphin upper frequency of hearing (Johnson, 1968). The lowest frequency was chosen so that $f_0$ ≥ 62 kHz. For frequencies ≤ 62 kHz, the energy in an echo was at least 30 dB down from the peak. For a desired Q, N was chosen so that $f_0$ was as close to 62 kHz as possible.

The counterpropagation neural network was simulated by the Neural Works Professional II Plus program from Neural Ware, Inc. The network consisted of an input layer of N elements, a normalizing layer of N+1 elements, a Kohonen layer of N elements, and an output layer of two elements. Echoes associated with the standard target were paired with echoes from each comparison target. Twenty echoes from each target were used for the training set and 10 echoes from the remaining thirty echoes were used for the test set. The network’s capability to discriminate the standard from each comparison target was determined for different values of Q and N.

The performance of the network with noisy data was determined by first adding a different burst of Gaussian random noise to each target echo. The noisy echo was then passed through the constant-Q filter bank. A noise burst was created by passing white noise through a cosine taper window having a half-power width of 264 μs. A width of 264 μs was chosen to correspond to the dolphin’s integration time of 264 μs determined by Au et al. (1988).

II. RESULTS AND DISCUSSION

The performance of the counterpropagation network for the free-field echoes and Q values of 4 and 5 are shown in Fig. 1, along with the dolphin’s performance. The value of N was equal to Q-1. The network’s performance for a Q of 4 was not as good as the dolphin for most of the comparison targets. However, for a Q of 5 the network’s performance was better than the dolphin for most of the comparison.

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Fig. 1. Results of network in discriminating the standard and comparison targets.
targets. Therefore, a constant-Q filter bank consisting of 4 filters each having a Q of 5 produced target features that allowed the network to perform better than the dolphin.

The dolphin's critical bandwidth at 120 kHz is approximately 43 kHz (Au and Moore, 1990) indicating a Q of 2.8. Although the counterpropagation network could perform better than the dolphin, the filter bank had to have a higher Q than the dolphin. Furthermore, the network had a relatively simple task of sorting 50 echo pairs in files that were already time-windowed so that only target echoes were present. The dolphin had a more complex task of echo-locating the targets, ignoring irrelevant echoes, determining the proper time window, remembering echo characteristics, as well as report to the experimenter.

Typical echoes from the standard and the comparison target having the closest wall thickness to the standard are shown in Fig. 2. Small differences in the spectrum for the comparison target (dash line) compared to the standard target can be seen. The spectrum of the comparison target is shifted slightly toward the left of the spectrum for the standard target. Note that echoes from the same target may be slightly different as a result of wave and wind induced motion on the target and test pen.

The performance of the network for the noisy echoes is shown in Fig. 3. For signal-to-noise ratios of 15 and 10 dB, the network performed better than the dolphin with most of the comparison targets when the Q was equal to 8. For a Q of 7, the network was worst than the dolphin for a signal-to-noise ratio of 15 dB and slightly better than the dolphin a signal-to-noise ratio of 10 dB.

The similarity between the standard and comparison targets can be expressed by a Euclidean distance measure. Let \( E_i(f) \) equal the normalized energy in the i-th filter averaged over all the standard target echoes in the training set, where \( i = 1 \) to \( N \). Let \( E_j(f) \) be the corresponding energy for a comparison target in the test set and \( j \) be the index of the j-th echo in the test set of M echoes. The Euclidean distance \( d_{ij} \) of the j-th comparison echo is

\[
d_{ij} = \left( \frac{1}{N} \sum_{f=1}^{N} (E_i(f) - E_j(f))^2 \right)^{1/2}
\]

The similarity measure averaged over the M test set echoes is shown in Table 1 for a Q of 5 and no noise, and a Q of 8 with a 10-dB signal-to-noise ratio. The results indicate that the standard target can be differentiated from the comparison targets provided a good threshold value of \( d \) is chosen.

Table 1. Euclidean distance measure of similarity

<table>
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<tr>
<th></th>
<th>THINNER COMPARISON TARGET</th>
<th>THICKER COMPARISON TARGET</th>
<th>THINNER COMPARISON TARGET</th>
</tr>
</thead>
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<tr>
<td></td>
<td>No Noise</td>
<td>Q = 5</td>
<td>Q = 8</td>
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<td>Standard (test)</td>
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<td>.2 mm</td>
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<td></td>
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<td>.3 mm</td>
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<td></td>
<td></td>
<td>.5 mm</td>
<td>.5 mm</td>
</tr>
<tr>
<td>Standard (learn)</td>
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<td>.2 mm</td>
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<td></td>
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CONCLUSIONS

The results suggest that a counterpropagation network can discriminate target echoes as well or better than a dolphin by preprocessing the echoes with a bank of constant-Q filters. However the filters must have a higher Q than the Q of 2.8 for the dolphin. A Q of 5 \((N = 4)\) produced results that enable the network to perform better than the dolphin in the noise-free condition. A Q of 8 \((N = 7)\) was needed when the echoes were mixed with white noise. The Euclidean distance measure indicated that the standard target echoes could be classified if the appropriate threshold value is chosen. Nevertheless, use of a neural network is a simple way of discriminating targets.

REFERENCES

NEW DEVELOPMENT IN ARRAY SIGNAL PROCESSING

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Array processing deals with methods for processing the output data of an array of sensors located at different points in space in a wavefield. Array signal processing plays an important role not only in sonar, but also in radar, geophysics, radio astronomy, biomedical engineering, etc.[1]. In recent years it makes a big progress along with the fast development of many related sciences and technologies[2,3]. In this paper we introduce some advanced topics on array signal processing mainly in underwater acoustic area, and show the improved research results, most of them have been done by authors' research groups.

HIGH–RESOLUTION TECHNIQUE (HRT)

The classical methods in array signal processing are the conventional beamformers which are well used to estimate the direction of arrival (DOA) for a single source. However it is difficult to show the true angles of multiple sources which are located inside the main beam, resulting in loss of resolution. A new sort of methods has been proposed in recent twenty years more and has shown the improved performance over classical methods. In general they are called as high–resolution technique. Capon’s minimum variance estimator [Capon,1969] and linear prediction (or maximum entropy) method[Burg,1967] were the well–known methods. The another set of approaches is based on the subspace orthogonality principle by using the eigen–decomposition of the array output covariance matrix to achieve the promising performance. Typically it contains the Pisarenko harmonic decomposition (PHD) estimator[ Pisarenko ,1973], multiple signal classification (MUSIC) estimator [Schmidt,1981], minimum norm or KT estimator [Kumaresan and Tufts, 1983] and many modifications. Principal component technique is useful to reduce the effect of noise in eigenstructure–based methods. The estimation of signal parameters via rotational invariance techniques (ESPRIT) [Roy etc. ,1986] utilizes an underlying rotational invariance among signal subspace induced by pairwise subsets of an array of sensors. Some advantages for less computation, omitting the search in space, and robustness have been achieved by ESPRIT and other extensions of TLS–ESPRIT [Roy and Kailath, 1989], PRO–ESPRIT [Zoltowski, 1989]. The direction finding in a coherent environment and for wideband sources are also important aspects in array signal processing.

Statistical efficiency study of high resolution methods is being conducted to reveal the essentials and develop the new approaches by many researchers. Finding the new high resolution methods which work well in low signal to noise ratio(SNR) is also an interested problem.

A new modified covariance matrix decomposition method is proposed[4]. It is based on forward–backward smoothing and increase the accuracy of estimated matrix, thereby lowering the resolution threshold for most of the eigenstructure–based methods. Simulation results show ,for example, that for high resolution frequency estimation of two sinusoidal signals with much closed frequencies by using 25 data points, it makes the resolution threshold reducing 10dB, from 12dB down to 2dB.

A new high resolution method –dynamic programming–type algorithm (DPA) has been presented[5]. It is motivated by the dynamic programming theory to obtain the improved high resolution performance not only in frequency estimation , but also in DOA estimation. For two narrow band signals with frequencies spaced to equal the Fourier resolution, the DPA method can decouple these two signals and hence get very accurate frequency / DOA estimation.

High–accuracy DOA / frequency estimation is also an important topic, especially in the direction finding for principal source or estimation of principal frequency. In the multiple beamformer system or digital spectral estimation of digital sequences, estimation errors always exist because the directions corresponding to the maximum of mainlobe in the multiple beamformers or the digital frequencies calculated by computer are all discrete. Differing with the conventional parabolic interpolation method to decrease the discrete errors, a new mapping interpolation method(MIM) is proposed[4]. Simulation results show that for the principal frequency estimation the MIM can get very accurate frequency estimates. When the Hamming window is used, the averaged mean square error (MSE) of frequency estimation can be reduced to 0.08% without noise, and down to 5% even though with SNR = 5dB. The MIM has been extended to the case of DOA estimation for principal source[4]. Similar results are also obtained.
ADAPTIVE PROCESSING (AP)

Adaptive signal processing becomes popular in array processing, especially in the case of unknown or time-variant environment. Adaptive theory and algorithms have been studied intensively.

For active sonar, reverberation is a strong interference. If conventional adaptive notch filter is used it will cancel the reverberation and signal together. An adaptive reverberation canceller (ARC) is formed to reduce the reverberation. Experiment results from a real-time VLSI ARC system show that if the reverberation dominates the background in the input of receiver the output of reverberation will be cancelled down to 40dB more.

Mainlobe constrained adaptive array (MCAA) is very useful in variant underwater environment, particularly in the case of multiple coherent sources. The MCAA with notch pattern transformation (NPT) was presented [6]. The preadaptive transformation is used. The primary input of the following noise cancellation loop is provided by a conventional beamformer which forms a quiescent pattern. The reference inputs generated by the transformation possess notch patterns in which the notch region coincides with the mainlobe of the quiescent pattern. Combining with wideband beamforming, a wideband adaptive array with mainlobe constrained can be implemented by simply using FIR filtering and unconstrained noise cancellation algorithms. It makes adaptive array more feasible in underwater measurement applications.

Accurate phase shift and time delay are key techniques for DSP micro-based beamforming. A constant phase-shift FIR filter for a narrow band centered at a given frequency has been designed successfully by using adaptive modeling technique[7]. For example, a 4-tap digital filter can provide arbitrary phase shift within ±0.04 bandwidth centered at frequency 0.2 with maximum error less than 1°.

ARTIFICIAL NEURAL NETWORK (ANN)

After about 20 years of silence, ANN research has been springing up worldwide. Processing features of ANN different from traditional signal and information processing are: learning by example, distributed and associative memory, abstraction and generalization ability, fault tolerance etc. Obviously ANN is very suitable for array signal processing.

DOA estimation can be conducted by ANN. A special ANN architecture with a Devis type pyramid net for DOA estimation of multiple correlated sources is proposed[8]. It provides very high resolution and accuracy of approaching the CR bound. The network is especially ideal for separating 2 coherent and narrowband plane-wave sources. For a 4-element linear sensor array, as an example, if the input SNR = 0, 14, 20, and 30dB, the ANN can resolve 2 coherent sources within 1/3, 1/5, 1/10, and 1/20 conventional beamwidth respectively. This is remarkable to compare with the other existing high resolution methods.

Wideband noise cancellation by ANN has achieved good results by using a 2-layer back-propagation network[9]. It was shown that the canceller can outperform its adaptive counterpart by 10-20dB noise suppression. The reason for improvement lies in better modeling ability obtained by the hidden layer connections and nonlinear properties of neuron function.

The research results above presented in this short paper are a part of work we have been done. The new development in array signal processing are also showing in many other aspects[2,3], and can be expected that it will be advanced faster in the next decade.

REFERENCES

CLASSIFICATION OF LAKE BOTTOM ECHO SIGNALS USING A WIDEBAND SONAR SYSTEM

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1. Introduction

Many studies are devoted to the detection and the classification of targets lying on the sea floor. A good knowledge of the signal behaviour of bottoms can improve the detection performances in reverberating medium. Relatively little research has been directed to bottom type classification using the acoustic echo signals.

The purpose of this paper is to emphasize the particular interest of bottom type classification realized, by using a wideband constant beamwidth transducer (20-140 kHz). Experimental data give a set of representative observations corresponding to acoustic echo signals backscattered from three different bottom types. Various parameters extracted from each echo signals, its envelope and its power spectrum are chosen. The descriptive analysis is used to study the correlation between variables and to discard redundant information.

A further approach consists of classifying the observation set by means of the linear discriminant analysis. The parameter discrimination phase yields the most pertinent variables for the classification phase. Finally, the classification findings are interpreted and the dominance of few spectral features is confirmed.

2. Collected data and feature extraction

Processed data, corresponding to three different bottom types (silt, pebbles and sand), were collected during an experimental survey in lake Leman. This survey was performed in shallow water and the average distance between sonar and each lake bottom was 11 meters.

The transmitted signal, as indicated in fig. 1 and 2, is wideband and linearly frequency modulated on the whole frequency band of the sonar system. The backscattered signals are digitized through a scope and then stored onto a personal computer.

From a preliminary study, high variability has been observed when various acoustic backscattered signals of the same bottom are processed. It has therefore been necessary to select representative echoes in relation to signal to noise ratio and energy.

A set of sixty observations is selected (twenty for each of the three bottoms) among the whole data. Forty-two parameters are extracted from each signal, its envelope and its power spectrum. Each spectrum has a 120 kHz bandwidth and is splitted into six 20 kHz slices. Features extracted from every 20-140 kHz spectrum are indexed by number one and those extracted from each 20 kHz bands are indexed by number two to seven.

The following parameters are computed from the signal and its envelope.

- \( E_{\text{ne}} \): energy of the signal
- \( M_{\text{ea}} \): mean value of the envelope
- \( S_{\text{de}} \): standard deviation value of the envelope
- \( M_{\text{ax}} \): maximum value of the envelope
- \( E_{\text{nc}} \): energy of the echo computed from the envelope
- \( D_{\text{uc}} \): duration of the echo computed from the envelope

The following parameters are computed from the Nth frequency band of the spectrum, for N varying from one to seven.

- \( E_{\text{en}} \): fraction of energy contained in the Nth frequency band
- \( M_{\text{ax}} \): maximum value of the Nth frequency band
- \( F_{\text{en}} \): frequency of the highest peak of the Nth frequency band
- \( M_{\text{ea}} \): mean value of the Nth frequency band
- \( S_{\text{de}} \): standard deviation value of the Nth frequency band

![Amplitude vs Time](image1.png)

**FIG. 1.** Transmitted chirp signal.

![Amplitude vs Frequency](image2.png)

**FIG. 2.** Power spectrum of the transmitted chirp signal.

3. Data analysis

Before implementing the classification stage, a descriptive approach can reveal additional information about variables. Two correlation matrices are computed from all forty-two parameters and from the ten parameters selected during the feature extraction stage of the linear discriminant analysis (\( S_{\text{de}}, M_{\text{ax}}, M_{\text{ex}}, F_{\text{e}}, E_{\text{ne}}, M_{\text{ax}}, M_{\text{ax}}, F_{\text{e}}, E_{\text{ne}}, E_{\text{ne}} \)). They provide information about correlation between variables and give the features to be discarded.
We feel interesting to implement the principal component analysis in view to visualize the correlation between various variables. The results of this method point out that in the case of forty-two features (respectively ten features), the two first eigenvalues from each of the two correlation matrices contain 78% of the overall information (respectively 87% of the overall information). Thus, the projection of the initial point cluster on the inertia principal plane, made of the two first inertia axis, provides a satisfactory qualitative representation.

In order to realize a bottom classification, the algorithm of linear discriminant analysis has been implemented. During the feature extraction stage, the most pertinent variables are selected by a step-by-step method, which singles out the most discriminating variable of the parameter set at each step.

To make a representative classification with twenty observations per class, it is necessary to interpret the average result of several classifications. The learning procedure provides the discriminant linear function and the test samples are randomly chosen so that they represent 50% of the total number of observations.

The result of one classification is then obtained and the process is applied forty times. The results of the averaging forty classifications are summarized as follows:

<table>
<thead>
<tr>
<th>Bottom types</th>
<th>Well-classified echoes after averaging 40 classifications</th>
<th>Mean value</th>
<th>Standard deviation</th>
</tr>
</thead>
<tbody>
<tr>
<td>silt</td>
<td>85.2%</td>
<td>9.91%</td>
<td></td>
</tr>
<tr>
<td>pebbles</td>
<td>52.8%</td>
<td>13.7%</td>
<td></td>
</tr>
<tr>
<td>sand</td>
<td>93.0%</td>
<td>7.01%</td>
<td></td>
</tr>
</tbody>
</table>

The classification results concerning the pebble bottom may be due to a non-optimal choice of selected features or be, more probably, imputable to experimental problem during the data acquisition.

The histogram of the forty classifications obtained for all classes is shown in fig. 3. The estimated mean value and standard deviation value are 76.7% and 6.25% respectively. The results are sufficiently focused around the mean value, so that they provide a fairly good classification.

4. Conclusion

In spite of a low number of observations, data analysis methods used in this paper show the ability of a discrimination of acoustic bottom echo signals with only a small number of features.

It is noticeable that 80% of the selected parameters are spectral features and that these parameters lead to a fairly good classification. Consequently, we feel that using a wider bandwidth for the transmitted signal will improve classification results.

From the result provided by the feature extraction stage, it has also been observed that the discriminant spectral information is focused on given frequency domain subband and this is really interesting for bottom classification purpose.

In the near future, the methods presented in this paper will be applied to other surveys where insonified bottom types are better known in order to establish a pertinent set of parameters and to improve the classification performances.

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A Neural Network Approach for Passive Source Localization and Tracking in Underwater Acoustics

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ABSTRACT

It is well known that the variance in Passive Source Localization Tracking (PSLT) in underwater acoustics is very large. This paper outlines a method of PSLT for underwater moving object based on neural network. The classical methods, such as moving average and Kalman filter, often give weak results because of unknown noise distribution. Contrast to classical methods, the present method, with neural network (NN) algorithms, can work well to some extent, especially in real environment. The performance measured in sea and comparison to Kalman method is also shown in this paper.

INTRODUCTION

For nearly two decades the PSLT has been a difficult one due to complex conditions in sea. The underwater moving object states, based on the sensor measurement, however, often have untolerant outlier and bias. Traditionally, either the Kalman filters or some other filtering techniques are used to modify the primitive target states. The major drawbacks of these methods are: (1) unaccuracy of models used for describing the states of the underwater moving target; (2) unknown priori knowledge; (3) In the real situation, the farther the moving target from the hydrophones, the bigger variance in time delay signal, which makes primitively estimated range biased and invalid. All these factors result in the poor performance of traditional methods.

The modification of target states can be thought of as finding reasonable elements of trajectory to replace biased one over a short successive time. This is a suitable task for the NN since it has been successfully used in pattern recognition problems[1] and can be used as filters[2]. In this paper, we overcame the drawbacks of traditional methods by using NN and accomplished our modification task quite well.

This paper is organized as follow: in part 1 system is briefly discussed; in part 2 basic formulas are given, in part 3 we report data arrangement and training, and we give test results in part 4; the final part is conclusion.

1 SYSTEM BLOCK

The system block are shown in Figure 1. The inputs of the system are time delay (TD) signals and the outputs are target states including range r(t), elevation, and azimuth angle θ(t). Both Kalman filter (KF) and NN are employed to modify these primitive states. For the system can estimate bearing information relatively well, we use the NN to modify the range information.

![Figure 1. system block](image)

2 ARTIFICIAL NEURAL NETWORK

Extended Back-propagation (BP) model has been used as basic modifical subsystem. This model is similar to the standard BP model except two differences: (1) The transfer function of the output layer neurons is linear rather than sigmoid so that the output is not restricted to (0-1). (2) The learning gain is changeable according to the sign of increments of the weight. The input signals are the data of range obtained from conventional network. The system adopt three layers (1x50x20) and sigmoidal (except for output layer) function as NN's parameters. To be more specific, some formulas are given below:

At time t, the input Nxl received primitive range vector:

\[
v_h = [y_1, y_2, \ldots, y_n]^T
\]

the Nxl vector in the output of hidden layer as:

\[
v_h = [v_{h1}, v_{h2}, \ldots, v_{hn}]^T
\]

where \( v_j = f(\sum_{i=0}^{N-1} w_{ij} y_{i} + \theta_{j}) \), \( w_{ij} \) is the weight from input neuron i to hidden layer neuron j; \( f(x) = 1/(1+\exp(-cx)) \) and c is a slant parameter; \( \theta_j \) is the threshold of neuron j.

In the output layer:

\[
v_o = \sum_{j=1}^{n} v_{hj} w_{oj} + \theta_o
\]

the output \( v_o \) is compared with the desired output the error is fed back through the network, and weights are adjusted to minimize this error.

The increments used in updating of weights, \( \Delta w \) and threshold levels \( \Delta \theta \), can be accomplished by the following rules:
\[ \Delta w_{ij}(n+1) = \eta \delta_j v_i w_{ij}(n) \]

\[ \Delta \theta_j(n+1) = \eta \delta_j + \alpha \theta_j(n) \]

where \( \eta_{ij} \) is the learning gain corresponds with \( \Delta w_{ij} \), \( \alpha \) is the momentum parameter. The error \( \delta_j \) for the output layer is calculated from:

\[ \delta_j = (t_k - v_j) \]

where \( v_j \) is the desired signal of the jth output neuron, and \( \delta_j \) is recursively back propagating the error to lower layers:

\[ \delta_j v_j (1 - v_j) \sum_k \delta_k w_{kj} \]

To speed training \( \eta_{ij} \) is changeable:

1. choose an initial step size \( \eta \) for whole network.
2. as long as the weight gradient has the same sign increase \( \eta \) by adding a small number.
3. if the weight gradient changes sign, reset \( \eta \) to its initial value.

At time \( \eta \), we move input sequence and desired signal one time interval along time axis, in this way the output signal has smaller variance.

3 DATA AND TRAINING

Before putting NN into use, it should be provided the trajectory rules as many as possible so that NN can learn and store these rules itself. According to these principle we give the NN simple rules such as line and circle as the elements of complex trace. The complex element is also given in the training data for example peaks, valleys and roll-offs. The input vectors are obtained by corrupting the desired signal with different noise. Some of them are shown in Figure 2. It is worthy to note that one of aims of using NN is to correct the bias in far range estimation, the training data should be arranged with this purpose. During training we train network with different initial weights to avoid local optima. The trained network with the smallest error is ready to use.

To test our NN has learned enough rules before using it in the system, first, we give it a simulated signal (Figure 3), the response of it is also shown in Figure 3 from which we can see NN has learned the ability of generalization to new case and reduced the bias when target is at far range. Second, we input time delays measured in sea to the conventional network and obtained primitive trajectory (Figure 4). With NN's modification we show the result in the Figure 5. To compare with this result we also give the trajectory of the same target by using Kalman method in Figure 6.

5 CONCLUSION

The neural network method combined pattern recognition and filtering techniques in the underwater moving target tracking problem successfully. It works better than Kalman filter in that it can repair bias and give smoother trace.

References

TRAINS OF TRANSIENTS: WHAT TO EXPECT
AND HOW TO MEASURE SIMILARITY

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INTRODUCTION

This work is concerned with the application of adaptive pattern recognition and neural networks to the field of transient signal recognition. Identification of transients can be an extremely difficult task, due in part to the variability and potentially large number of signals. These properties coupled with intra-class variability render most conventional methods of signal classification unsuitable. An additional consideration is the choice of the similarity measurement. An alternative metric for determining similarity is introduced along with some results of its implementation.

SYSTEM DESCRIPTION

Our approach to the problem of transient classification is a combination of conventional signal processing and neural network techniques which combine into a robust system capable of classifying a wide variety of signals. This methodology is based on a comparison between a temporal trajectory of input features and a predetermined trajectory of prototypical features representing a specific signal class. These prototypes are derived from all available class examples by a self-organization algorithm based on Hausdorff measure (discussed below), and provide a method of representing each class by a subset which effectively spans feature space, affording a substantial reduction in data overhead. The trajectory provides a temporal ordering of prototypes, thus furnishing identifiable signatures for class examples called episodes, which can be compared with input sequences and used to train neural networks. This concept is displayed graphically in figure 1.

![Figure 1: System architecture of neural-net classifier. Early processing is performed on the input data to provide necessary pattern structure. Background effects reduced, followed by extraction of the important features. A bank of neural networks compare the resulting feature space trajectories with those created by class examples. The network with the greatest value exceeding its threshold supplies the decision.](image)

The attributes used for decision making are functions of distances, indicating similarity from the unknown sequences to the prototypical sequences on a point-to-point basis (see fig. 2). These parameters are input to neural networks in order to provide a decision as to whether a class member is present.

![Figure 2: Temporal trajectories through feature space. The x’s symbolize ordered prototypes created during organization, the dots represent a group of ordered unknown patterns from an input sequence, the dotted lines convey the distances, as defined by a suitable metric, from the components of the input sequence to the prototypes at the time of observation. Referring to the dotted lines, sequence (62, 63, 64, 65) is visibly similar to the prototype sequence (1, 2, 4, 3) representing a class example.](image)

For example, with reference to figure 2 let there be 1000 patterns (numbered 1-1000) designating an unknown input containing noise and possibly a short signal sequence from a known class. Prior to run time, a known class has been examined and it has been determined that its four major sequential features are closest (as measured by a metric) to the prototypes #1, #2, #4, and #3, in that order. The goal is to determine if four sequential members from the input data are similar to the ordered prototypes. Suppose observations are collected at times t-3, t-2, t-1, and t, yielding patterns #61, #62, #63, and #64 respectively. Each of the distances from the aforementioned patterns to prototypes #1, #2, #4, and #3 are computed separately at time t, transformed by the similarity function f(t), and sent to the network. It is noted that temporal misalignment will occur because of the uncertainty of a signal’s starting point, therefore, at time (t+1) the distances from input patterns #62, #63, #64, and #65 to prototypes #1, #2, #4, and #3 are now computed, transformed, and sent to the network. During each data acquisition one new pattern is added to the input sequence while the oldest is removed and all of the appropriate similarities are again input to the network.

FEATURE SELECTION

After careful consideration we decided to utilize the STFT for the preprocessing stage. This method was incorporated for two major reasons. (1) It localizes the spectra, thus enhancing short term spectral variation. (2) It is easy to compute.

It is tempting to consider the power spectra as a vector, and by doing so, employ a metric such as a 2-norm in the algorithm above for determining spectral similarity. This is not a good choice for this type of data since vector norms do not consider adjacency of the dimensions, i.e., no dimension is any closer to any other than it is to all the rest. This can be circumvented by employing the Hausdorff metric which has the ability to incorporate adjacency in its measurements, allowing spectrally similar signals to be encompassed within the same cluster. The Hausdorff metric (eqn. 1) is sensitive to deviations from the normal, called singularities, which are useful in determining cluster boundaries.
\[ h(A, B) = \max \{ d^*(A, B), d^*(B, A) \} \]  
(1)

\[ d^*(A, B) = \max \{ d^*(x, B) : x \in A \} \]  
(2)

\[ d^*(x, B) = \min \{ d(x, y) : y \in B \} \]  
(3)

In words, eqn. 3 instructs one to find the distance from a point \( x \) in the set \( A \) to all points \( y \) in set \( B \), retaining the smallest.

Egn. 2 states that this same operation is to be performed for all of the points \( x \) in \( A \) and that the largest of these minima be retained. At this point the same operations are to be performed from the opposite set and the largest value also reserved.

Finally, eqn. 1 says to keep the largest of the two remaining values. This is the Hausdorff distance between the sets, characterized as a method that looks for the greatest dissimilarity between two sets, known as singularities. The measurement \( d \) in the equations above represents some suitable metric. Since straight-line distances on a 2-dimensional plot are being used in this work, a 2-norm is appropriate for the point-to-point measurements.

**NEURAL NETWORKS**

This work utilizes a network proposed by Pao [89a], which performs enhancements to the input patterns before they are passed to higher layers of the network. This configuration is called the Functional-Link Net (FLN), and it can, depending on the circumstances, be utilized as a flat net, obviating the need for hidden layers [Pao 89b] and allowing delta rule training to be employed. There are three primary models for the design of the FLN. Two earlier methods that we have used very successfully in the past are called functional enhancement and joint activation. The third form is called the random vector method, which is a new approach derived from the realization that the lower set of weights, in a network with two layers of weights, are not fixed by a particular set of patterns, but can vary depending on their initial values and the training method. This means that there are several equally suitable sets of weights at this level that can provide useful enhancements to the top layer. With this as motivation we can set the lower weights in a random fashion and simply learn those weights in the upper layer by using the delta rule. This is illustrated in figure 3 where it is shown that the hidden layer of activation functions can effectively be brought down to the input layer. The only stringent rule which needs to be followed is to choose random weights which inhibit saturation of the activation functions.

**EXPERIMENTAL RESULTS**

The class examples described here were generated by computer, and exhibit results similar to those we have achieved on measured data [Pao, Hemminger 91]. The first step in implementing this system is to choose an appropriate set of class examples and form a prototypical set of cluster centers. For this work, ten examples from each class were used for training the networks, along with ten additional examples needed for testing. The training examples from all classes are clustered with a Euclidean based algorithm using the Hausdorff metric. Class episodes are observed and Hausdorff measures accumulated for the purpose of training the neural networks. Prior to training a large collection of noise trajectories are accumulated and utilized to provide the system with a background group to use for counter training (training for non-target outputs). There is a separate network for each signal class. Four signal classes were chosen for this paper, one of which is shown below at its highest SNR (Fig. 4). SNRs for all signals were varied from 0 dB to 10 dB. The results from processing the four signal types are shown in the table of fig. 5.

**CONCLUSION**

In summary, the system described here is based on two major neural network paradigms, unsupervised and supervised learning including an alternative similarity measurement based on Hausdorff distance. This approach has demonstrated robust and reliable qualities in that it is not overly sensitive to parameter changes and has exhibited consistent behavior in its treatment of a wide range of class members. The architecture requires less processing than many current methods and lends itself well to real-time applications.

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KING-RESOLUTION MODE ESTIMATION AND MODAL DOMAIN BEAMFORMING
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INTRODUCTION

Array processing based on the plane wave model is only fit to limitless and homogeneous propagation media. Dispersion should be taken into account when signal propagation in a waveguide is studied, and corresponding array processing must be based on an appropriate propagation model. Array steering and processing as a spatial filtering in a layered waveguide was initially studied as early as the 1980s[12-14]. In the layered medium such as ocean the theory of normal mode is suitable for establishing the model. Recently, matched field processing and matched mode processing (modal beamforming) have received wide attention[15-19], their aim being 2-D or 3-D search for and localization of source in a waveguide. One of studied emphasis is exploring modern high-resolution signal processing methods to this kind of processing and explores possible paths for enhancing adaptability and robustness of processing (relative to incompleteness of aseasomal knowledge)[20].

The article has investigated the parameter estimations of normal mode propagation model and corresponding mode domain beamforming for a horizontal array and compared them with for a vertical array based on the previous papers[9,10]. Its emphasis is placed on the eigenvector method applied to the high-resolution estimates of mode wavenumbers and the mode filtering for a horizontal array. From a view of spatial processing of array, since horizontal array (with a source in its endfire direction) can not provide as best chance of resolving normal mode as vertical array, it is necessary to increase array aperture in order to obtain accurate model parameters such as mode wavenumbers. Generally speaking, to achieve as same results as vertical array, the aperture of horizontal array is at least ten times as long as vertical array.

1 MODE WAVENUMBER ESTIMATION FOR A HORIZONTAL ARRAY

The pressure at element $i$ (depth $z_i$) originating from a source (range $r_s$, depth $z_s$) can be expressed as sum of $N$ modes

$$p(r, z_i) = \sum_{n=1}^{N} \Phi_n(z_i) e^{j\omega_n r_s}$$

(1)

where $\Phi_n(z_i)$ denotes the depth function of $n^{th}$ mode, $r_s$ and $z_s$ denotes horizontal wavenumber and depth of $n^{th}$ mode, respectively. $r_s$ denotes the range from element $i$ to first element (nearest to source in endfire direction of array) (Fig. 1). $N$ denotes number of elements.

The data vector of array is the expression (2) (omitted) i.e.,

$$\mathbf{p}(r, z_i) = \mathbf{D}(r) \mathbf{\Phi}(r, z_i)$$

(3)

Considering noise results in

$$\mathbf{p}(r, z_i) = \mathbf{D}(r) \mathbf{\Phi}(r, z_i) + \mathbf{n}(r, z_i)$$

(4)

Covariance matrix of array data can be written as

$$\mathbf{R}_p = \mathbf{D}(r)^\mathbf{H} \mathbf{D}(r) + \mathbf{R}_n$$

(5)

So far we can see that $\Phi_n(r, z_i)$ denotes $N$ normal modes with different wavenumbers received by first element (at depth $z_s$) and generated by a source (depth $z_s$ and range $r_s$) in a waveguide, its corresponding $NN$-wavenumber matrix being $\mathbf{D}(r)$. This is a counterpart of situation of $N$ independent sources in homogeneous medium field and with different DD's corresponding to $NN$-wavenumber matrix. In the later the various DD's want to be estimated, while the wavenumbers in the former. Eigenvector or signal subspace approach developed recently to the kind of high resolution estimations is successful. Eigenvector decomposition of covariance matrix of array data results in

$$\mathbf{R}_p = \mathbf{A} \mathbf{A}^\mathbf{H}$$

(6)

where $\mathbf{A}$ denotes the diagonal matrix consisting of $N$ eigen-values and $\mathbf{A}$ denotes the matrix consisting of corresponding eigenvectors. $\mathbf{A}^\mathbf{H}$ can be written as

$$\mathbf{A}^\mathbf{H} = \mathbf{A}_0 e^{-j \theta}$$

(7)

assuming that $\mathbf{R}_n = 0$. $\mathbf{A}^\mathbf{H}$ and $\mathbf{A}_0$ has $N$ non-zero eigen-values.

High resolution eigenvector method is based on two subspaces orthogonal each other, signal subspace spanned by the eigenvectors corresponding to the $N$ largest eigenvalues of covariance matrix $\mathbf{R}_p$ of array data and noise subspace spanned by the eigenvectors corresponding to the $N-N$ smallest eigenvalues of $\mathbf{R}_p$. According to this principle, the various mode wavenumbers (phases) can be estimated by linear products of $N$ wavenumber vectors belonging to signal subspace and $N-N$ eigenvectors belonging to noise subspace and corresponding to the smallest eigenvalues. Besides, $N$ eigenvectors belonging to signal subspace and corresponding to the largest eigenvalues, i.e., mode functions (amplitudes) can also be estimated.

Eigenvector beamforming resolving mode wavenumbers can be expressed as follows:

$$\mathbf{F}(\omega_k) = \mathbf{F}(0) \mathbf{E}(\omega_k) \mathbf{A}^{-1} (\mathbf{E}(\omega_k))^{-1}$$

(8)

where

$$\mathbf{F}(\omega_k) = \left[ e^{j \omega_k r_s}, ..., e^{j \omega_k N r_s} \right]$$

$$\mathbf{A}^{-1}$$ is $ NN$-wavenumber $NN N \times NN N$ matrix

$NN N$ denotes $NN$ normal vectors composing wavenumber matrix $A$, and $\omega_k$ denotes $NN$-th eigenvalues and corresponding eigenvector, respectively.

Next section discusses another method of estimating mode amplitude (amplitude and phase).

2 MODE FILTERING AND MODAL BEAMFORMING FOR A HORIZONTAL ARRAY

Expression (1) can be rewritten as

$$\mathbf{p}(r, z_i) = \sum_{n=1}^{N} \Phi_n(z_i) e^{j\omega_n r_s}$$

(9)

Theoretically,

$$\mathbf{p}(r, z_i) = \sum_{n=1}^{N} \Phi_n(z_i) e^{j\omega_n r_s}$$

(10)

Array data vector is expressed in terms of matrices

$$\mathbf{p} = \mathbf{F} \mathbf{\Phi}$$

(11)

Eigenvector algorithm of mode filtering for estimating mode amplitudes $\mathbf{\Phi}$ from array data is

$$\mathbf{\Phi} = \mathbf{F}^{-1} \mathbf{F}^{-1} \mathbf{p} (\mathbf{1} / \mathbf{F}^{-1} \mathbf{F}^{-1} \mathbf{p}) \mathbf{F}^{-1} \mathbf{\Phi}$$

(12)

$\mathbf{p}$
\[ R = \phi(z, \omega) e^{i \omega z} - \phi(z, \omega) e^{i \omega z} \]

\[ F = L \phi(z, \omega) e^{i \omega z} - \phi(z, \omega) e^{i \omega z} \]

(13)

Obviously, the elements in \( \phi \) are complex including magnitudes and phases.

The expression (12) also indicates the mode filtering for a vertical array, but \( F \) should be rewritten as

\[ R = \phi(z, \omega) - \phi(z, \omega) \]

\[ F = L \phi(z, \omega) - \phi(z, \omega) \]

(14)

After data array is transformed into mode space from element space through mode filtering, mode domain beamforming essentially is an extension of classic beamforming in element domain to mode domain. Co-polar or steering vectors are mode functions \( \phi(z, \omega) \) and horizontal wavenumbers \( \omega_k \) deduced from mode programs according to basis cost functions parameters of medium and mutual-correlated with the transformed array data. The beam responses can be written as follows:

\[ B_{\text{beam}}(r, \omega) = \sum_{\omega_k} \left| \phi(z, \omega) \right|^2 \]

(15)

3 COMPUTER SIMULATIONS

We consider the field conditions shown in Fig.2. Range from source to vertical array or first element of horizontal array \( r = 4000 \text{ m} \), source depth \( z = 50 \text{ m} \), number \( N \) of elements of horizontal or vertical array \( = 16 \), separation of adjacent element \( = 5 \text{ m} \), aperture of array \( = 45 \text{ m} \), source frequency \( = 245 \text{ Hz} \). Scatter of eigenvalues of matrix \( F^*F \) for vertical and horizontal arrays is shown in Tab.1 when mode filtering is made by eigenvector method. It is seen from scattering of eigenvalues shown in Tab.1 that rank of \( F^*F \) for horizontal array drops much faster than for vertical array and resolution of decomposition and accuracy of estimation of normal modes for horizontal array are lower than for vertical array under conditions of same element numbers and array aperture. The estimated mode magnitude to theoretical value ratios for two arrays listed in Tab.2. If \( [1.1 \leq 1.0 \leq 1.1] \leq 30 \% \) is regarded to be accurate, then estimated mode numbers using vertical or horizontal array are 10 or 3.

Based on mode filtering, mode domain beamforming has been made according to (15). Beam patterns for horizontal and vertical arrays are shown in Fig.3 and 4. Table.1 Scatters of eigenvalues of \( F^*F \) for horizontal and vertical array (omitted). Search scopes for depth and range are \( 0 \leq 50 \text{ m} \) and \( 1 \text{ km} - 6 \text{ km} \), and separations of adjacent points at depth and range are 1.0 m and 1.5 m, respectively. It is seen from Fig.3 and 4 that horizontal array cannot realize two dimensional localization of source as vertical array can do, because of its poor resolution and accuracy of mode estimation.

We have also simulated mode filtering and model beamforming for horizontal enlarged array, making element number \( N \) constant ( \( N = 16 \) ) and taking array apertures as \( 90 \text{ m}, 150 \text{ m}, 300 \text{ m}, 450 \text{ m}, 600 \text{ m}, 750 \text{ m}, \) respectively. Scatters of eigenvalues of \( F^*F \) and estimated mode magnitude to theoretical value ratios are shown in Tab.3 and 4, respectively, for horizontal arrays with aperture \( 750 \text{ m} \) and \( 300 \text{ m} \). Tab.3 Scattering of eigenvalues of \( F^*F \) for horizontal arrays (omitted). Comparing Tab.3 and 4 with Tab.1 and 2, it is seen that rank of \( F^*F \) gradually increases along with enlargement of horizontal array aperture, and resolution and estimation accuracy of modes are also enhanced. If relative magnitude error of mode estimate \( < 30 \% \) is still taken as criteria, estimated mode number using the horizontal array with aperture \( 300 \text{ m} \) is 9, while using the horizontal array with aperture \( 750 \text{ m} \) comes up to 13.

The beam patterns of mode beamforming for horizontal arrays with enlarged aperture are shown in Fig.5-10. (Fig.5-7, 9-10 are omitted). It is seen from the figures that when the aperture is enlarged up to \( 90 \text{ m} \) there is no peak of sidelobe larger than mainlobe's, and sidelobes are further suppressed along with enlargement of aperture. As comparing with pattern of beam response for vertical array shown in Fig.4, it is also seen that in order to achieve as vertical results as vertical array, the aperture of horizontal array is at least ten times as long as vertical array.

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An Underwater Acoustic Array Signal Processing System (UWASPS)

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[Abstract] In this paper, an underwater acoustic array signal processing system (UWASPS) is presented, which includes two reconfigurable subsystems; a fixed-point array signal processing system (ASP) [1] with 16 TMS320C25 DSP Modules (DSPMs), and a floating-point array signal processing system (FASP) [2] with 10 TMS320C30 DSPMs. Flexible system management and high data processing speed up to 1330 MOPS and 330 MFLOPS are achieved through the multiprocessors-based hierarchical architecture, the multilevel-data-processing-buses organization based on VMEbus, the macro-pipeline modular architecture of the subsystems based on C25 and C30 DSPMs. The advance architecture and high performance makes the UWASPS ideal for optimum space-time and real time signal processing in modern digital sonar system.

1. INTRODUCTION

With the continuing advances in VLSI technology, digital techniques have found widespread applications in the fields of underwater acoustic signal processing. Today, in order to obtain optimum space-time processing, a modern digital sonar system entails self-adaptive match technique and multibeam, multichannel signal processing technique, which demand more powerful digital processors to fulfill real-time signal processing with higher speed and accuracy. Therefore, it is very important to rapidly develop underwater acoustic array signal processing systems by means of advance VLSI technology.

The 2nd and 3rd generation products of TMS320 family have typical architectures adaptable for various digital signal processing algorithms [3, 4]. Their multiple internal buses, powerful instruction set and dedicated devices on chip make them easy to be designed into modular functional units of an array signal processing system.

2. ARCHITECTURE OF UWASPS

As illustrated in Fig. 1, the UWASPS has a hierarchical architecture with three levels: data processing level (DPL), subsystem control level (SCL) and system management level (SML).

The DPL is the lowest level. It includes two parts of high data processing and communication processing speed. One part drives 16 cascaded C25 DSPMs in the ASP and the other 10 cascaded C30 DSPMs in the FASP. All of the DSPMs are configured in a macro pipeline structure. Since time succession behavior of most algorithms, such as wavebeam forming and parameter detection, this architecture is helpful to sufficiently exploit the parallelism of application algorithms.

The middle level (the SCL) is based on MC68010 single computer board in the ASP and on MC68020 in the FASP. The two CPU boards can not only independently manage a D16 VMEbus of the ASP and a D32 VMEbus of the FASP, but also dispatch and schedule jobs of a batched task among DSPMs of the two subsystems, respectively. A DMA/SC module [5], containing a DMA controller (DMAC) and a system controller (SC) specially devised for the ASP, mainly provides the ASP interrupts management and DMA data transfers. Additionally, it also supplies both the ASP and the FASP a mutual data exchange channel via a dual data parallel port on the board.

The SML is the uppermost level based on a host microcomputer or a powerful workstation. The host computer should have two parallel or serial interfaces to communicate directly with the two lower control CPUs. Its rich management proceedings and program resources make it easy for an user to flexible call any algorithm programs, efficiently dispatch any jobs, and reasonably control concurrent process of a batched task.

3. CONFIGURATION OF SUBSYSTEMS AND COMMUNICATION MODE

(1) Fixed-point Subsystem (ASP) [1]

Usually a front-end signal processing unit of a sonar system requires high speed in real time data processing. As one of the most efficient parallel structure of exploiting parallelism of an algorithm, a systolic array structure is reasonably used in the ASP subsystem to meet the demand of modern underwater signal processing. With this advancing architecture, the ASP can operate up to 9.6 MOPS under drive of a global clock of 40MHz.

The organization of DSPMs is shown in Fig. 2. Each DSPM has three C25 DSP chips to form a one-host-and-two-slaves structure. The host C25 exchanges information with the control CPU through its dual-port memory and communication control unit (DM/CC), and accesses the two global memory units (GM1, GM2) of next DSPM via its macro-pipeline communication interface unit (MC1). The two slaves can in parallel access I/O peripherals through their own I/O channel switch units (K1, K2). One ROM on board stores the host C25's initializing program, monitor and control program, and some kernel programs for algorithms. Six data or program local memory units have total capacity ranging from 80K words to 200K words.

(2) Floating-point Subsystem (FASP)

The FASP has an asynchronous data-driven pipeline architecture. It incorporates ten C30 DSPMs, each with its own 640Kbytes of memory and capable of 33. 3MFLOPS, giving the FASP 6400Kbytes and a peak pro-
processing speed of 333 MFLOPS. The communicating speed between every two DSPMs is up to 33.3 MBytes/sec.

Each DSPM has only one C30 DSP chip as the floating-point processing unit and three parallel data interfaces; VSI, FSI and BSI (shown in Fig. 3). The VSI provides C30 data exchange channel with the control CPU on VMEbus, the FSI does C30 with the front-end peripherals and the BSI does C30 with the back-end peripherals. The communication mechanism among the DSPMs is simple and asynchronous. It includes two essential modes; interrupt mode and polling mode. Additionally, there are two standard DSP serial ports of TMS320 family (SPI, SPI1) with data transfer speed up to 8 MBit/sec. Four memory units, ROM, common memory (CM), local memory (LM) and expansion memory (XM), are based on prime-bus (Pbus) and expansion-bus (Xbus) of C30, respectively.

(3) Communication Channels Between Subsystems

Two ways can be used to exchange data blocks between the two subsystems. The one is a dual data transfer channel based on the parallel port of DMA/SC module, working in the single address DMA communication mode of MC68450. The other is based on the D16/D32 VMEbus compatibility of C30 DSPMs when one C30 DSPM is prefixed on the D16 VMEbus in the ASP so that data block can be transferred from its BSI to the FSI of the 1st C30 DSPM of FASP. Since both FSI and BSI are dual parallel data ports, it is possible to transfer data blocks from the FASP to the ASP.

4. CONCLUSION

The UWASPS is newly developed in the Institute of Acoustics, Academia of Sinica. In this paper, we have presented the architectures of UWASPS and its two subsystems, discussed their performance and operation speed. For the modular and flexible structure of DSPMs, the UWASPS can be physically configured as a parallel pipeline oriented SIMD or MIMD system so as to adapt to different parallel structures of algorithms. Both ASP and FASP can also be used independently in various fields of signal processing as required.

REFERENCES


WIDE BAND SONDER FOR FISHERIES
ECHO CLASSIFICATION AT SEA

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INTRODUCTION

We have shown, in previous work, the interest of using wide band systems for fisheries for a better fish size estimation [1], [2], [4] as well as for fish behaviour description [8] using statistical considerations. These previous research and considerations are mainly based on tank experiment results.

One major interest of such wide band soniders is the performance that could be obtained at sea and particularly for species discrimination.

For this purpose, we have developed a wide band sounder system, in collaboration with IFREMER. This system is now regularly operated for sea experiments on known areas. Species could be identified by trawling.

A special purpose system has been developed for data storage and display. Echoes are then processed in the laboratory. We will show preliminary results obtained using a "blind" time-frequency approach. Other approaches using a priori target models are also under investigation.

PROTOTYPE DESCRIPTION

The wide band sounder prototype developed covers a frequency bandwidth ranging from 20 kHz to 80 kHz.

It uses a square transducer (50 cm x 50 cm) of 100 elements divided into 5 square rings (for beam pattern switching).

The bandwidth has been split into 3 sub-bands. The number of rings used for each sub-band has been optimized in order to reduce beam pattern variations with frequency:

- Low Frequency (L.F.) 20 kHz to 40 kHz
- Medium Frequency (M.F.) 40 kHz to 60 kHz
- High Frequency (H.F.) 60 kHz to 80 kHz

As no frequency control of the beam pattern has been used, it may vary in each sub-band. Nevertheless the experimental results obtained on a spherical target show that the global energy angular distribution is quite the same in each sub-range (main lobe at -3 dB): L.F.: 6°, M.F.: 6.5°, H.F.: 5.5°.

The transmitted signal is a chirp signal, the expression of which is:

\[ S(t) = \exp \left( \frac{-\log^2(t/\tau_0)}{2 \log^2(g)} \right) \cos \frac{\log(t/\tau_0)}{\log(g)} \]

where:
- \( \tau_0 \) is the central date (peak of the envelope)
- \( b \) is a parameter defining the signal duration
- \( g \) is a parameter defining the signal bandwidth

This signal has been chosen for future use in processing and classification schemes based on a target model [10].

A real-time digital matched filter is used to separate between signal duration and time resolution [3].

SEA DATA STORAGE AND DISPLAY

The echoes obtained during sea trials are recorded on a multi-track wide band (\(< 300 \) kHz) analog tape recorder. The recorded data are then digitized and processed in the laboratory using a P.C. based multi-channel system.

The tape speed could be reduced (if needed) in order to storage all digitized data on a hard disk. Data are then compacted and stored on optical disks.

In order to establish a correlation between the digitized data and the paper charts obtained during sea cruises, a special purpose software which could display data in both "echogram" mode (figure 1) and "curve" mode (oscilloscope like) has been developed. Any particular ping can thus be selected from a set of pings; then, for this particular ping, any echo can be selected and isolated (by gating) for further processing (figure 2). It is important to note that the system stores and displays the whole echoes and not only the echo envelope as it is commonly achieved for data loggers used in fisheries.

BLIND APPROACH: WIGNER-VILLE DISTRIBUTION

The first approach we followed for echo processing was a blind one that did not use any a priori information on the echo to analyze. It has been shown that the so-called Wigner-Ville distribution [2], [6] is the "blind" time-frequency approach that could be used to look for characteristic patterns in the time-frequency plane [7]:

\[ W_z(t,f) = \int_{-\infty}^{+\infty} x(t+\tau/2) x^*(t-\tau/2) e^{-i2\pi\tau f} d\tau \]

where \( z \) is the analytic signal associated to the signal under investigation.

We did use, in fact, a digitized and smoothed version of this representation, i.e. the Smoothed Pseudo Wigner-Ville Distribution [6]:

\[ SPW_z(t,f) = \sum_{n=N-1}^{N-M} |h(n)|^2 K_M(t,n) e^{-i4\pi f n} \]

with:

\[ K_M(t,n) = \sum_{m=M-1}^{M-1} h(m) x(t+m+n) x^*(t+m-n) \]

where \( h \) is a short-time analyzing window (defined on 2N-1 samples), \( g \) is a smoothing window (defined on 2M-1 samples) and \( f \) is expressed as a reduced frequency \((-1/2c \leq f \leq 1/2)\).\n
Figures 3 and 4 display the time-frequency plots for two typical echoes of isolated fishes of two different species (species A and B). One can note some differences between these two patterns in several domains: time domain, frequency domain or time frequency plane. The common curved pattern is in fact due to the chirp nature of the transmitted signal. Only the variations around this pattern are due to the echo signature.

Other echoes from the same sequence show a similar behaviour.

For both figures, the time-frequency distribution is displayed together with the processed signal envelope and with its power spectrum.

These preliminary results point out the difficulty of finding objective criteria to quantify echo resemblance for classification purposes. The use of a time-frequency approach is very interesting to understand the physical phenomena but does not provide any information reduction (as an echo is replaced by an image). Such a blind approach has to be associated to a priori models for reducing the amount of information needed to describe an echo.

PROCESSING USING A PRIORI MODELS

For an effective information reduction, we will need an a priori echo model that is matched as much as possible to the echo formation mechanisms. Some mathematical models can be very attractive as they may be associated to several extraction algorithms that can be found in the literature [8]; but they can hardly be associated to echo formation mechanisms.

Other models are more matched to this problem. But, in this case, the literature review is not very wide. We have chosen a model based on bright spot structure known as the generalized transversal filter description [10]: the target is
replaced by a set of bright spots reflecting a weighted sum of the signal, its derivatives and integrals. The weights and the time delays are used to describe the echo. Such a model leads to the use of a transmitted signal such as the one transmitted by the sounder prototype [3] it can however be extended to a general case [9]. The sets of parameters can be used for target classification. The echo processing has to be done very carefully as it is faced to a lot of numerical problems. We are now investigating these problems and mainly the sensitivity of echo reconstruction to numerical errors.

CONCLUSION AND PERSPECTIVES

In addition to the advantages already mentioned, the use of wide band systems for fisheries seems to be a promising way to select fish species from spectral signatures. The first results obtained at sea seem promising. The real performance have to be expressed in terms of probability of correct classification and probability of miss. The final performance are strongly connected to the robustness of the parameter extraction algorithms and to the classification schemes that will be used. It is important to note that the work already achieved mainly concerns isolated fishes. It can also be extended to fish schools. As echoes may probably overlap, the processing will preferably be achieved, in this case, after pulse compression as the range resolution will be improved. In this case, the spectral signature will have to be associated to geometrical parameters extracted from the school shape in order to increase the selection performance.

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REFERENCES

BERGEN ECHO INTEGRATOR: A SYSTEM FOR POSTPROCESSING ECHO SOUNDER DATA

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INTRODUCTION

Echo sounding is a basic technique for surveying the water column, especially for determining biological scatterers. Echo integration is a method for quantifying the concentration of scatterers \(<1\). It does this by integrating the intensity of the calibrated, detected echo signal after compensation of this for absorption and geometrical spreading, as though the scatterers were distributed uniformly in the plane transverse to the acoustic axis at each range. Use of a directional transducer is presumed. Typical echo sounder frequencies are in the range from 20 to 500 kHz. Transducer beamwidths are typically in the range from 2 to 10 deg, measured between opposite half-power points of the one-way beam pattern.

Echo integration is thus seen to involve a series of operations on a calibrated echo sounder signal. Since the objective of echo integration is generally quantification of scatterers in the water column, the process can be imagined to be considerably more complicated. In fact, the echo signal must also be classified as to scatterer species and size distribution. Given movement of fish within the water column, including vertical migration and avoidance reactions, decisions about the nature of observed fish cannot properly be taken in advance of surveying. Thus, the process of echo integration, in the wider operational sense, also involves postprocessing of data.

It is clear that a postprocessing system, unlike an echo sounder, must be more than a black box. Important decisions as to scatterer identity and occurrence are expected to depth and salinity distance can only be taken by the trained professional. To aid in this process, the Bergen Echo Integrator has been developed. This has previously been described \(<2,3\>. The aim here, following an overview of the system, is to call attention to a number of its uses and to indicate what extensions might be necessary to prepare the system for these.

SYSTEM OVERVIEW

Preprocessed data

The Bergen Echo Integrator (BEI) expects to receive high-volume acoustic data from an echo sounder preprocessor together with such low-volume auxiliary data as detected bottom depth, time, position, and ship's log, in addition to data from other sensors. The acoustic data consist principally of values of mean volume backscattering strength \( S_v \) in the water column and about the detected bottom. If the transducer is a phased array, the acoustic data may also include values of single-scatterer target strength. The various data are arranged in formatted blocks called datagrams.

Echogram display

When using BEI, \( S_v \)-values are typically presented in echograms consisting of 1000 pings. Each ping is represented by 500 \( S_v \)-values in the water column and 150 generally high-resolution \( S_v \)-values in the region about the bottom as detected by the echo sounder preprocessor. Displayed \( S_v \)-values are color-coded according to each of several schemes. The default scheme, dark-red-light blue, attempts to combine the advantages of the gray scale for shape recognition and of the red-blue color scale for judging echo strength. When interpreting the echogram, it is possible for the operator to indicate occurrences of specific scatterer classes by encircling these on the screen, and to correct erroneous bottom detections by redrawing the bottom line.

Data storage

High-volume, acoustic data are stored in BEI in files. Low-volume data from the echo sounder preprocessor and other sensors are stored directly in the database. The results of data interpretation are also stored in the database.

Hardware

Loading echograms, performing operations on these connected with, for example, interactive thresholding and summation of echo data in subregions of the echogram, and storing the resulting data, while recording new high- and low-volume data that are being continuously sent to BEI, define the minimal hardware capacity. To attain this, while allowing for future expansion of the system, a workstation-level computer, with RISC architecture, is required that meets the following response time constraints: 10 MIPS, 2 MFLOPS, and 2-3 Myt byte/s 10 bandwidth.

Standard software

BEI itself consists of software to be run with the following standard systems: UNIX operating system, C-programming language, X Window System, and INGRES SQL-relational database. As much as possible, BEI has adhered to the principle of using non-proprietary, standards-level software, to ensure transportability without regard to machine manufacture.

BEI processes

A radical departure from traditional echo integrator design is based on a monolithic program run on a CPU, is the organization of BEI as a set of loosely connected processes. These intercommunicate through the message mechanisms implicit in X-Windows. Represented processes include those for (1) Configuration, by which key data are set in the database, (2) Survey grid, to display the location of data collection stations, as for use in selective data to be retrieved and displayed, and (3) Scrutinizing, including control of windows containing echogram data, target strength data, and table of potential scatterer classes, for assignment of echo integral values. These values are computed by integration of the antilogarithm of \( S_v \) over depth, summing this result over series of pings, and normalizing this to a unit area, often one square nautical mile, to determine the absolute physical quantity of area backscattering coefficient \( g_a \) \(<3\>. The same process also supports operator definition of the region of scatterer occurrence by encircling this with a mouse-drawn line, as well as redefinition of the detected bottom line. Completion of the scrutinizing process for each echogram is accompanied by transfer of the resulting \( g_a \)-values to the database. Two further important processes are those of (4) STD station, and (5) Fish station, which organize the respective data for easy retrieval and display during the scrutinizing process, for example.

External communications

All external communications occur over a LAN. The particular form of this is Ethernet, with TCP/IP and UDP/IP protocols.

APPLICATIONS
Fish abundance estimation

The particular use of BEI in this major current application is in measurement of the acoustic density of fish along line transects. The number density of fish is derived by dividing the area backscattering coefficient by the mean backscattering cross section of subject fish species.

Plankton-surveying

Plankton, such as krill, may also be surveyed by BEI, but higher-frequency data may be required from the echo sounder preprocessor because of the known weakness of scattering by such small animals. This use is in a primitive state of development, mainly for lack of knowledge of the scattering properties of plankton and scarcity of sufficiently high-frequency, high-performance echo sounders.

Fish-plankton discrimination

Discrimination of fish and large plankton, at least, may be aided by BEI, as indeed it is. At present, data collected simultaneously or concurrently at different frequencies, say at 38 and 120 kHz, are visually compared through their corresponding echograms. The dramatic increase in small-scatterer target strength, hence echo strength, with frequency will be evident in such a process. Automation, as in computation and display of a difference-signal echogram, is being discussed.

Acoustic classification of fish

This is also eminently feasible with BEI, although not yet attempted. Three methods are mentioned here. (1) Discriminant analysis based on single-frequency data may allow automatic identification of scatterer type by virtue of visual echogram features, such as echo strength and proximity of echoes to surface or bottom <5>. Rather simple algorithms or algorithm classes could be readily investigated by BEI. (2) Wideband acoustic data could be analyzed, as by spectral analysis, but exploitation of wideband sources as currently proposed or under development, as in <6>, will still require fundamental knowledge of the frequency dependence of target strength, which is currently lacking. It will also require development of suitable algorithms for both data display and analysis in BEI, although preliminary work could be performed at multiple frequencies, as these are available from the SIMRAD EK500 echo sounder system, a convenient preprocessor <7>. (3) Pattern recognition algorithms or neural-network machines could be applied to single- or multiple-frequency echo sounder data. Such analyses could be facilitated by BEI, especially by addition of commercially available software or use of other processors attached to the LAN.

Fish target strength

Target strength studies of fish are being undertaken by means of BEI, when the preprocessor provides single-fish data. These could also be extended to plankton, although the problem of ensuring sufficient signal strength may require processing of the original digitized, detected echo signal before Sp-computation by the preprocessor. This is possible with BEI because of provision for storing high-volume data in files.

Fish behavior

Fish behavior is being studied by means of data on target position that accompany single-fish target strength data. E. Oma has developed a target-tracking program that indicates relative or absolute movement of fish in the transducer beam. Studies like this contribute generally to knowledge of fish behavior, including that of avoidance reactions to the passage of a vessel, whether sampling fish by acoustic means or by trawl.

Extinction correction

Correction of underestimates of fish density in highly concentrated and extended fish aggregations can be effected in BEI. An algorithm for this is available <8>.

Bottom-typing

Acoustic classification of the sea bed may be effected easily with BEI because of the possibility of integrating the bottom echo and the bottom-surface-bottom echo. By comparing the energy in the two, a measure of bottom hardness can be derived <9>. Certainly this idea has been proved elsewhere, but without the flexibility afforded by the new system. Still, an investigation must be performed to determine the best algorithm, and, for example, whether this is to operate ping-by-ping or with respect to averages over sets of pings.

SUMMARY

The Bergen Echo Integrator is by design a powerful, flexible, and user-friendly software system for the post-processing of echo sounder data. Applications for this in fisheries and plankton research are many. Some have been realized, others are planned, while some belong to the future.

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USE OF SONAR MODELLING TO CONSIDER THE EFFECT OF FISH MOVEMENT IN FISH ABUNDANCE MEASUREMENT

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INTRODUCTION

For some years estimates of fish abundance have been made using the output of an echo-sounder system. By considering the integrated energy of the echo of a shoal and with knowledge of the sonar parameters and the target strength of the species of fish involved, a crude estimate of the biomass can be made. Some years ago one of the present authors published some work in which a computer model of a typical sonar system was used to try to evaluate the accuracy with which the number of targets could be estimated from the integrated energy [ref. 1].

However the model used was fairly simple and did not take into account some important effects such as the directivity of the target strength of the fish which may have a significant effect on the estimates. The present work is an attempt to modify the computer model by considering these factors and to try to approach more nearly the practical situation.

BRIEF DESCRIPTION OF MODEL

In reference [1] it is shown that the output of a simple echo-sounder can be represented by the convolution of two functions

\[ v(R) = [h(R) \ast T(R)] \]

where \( h(R) \) is the transmitted pulse shape, \( R \) is the range of the ith target and \( T \) is its amplitude taking into account a number of factors which include propagation losses, the target strength of the fish, the sensitivity of the receiver and the position of the target in the beam etc. It should be noted that the output waveform of the echo-sounder receiver is expressed as a function of range rather than time, since the signal is normally displayed using range (depth) as the variable in an echo-sounder.

\( T \) will not in general be constant from pulse to pulse because of the variability of the many factors which contribute to its magnitude. One of the major factors will be due to the movement of the fish itself. Many published papers have shown that the target strength of a fish depends strongly on the orientation of the fish as well as on its size and on the frequency of the acoustic signal. In the earlier modelling work referred to, it was assumed that the target strength of the fish varied according to an arbitrary probability distribution and for simplicity the Gaussian distribution was chosen. This paper takes a closer look at this choice.

EFFECT OF FISH MOVEMENT

The orientation of the fish will include tilt, roll and pitch, but when the fish is viewed in the dorsal aspect it has been shown that the significant factor in the determination of the target strength is the tilt angle [ref. 2]. In ref. [3] it is shown that the shape of the curve relating the target strength to tilt angle varies considerably with species. Figure 1 shows some typical measurements due to Foote and it can be seen, as might be expected, that there is a strong inverse correlation between the width of the main lobe and the length of the fish. In fact the curves are not unlike those that might be expected from a crude approximation of a fish as a reflector of an appropriate length.

In order to make an estimate of the effect of this variation of target strength with tilt angle, it is necessary also to know, or to assume, the probability distribution that represents the way in which the tilt angle will vary in practice. In reference [3] and [4] measurements are reported on different species and it is shown that this distribution is close to normal. Using these two functions it is then possible to predict the probability distribution for the variation of the target strength and hence to calculate the effect of this probability distribution on the estimation of biomass.

In our model we have assumed, again for simplicity, either a uniform or a normal distribution for the probability distribution of tilt angle. For the variation of target strength with tilt angle we have used one of the two functions shown in figure 2. These curves are not unlike the measured curves and have been chosen to represent two substantially different fish lengths. Obviously the result will also be dependent on the assumed mean tilt angle. In figure 3 we see the various probability distributions calculated from the above assumptions. They are shown in four groups one for each combination of the two directivity functions and the two probability distributions of angle, and for each case four different mean angles of tilt. It is fairly clear from the nature of the problem that the probability dis-
tribution of target strength of the fish derived will not be symmetrical and indeed we see in figure 3 that this is so. When the mean angle of tilt of the fish is not zero then the possibility of a double moded distribution arises and this can be seen clearly in the figures.

Using these distributions it is possible to calculate the mean and variance of the target strength of the fish and in figure 4 we see some typical results and their dependence on the various parameters.

It is also possible to use these results in the simulation of the estimate of fish abundance as was done in the original work but with this better approximation to the target strength variation of the fish. Basically a random distribution of fish is assumed and the sonar model determines the probable sonar trace that will be obtained. From this the number of fish echoes can be obtained using the standard methods and the result can be compared with the original number simulated. Such an experiment can be repeated many times on a computer giving plots like those shown in figure 5. Ideally we would like these results to lie on a straight line with a slope of unity but it can be seen that there is considerable scatter and that the scatter increases the greater the movement of the fish.

CONCLUSIONS

This paper summarises part of the work which was contained in an M. Phil. thesis written by one of the authors. It shows that the simulation method can help considerably in determining the effects of the variation of different parameters in the problem of fish abundance measurement by acoustic means.

REFERENCES

INTRATHERMOCLINE LENSES AND THEIR EFFECT ON ACOUSTIC FIELDS IN THE OCEAN.

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In the last decade very intense intrathermocline eddy lenses have been discovered in the ocean. Their hydrological characteristics strongly differ from those typical for surrounding waters. So, for a lens found in 1989 in the Canarian basin this difference in temperature at depths of 900-1000 meters makes up to 4.6°C for the temperature, 1.1 ‰ for salinity, and 18.6 m/s for sound speed. It seems to be the most great anomalies of all oceanological characteristics of lenses. Now more than 300 lenses have been found in very different geographic zones from the tropical zones to the Arctic and Antarctic regions, at depths from 50 down to 4000 meters. Parameters of lenses fall inside a rather wide range — the diameter of 1-100 km or even more, vertical size from tens of meters up to several kilometers, eddy current velocity up to tens cm/s, and lifetime of 1-16 years depending on sizes thereof and the other conditions. Field observations show that the most pronounced lenses are generated at boundaries of different water masses. The typical example is intrusion of the Mediterranean waters at intermediate depths into the Atlantic ocean in the vicinity of the Gibraltar strait. Lenses of this kind ("meddies") have been studied best of all.

MECHANISM OF LENSES GENERATION.

One of the most probable mechanisms of generation of anticyclonic lenses dominating in the ocean is associated with the Coriolis force effect on spreading waters. It allows a very simple physical interpretation. Consider a water volume of increased density invaded into the stratified ocean at some depth. Having lowered down to the depth corresponding to its density, this water mass begins to spread in lateral directions. Under the action of the Coriolis force the water mass considered turns to the right side (in the northern hemisphere) and eventually gives rise to anticyclonic circulation around its core. As a result, spreading waters form the quasigeostrophic current. Neglecting by the friction and centrifugal force, we obtain

\[ v(z) = \frac{\rho f}{\partial \rho / \partial \sigma}, \]

where \( v(z) \) is the horizontal current velocity, \( f = 2qg \sin \phi \) is the Coriolis parameter, \( \phi \) is Earth's angular frequency, \( g \) is the geographic latitude, \( \rho \) is the density of the radic pressure gradient. For the estimation of the magnitude of \( \partial \rho / \partial \sigma \) we suppose that the buoyancy frequency \( N \) of surrounding waters is a constant. In this case it can be shown that the pressure difference \( \Delta P \) at the low boundary \( z = 0 \) of the water mass of thickness \( h \) and at the same depth in surrounding waters is

\[ \Delta P = \rho h N^2/2 \]

Then \( \Delta P/\sigma r \) can be estimated as \( \Delta P/L \), where \( L \) is horizontal width of the geostrophic current zone. From (1) we have

\[ v(z) = \frac{(\rho h)^2/(2\sigma)}{L}. \]

On the other side, the current velocity \( v_s \) of spreading waters can be estimated as the ratio \( L / R_s \), where \( R_s \) is any characteristic temporal scale, \( \tau = f/L \) and, therefore, \( v_s = fL \). Since the Coriolis force changes only the current direction but not its velocity magnitude, then \( v = v_s \). It gives

\[ L = \frac{N h^2}{2} = \frac{R_s}{\sqrt{\tau}}. \]

where \( R_s = N h/\tau \) is the local Rossby radius determined by the value of lens thickness \( h \) and the buoyancy frequency of surrounding waters. The expression (3) gives the zone width \( L \) of geostrophic current generated under the action of the Coriolis force around limited volume of the water mass invaded into the stratified ocean. It is evident, that the geostrophic current of the same width will be generated at the edge of a water flow running out through the Gibraltar strait into the Atlantic ocean.

Due to baroclinic instability eddies (lenses) of radius \( L \) are took off from the edge of this geostrophic current. This conclusion is confirmed by results of comparison of values of anticyclonic lenses radius \( R \), with local Rossby radius. The data for about 20 lenses (the most of them are meddies) observed in the northern Atlantic for which both parameters \( h \) and \( R \) were measured, have been analyzed. It is shown that near the Gibraltar strait the spread of Burger number values \( \text{Bus}(R/R_s)^2 \), is rather small. The ratio \( R/R_s \) varies from 0.3 up to 0.8 with the mean value of 0.7. While drifting in the anticyclonic gyre of waters in the northern Atlantic lenses undergone some evolution. They became more "flat" due to viscosity effect and energy loss to generation of internal waves. The mean ratio between their thickness \( h \) and diameter varies in a wide range from 10^-3 up to 10^-2. The ratio \( R/R_s \) changes also. At first it increases up to 0.9 at a range 2000 miles from a place of their formation, and then it decreases down to 0.1-0.2 at a range of 5000-6000 miles, where lenses are not intense because they are already at a final stage of their existence. Considerations developed above allow to look on the fine structure in the stratified ocean from a new point of view. Let us estimate thickness of a
water volume for which it is possible to neglect by the viscosity effect. This effect is small as compared with that of Coriolis force if the Ekman number $E=2u/\tau^2$ is small relative to unity. Here $u$ is the viscosity of water. Assuming $E=0.1$ and $v=1$ cm/s, $f=10^{-4}/s$, we obtain $h>4.5$ m. Thus, if thickness of a protruding water volume in the stratified ocean exceeds some meters, lenses are formed which can exist for a long time. Such intrusions of diverse waters are frequently occurred in the ocean. Therefore, it can be suggested that a considerable part of the ocean stratification is formed by anticyclonic eddies-lenses of various scales. For the smaller thicknesses of the lens, viscosity effects begin to play a dominating role, but the Coriolis force influence decreases. As a consequence, thin intrusions of diverse waters do not live for a long time but spread to great distances transforming into very thin sheets.

**ACOUSTIC EFFECTS OF LENSES**

The disturbances of a sound field due to lenses result in changing of spatial structure of the acoustic field. Below some results of computer calculations of the acoustic intensity are discussed. The normal modes and ray methods have been used for the acoustic field calculation in range-dependent ocean. Calculations of the horizontal wave numbers are carried out in the adiabatic approximation, and those of amplitude of modes are made taking into account, in part, the interaction between modes. As a rule, coupling of each mode with a group of 6-8 neighboring modes are taken into account. The energy exchange for modes are calculated by standard method. Along the propagation path a number of modes effectively generated vary within a zone occupied by lens. Their amplitudes are found on the basis of assumption that the energy flux within this group of modes is constant. The spatial structure of the acoustic field is calculated by ray method. The main results have been obtained for two meddies. The first one was discovered in tropical part of the North Atlantic(1965). The maximum excess of the sound speed at the anti-channel axis of the lens is equal to about 16.5 m/s. The antichannel due the lens encompasses a range of depths 850-1200 m, its horizontal extent is 60 km. The second lens was found in Canarian basin, and its characteristics were described formerly. Computer results show that acoustic effects of lens is quite different in the case when there is zonal structure of acoustic field and when it is absent. The most striking effect provided that the zonal structure exist is insonification of the shadow zones. The ray diagram usually differs from that which is typical for the stratified ocean. The other effect associated with lenses is the displacement of the convergence zones boundaries. In the presence of lenses the convergence zones occur at more far from the sound source than in the case of the stratified ocean. The greater the zone number is, the greater the displacement of its boundaries. The displacement of the sixth zone is equal to 7-8 km. It should be noted that the sound intensity level is approximately the same whether a lens exist or not. It is accounted for by modes of the first numbers propagating under small grazing angles are refracted more strongly and leak from the lens zone. In this particular case leaking modes have relatively small excitation coefficients and, therefore, their contribution to the total acoustic field is negligible. One else effect lies in non-monotonic changing of the distance between convergence zones and sound source when the latter moves apart. There are some spatial intervals where this distance varies the most quickly. Situation will be quite different if a sound source is near the axis of the underwater sound channel in the stratified ocean and the zonal structure is absent. In this case a very deep new shadow zone is formed within a region occupied by lens. For parameters of the lens considered the transmission loss is about 25 dB. The reason for such strong decreasing of the sound intensity level in this shadow zone is that modes with high excitation coefficients due to refraction by the lens go out from it.

**EXPERIMENTAL STUDIES**

Measurements of the acoustic field in the presence of the lens discovered in the Canarian basin have been carried out. Experimental data obtained show good agreement with theoretical predictions. In the shadow zone a small scale fine structure of the acoustic field was discovered. The experiments carried out with very high sensitive deep water vertical receiving array (32 omnidirectional transducers) reveal the influence of the lens on the characteristics of the dynamic ambient noise in the ocean. The measurements were made at three depth 0.6, 1.6 and 1.0 km, i.e. above the lens, under it and in the core of the lens, in the frequency range from 0.5 to 3 kHz. At the depth 1 km the angular half width of the refraction minimum decreases to 3° (instead of 11° for the stratified ocean) due to the lens. As a result, the mean level of acoustic field within the main lobe of the receiving array increases. At the depths 0.6 and 1.6 km the angular half width changes only slightly and the acoustic field level is approximately the same as in the absence of the lens.
THE WKBZ ADIABATIC MODE APPROACH TO SOUND PROPAGATION IN GRADUALLY RANGE-DEPENDENT CHANNELS

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INTRODUCTION

In most regions of the ocean the vertical gradient of sound velocity is about a thousand times the horizontal one, so the ocean is often considered as a plane-stratified medium\(^1\). However, in the area of the convergence of cold and warm currents or in the regions with synoptic eddies the horizontal variation of sound velocity cannot be neglected. Even in the area of the ocean with a little horizontal variation the influence of the horizontal variation on the long-distance propagation must be taken into account due to accumulative effect.

Many methods for calculating the acoustic field in the ocean with horizontal variation have been developed, such as parabolic equation method\(^2\), horizontal ray theory\(^3\), Gaussian beam approach\(^4\) and so on. A WKBZ mode approach is proposed in Ref.5, and it can be used for fast and accurate calculation of the field of waveguide modes in a horizontal channel. In this paper, a WKBZ adiabatic mode approach is developed, and the approach may be used for gradually range-dependent channels and it takes the bottom interaction into account.

WKBZ ADIABATIC THEORY

Assume that the ocean medium varies with horizontal range so gradually that the mode coupling may be neglected, resulting in the "adiabatic approximation". Under the condition of adiabatic approximation, the acoustic field of a harmonic point source may be expressed as\(^5\)

\[
P(r,z,x_0) = \frac{\mathbf{v}_0}{r^{\nu_0}} \sum \psi(x_0,0) \psi(x_0)^* \left( \frac{v_0}{r} \right) \exp \left( i \phi_0 (x_0) \right)
\]

where \(v_0(0)\) and \(v_0(r)\) are the "local" eigenvalues of modes, \(\psi(x_0,0)\) and \(\psi(x_0)^*\) are the "local" eigenfunctions of modes, respectively at the source and receiver positions.

For definiteness, we suppose that the velocity near the sea-bottom is greater than that near the surface. In a ocean channel there are two types of normal modes: the waveguide modes and bottom-reflection modes. For the waveguide modes, the eigenrays do not touch the sea-bottom, and when the seawater absorption is not taken into account, the eigenvalues are real and determined by the following equation

\[
\int_{\frac{N}{2}}^{\frac{N}{2}} \left( k_0 (x) - \mu_0 (x) \right) dy = - \frac{\pi}{2} + \pi \frac{\theta_1}{2} + \pi \frac{\theta_2}{2}, \quad l = 0,1, \ldots, L
\]

where \(\phi_1\) and \(\phi_2\) are the phase shifts at upper and lower turning (or reflecting) depths \(\eta_1\) and \(\eta_2\), respectively. For the bottom-reflection modes, the eigenrays reflect from the sea-bottom, and the eigenvalues are generally complex, i.e., \(\nu_0 = \mu_0 + i \beta_0\). The horizontal wavenumber \(\mu_0\) and attenuation coefficient \(\beta_0\) are determined by

\[
\int_{\frac{N}{2}}^{\frac{N}{2}} k_0 (x) - \mu_0 (x) dy = (l + \frac{1}{2}) \pi + \pi \frac{\theta_1}{2},
\]

\(l = L, L+1, \ldots\)

where \(\nu_0\) is the bottom-reflection coefficient, \(S_1\) is the cycle distance of an eigenray.

In this paper, the WKBZ approximation\(^6\) is used to calculate the "local" eigenfunctions of modes. The characteristics of long-range propagation, especially the fields in convergence zones are mainly determined by the waveguide modes, while the fields in the deep shadow zones are mainly contributed by the bottom-reflection modes. Since there are a great number of bottom-reflection modes whose interference distances are much shorter, the field of bottom-reflection modes is suitable for smooth-averaging\(^7\). We then express the whole intensity as

\[
I = I_w + I_s
\]

where \(I_w\) is the intensity of waveguide modes by coherent superposition, \(I_s\) is the smooth-averaged intensity of bottom-reflection modes.

NUMERICAL SIMULATION AND EXPERIMENTAL RESULTS IN THE PHILIPPINE SEA

We apply the WKBZ adiabatic mode approach to the ocean channel in the Philippines sea. Figure 1 shows the velocity profiles\(^8\), in which the curves 1 and 2 are the profiles measured at 0 and 250 km, respectively. It can be seen that the two profiles have obviously difference above the depth of 500 m.

For the frequency of 109 Hz, the source depth of 100 m and receiver depth of 107 m, ignoring the horizontal variation of velocity profile, the transmission-loss curves of waveguide modes calculated according to the profile 1 and 2 are shown in Fig.2(a) and (b), respectively. Taking account of the horizontal variation from profile 1 to 2, the transmission-loss curves corresponding to the intensity \(I_w\) of waveguide modes and the intensity \(I_s\) of bottom-reflection modes are shown in Fig.2(c), where the solid and dashed curves denote \(I_w\) and \(I_s\), respectively. It can be seen from Fig.3 that the horizontal variation of velocity profile affects the positions and forms of convergence zones obviously.

In Fig.3 is shown the comparison between the calculated and measured results, where the curve in Fig.3(a) is the calculated one including \(I_w\) and \(I_s\), and the curve in
Fig. 3(b) is the measured one from the research report of Akulichev [6]. It may be seen from Fig. 3 that the positions and forms of convergence zones calculated and measured are coincident well, and the fields in the deep shallow zones less than 100 km are also consistent. For the fields in the deep shadow zones greater than 100 km, the measured curve is higher than the calculated one due to the noise background.

References

THE EFFECTS OF INTERNAL SOLITONS OR BOTTOM RELIEF ON SOUND PROPAGATION IN SHALLOW WATER

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Introduction

Strong bottom interaction, variable media and multipath propagation make shallow water acoustics very challenging. One of unexplained phenomenon for shallow water sound propagation, in the contrast with the so-called the optimum propagation frequency, is that high acoustic propagation loss over some frequency range is frequently reported\(^1\) - 3. Weston attributes such a phenomenon to fish activities\(^1\). In some case it is associated with sediment shear-wave resonances within the layer of sediment\(^2\). Fig. 1 shows a frequency response for shallow-water sound propagation in the summer (with a strong thermocline shown in Fig. 2) in the Yellow Sea off China, obtained by Zhou and his group at the IAAS\(^3\). Around 600 Hz and above 1100 Hz the transmission loss is abnormally large. Several years of observations, all made in August at the same area with a similar strong thermocline, have shown that the frequency responses of sound propagation is often a strong function of both time and propagation direction. The variability of the high loss frequency range and the apparent anisotropic characteristics of sound propagation occur with no obvious explanation in terms of fish swim bladder resonance. For the strong thermocline shown in Fig. 2, if both sound source and receiver are located beneath the thermocline, the sea surface influence on the long-range sound field (or lower modes) is negligible. While it might be possible that one could use the Weston fish model (or some other model) to quantitatively match the experimental data, we consider the possibility that the acoustic normal-mode conversion caused by internal soliton or bottom relief packets may also explain the data.

Ocean models

1. Internal soliton packets. Our numerical results have shown\(^5\) that the acoustic mode-coupling induced by internal wave packets with a gated sine waveform (model I) exhibits frequency, soliton wavelength and packet length resonances. The interaction between the acoustic waves and internal wave packets could provide a plausible explanation for the observed anomalous propagation in the summer. Now we extend the model to consider the influence of the characteristic properties of more realistic oceanic soliton packets.

Solitons in shallow water (\(h/\lambda_s << 1\)) are described to first order in wave amplitude by the Korteweg-de Vries (KdV) equation: \(A_t + C_kA_x + \mu A_A_x + \delta A_{xx} = 0\). A solution to this equation is \(A(x - ct) = \text{sech}^2(\frac{\alpha x}{2})\). The density (temperature) profiles corresponded to Fig. 2 is very reasonably described by a two-layer fluid model. The soliton phase speed \(c\) and wavelength \(\Lambda\) (and scale parameter \(L\)) are then easy to obtain. When the initial solitary wave propagates, it often evolves more regular solitons and exhibits clear nonlinear features. Instead of the previous simple gated sine function, three packets with typical characteristics of internal solitons are used: the classical \(\text{sech}^2\) soliton shape which decreases in wavelength and amplitude toward the rear of the packet (Model II, shown in Fig. 2).

2. Bottom relief group. For simplicity, we assume that an undulating seabed can be expressed by a half sine function, shown in Fig. 2 (Model III).

Approach

The presence of internal soliton or bottom relief packets makes the environmental parameters range dependent. The parabolic equation (PE) model is used to numerically simulate the effect of internal soliton packets or bottom relief.
groups on sound propagation in shallow water. We arbitrarily put three packets of solitons or bottom relief located at 5 km, 15 km and 25 km along propagation track. The PE method is used to calculate the frequency response of acoustic transmission loss. When analyzing the characteristics of acoustic mode-coupling, we put just a single group of solitons or bottom relief at 15 km and use the first normal mode alone as the initial input field to the PE code (IFD). The acoustic field obtained using PE model is decomposed into the normal modes. The results are compared with mode-coupling theory in order to improve our understanding of how mode conversion acts to become an important mechanism of acoustic attenuation in shallow water.

Results and discussion

1. Similar to model I and model III, internal wave packets with the classic characteristic properties of soliton could also cause abnormally large attenuation for acoustic propagation in shallow water. The differences between the results with and without the soliton packets is shown in Fig. 1 by * * * * .

2. The interaction of sound waves with soliton or bottom relief group exhibits (signal) frequency, (radial inhomogeneity) wavelength and packet length resonances. As an example, the bottom relief wavelength and group length resonances are shown in Fig. 3 \( f = 630Hz, r = 30km \).

3. For a shallow water with a strong thermocline the apparent resonance interaction of sound waves with three ocean models could only occur under a specific circumstance: when the acoustic mode coupling caused by environmental parameter variation transfers a significant amount energy from lower mode into higher-order modes with much larger attenuation rate. For example, the PE field at 18 km is decomposed to normal modes with the results shown in Fig. 4 \( f = 630Hz, r = 18km \). Apparently, after the interaction of the first mode with a packet at 15 km a significant amount energy has been coupled into higher-order modes.

4. The "resonancelike" behavior of transmission loss predicted by the PE analysis is consistent with mode coupling theory. Significant energy transfer will occur between mode \( m \) and \( n \) if \( k_{\text{coh}} \approx k_m - k_n \), here \( k_{\text{coh}} \) is the wave number of the spectrum peak of the radial inhomogeneity. The mode coupling is a periodic (resonance) function of the soliton (or relief) group length.

5. The main intention of this work is to show what would happen to sound propagation if soliton or bottom relief packets are present, not to solve an inversion problem for internal waves or to determine sound propagation for a specific area. It is expected that any abnormal attenuation for sound propagation in a real shallow water environment might be caused by any combination of several different mechanisms, internal wave/or bottom relief) interaction must be considered as one such mechanism.

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THREE-DIMENSIONAL ACOUSTIC RADIATION FROM A VOLUME ARRAY AND PROPAGATION IN LATERALLY STRATIFIED MEDIA WITH IMPROVED NUMERICAL EFFICIENCY BASED ON GLOBAL MATRIX METHOD

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Introduction
The propagation and radiation in the laterally stratified media has been a classical problem in the analysis and prediction of ocean acoustic pressure field. Among various methods available nowadays, Global matrix method\[1,3\] gives the full wave solution for acoustic waves as well as elastic waves of P, SV and SH, so as to treat ocean bottom interactions. Another advantage comes from its ease of handling the directional sources such as couplets\[4\], seismic sources\[4,5\] and line arrays\[2,3,5\].

The radiation from a horizontal array of sources has been treated in the paper by Schmidt and Glattet\[4\]. However, the method is computationally intensive due to the higher orders of Bessel functions required to expand the range-direction field for each Fourier orders in azimuthal angle. In this article, a superposition method, which eliminates Fourier expansion in azimuthal angle, is discussed and numerical examples will be given with interpretations.

Theory
The superposition method, literally, utilizes the superposition of radiated field by a single source placed on axis to represent the field caused by a source displaced from the z-axis (i.e. r=0 in cylindrical coord.). Figure(1) shows the horizontal array with sources off the z-axis. The field caused by this array can be represented by summation of the field in the receiver positions in Figure(2). For the horizontal displacements and shear stresses, the field parameters need to be transformed in the proper coordinate system, which can be either rectangular or cylindrical. Hence, the directivity caused by the horizontal line array can treat the horizontally or vertically polarized shear waves.

![Figure 1: Array of horizontally distributed sources.](image)

![Figure 2: Reduced numerical Model.](image)

Numerical Examples
Two examples are discussed to demonstrate the applicability of the superposition method. The first example is the case of a vertical line array, which can be solved by 2-dimensional version of Global Matrix Method\[1\]. The layer consists of 3 layers. The uppermost layer is a fluid half space with sound speed 1500 m/sec. The second layer is the ocean bottom with compressional wave speed of 1600 m/sec. and shear wave speed of 400 m/sec., respectively. The third and last layer is the sub-bottom half space with compressional wave speed 1800 m/sec., and shear wave speed 600 m/sec., respectively. Figure(3) is the result by 2-dimensional version for the source array of 41 elements centered at 50 m steered to 25° downward and
reflected from the ocean bottom. The y-axis is the depth axis from 50 m to 125 m. The x-axis is the range from 0 m to 300 m. Figure(4) is the result by 3-dimensional version using superposition method for the same geometric and source configurations. Figure(3) and Figure(4) show good agreement.

The second example is the case of a horizontal array which can only be treated in 3-dimensional version. Figure(5) shows the acoustic field from a horizontal array with 41 elements of 0.75 m apart steered to 65° (grazing angle 25°) in the vertical direction. This steering direction has been used to observe the same radiation direction as the previous example. As expected, the beam is broader than that of the first example, since the beam steering direction is toward the end-fire. The strong field in the near range is considered to be caused by interpolation error due to spatial sampling, which is required for field parameter transformation. This problem remains to be further studied and resolved.

Summary
The superposition method is shown to be numerically efficient without convergence problem. Also, the method describes the 3-dimensional characteristics of the elastic waves accordingly.

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SPATIAL COHERENCE IN MULTIPATH UNDERWATER ACOUSTIC CHANNELS

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INTRODUCTION

Whether the sound wave propagates in the deep sea or in shallow water, there are always many paths from the source to receiver because of the refraction caused by the inhomogeneity of sea—water and the reflection from boundaries. Therefore, the multipath transmission is an important characteristic of underwater acoustic channels[1,2].

The multipath transmission causes the waveform distortion and time spread of signals, and then results in coherence loss of acoustic field[3,4]. For a moving source or receiver, since the signals through different paths have different Doppler frequency—shifts, the multipath propagation makes the time and frequency coherence decrease[5,6]. As the phase and group velocities for different paths are different, the multipath propagation may cause severe bearing error for a time—delay compensated receiver array[4,7].

In general, the horizontal transverse correlation of acoustic field is mainly determined by the random fluctuation of medium, the vertical correlation is chiefly determined by the multipath transmission, and the horizontal longitudinal correlation is dependent on both the causes mentioned—above. This paper mainly discusses the vertical and longitudinal correlations due to the multipath transmission.

SPATIAL COHERENCE THEORY

Assuming that the ocean channel is a horizontally stratified medium, the acoustic field of a harmonic point source may be approximately expressed as

\[ P(x_1,x_2,t) = \sqrt{8 \pi} \sum x_i \Psi(x_i) \overline{\psi(x_i)} \sqrt{\mu_i} \exp(i \mu_i x + n \pi/4), \]

where \( x_i \) and \( \mu_i \) are the source and receiver depths, \( \beta_i \) are the real and imaginary parts of a mode eigenvalue respectively, and \( \psi(x) \) is the eigenfunction. In this paper, the generalized phase—integral (WKBJ) approximation[8] is used to calculate the eigenfunction, i.e.,

\[ \Psi(x) = \left( \frac{-1}{\sqrt{2 \pi}} \int_{-\infty}^{\infty} e^{i \mu z} \exp \left( \int_0^z \sqrt{\mu^2 - k^2(y)} \, dy \right) \right) \left( \frac{\sqrt{2 \pi \text{erf}(z)}}{\sqrt{\text{erf}(z)}} \right)^{1/2}, \]

where \( k(z) = \omega / c(z), E = 0.875, b(z) = |d k^2(z) / dz|, S_i \) is the cycle distance of an eigenray, and \( \eta_i \) and \( \zeta_i \) are the upper and lower turning depths of an eigenray[9,10].

The spatial correlation \( \Gamma(x_1,x_2,r,d,l) \) and normalized correlation coefficient \( \gamma(x_1,x_2,r,d,l) \) of acoustic field respectively are defined as

\[ \Gamma(x_1,x_2,r,d,l) = \frac{\int P(x_1,x_2,t) P^*(x_1,x_2 + d + l) \, dt}{\left( \int |P(x_1,x_2,t)|^2 \, dt \right)^{1/2}} \]

where \( d \) and \( l \) are the vertical and longitudinal separations between two receivers respectively, and the overbar \( \overline{\cdots} \) denotes the range and depth—average[11].

Substituting Eqs.(1) and (2) into Eqs.(3) and (4), one get the integral expressions of spatial correlation and correlation coefficient as follows:

\[ \Gamma(x_1,x_2,r,d,l) = \overline{\int G(x) \cos(k_1 \sinh(r / n \pi)) \exp(-x_1^2 / 2 \sigma_x^2) \, dx_1} \]

where \( a_\omega = \cos^{-1}(\max(k_1,k_2) / k_3) \) is the grazing angle at the depth with the minimum velocity, \( a_1 \) and \( a_2 \) are the grazing angles at the source and receiver depths, respectively. The function \( G(x) \) is given by

\[ G(x) = \exp(-2k_x^2 \frac{x^2}{2 \sigma_x^2} - 2(x_1^2 + x_2^2 + x_3^2)), \]

where \( D(x;z) = 0.875 \left( \frac{1}{n \pi} \right) \frac{|\exp(z_1^2 / 2 \sigma_x^2) \exp(-2(x_1^2 + x_2^2 + x_3^2))|^{1/2}}{|\exp(z_2^2 / 2 \sigma_x^2) \exp(-2(x_1^2 + x_2^2 + x_3^2))|^{1/2}} \)

NUMERICAL SIMULATIONS AND EXPERIMENTAL RESULTS

We have used Eq.(6) to make a great number of numerical simulations, Fig. 1 is an example. In Fig.1 are shown the vertical and longitudinal correlation coefficients versus the receiver depth in a shallow water with thermocline, where \( f = 1 kHz, x_1 = 10m, d = 1.5m, l = 50m, and r = 5, 10, 20, 30, 40Km \). It can be seen from Fig.1 that (1) the correlations have obvious depth and range dependence, (2) the correlations received above the thermoline are obvi-
ously greater than those below the thermocline when the source is above the thermocline, (3) the more distant the range, the stronger the correlation, and (4) the longitudinal correlation is much stronger than the vertical one.

In Fig. 2 are shown the measured results of longitudinal correlations in a shallow water with thermocline, where the source depth is 7m, the receiver depths respectively are 7 and 25 m, the range is 7.6 km, the longitudinal separation is 50m, the signal is the narrow-band one with the carrier frequency of 630Hz and the band of 1/3 Oct. The measured results show that the correlation coefficient for the source and receiver depths of 7m is 0.86, while that for the source depth of 7m and the receiver depth of 25m is 0.75, the former is obviously greater than the latter.

CONCLUSIONS

(1) The longitudinal correlation in multipath channels has reciprocity, while the vertical one has not.

(2) The longitudinal correlation is much stronger than the vertical one.

(3) In shallow water with thermocline the spatial correlations for both the source and receiver above the thermocline are greater than those for the source above and the receiver below the thermocline.

(4) The longitudinal and vertical correlations increase with increase of range.

REFERENCE

SPATIAL COHERENCE OF SOUND IN SHALLOW WATER

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In this paper, a method for the problem of horizontal coherence of sound in shallow water is developed by using the concept of horizontal rays, vertical modes (HRVM)$^1$, and adiabatic approximation. Analytical and numerical results were compared with those of measurements.

THEORETICAL MODEL

We consider an ocean model with N layers and a half space fluid basement. The pressure release surface $H_0$ and the layer boundaries $H_i$, as well as the sound speed at each layer $c_\nu$ can all be slowly varying function of horizontal position. In this case, the potential functions $U(r,s)$ is governed by Helmholtz equation and boundary conditions:

$$\nabla^2 U + \frac{\beta^2}{c^2} U + \mu^2(x,y,z) U = 0 .$$

According to HRVM theory$^1$, if we neglect the mode coupling effects, an approximate solution for $U$ can be written as

$$U = Q \sum_{i=0}^{\infty} \left\{ \frac{\rho_i(r,s)}{\mu_i} \exp\left\{ -\frac{\beta}{c} r + i \int_{0}^{r} \mu_i (t') \exp \left\{ -\frac{\beta}{c} t' \right\} dt' \right\} \right\}$$

where the integrals in exponent are along the horizontal eigenrays of mode number $m$, $\rho_i$ and $\mu_i$ are normalized eigenvalues and eigenvalues of the unperturbed eigenvalue problem:

$$\frac{\partial^2}{\partial z^2} \rho_i + (k^2(r,s) - \mu^2_i) \rho_i = 0 .$$

For small disturbance, we rewrite:

$$k(r,s) = k_0(1 + \sigma_0 r(s)),\quad H_i(r) = H_0 + \sigma_i H_f,\quad \sigma_i \equiv \frac{\sigma_i}{c_i},\quad \sigma_i = \frac{\partial H_i}{\partial s},\quad \sigma_i = \frac{\partial H_i}{\partial r},$$

By inserting Eqs.(4)-(7) into Eq.(2) we can get equation and boundary conditions for $U_i(r,s)$, from which:

$$\beta \mu_i \rho_i = - \left[ \int_{0}^{s} \rho_i \frac{\partial \sigma_i}{\partial s} ds - \sum_{j=1}^{m} \rho_{ij} U_{ij} \right]$$

where

$$\rho_{ij} = \frac{1}{2 \mu_i} \left\{ \left( \frac{\sigma_i}{\beta} \mu_i \right)^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

and

$$A_i = A_i H_i .$$

We assume $\eta$ and $h$ to be statistically independent with each other. The effects of $\eta$ and $h$ can be treated separately.

We assume that $h(x,y)$ is homogeneous random function of Gaussian type. The spatial coherence function of two points $r_i(x_1,y_1)$ and $r_j(x_2,y_2)$ can be expressed as

$$C(r_i,r_j) = \langle U(r_i,s) U^*(r_j,s) \rangle = \sum_{m=0}^{\infty} \left\{ \frac{\beta^2}{c^2} \left[ B_{mm} + i \int_{0}^{r_i} \rho_i \rho_j \sigma_i \sigma_j d^2 \right] \right\}$$

where

$$B_{mm} = \frac{1}{2} \left\{ \left[ \frac{\sigma_i}{\beta} \mu_i \right]^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

and

$$\rho_i = \rho_i(r,s),\quad \rho_j = \rho_j(r,s),$$

Therefore, the integral terms in the exponent of Eq.(10) can be expressed in the generalised form

$$\Gamma_{mm} \Gamma_{mm} = \int_{0}^{r_i} \int_{0}^{r_j} d^2 \rho \rho \sigma \sigma W_{ij}(\sigma,\tau)$$

where

$$\Gamma_{mm} = \frac{1}{2} \left\{ \left[ \frac{\sigma_i}{\beta} \mu_i \right]^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

and

$$\Gamma_{mm} = \frac{1}{2} \left\{ \left[ \frac{\sigma_i}{\beta} \mu_i \right]^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

for water mass, and boundary disturbance respectively.

We assume the horizontal correlation function in the form

$$W_{ij}(\sigma,\tau) = \exp \left\{ - \frac{1}{2} \left[ (x_1 - x_2)^2 + (y_1 - y_2)^2 \right] \right\}$$

and

$$\Gamma_{mm} \Gamma_{mm} = \int_{0}^{r_i} \int_{0}^{r_j} d^2 \rho \rho \sigma \sigma W_{ij}(\sigma,\tau)$$

We can get the transverse horizontal coherence of two receivers separated by the distance $d$:

$$C(r,s) = \sum_{m=0}^{\infty} \left\{ \sum_{n=0}^{\infty} \left[ B_{mm} + i \int_{0}^{r_i} \rho_i \rho_j \sigma_i \sigma_j d^2 \right] \right\} \exp \left\{ -2 \delta \rho \sigma \sigma W_{ij}(\sigma,\tau) \right\}$$

where

$$\delta \rho \sigma \sigma = \frac{\sqrt{\pi} \sigma \rho}{2 \sqrt{2}} \int_{0}^{\infty} x^3 e^{-x^2} dx,$$

and

$$\sigma(0) = 1 .$$

Finally, the transverse horizontal coherence function:

$$C(r,s) = \sum_{m=0}^{\infty} \left\{ \sum_{n=0}^{\infty} \left[ B_{mm} + i \int_{0}^{r_i} \rho_i \rho_j \sigma_i \sigma_j d^2 \right] \right\} \exp \left\{ -2 \delta \rho \sigma \sigma W_{ij}(\sigma,\tau) \right\}$$

NUMERICAL ANALYSIS

Spatial coherence in winter condition

In winter, most shallow sea is isothermal. If the frequency is high enough, we have the approximate results:

$$\Gamma_{mm} = \frac{1}{2} \left\{ \left[ \frac{\sigma_i}{\beta} \mu_i \right]^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

for surface and bottom, (18)

$$\Gamma_{mm} = \frac{1}{2} \left\{ \left[ \frac{\sigma_i}{\beta} \mu_i \right]^2 - \frac{\mu_i}{\rho_i} \sigma_i \right\}$$

for water mass, (19)

$$\delta \rho \sigma \sigma = \frac{\sqrt{\pi} \sigma \rho}{2 \sqrt{2}} \int_{0}^{\infty} x^3 e^{-x^2} dx,$$

for water attenuation, (20)

$$\delta \rho \sigma \sigma = \frac{\sqrt{\pi} \sigma \rho}{2 \sqrt{2}} \int_{0}^{\infty} x^3 e^{-x^2} dx,$$

for bottom attenuation. (21)
It can be seen that the influence of bottom and surface disturbance and that of water mass on sound field coherence have significant difference. The $\Gamma_{\infty}$ of surface disturbance increases with $n^4$, but the water mass induced $\Gamma_{\infty}$ is almost constant for all $n$. That is, the coherence of higher modes is more sensitive to the bottom and surface disturbance. Figure 1 gives an example of $\Gamma_{\infty}$ and $\delta_n$. We can see from these curves that for the first few modes, the influence of water mass is the major cause of coherence decrease; the boundaries, however, dominate the higher modes. The range dependence of spatial coherence of sound can be divided approximately into: "ray acoustic region", "boundary region" and "water mass region". In the "ray acoustic region", the higher modes dominate the field and the coherence decreases rapidly with range. In the "boundary region" the highest modes are severely attenuated. The boundary disturbance is the major cause of the decoherence. Because of the approximate $n^4$ increase of the mode attenuation ratio, the lower modes, with greater coherence, play a more and more important part in the total field. As a result, the coherence of the field increases with range. In the "water mass region", the field is dominated by the sound speed disturbance in water layer. The coherence decreases exponentially. Figure 2 gives a comparison of the results of theory and measurement in winter condition. The measurement was carried out in December, 1987, at the Yellow Sea.

Spatial coherence in summer condition

In northern areas, the major characteristic of shallow sea in summer is the thermocline. The surface randomness is the most important source of the fluctuation of the highest modes. The bottom and internal waves determine the behavior of the lower modes. Most of the energy of the lowest modes are restricted under the thermocline and the higher modes have homogeneous contribution in water column. Therefore, if both the source and receivers are under the thermocline, so that the lower modes dominate the field, the coherence must be greater than in other cases. Figure 3 gives curves of coherence at different receiver depths. These results are very similar to the measurement results (Scholtz and Zhang, et al). Figure 4 gives examples of coherence of different ranges and frequencies. For low frequency the number of modes is small, the coherence decreases in range increase. For higher frequency, no simple relation is found between coherence and range.

REFERENCES
The sound field in the layered or nearly layered media far enough from the source can be represented as a finite sum of propagating normal modes. Modal description is especially convenient for low frequencies when number of propagating modes is small. Therefore \( N \) complex amplitudes \( a_n \) of these modes contain complete information about the field.

For this reason, experimental measurement of \( a_n \) is of interest. To measure \( a_n \) long enough vertical array can be used which spans the most part of the waveguide. Now only a few publications concerning measurements of this sort are available. In particular, such experiments were carried out in shallow sea [1,2], in Arctic [3] and in Pacific Ocean [4]. One of the main difficulties of such experiment consists in controlling with high accuracy the space configuration of the array. If this problem is solved then for cw-signal the following linear set of equations arises:

\[
P_n'(z) = \sum_{n=1}^{N} a_n \{ \mathbf{u}_n(z) \exp(\xi_n x_n) \}, \quad n = 1, 2, ..., N
\]

Here \( N \) is total number of hydrophones, \( P_n'(z) \) is theoretically measured signal at the \( n \)-th hydrophone, \( x_n, z_n \) - its horizontal and vertical coordinates (in the plane of sound propagation), \( \mathbf{u}_n(z) \)-profile of the \( n \)-th mode and \( \xi_n \)-corresponding wave number. The solution of this set of equations was discussed in detail in [5].

Our experiment was made in Norway sea in 1990 [5] with the help of 560 m long vertical array with 29 equally spaced hydrophones (arrangement of this array is described in [6]). The array was deployed from the vessel and its space configuration was controlled with special acoustical system. The 105 Hz transmitter of cw-signal was set at the depth \( z = 550 \) m and \( 105.5 \) km far from the array. The propagation conditions between the source and the array were inhomogeneous. Fig.1 shows isolines of sound speed, bottom profile along the acoustic path, array and source positions.

Mode amplitudes were calculated according to the procedure described in [5]. Arbitrarily scaled intensities \( |a_n|^2 \) for the first 11 modes (theoretically predicted number of water-borne ones) are shown in Fig.2 by solid vertical lines. The dashed lines correspond to mode intensities theoretically calculated from adiabatic mode theory (AMT).

![Fig.2](image)

**Fig.2** Measured (solid lines) and theoretically predicted (dashed lines) intensities of modes.

For given complex mode amplitudes \( a_n \), we can attempt to reconstruct position of the source "emitting" the sound field backward numerically. According to AMT the modes generated at the point \( x_s, z_s \) have the following amplitudes at the point \( x, z \):

\[
b_i = u_i(x_s, z_s) \exp\left( \int_{x_s}^{x} \xi_i dx \right)
\]

Thus the "angle" included between two vectors \( a_i \) and \( b_i \) depends on assumed coordinates of the source:

\[
C(x_s, z_s) = \left( \sum_{i=1}^{N} a_i b_i^* \right) \left( \sum_{i=1}^{N} |a_i|^2 \right)^{-1/2} \left( \sum_{i=1}^{N} |b_i|^2 \right)^{-1/2}
\]

Dashed lines in Fig.3 demonstrate the following dependence

\[
C_{\text{max}} = \max_{x_s, z_s} |C(x_s, z_s)|
\]

![Fig.3](image)

**Fig.3** Function \( C(x_s, z_s) \) determining source location.
One can see that $C(x)$ is maximum at $x_0 = 109$ km that differs by 4.5 km from the real position of the source.

We believe that this discrepancy is owing to lack of information on propagation conditions. The sound speed profiles (SSP) were measured at five points: two terminal (A and S) and three additional points $P_1, P_2, P_3$ (see Fig.1). Due to drift these points appeared to be moved about 6 km apart from the actual acoustic path. Basing on the obtained acoustical data we attempted to reconstruct SSP along the path using the method called the matched field tomography. In our experiment appropriate procedure was as follows. Proceeding from five available SSP we calculated the average profile $C(z)$. Deviations $\Delta C(z) = C(z) - C(z)$, $k = 1, 2, \ldots, 5$, were approximated by empirical orthogonal functions (EOF). Only two of EOF were found to ensure good approximation for all $\Delta C(z)$. Sound speed between points $A, P_1, P_2, P_3, S$ was interpolated linearly (SSP at the terminal points A and S were assumed as given). Hence certain vector $q$ with 25 components described completely sound speed field along the path. For given $q$ the signal $p_m$ at the $n$-th hydrophone was calculated theoretically according to AMT. It is worth mentioning that the quantities $p_m$ cannot be generally considered as linear functions of $q$ for practically possible variations of $q$. The tomography procedure consisted in searching out the vector $q$ which ensures maximum of the following function:

$$K(q) = \left( \sum_{n=1}^{N} \left( \sum_{n=1}^{M} p_n(q) \right)^2 \right)^{1/2}$$

For this we used the gradient descent method with starting point $q = 0$ (i.e. layered medium with SSP $C(z)$).

Measured (dashed lines) and corresponding to $q$ SSP $\Delta C(z)$ (solid lines) are represented in Fig.4.

One can see that the reconstructed medium appeared to be close enough to the measured one and maximum deviation is of the order of 1 m/s.

Then locating of the source position was accomplished in tomographyally reconstructed medium. The result is shown in Fig.3 by solid line. The new position of the source was found to be 106 km from the array what deviated only by 0.5 km from its real position.

Mention that the curve become more symmetric with respect to vertical line $x_{max}$ as it should be.

Literature:
COHERENCE OF SURFACE REVERBERATIONS OF DIFFERENT BEAMS

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The study of horizontal and vertical coherence of sound scattered from rough surfaces has been the subject of many papers and monographs (e.g., [1-3]). This presentation gives some results of theoretical and experimental studies on the coherence of surface backscattering of sound received by two non-collinear receiving beams. As was limited by the recording system, only the envelopes of the reverberation signals were recorded and analyzed.

I. Theory

As is shown in Fig. 1, the average plane of the rough surface is placed on the X-Y plane and the source and receivers are on the Z-axis. It is assumed that the surface is slightly rough and can be described by

\[ z = \sigma \xi(x, y), \quad \text{with} \quad \langle \xi \rangle = 0, \quad \langle \xi^2 \rangle = 1 \quad \text{and} \quad 2k \sin \psi \approx \pi \sigma \xi(x, y) \]

where \( \sigma \) is the rms value of the surface roughness, \( \langle \cdot \rangle \) means ensemble average, \( \sin \psi = z_0/R \), and \( k \) is the wave number of the incident wave. By using perturbation methods to the integral expression for the sound pressure and ignoring all the terms of order \( (k\sigma)^2 \) and \( 1/R^2 \), the following expression for the backscattered signal (reverberation) can be derived [3]:

\[ u = \frac{k^2}{R} \int \int \frac{e^{i\frac{4\pi}{\lambda} D(\psi) \sin \psi}}{R^2} r^{2 \psi} \zeta(r') \, ds \]

where \( S \) is the area illuminated by the incident beam and \( D(\psi) \) is the directivity pattern of the receiver, the beam width of which is assumed to be much less than that of the transmitter. If the scattering area is located in the far-field of the source, i.e., \( d \ll R \) where \( d \) is the distance of the illuminated area, by replacing \( R = R_0 \cos \psi \), where \( r' \) varies from 0 to \( d \), in the exponent and \( R = R_0 \) elsewhere, Eq.(2) can be simplified further:

\[ u = \frac{k^2}{R_0} P(\alpha) \int \int D(\psi) \exp[-i2k \cos \psi] \zeta(x, y) \, ds \]

where \( P(\alpha) \) is a function of incident angle and the parameters of the bottom materials.

The coherence function of the reverberation signals received by two directional receivers can thus be expressed as

\[ C(\Delta \phi) = u^* u(\Delta \phi) \int \int D(\psi) \exp[-i2k \cos \psi] \zeta(x, y) \, ds \]

where \( u^* \) is the complex conjugate of \( u \), \( W(\xi', \psi') \) is the correlation function of the surface roughness and is assumed to be related only to the distance between the two points, and where \( r_1, r_2 \) is changed back to \( r_1, r_2 \). In this presentation, the Gaussian type correlation function will be used for the rough surface:

\[ W(\xi', \psi') = \exp[-(r_1^2 + r_2^2 - 2r_1 r_2 \cos(\psi_1 - \psi_2))/a^2] \]

where \( a \) is the correlation length of the surface roughness. The directivity pattern of the receivers is assumed to have the form

\[ D(\psi) = \frac{\cos(\pi \psi/2 \beta)}{\cos \psi} \]

where \( \beta \) is the half-power beamwidth of the receivers.

If a pulsed signal is transmitted by a point source and two directional receivers with the sound axis direction differed by the angle \( 2 \Delta \phi \), as was in Fig. 1, Eq.(4) can be written as

\[ C(\Delta \phi) = B \int \int D(\psi) \exp[-i2k \cos \psi] \zeta(x, y) \, ds \]

where \( d = cT/\cos \psi \), \( T \) is the pulse-length of the incident wave and \( B \) is a factor independent of \( \Delta \phi \).

Changing variables in Eq.(7) to

\[ \xi = \phi_1 - \phi_2, \quad \eta = \phi_1 + \phi_2, \quad z = r_1 - r \]

performing the integration over \( \eta \) and \( t \), using the assumption of \( r \approx d \alpha \), and ignoring all the items of \( O(1/r) \), then the following two-dimensional integral expression can be derived:

\[ C(\Delta \phi) = C_1 \int_0^\infty A(\xi) J(\phi_1) J(\phi_2) \, d\xi \]

where \( \phi_1 = \xi - 2\phi \pm \Delta \phi, \ K = 2k \cos \psi \),

\[ A(\xi) = \sin \frac{\xi}{2 \beta} \frac{\pi}{2 \beta} \]

\[ J(\phi) = \cos \phi e^{-\left(\frac{\phi}{\beta} - i \phi \sin \phi \cos \phi \right)} \int_0^\infty e^{-i2k \cos \phi} \cos \frac{\phi}{\beta} \, ds \]

where \( \text{erfc}(z) = \frac{2}{\sqrt{\pi}} \int_z^\infty e^{-t^2} \, dt \) is the complementary error function and \( C_1 \) is a factor independent of \( \Delta \phi \).

No simple analytical expression can be derived from Eq.(9), therefore, a numerical integration method was used to get the numerical results.

Fig. 2 gives some examples of the calculated results, where the abscissa is the half of the normalized angle between the axes of the two beams, \( \Delta \phi/\phi \), and the ordinate is the normalized cross coherence of the signals in the two receiving
beams, $C(\Delta \phi)/C(0)$. The results for three correlation distance of the surface roughness $a/\lambda=0.25$, 2 and 5 are presented. It may be seen that the normalized coherence increases with the correlation distance $a$. The other calculations (not presented here) showed that the coherence changes only slightly when the receiving beamwidth changes. The results are similar with what was obtained by other methods [4].

II. Experiment

The experiment was performed in a water tank. Sound beams were transmitted over a water-sand interface on the bottom of the tank and the reverberation signal was received by receiving beams of different directions. The half-power beamwidth is about 25° for the projector and 10° for the receivers with sidelobes that were 25 dB down from the mainlobe and located about 35° off the mainlobe. The sand surface was pressed by specially made mould plates to form rough surfaces of certain statistical characteristics. The transmitted signals are 0.1 and 0.2 ms impulses with the carrier frequency being 110 kHz. Fig. 3 gives the block diagram of the experimental arrangement. The gating system is to prevent the receiving system from overloaded by the direct signals and the signals reflected from the walls of the tank and the water surface.

Fig. 4 gives some experimental results and two theoretical curves as comparison. As the experiment recorded only the envelopes of the signals, some additional calculations are needed for the experimental results. Assuming the reverberations to be narrow band Gaussian type stochastic processes:

$$\xi(t) = E(t)e^{j\phi(t)} , \quad <\xi> = 0, \quad <\xi^2> = \sigma^2_x, \quad n = 1, 2$$ (10)

where $E(t)$ is the envelope, $\phi(t)$ is the phase, then the cross correlation coefficient for the envelopes of the signal can be related to that for the signals by [5]:

$$\rho_\phi = 0.91 \rho_x$$ (11)

where $\rho_x = \frac{<E_x E_y> - <E_x> <E_y>}{\sqrt{(<E_x^2> - <E_x>^2)(<E_y^2> - <E_y>^2)}}$ are the correlation coefficients for the signals and the envelopes, respectively.

III. Conclusions

The results of theory and experiments showed that the normalized coherence of boundary reverberations received by non-collinear receiving beams is approximately proportional to the ratio of the overlapped part of the two beams and their effective beamwidth, when the correlation length of the surface roughness $a$ is less than the wavelength of the signals, and increases slightly as $a$ becomes larger.

Acknowledgments

The author is indebted to Professor Shi'e Yang for providing the initial motivation of this work and for many helpful discussions. The experiment of this work was performed at the Underwater Acoustics Laboratory of Harbin Shipbuilding Engineering Institute, Harbin 150001, China.

References


Fig. 1 Geometry of the scattering problems.

Fig. 2 Correlation versus angle between beams, theoretical.

Fig. 3 Block diagram for the experiment.

Fig. 4 Correlation versus angle between beams, dots -- experiment, lines -- theory.
WAVE FIELD SPLITTING, INVARIANT IMBEDDING, AND PHASE SPACE ANALYSIS APPLIED TO OCEAN ACOUSTIC WAVE PROPAGATION MODELING

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INTRODUCTION

Due to the extremely large size of realistic three-dimensional ocean propagation problems, there is an understandable desire to incorporate marching (one-way) methods into the solution algorithm for the inherently elliptic, frequency-domain formulation. The geometry of the problem provides the key insight. For the purpose of illustration, imagine a two-dimensional ocean waveguide with flat top and bottom surfaces divided into three distinct regions: left- and right-hand half-spaces separated by a transition region of arbitrary length. The two half-spaces are taken to be transversely inhomogeneous, while the transition region is both depth and range dependent. In the ocean problem, boundary data is given on the top and bottom boundaries; however, no total wave field (or derivative) values are known on the vertical transition-region boundaries. Only sources in the left- and right-hand half-spaces are prescribed. The problem is a scattering problem. Recognizing that typical ocean propagation problems are essentially scattering problems in terms of a transition region and transversely inhomogeneous half-spaces, wave field splitting, invariant imbedding, and phase space methods [1,2] reformulate the problem in terms of an operator scattering matrix characteristic of the transition region. The subsequent equations for the reflection and transmission operators are first-order in range, nonlinear (Riccati-like), and, in general, nonlocal.

MATHEMATICAL DEVELOPMENT

For sound propagation in the ocean, the initial modeling is provided by the scalar Helmholtz equation,

$$\left( \nabla^2 + k^2 K^2(\vec{x}) \right) \phi(\vec{x}) = 0 ,$$

(1)

where $K(\vec{x})$ is the refractive index field and $k$ is a reference wave number. The environment can be characterized by a refractive index field with a compact region of arbitrary variability superimposed upon a transversely inhomogeneous background profile. While the model described is two-dimensional, the analysis is n-dimensional.

The geometry of the formulation suggests the following wave field decomposition. Split the wave field $\phi(\vec{x})$ into two components, $\phi^+(\vec{x})$ and $\phi^-(\vec{x})$, via the transformation

$$\begin{pmatrix} \phi^+(\vec{x}) \\ \phi^-(\vec{x}) \end{pmatrix} = \begin{pmatrix} 1 & 1 \\ iB_1 & iB_2 \end{pmatrix} \begin{pmatrix} \phi^+(\vec{x}) \\ \phi^-(\vec{x}) \end{pmatrix} ,$$

(2)

where $B_1$ and $B_2$ are the two operator solutions of the simple quadratic operator equation

$$B^2 - \left( k^2 K^2(\vec{x}) + \nabla_k^2 \right) = 0 .$$

(3)

This results in the equivalent formulation [1,2]

$$\begin{pmatrix} \partial_x \left( \phi^+(\vec{x}) \right) \\ \partial_x \left( \phi^-(\vec{x}) \right) \end{pmatrix} =$$

$$\begin{pmatrix} -\frac{i}{2} \partial_x B_1^{-1} + 1 & iB_1 \\ \frac{i}{2} \partial_x B_2^{-1} & -\frac{i}{2} \partial_x B_2^{-1} + 1 \end{pmatrix} \begin{pmatrix} \phi^+(\vec{x}) \\ \phi^-(\vec{x}) \end{pmatrix} ,$$

(4)

which holds in all three regions. Choosing $B_1$ to correspond to the forward (outgoing) wave radiation condition and $B_2$ to correspond to the backward wave radiation condition completes the identification. For the generally inhomogeneous transition region, $\phi^+(\vec{x})$ and $\phi^-(\vec{x})$ do not have a special physical interpretation. For the two simply inhomogeneous half-spaces, $K^2(\vec{x}) = K^2(\vec{z})$ and (4) is diagonal. The diagonal system represents the exact decoupling (splitting) of the total wave field $\phi(\vec{x})$ into physically identifiable forward $\phi^+(x,z)$ and backward $\phi^-(x,z)$ wave field components in the transversely inhomogeneous environments. The forward evolution (one-way) equation,

$$\left( i/k \right) \partial_x \phi^+(x,z) + \left( K^2(\vec{z}) + (1/k^2)\nabla_x^2 \right)^{1/2} \phi^+(x,z) = 0 ,$$

(5)

with $B \equiv \left( K^2(\vec{z}) + (1/k^2)\nabla_x^2 \right)^{1/2}$ is the formally exact wave equation for propagation in a transversely inhomogeneous half-space supplemented with appropriate outgoing wave radiation and initial-value conditions [1,2].

For wave propagation problems in the presence of two (generally different) transversely inhomogeneous half-spaces separated by a planar transition region of arbitrary length and inhomogeneity, (4) represents a type of two-point boundary-value problem. For the Helmholtz equation, designating the left and right boundaries of the transition region at $x = a$ and $x = b$, respectively, and generally locating a source in each half-space, the incident wave fields are connected to the scattered wave fields through the operator-valued scattering matrix $\mathbb{S}(a,b)$,

$$\begin{pmatrix} \phi^+(b,z) \\ \phi^-(b,z) \end{pmatrix} = \begin{pmatrix} S(a,b) & \phi^+(a,z) \\ \phi^-(a,z) & \phi^-(b,z) \end{pmatrix} R^+(b,z) ,$$

(6)

The scattering matrix is defined in terms of the appropriate forward (right-traveling) and backward (left-traveling) reflection and transmission operators associated with the transition region. Invariant imbedding methods [1,2] enable the construction of the operator-valued scattering matrix $\mathbb{S}(a,b)$, which then immediately allows for the construction of the reflected and transmitted wave fields by the one-way marching algorithms.

Invariant imbedding intuitively views the scattering matrix for a finite region as being composed of scattering matrices of a large number of contiguous subregions, and thus computes the effect of adjoining a very thin slab to the left-hand side of the transition region. This results in, for example, [1,2]

$$\left( i/k \right) \partial_x R^+(x,b) = \gamma(x) + \delta(x) R^+(x,b)$$

(7)

and

$$\left( i/k \right) \partial_x T^+(x,b) = -T^+(x,b) \alpha(x) + T^+(x,b) \beta(x) R^+(x,b) ,$$

(8)

with the initial conditions $R^+(b,b) = 0$ and $T^+(b,b) = 1$. In (7)
and (8), \(\alpha(x), \beta(x), \gamma(x),\) and \(\delta(x)\) are related to the matrix operators in (4) [1].

In splitting the wave field in a manner which corresponds to the physical experiment and applying invariant imbedding to transform the scattering problem in (4) into an initial-value problem for the scattering operators, it is clear that an explicit representation of the square root Helmholtz operator \(B\) (and \(B^{-1}\), and derivative) and a one-way propagation theory are necessary.

The formal one-way Helmholtz wave equation (5) can be recast and analyzed within the phase space framework as a Weyl pseudo-differential equation in the form [1,2]

\[
(i/k) \partial_x \phi^+(x,z) + (k/2\pi)^{n-1} \int_{R^{n-1}} d\xi d\eta \frac{\partial_x \partial_\xi}{4\pi} \frac{\partial_x \partial_\eta}{4\pi} 
\]

\[
- \Omega_B \left( \psi, (z + z'_z)/2 \right) \exp \left( ik \|z - z'_z\| \right) \phi^+(x,z'_z) = 0
\]

where \(\Omega_B(p,q)\) is the symbol associated with the square root Helmholtz operator \(B = (\mathbf{K}^2(x) + (1/2\pi)^{n-1} \mathbf{I})^{1/2}\). In the Weyl pseudo-differential operator calculus, \(\Omega_B(p,q)\), the operator symbol, is defined through the Weyl composition equation

\[
\Omega_B(p,q) = \mathbf{K}^2(x) - p^2 = 
\]

\[
(k/x)^{2n-2} \int_{R^{n-1}} d\xi d\eta d\zeta \Omega_B(1 + \xi + \eta + \zeta)
\]

\[
\cdot \Omega_B(2ik \cdot z' \cdot q + \zeta) \exp \left( 2ik \cdot z \cdot x \cdot q \right)
\]

(10)

with \(\Omega_B(p,q)\) the symbol associated with the square of \(B, B^2 = (\mathbf{K}^2(x) + (1/2\pi)^{n-1} \mathbf{I})^{1/2}\).

Such solutions representations for pseudo-differential equations as (9) can be directly expressed in terms of infinite-dimensional functional, or path, integrals [1,2], following from the Markov, or semigroup, property of the propagator. This is detailed in [2]. The one-way marching algorithm is based on (1) the matching range step (following from the path integral), (2) a sophisticated symbol analysis (reflecting the detailed study of the (Helmholtz) Weyl composition equation (10)), and (3) Fourier component, or wave number, filtering in phase space (for increased efficiency, decreased computational time, and reduced error).

The detailed numerical algorithm is discussed in [2]. Sufficiency approximates accurately approximations of the square root PDO symbol over the relevant region of phase space result in very accurate numerical wave field calculations [2].

**ALGORITHM SYNTHESIS**

The wave field splitting, invariant imbedding, and phase space methods can now be synthesized into an explicit solution method. As an example, for the limiting case of a transversely homogeneous environment, the equation for the reflection operator symbol takes the form

\[
(i\lambda) \partial_x \Omega_R = \Omega_\gamma + (\Omega_\delta - \Omega_\alpha) \Omega_R + \Omega_\beta \Omega_R^*, \quad \text{with} \quad (11)
\]

\[
\Omega_\lambda \bigg|_{x=a} = 0, \quad \text{and} \quad (12)
\]

\[
\Omega_\gamma = (i/\alpha) \left( \frac{\mathbf{K}(x) \mathbf{K}(z)}{\mathbf{K}(x) - p^2} \right), \quad \Omega_\beta = -\Omega_\gamma, \quad \text{and} \quad (13)
\]

\[
\Omega_\delta - \Omega_\alpha = 2 \left( \mathbf{K}(x) - p^2 \right)^{1/2}, \quad (14)
\]

In (13) and (14), an appropriate outgoing (forward) wave radiation condition is understood. Equation (11) is first-order in range, nonlinear (Riccati-like), and, in the general (transversely inhomogeneous) case, nonlocal. The system is well-posed, but stiff [1]. The stiffness is seen to arise from the local reflections and transmissions.

The algorithmic application of the solution method proceeds in the following manner. The basic idea is to compute the reflection and transmission operator symbols, and then, in the manner prescribed by the Weyl calculus, apply them to the appropriate incoming fields to produce the appropriate initial data for well-posed, one-way marching. Assuming a source in the left half-space only, the reflection and transmission equations (given symbolically in (7) and (8)) can be simultaneously solved in an efficient manner owing to their marching (one-way) nature. The equations are marched from \(x = b\) in incremental steps \(\Delta x\) to \(x = a\), which then regains the physical medium. Each incremental step is transversely inhomogeneous in general. If only reflected and/or transmitted fields are desired, the storage requirements are minimal, since only the symbols (operators) for the physical medium are required. If the wave field in the transition region is also desired, then the incremental reflection symbols (operators) corresponding to the successive imbedding problems must be stored. Starting from the physical problem at \(x = a, \phi^-(a^-, z)\) is determined from the given \(\phi^+(a^+, z)\), the determined \(\Omega_{R^-}(a,b)\), and the Weyl symbol calculus. The continuous \(\phi(a, z)\), and \(\delta_{\phi}(a, z)\) then follow from (2). \(\phi(a^+, z)\) now follows from the inverse of (2) [1]. The Weyl symbol calculus and the approximate constructions of \(\Omega_B [1,2]\) are used in the applications of (2) and its inverse. In the transversely inhomogeneous step, \(\phi(a^+^+, z)\) is now one-way marched to \(x = (a + \Delta x)^+\), producing \(\phi^+(a + \Delta x^+, z)\). One now has the problem of an incoming wave field incident on a new slab with a known (previously calculated and stored in the invariant imbedding procedure) reflection symbol (operator). The above outlined procedure is then repeated, in effect, in a layer-stripping manner to \(x = b\). Right- and left-moving wave fields are always propagated in their natural (well-posed) directions. The two-way nature of the elliptic problem in the transition region is accounted for by first marching the reflection symbol (operator) equation from \(x = b\) to \(x = a\), and then successively calculating the field by marching back from \(x = a\) to \(x = b\).

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BISTATIC REVERBERATION IN THE SEA

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In some case, a separated source and receiver are employed and the arrangement is said to be bistatic and reverberation is said bistatic reverberation. In a sense, the reverberation law is different for a bistatic reverberation than monostatic reverberation.

The assumption necessary are the following:

1 A density of scatterers so large that a large of scatterers occur in an elemental volume.
2 A random, homogeneous distribution of scatterers throughout the volume.
3 An absence of multiple scattering.
4 A pulse duration short enough for propagation effects over the range extension of the elemental volume to be neglected.

See Fig 1. Let A is source of sound, B is point of receiver, A and B is the focus of the prolate spheroidal, where d is distance from the focus, if and YA and YB are distances from the focus to some scattered object.

Quoted the prolate spheroidal coordinates parameter \( \xi \) and \( \eta \) is defined as follows:

\[
 \begin{align*}
 \gamma_A + \gamma_B &= \xi d \\
 \gamma_A - \gamma_B &= \eta d
\end{align*}
\]

Where \( \xi \) goes from 1 to \( \eta \) goes from -1 to +1.

The Surface constant is a prolate spheroid with interfocal distance, because + = Constant.

So scattering signal produced by all elemental volume of surfaces prolate spheroidal arrives back the receiver at the same instant of time.

If source A projected a pulse signal.

The incident intensity \( I_0 \) will be by assumption above \( I_0 \) and \( I_0 \) intensity of source arrives back the receiver B scattering intensity we obtain \( \frac{1}{\gamma_A^2} \), \( S_y \) - Scattering coefficient. According to characteristic of prolate spheroidal, by reverberation is produced by scattered distributed over the (thickness \( \frac{S_y}{2} \)) surface \( \xi \) - constant of prolate spheroidal.

Where \( \xi \) is the pulse duration and \( c \) the velocity of sound. We obtain scattering intensity is

\[
 I_\xi = \frac{I_0}{\gamma_A^2} S_y \frac{c_x^2}{2} \int dS
\]

replace \( \gamma \) and \( \eta \) go on integral.

We obtain reverberation level \( (\text{in decibels unit)\quad R_\nu = \frac{10}{10} \log \frac{I_\nu}{I_0} \quad S_y \quad c \quad \frac{j}{d} - \frac{1}{4} \pm \frac{1}{4} \quad \text{const.} \quad - \frac{1}{3} \quad \left[ \frac{1}{n - \frac{1}{2}} + \frac{1}{2} \right] \quad (6)} \)

Fig 1. prolate spheroid

When distance is very large that is equation 2 become,

\[
 R_\nu = \frac{10}{10} \log \frac{I_0 \gamma_y \gamma_y S_y \frac{c_x^2}{2} \int dS}{\gamma_x^2} = 2 \gamma \quad \text{equation (8) become,}
\]

\[
 R_\nu = \frac{10}{10} \log \frac{2 \gamma \gamma_y S_y \frac{c_x^2}{2}}{\gamma_x^2} \quad (8)
\]

This is general monotonic riverberation formula (nondirectional source).

The experiment were made at the south china sea.
The projection source of sound is the explosion of a 1 kilogram TNT set up on the point were is 60 metres deep below the sea surface. The depth of receiver is 60 metres, it is 1.45 Km from the projector to receiver. The sound signal received is recorded by using the tape recorder. We dealt with experiment data by 1/3oct filter, one of the central frequency is 0.8Kc and 2.6Kc.

The fig(2) and fig(3) show the variation of time with the reverberation intensity, the real line comes from the formula (2). It's obvious that a relation between t and is

According to Fig(2) and Fig(3), we conclude that bistatic reverberation is different from monostatic reverberation. For the monostatic reverberation it is monoton decreasing when the value of time(or distance) increases.

With time increasing bistatic reverberation firstly increases to reach the maximum value (where is 1.33 and t is 500ms) and finally falls down. But the maximum value worked out in the experiment is at the point where t equals 270ms. So the result from experiment is basically consistent with that from theory. During the experiment, we also discovered that there is influence of source near the surface. We would leave it for the following study.

After all that there exists undoubtedly a maximum value for bistatic reverberation is the important difference from monostatic reverberation.

References
ACOUSTICAL INTENSITY STATISTICS
OF BREAKING WAVES

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INTRODUCTION

Breaking surface waves have been identified as the main source of wind-generated ambient sound in the ocean\(^1\). However, previous studies of ambient noise have generally been based on the spatial average of contributions from breaking waves uniformly distributed over the infinite surface. No prior work has thus far been done on the acoustic intensity of individual breaking waves. Characterization of the sound radiated by breaking waves is important not only to understanding the sound generation mechanism, but also to remote sensing of the wave breaking process which is crucial in air-sea interaction. We have recently developed a novel acoustical instrument that tracks breaking surface waves using ambient sound, thus providing an opportunity to investigate the acoustic radiation from breaking events. In this paper, we shall focus on statistical analysis of the acoustic intensity of breaking waves and discuss its potential use in remote measurement of wave energy dissipation.

TECHNICAL APPROACH

Our acoustic observations were made with a small broadband hydrophone array of span 8.5 m and 5 kHz bandwidth. The acoustic instrument is designed as a self-contained package and deployed at a depth of 25 m beneath the ocean surface (Fig. 1). Omnidirectional hydrophones are placed at the ends of the arms. Such a hydrophone array allows us to track each individual breaking wave and determine its spatial and temporal statistics, such as duration and velocity.\(^2\)

Figure 2 shows the trajectories of tracked breaking waves in a 45 s period, based on a recent open ocean experiment. The wind at this time was 11 ms\(^{-1}\) from the north. These events have a mean speed of 5.25 ms\(^{-1}\) relative to the instrument, in the downwind direction. The corresponding wave spectrum (not shown here) has a dominant wind wave component of period 6.6 s, with corresponding phase speed of 10 ms\(^{-1}\). It is observed that the mean event speed is close to the group velocity of the wind waves.

OBSERVATION OF ACOUSTICAL INTENSITY

By using the technique described above, we can isolate each individual breaking event and analyse its acoustical radiation properties. Previous analysis\(^3\) of the power spectrum and coherence of acoustic signals from breaking waves has suggested that breaking waves radiate sound predominantly over the range 100 – 500 Hz. We therefore calculated a time series of acoustic power in this frequency band from one hydrophone. Locations of breaking waves during this period were simultaneously tracked. The received intensity of each breaking event was then found in the series, and the source intensity at 1 m distance was determined by assuming a dipole radiation pattern and neglecting absorption. Such calculations were performed over a period of 30 min in an area of radius 40 m. The estimated distribution of the source intensity (in dB) based on all the tracked breaking waves is given in Fig. 3. The background noise received at the instrument depth is estimated to be 46 dB for these data.

DISCUSSION

Laboratory work using colliding plane waves has demonstrated that the acoustic power radiated by a breaking wave is proportional to the energy dissipated due to breaking.\(^4\) The dissipated energy is proportional to the difference of the upstream and downstream surface displacement variances (\(a_1^2\) and \(a_2^2\) respectively) of the breaking wave, and therefore,

\[ I_0 \sim E_{dis} \sim a_1^2 - a_2^2 \]

where \(I_0\) is the source intensity at 1 m. It is well known that the distribution of wave amplitude is closely Rayleigh, and hence \(a^2\) is found to be exponentially distributed. Assuming \(a_1^2\) and \(a_2^2\) are independent with the same exponential distribution,
it can be shown that \( z = (a_1^2 - a_2^2) \geq 0 \) is also exponentially distributed. Therefore we expect that \( I_0 \) has an exponential distribution. In order to facilitate comparison between the data and this model, we make a transformation \( x = 10 \log_{10} I_0 \) to reduce the dynamic range of \( I_0 \). The resulting distribution of \( x \) is found to be

\[
I_x(x) = \alpha e^{-\alpha x} \exp(\alpha x - e^{\alpha(x-\beta)})
\]

where \( \alpha = \ln 10/10 \) and \( \beta = 10 \log_{10} I_0 \).

Equation (1) is then fitted to the data in Fig.3, where the resulting curve is also plotted. It can be seen that the fit is fairly good for intensities greater than 86 dB. Significant deviation from the curve below \( I_0 = 86 \) dB can be explained on the basis that only those events with intensity above a certain threshold can be detected and thus the statistics at lower intensities are biased by the background noise. However, the fact that \( I_0 \) closely follows the Rayleigh distribution at higher values supports the hypothesis that the acoustic power released by breaking waves is proportional to the dissipated wave energy.

In summary, observations of the acoustic intensity statistics of individual breaking waves in the ocean using the novel instrument appear to be consistent with the dependence of the radiated acoustic power on the dissipated wave energy obtained for the special case of colliding plane waves in the laboratory. This suggests that energy dissipation by wave breaking at the ocean surface may be probed by using ambient sound.

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References

Figure 1: Acoustical instrument for passive detection of breaking waves.

Figure 2: Plan view of breaking wave trajectories for a 45 s time series. Wind direction is from the north. Relative positions of hydrophone locations are also shown.

Figure 3: Estimated probability density function of the source intensity (in dB) of breaking waves at 1 m. The curve is from Eq.(1) in the text with \( \beta = 85.41 \) dB.
SIMULATION OF MULTICHANNEL NARROWBAND OCEANIC REVERBERATION WITH SPECIFIED TIME–SPACE CORRELATION FUNCTION

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INTRODUCTION

Synthesis of multichannel narrowband oceanic reverberation sequences with specified time–space i.e. auto–cross correlation function is of considerable importance in active sonar testing and designing. The recent basic technique in obtaining a reverberation complex envelope time series with a specified time–varying auto correlation function is the piecewise AR or time–varying AR model [1,2,3,4]. In this paper, we have proved that if we first use a time–varying AR filter to generate multichannel reverberation complex envelope time series, each of which has the same specified time–varying auto correlation and is independent of the others, and then use a MA filter to process them properly, the resulting multichannel time series will have the specified auto–cross correlation function. By modulating the complex envelope time series with the active sonar carrier, we get the desired reverberation sequences.

THEORY AND EXPERIMENTAL RESULTS

Denote the M–channel narrowband oceanic reverberation sequences as

\[ X \psi = Re[X \psi e^{j2\pi n \nu}], \quad i = 1,...,M. \]  

(1)

where \( X \psi \) is the complex envelope time series of the \( i \)-th channel reverberation sequence, \( f_0 \) is the carrier frequency of the sonar transmitter. According O'lehesyki's theory [9] and by making use of the relations between the correlation functions of narrowband sequences and the ones of their complex envelope time series, we can obtain the auto–cross correlation function of multichannel reverberation complex envelope time series

\[ R_{X \psi X \psi} (n, n - m) = E[X \psi \overline{X \psi}] = R_{XX} (n, n - m) R_s(d_\psi), \]  

(2)

where

\[ R_{XX} (n, n - m) = E[X \psi \overline{X \psi}], \quad i = 1,...,M \]  

(3)

is the time–varying auto correlation function of the reverberation complex envelope time series which is the same for each channel. \( R_s(d_\psi) \) is the normalized cross correlation function between the two reverberation sequences in the \( i \)-th and \( j \)-th channel. \( d_\psi \) is the distance between the \( i \)-th and \( j \)-th element of the sonar array. When \( n = m, \) we have

\[ R_{X \psi X \psi} (n, 0) = R_{XX} (n, 0) R_s(d_\psi). \]  

(4)

Our aim is to generate \( M \)-channel reverberation complex envelope time series which has the auto–cross correlation function specified by (2). It can be achieved by two steps. In the first step, we use the conventional AR filtering method [1,2,3,4] to generate \( M+q \)-channel reverberation complex envelope time series \( Y \psi \), \( i = -(q-1),...,0,1,...,(M-1), \) each of which is independent of the others and has the same time–varying auto correlation function specified by (3), i.e.

\[ R_{Y \psi Y \psi} (n, n - m) = E[Y \psi \overline{Y \psi}] = R_{XX} (n, n - m). \]

In the second step, we use a \( q \)-order MA filter to process the set of \( M+q \)-channel sampling values \( Y \psi \) at every fixed time index \( n \), and get \( M \) complex sampling values \( \tilde{X} \psi \) at the same time index:

\[ \tilde{X} \psi = \sum_{k=-q}^{q} a_k \overline{Y \psi}^{-k}, i = 1,...,M, \]  

(5)

where \( a_k, \ k = 1,...,q, \) are the MA coefficients. We now prove that under certain conditions we can choose suitable \( a_k, \ k = 1,...,q, \) to make cross correlation between \( \tilde{X} \psi \) and \( \tilde{X} \psi \) equal the one specified by (4), that is

\[ E[\tilde{X} \psi \overline{\tilde{X} \psi}] = E[\sum_{k=-q}^{q} a_k \overline{Y \psi}^{-k} \sum_{k=-q}^{q} a_k \overline{Y \psi}^{-k}] = R_{XX} (n, 0). \]  

(6)

The fact that \( \tilde{Y} \psi \) is independent of \( Y \psi \) when \( i \neq j \), and

\[ E[|\tilde{Y} \psi|^2] = E[|\tilde{X} \psi|^2] = R_{XX} (n, 0) \]

enable us to get the below equation from (6)

\[ \sum_{k=-q}^{q} a_k a_{-k} \delta_{i-j} = R_s(d_\psi), \quad i,j = 1,...,M. \]  

(7)

This is a nonlinear algebraic equation. Because at the right side \( R_s(d_\psi) \) only takes real value, so does every \( a_k \). Hence the superscript * can be omitted in (7). In general, the equation (7) has no solution. But when a sonar array takes some uniform symmetry (for example, a line or circular array of equally spaced elements), \( d_\psi \) is determined only by \( |i - j| \), and \( R_s(d_\psi) \) can be replaced by \( R_s(|i - j|) \). If we set \( v = |i - j| \), \( q = M \), the equation (7) has the form:

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**B5-8**
The nonlinear equation (8) is usually soluble. The solution for M real MA coefficients \(a_k\), \(k = 1, \ldots, M\), can be obtained by using conventional approaches (for example, the Newton-Raphson method).

It is easy to identify that the multichannel time series \(X^M\) synthesized by this two-step method has exactly the same auto-cross correlation function as specified by (2).

Fig. 1 shows the narrowband oceanic reverberation synthesis system which is constructed by applying the theory described above. The sonar array is a circular arc of \(M\) equally spaced elements, where \(M = 28\). The order of the time-varying AR filter is \(p = 20\). The order of the MA filter is \(q = M = 28\). The system uses two TMS320C25 and two INMOS A100 chips to do all the filtering and modulating functions. The single channel reverberation spectrum and its time-evolution are shown respectively in Fig. 2 and Fig. 3. Fig. 4 presents the comparison between the measured and the specified values of the normalized cross correlation of the multichannel reverberation complex envelope time series. The agreement is satisfactory.

CONCLUSION

The proposed two-step (AR and MA filtering) method for synthesizing multichannel narrowband oceanic reverberation sequences has some advantages such as being simple, compute-saving and of good fidelity. It can be used for testing active sonar with a linear or circular array of equally spaced elements.

REFERENCES

ENHANCEMENT OF HYDRODYNAMIC FLOW NOISE RADIATION BY THE REGULATION OF AIR BUBBLES IN A TURBULENT WATER JET

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An experiment is performed to show that the near-field hydrodynamic radiation of flow noise (generated by a turbulent submerged water jet) is enhanced when the turbulent flow is modified to become a two-phase flow containing air bubbles. Acoustic intensity spectra, in the frequency band between 5Hz and 7000Hz, are measured using a digital spectrum analyzer from signals generated by a hydrophone placed at the position of Z=4D and R=4D. Here, Z and R are the axial and radial positions from the nozzle exit, respectively. The water velocity is 12m/s at the nozzle exit diameter D=0.635cm. An amplification factor defined by the ratio of intensities

\[ \frac{I_{\text{air}}}{I_{\text{fluid flow}}} \]

is measured as a function of the void fraction \( \beta \) of the air bubbles.

I. INTRODUCTION

Using Lighthill's\(^1\) aerodynamic theory as a starting point, Crighton and Ffowcs Williams\(^2\) showed (under reasonable hypotheses) that the effect of bubbles in a turbulent flow is to increase the acoustic power output of radiated noise by the factor \( \left( \frac{c_{\text{fluid}}}{c_{\text{mixture}}} \right)^2 \), where \( c_{\text{fluid}} \) and \( c_{\text{mixture}} \) are the low frequency sound speeds in the fluid alone and in the fluid-air bubble mixture, respectively. More recently, Prosperetti\(^3\) investigated this amplification mechanism and was able to show that the low frequency ambient noise in the ocean might be explained by the amplification effects of bubble layers in a turbulent ocean - caused by breaking waves.

Measurements of the acoustic intensity are made as a function of void fraction \( \beta \) (the ratio of air volume to total volume) in an effort to verify the theoretical amplification predictions made in Ref 1 and 2. For low frequencies, the sound speed in a mixture can be expressed by \( c_{\text{mixture}}^2 = \frac{1}{\beta} (c_{\text{air}}^2 + (1-\beta)c_{\text{water}}^2) \). Define \( \epsilon = \rho_{\text{air}} / \rho_{\text{water}} \) and \( \chi = \rho_{\text{air}} c_{\text{air}}^2 / \rho_{\text{water}} c_{\text{water}}^2 \), which involve the densities \( \rho \) and sound speeds \( c \) of each phase in the mixture. Here \( \beta < 1 \) and let \( f \) and \( g \) represent the fluid and gas phases, respectively.

II. EXPERIMENTAL SETUP

The experiment is performed in a 6.5m x 6.5 m x 5m deep section of the U.S. Naval Academy Hydrodynamics Tow Tank. An apparatus is constructed to produce a turbulent shear flow that is generated by a submerged circular jet which is arranged to flow in an upright position. See Fig.1. The submerged water jet apparatus has a plenum section followed by a conical section that is joined to a circular nozzle (which tapers from a 7.62cm diameter to a 0.635cm diameter at the exit). The bubbles are generated in the nozzle throat by a fritted ceramic disk (1cm diameter, 5μm pore size) that is housed in a glass Buchner funnel. Compressed nitrogen gas is regulated and metered from a copper pipeline above the surface.

Measurements of the near-field acoustic pressure spectrum are made from a hydrophone that is shown in Fig.1. Voltage signals from the hydrophone are band pass filtered between 5Hz and 7000Hz and amplified. A digital oscilloscope, with a Fast Fourier Transform algorithm, computes the average acoustic pressure spectrum of 50 trials.

III. EXPERIMENTAL RESULTS

In Fig.2, the average of 50 trials of the squared rms acoustic pressure spectrum is shown for the following two cases: flow noise generated by (1) turbulence alone and by (2) turbulence with bubbles. (We define a power spectrum to be proportional to the squared rms pressure spectrum.) The volume flow rate of \( N_2 \) gas and water are measured to be 19±3 cc / min and 2.5 x 10^4 cc / min respectively. From these measurements, a void fraction \( \beta \) was estimated to be 7.6 x 10^-4.

**Fig.1.** Experimental setup for measuring flow noise from a two-phase turbulent flow containing air bubbles.

**Fig.2.** Flow noise power spectrum. Turbulent flow containing bubbles, without bubbles. (Linear vertical scale is in relative units.)
Since the pressure to voltage sensitivity of the omni-directional hydrophone did not change from case (1) to case (2) an amplification or gain factor G can be computed from the results in Fig.2. Define \( G = \frac{L_{2}}{L_{1}} \) where the intensity \( I \propto \int_{f}^{f} (\text{rms pressure spectrum})^{2} df \) corresponds to either case (1) or (2). The amplification factor \( G \) was computed to be \( G=2.76 \) when \( f_{1}=5\text{Hz} \) and \( f_{2}=4000\text{Hz} \).

**Amplification of Hydrodynamic Flow Noise**

![Amplification Graph](image)

**Fig. 3.** Intensity gain of turbulent (two-phase) flow noise vs void fraction.

Amplification measurements of the hydrodynamic flow noise are presented in Fig.3. Here, the gain \( G \) is measured as a function of void fraction \( \beta \). The volume flow rate of the water is kept constant for all data runs while the volume flow rate of \( N_{2} \) gas going to the bubbler is varied. The captured pressure signal, shown in Fig.4, is a single transient burst (labeled trial 1). For this case the turbulence is generated with a bubble void fraction of \( \beta = 4.9 \times 10^{-4} \).

**Fig. 4.** Transient pressure signal vs time of turbulent two-phase flow noise.

Figure 5. shows the corresponding power spectrum of the burst. The spectral peaks between 1 and 2kHz are significant when compared with the relatively low frequency flow noise spectrum of the turbulent jet (in the absence of bubbles). These peaks rapidly increase in magnitude with increasing \( \beta \). This effect might be an extremely large amplification of the minute tail end of the flow noise spectrum. More likely it may be due to collective oscillations of the bubble cloud in the jet plume.

**Fig. 5.** Power spectrum (for the transient pressure signal in Fig.4.) of the turbulent two-phase flow noise.

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**References**

SOUND ABSORPTION IN THE OCEAN: PH EFFECTS

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INTRODUCTION

Sound attenuation in the sea has a long and interesting history. Early ocean experiments in the 10 kHz range showed excess losses an order of magnitude greater than in fresh water. The mechanism in this frequency range was found to be absorption by ionic relaxation of magnesium sulfate (MgSO₄) with a relaxation frequency near 100 kHz. Later ocean measurements revealed an excess low-frequency loss indicating a 1 kHz relaxation. Measurements by the T-jump method showed that the relaxation involved boric acid (B(OH)₃), a minor constituent of sea water. Resonator experiments showed the ionization mechanism to be an acid/base exchange with the bicarbonate/carbonate equilibrium. Unlike the magnesium sulfate case, absorption by the boric acid/carbonate system depends on pH. Other pH dependent relaxations were also discovered in the resonator studies; however, only the magnesium carbonate (MgCO₃) plays a significant role.

As result of this work, a simple formula was developed that includes three relaxations with pH, temperature, salinity and depth factors:

\[ A = A_1 \cdot (\text{MgSO}_4) + A_2 (\text{B(OH)}_3) + A_3 (\text{MgCO}_3) \]

\[ A_n (\text{MgSO}_4) = a_1 \cdot P \cdot Q \cdot F \]

\[ a_1 = 0.5 \times 10^{-10} \cdot (\text{mm}^2) \]

\[ F_1 = 5 \times 10^7 \text{m} \]

\[ a_2 = 0.1 \times 10^{-10} \cdot (\text{mm}^2) \]

\[ F_2 = 5 \times 10^7 \text{m} \]

\[ a_3 = 0.8 \times 10^{-10} \cdot (\text{mm}^2) \]

\[ F_3 = 5 \times 10^7 \text{m} \]

(1)

Values are: A, dBr/km; frequency f, kHz; relaxation frequency f₀, kHz; T, °C; salinity S, %. The mean value pH=8.0 is used as the reference. The H₂O absorption term is negligible below 100 kHz and is omitted. World Ocean pH varies roughly from 7.7 to 8.3, corresponding to an absorption ratio of nearly a factor of four at low frequencies. The magnesium sulfate term (A₁) includes the depth factor d(km). Those of the other relaxations are evidently small and are not considered.

More recent resonator investigations have confirmed the mechanisms and have extended the accuracy of measurements by a large factor. These and other theoretical studies now allow the absorption formula to be based on purely chemical grounds. However, pH must be known to within ±0.05 units to keep prediction errors within the desired limits of ±15% and the simplified formula is adequate for the present discussion because of the pH data base limitations.

SEA MEASUREMENTS

Low-frequency attenuation experiments are usually carried out in sound channels in order to avoid any extra losses due to surface and bottom effects. Explosive sources allow measurements over a wide frequency range. With both source and receiver near the channel axis, propagation loss is measured vs. range. A simple way of analyzing the data is to subtract cylindrical spreading loss from propagation loss and obtain the attenuation coefficient (dBr/km) at selected frequencies by linear regression.

Fig. 1: Sound speed profiles.

Sound speed profiles typical of mid-latitudes are shown in Fig. 1. The sound channel is formed by downward refraction in the thermocline near the surface and upward refraction by the pressure gradient at greater depth. The axis in these cases is near 1 km depth.

Fig. 2: Attenuation Measurements

Results of three experiments are compared in Fig. 2. The axis depth of the E. Med. sound-speed profile (not shown) was roughly 150 m. The solid lines computed from Eq. 1 for these cases show that the absorption model accounts for the data quite well.
CHEMICAL FACTORS

For simplicity, consider only the pH effect on the boric acid process, the main equilibria being:

\[
\begin{align*}
\text{B(OH)}_3 + \text{CO}_2^2- &\rightleftharpoons \text{B(OH)}_4^- + \text{HCO}_3^- \\
\text{Ca}^{2+} + \text{CO}_2^2- &\rightleftharpoons \text{CaCO}_3
\end{align*}
\]  

(3)

Effectively, the top one controls the relaxation frequency and the bottom one controls magnitude.

![Concentrations vs. pH](image)

Concentrations of the components involved vs. pH are shown in Fig. 3, where the ionic terms include all components associated with Na\(^+\), Mg\(^{2+}\), and Ca\(^{2+}\). Magnitude depends on the concentration product [B(OH)]\([\Sigma \text{CO}_3^-]\), where \([\Sigma \text{CO}_3^-]\) includes all species. The equilibrium \(\text{HCO}_3^- \rightleftharpoons \text{CO}_2^2- + \text{H}^+\) is the control. Since total \(\text{CO}_2\) is nearly constant in the World Ocean, \(\text{pH} = -10 \log[\text{H}^+]\) controls absorption. From Fig. 3, it is clear that absorption varies nearly as \(10^\text{pH}\) in the sea water range.

![pH Profiles](image)

Profiles of pH for the locations of Fig. 1 and for the E. Med. are compared in Fig. 4, showing the extremes of variability.

WHAT CONTROLS pH?

Because profiles have regional and probably seasonal dependence, an archival data base may account for only gross features. Examination of the mechanisms controlling pH may therefore be useful in estimating temporal variability.

High surface pH indicates that atmospheric \(\text{CO}_2\) exchange is a critical part of the mechanism. In a lifeless ocean, partial pressure \(P_{\text{CO}_2}\) would govern equilibrium. However, biochemical factors change this simple picture dramatically.

The main part of the chain is photosynthesis. Phytoplankton consume \(\text{CO}_2\) and release \(\text{O}_2\), using nitrogen and phosphorus compounds as nutrients and the effect is to increase pH. Zooplankton and higher life forms consume \(\text{O}_2\), which tends to act in reverse. However, formation of skeletal \(\text{CaCO}_3\) causes much stronger decrease in pH.

In deep water, decay of organic matter tends to reverse the processes. Decay of photosynthetic products will decrease pH. However, effects of \(\text{CaCO}_3\) depend on how much dissolves and how much is buried in sediments.

Ocean circulation is obviously a critical part of the process. The N. Pacific minimum suggests that the pH reduction involves a long circulation time. The problem therefore is how to model the circulation and the \(\text{O}_2/\text{CO}_2/\text{CaCO}_3\) budget in order to account for both the pH trends and the nearly constant total \(\text{CO}_2\) in the World Ocean.

REFERENCES

A RAY-MODE THEORY OF SURFACE-GENERATED AMBIENT NOISE IN THE SEA

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The sea ambient noise is a kind of noise backgrounds in underwater acoustical signal detection, it can reduce the performance of information receiving in the sea. Otherwise, the principal oceanological, biological and meteorological informations are contained within ambient noise. Therefore, the environmental parameters of the sea may be abstracted by using inverse methods for research the ambient noise field1,2.

Interaction of wind and sea-surface is one of the principal causes in generating of the sea ambient noise. If we are not consider the mechanism of generation of the ambient noise, the investigation on the ambient noise may be summed up as a problem of the sound transmission in the sea caused by surface distributed noise source. The assumption are that:

1. The statistically independent directional sound sources are uniformly distributed on the surface3,4.
2. The sea is assumed to be a stratified medium.

INTENSITY, SPATIAL CORRELATION AND DIRECTIONAL DENSITY FUNCTION OF AMBIENT NOISE

According to the ray-mode theory of ocean acoustics, by using the WKBZ method to solve one-dimensional wave equation, taking the smooth-average over range and depth and by ensemble-averaging, the correlation function can be obtained as follows:

\[
< P(\delta)P^*(\delta) > = \frac{2\pi}{V_\delta} \left( \begin{array}{c} \frac{\cos\theta \left( k_{\text{sonar}} \right) \left( {k_0}^2 \cos\theta \right) \delta(k_{\text{sonar}}, -\ln V_\delta) \, \mathrm{d}k_{\text{sonar}}}{V_\delta^2 (\delta + \alpha) \delta(\delta + \sin^2 \alpha)^{1/2}} \end{array} \right)
\]

where

\[
k = k_{\text{sonar}}, \quad \alpha = \arccos\left( \frac{k_{\text{sonar}}}{k_0} \right), \quad a = \arccos \frac{k_{\text{sonar}}}{k_0}.
\]

The characteristics of ambient noise may be discussed from Eq.(1) as following.

DIRECTIONAL DENSITY FUNCTION N(\alpha,Z)

By using Euler's formula, equation (1) can be written as

\[
< P(\delta)P^*(\delta) > = \int_{-a_s}^{a_s} 2\pi \cos\theta \text{N}(\alpha,z) f_{\text{sonar}}(k_{\text{sonar}}) e^{i \theta \alpha \delta} \, \mathrm{d}k_{\text{sonar}}.
\]

in which \(N(\alpha,z)\) is the directional density function obtained by ray-mode theory, it is equal to

\[
N(\alpha,z) = \frac{\sin^2 \alpha}{2V_\delta^2 (\delta + \sin^2 \alpha)^{1/2}} \left( \begin{array}{c} \frac{1}{V_\delta^2} \end{array} \right) \quad \alpha < 0
\]

when the frequency is low and the grazing angle \(\alpha\) is great enough, Eq.(3) can be simplified as

\[
N(\alpha,z) = \frac{1 - \nu^2}{2V_\delta^2 (\ln V_\delta)^2} N'(\alpha,z),
\]

where the function \(N'(\alpha,z)\) is the directional density function deduced by the ray theory, it may be shown to be2

\[
N'(\alpha,z) = \frac{a}{(1 - \nu^2) \sin^4 \alpha} V_\delta^2 \quad \alpha > 0
\]

The factor \(1 - \nu^2 / 2V_\delta^2 (\ln V_\delta)^2\) in Eq.(4) is not great than 1 dB. It will be shown that the function \(N(\alpha,z)\) obtained by ray-mode theory can be applied to large grazing angles.

SPATIAL CORRELATION COEFFICIENT \(\Gamma (d,\rho,\zeta)\)

According to the definition of spatial correlation coefficient

\[
\Gamma (d,\rho,\zeta) = \frac{< P(\delta)P^*(\delta) >}{\int_0^\infty \text{N}(\alpha,z) f_{\text{sonar}}(k_{\text{sonar}}) e^{i \theta \alpha \delta} \, \mathrm{d}k_{\text{sonar}}}
\]

As a result, we get from Eq.(1)

\[
\Gamma (d,\rho,\zeta) = \int_{-\nu}^{\nu} \text{N}(\alpha,z) f_{\text{sonar}}(k_{\text{sonar}}) e^{i \theta \alpha \delta} \, \mathrm{d}k_{\text{sonar}}
\]

INTENSITY OF THE AMBIENT NOISE

Let \(d = \rho = 0\), we get the intensity from Eq.(1)

\[
\text{N}(\alpha,z) = \frac{2}{V_\delta^2 (\delta + \sin^2 \alpha)^{1/2}} \nu \text{N}(\alpha,z)
\]

NUMERICAL RESULTS

SHALLOW WATER WITH THERMOCLINE

The sound speed upper and lower the thermocline are taken as 1530 m/s and 1500 m/s, respectively. The thickness of the thermocline layer is 5 m. \(V_\delta = 1\), \(V_\delta = 0.4\), and \(\alpha = 0\). The noise sources distributed on the sea surface are assumed to be dipole, i.e., \(\text{N}(\alpha,z) = \sin^2 \alpha\). We calculate the directional density function \(N(\alpha,z)\) from Eq.(3). Figure 1 shows the dependence of the function \(N(\alpha,z)\) on the grazing angle \(\alpha\). It is seen from Fig.1 that the function \(N(\alpha,z)\) above and below the thermocline differ from each other, the value of function \(N(\alpha,z)\) lower the thermocline will equal to zero at \(|\alpha| < a_s\).

UNDERWATER SOUND CHANNEL WITH BILINEAR PROFILE

We assume that the axis of the sound channel is located at the depth above 1000 m, the gradients of sound speed above and below the axis are equal to \(-4 \times 10^{-3} \text{m}^{-1}\) and
$1 \times 10^{-7} \text{m}^{-1}$, respectively, $c(1000) = 1480 \text{ m/s}$. The water depth is 6000 m. The dependence of the noise intensity $I(z)$ on the depth $z$ is obtained from Eq.(8), it is shown in Fig.2.

REFERENCES


![Fig.2](image) Noise intensity $I(z)$ versus depth $z$ in a underwater sound channel with bilinear profile.

![Fig.1](image) Directional density function $N(\alpha,z)$ of surface-generated noise in shallow water with thermocline.

a. Above thermocline  
b. Below thermocline
A PROBABILITY EXPRESSION FOR SHIP NOISE POWER FLUCTUATION AND ITS EXPERIMENTAL CONFIRMATION

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INTRODUCTION

Recently, an increasing number of people enjoy marine sports in shallow water areas, such as skin diving, scuba diving etc. An underwater loudspeaker can be employed to inform divers, using acoustic signals, as a simple and effective way of preventing diving accidents. The use of an audio signal is very important, since the diver usually has no communication apparatus. It is fundamentally necessary to investigate the state of the underwater acoustical environment in the audio frequency bandwidth. Many investigators have already considered the underwater sound field from various viewpoints [1]. In most of the previous studies, however, theoretical and/or experimental considerations have been confined to sound waves containing high frequency components (such as ultrasonic waves) in deep water areas.

It is necessary to grasp the statistical properties of random noise radiated from a moving ship (which is an important sound source), in order to consider the stochastic acoustical environment. On the other hand, when one considers the statistical properties of ship noise power fluctuations measured at an observation point, the following should be taken into account: (A) The sound pressure waveform of ship noise is generally considered as the sum of broadband random noise (with a continuous spectrum) and periodic waves (with a line-component spectrum) [1]. (B) The sound propagation environment between the sound source and the observation point usually shows nonstationary properties, caused by temporal changes of the sound propagation path etc. (C) An analogy of fading in the field of radio wave propagation. Accordingly, the actual sound power of the periodic sound pressure wave (the line-component spectrum) at an observation point shows nonstationary fluctuation patterns. From the above practical viewpoints, in the previous paper [2], the statistical theory for the probability density function (pdf) of nonstationary ship noise power fluctuation has been proposed, by paying special attention to the standard shape of the sound pressure spectrum level. But, in the above study, the expression of the pdf has been derived in the form of infinite expansion series. Accordingly, in practical applications, the problem of the truncation error of pdf remains to be considered.

In this paper, for the purpose of solving this problem, an approximate expression of the pdf for nonstationary ship noise power fluctuation is first derived in the form of weighted sums of Gaussian distributions. The validity and the usefulness of the theoretical result has been experimentally confirmed using a digital simulation technique and actually observed ship noise data.

THEORETICAL CONSIDERATION

Mathematical model for ship noise power fluctuation

Let \( p(t) \) be the sound pressure wave of ship noise measured at an observation point. Suppose that \( p(t) \) is constructed as the sum of two different sound pressure waves \([1]\):

\[
p(t) = s(t) + n(t). \tag{1}
\]

Here, \( s(t) \) is an arbitrary periodic sound pressure wave and \( n(t) \) is a non-white Gaussian random noise with frequency bandwidth \( W \). In general, \( s(t) \) and \( n(t) \) can be expressed in the form of Fourier expansion series:

\[
s(t) = \sum_{n=0}^{\infty} (a_n \cos \omega_n t + b_n \sin \omega_n t), \tag{2}
\]

\[
n(t) = \sum_{n=0}^{\infty} (a_n \cos \omega_n t - b_n \sin \omega_n t), \tag{2}
\]

\[
n_0 = 2\pi s_0^2 \epsilon / W. \tag{2}
\]

The values of the Fourier coefficients \( a_n \) and \( b_n \) fluctuate with the lapse of time, since the sound propagation environment between the sound source (a moving ship) and the observation point shows nonstationary properties caused by temporal changes of the sound propagation path etc. And the Fourier coefficients \( a_n \) and \( b_n \) are random variables governed by the Gaussian distribution with mean zero and variance \( V_n \) \( (n=1,2,\ldots). \) When \( p(t) \) is passed through an indicating instrument containing a mean squaring circuit with averaging time \( T \), the output sound power is obtained as \([3]\):

\[
1 = \frac{1}{T} \int_0^T p^2(t) dt = \sum_{n=0}^N \left( V_n + \frac{(a_n^2 + b_n^2)}{2} \right)^{-\frac{1}{2}}, \tag{3}
\]

with

\[
V_n = \frac{1}{T} \left( a_n^2 + b_n^2 \right), \quad K = N + 1 - 2W. \tag{4}
\]

Here, the frequency band of \( n(t) \) is from \( f_0 \) to \( f_0. \)

\[ \mu(f) \] dB/Hz

Fig.1 Sound pressure level of ship noise.

It should be noticed that the value of \( K \) in Eq. (3) depends on the frequency bandwidth \( W \) of \( n(t) \). The mathematical expressions of Eqs. (3) and (4) are correct in principle. However, a normal frequency analyzer (such as FFT equipment) has a certain constant number for the frequency bandwidth \( W \) (Let \( N \) be a constant number. \( W = N \alpha \); see the solid line in Fig.1). Let us therefore introduce an approximate expression for Eqs. (3) and (4) (as a mathematical model) as follows:

\[
1 = \sum_{n=1}^N (c_n + c_n^2)^2 = \sum_{n=1}^N \epsilon_n, \quad (c_n^2 + c_n^2). \tag{5}
\]

Here, \( c_n \) is a value corresponding to the component existing in the frequency interval \( \Delta \) of the periodic sound pressure wave \( s(t) \), and \( \epsilon_n \) is a random variable.

[1]
governed by the Gaussian distribution with mean zero and variance $\sigma_n^2$ corresponding to the component existing in the frequency interval $\Delta$ of the Gaussian random sound pressure wave $n(t)$ (see Fig. 1).

In general, the random variables $c_n$ and $c_m$ $(n \neq m)$ are not statistically independent of each other, since the frequency characteristic of $n(t)$ is not white. However, we neglect here these correlation characteristics. $c_{n} \sim c_{m} \sim \mathcal{N}(0, \Delta \sigma_{n} ^2)$. We introduce the approximate assumption that $c_n$ and $c_m$ $(n \neq m)$ are statistically independent. Using the following facts:

1. The correlation characteristics between two different frequency components are not usually considered in actual experiments on the power spectrum analysis.

2. When one considers the above correlation characteristics in the statistical analysis, the mathematical treatment becomes very complex.

Using the above approximation, let us introduce a simplified pattern for the sound pressure spectrum level of $n(t)$, as shown by the dotted line in Fig. 1. In this case, the sound pressure spectrum level of non-white Gaussian noise is approximated by the combination of partial flat sound pressure spectrum levels. Thus, equation (5) can be rewritten as follows:

$$1 = \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} (c_{ij} + c_{ij})^2 = \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} E_{ij}$$

$$= E_{ij}(c_{ij} + c_{ij})^2$$

Here, $c_{ij}$ is a random variable governed by the Gaussian distribution with mean zero and variance $\sigma_{c_{ij}}^2$ $(i=1, 2, \ldots, A)$.

**Approximate probability expression for nonstationary ship noise power fluctuations**

By considering the fluctuation of $c(c_{ij}, c_{ij}, \ldots, c_{ij})$, the moment generating function (mgf) of the nonstationary sound power fluctuation $L_1$ over a long time interval can be given as [4]:

$$m(\theta) = \exp \left( \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} (c_{ij} + c_{ij})^2 \right)$$

$$= m(\theta; \Sigma)$$

where

$$m(\theta; \Sigma) = \exp \left( \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} \frac{(c_{ij} + c_{ij})^2}{\sigma_{c_{ij}}^2} \right)$$

From Eqs. (6) and (8), we can easily find

$$c_{ij} = \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} (c_{ij} + c_{ij})^2$$

By using this relation, we obtain

$$1 = \sum_{i=1}^{N_1} \sum_{j=1}^{N_1} (c_{ij} + c_{ij})^2$$

Accordingly, equation (8) becomes

$$m(\theta; \Sigma) = \exp \left( \sum_{i=1}^{N_1} \frac{(c_{ij} + c_{ij})^2}{\sigma_{c_{ij}}^2} \right)$$

**Experimental values**

**CONCLUSION**

In this study, an approximate expression of pdf for the sound pressure fluctuation generated by a simple moving vehicle has been derived. By paying special attention to the standard shape of sound pressure spectrum level, the validation and usefulness of the theoretical result has been confirmed experimentally by applying it to the data obtained by digital simulation and actually observed ship noise data.

**REFERENCES**


1. Introduction

As fluid passing through the cavity with a flow rate, the medium of mouth and inside of the cavity are oscillated and sound are radiated. Not only stochastic turbulence noise are possibly produced but also single frequency sound of specially designated pitch are possibly produced. And so spectrum of radiated sound is continuous spectrum plus characteristic line spectrum. Flow-induced oscillation are called by the case.

Unsteady flow of the interface and hydrodynamic self-excited oscillation is produced by flow-induced oscillation, at the same time stationary wave oscillation in the cavity is produced. They can cause elastic vibration of the cavity wall and sound radiation.

2. Experimental work

2.1 Water tunnel

The interior profile of the testing section of the water tunnel is a square sized 20 × 20cm². The length of the testing section of the water tunnel is 1.22 meters. The maximal flow rate is 22m/s.

2.2 Experimental model and cases

The protuberant special cavity with streamline is made of copper sheets and structural drawing see Fig.1. The model consists of up and down two parts. two parts of the model all are cavity.

![Fig.1. Structural drawing of experimental model](image)

Hydrophone $H_1$ is installed at the bottom of the special cavity in order to measure oscillating sound pressure of the cavity. Hydrophone $H_2$ is installed at the top surface of the wall of the testing section in order to measure radiated sound pressure.

The flow rate is from 3m/s to 12m/s.

3. The experimental results and analysis

Figures of power spectral density of accelerometer $H_1$ in the cavity see Fig.2.

![Fig.2. Figures of power spectral density of accelerometer $H_1$ in the cavity](image)

3.1 Natural oscillation of the cavity

Protuberant special cavity with streamline, may be regarded as a Helmholtz resonator from acoustics. Its resonance frequency is:

$$f_0 = \frac{1}{2\pi} \sqrt{\frac{1}{M_\ell C_\ell}}$$  \hspace{1cm} (1)

Where $M_\ell$ is acoustic mass of the cavity, $C_\ell$ is acoustic compliance of the cavity. According to geometric measurements of the model, may get resonance frequency of this cavity in the air:

$$f_0 = \frac{1}{2\pi} \sqrt{\frac{1}{M_\ell C_\ell}} = \frac{C}{2\pi \sqrt{\frac{\epsilon}{VL}}} = 221Hz$$  \hspace{1cm} (2)

When putting model into water, because characteristic impedance of aqueous medium is large than air, so acoustic mass $M_a$ of the cavity increase, thus resonance frequency of Helmholtz resonator in the water descend somewhat than air. A line spectrum of no changing with flow rate is seen in Fig.2, it is approximately 190Hz, it is namely resonance frequency of Helmholtz resonator in the water. Resonance frequency $f_0$ of no changing with flow rate see Fig.3.
3.2 Hydrodynamic oscillation of the cavity

From figures of power spectral density (PSD) in Fig. 2 is seen some resonance peaks in low-frequency range, they rise along with an increase in flow rate, these resonance peaks contain appearance of a bunch of line spectrum. Their numerical value of frequency are \( f_1 \) and \( f_2 \). Figure of frequency \( f_1 \), \( f_2 \) of hydrodynamic oscillation changing with flow rate see Fig. 4.

Relation of Strouhal number and Mach number of hydrodynamic oscillation and natural oscillation of the cavity (two modes) is drawn in same figure, in order to study coupling characteristic between two modes, see Fig. 5.

The relation between the Strouhal number and Mach number for hydrodynamic oscillation of the protuberant special cavity with streamline is given by the equation:

\[
S_m = \frac{m - 0.25}{M + 1.88m}
\]  

(3)

where \( S_m \) is the Strouhal number, \( M = V/C \) is the Mach number, \( V \) is the water velocity of mouth of the cavity, \( C \) is the acoustic speed, \( m \) is the order number of resonant mode, \( m \) is equal to positive integer.

3.3 Frequency spectrum level of third-octave filter of flow-induced oscillation

The frequency spectrum level analysis of third-octave filter is made for flow-induced oscillation of the model.

The figure of general sound level changing with flow rate see Fig. 6.

\[
P = 1.585 \times 10^7 V^{2.20} (\mu Pa)
\]  

(4)

The increment of noise spectrum level is directly proportional to square of flow rate in case low Mach number.

4. Conclusion

1. In water and low Mach number, there are two modes in the cavity for protuberant special cavity with streamline: natural oscillation and hydrodynamic oscillation of the cavity. Frequency of natural oscillation of the cavity does no changing with flow rate. Frequency of hydrodynamic oscillation of the cavity changes with size of mouth of cavity. When frequency of the two oscillation is very close or equal to each other, intense oscillation of resonant line spectrum is produced.

2. The relation between Strouhal number and Mach number of hydrodynamic oscillation of cavity in water is given by the equation (3).

3. In water and low Mach number, the increment of general sound level of flow-induced oscillation is directly proportional to square of flow rate. Theirs relation is given by the equation (4).

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References

SOUND INDUCED BY HYDRODYNAMIC FLOW OVER AN AXISYMMETRIC BODY

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INTRODUCTION

When an axisymmetric body moves in water at high speed, the interaction between the body and flow leads to the fluctuations of velocity and pressure on the surface, which are generally named as "flow noise". Flow noise induces sound through two approaches [1,2,3]: one is its direct radiation, the other is the sound radiation due to shell's vibration induced by flow. In general, the former one, which dominates at high frequencies, while the noise of low frequencies is attributed mainly to the latter one.

CHARACTERISTICS OF FLUCTUATING PRESSURE SPECTRUM AND FLOW INDUCED SOUND

Analysis of mechanism and results of model test show that the flow on body's surface belongs to that of axisymmetric turbulent boundary layer. Except transition and separation zone, fluctuations on the surface of axisymmetric body are almost the same as those on the surface of long cylindrical shell. This result has been verified by our tests for several times.

Test study has been carried out in water tunnel for the flow noise of several kinds of models whose length are 1500 mm and whose maximum diameters are 180 mm. Theoretical analysis and the results of model tests agree well as can be seen in Fig. 1, from which we know that the part of spectrum at low frequencies is even and that highest spectrum level at low frequencies and its variance with speed in different frequency ranges accord better with Corcos's model. If two materials [4], which are metal and glass-reinforced plastic, are used to make models with same outlines, the fluctuating pressure on their surfaces remain various, as is illustrated in Fig. 2.

The significance of shell's vibration induced by flow can been seen from Fig. 3. When speed gets from 6 m/sec to 12 m/sec, vibration acceleration increases by 20 dB, and many discrete spectra appear in low frequency range. Hence it is very possible for shell's vibration induced by flow to form one of the main sources of sound radiation at low frequencies.

At high frequencies, fluctuating pressure within turbulent boundary may remarkably radiate sound which produces great influence on the noise at head terminal. The head section being a flat terminal, account must be taken into sound diffraction at high frequencies.

According to Procos Williams-Hawings' wave equation [5], pressure in external field can be written as

\[ p(x,t) = -\frac{1}{4\pi} \int \frac{p(x,t)}{r^2} dA(y) \]

\[ + \frac{1}{4\pi} \int \frac{\partial p}{\partial r} dA(y) \]

\[ + \frac{1}{4\pi} \int \frac{\partial^2 p}{\partial r^2} dA(y) \]

(1)

Let the surface impedance of axisymmetric body \( Z = p_x U_x, a = \rho c / 2 \). If \( \Theta(w) \) denotes the self power spectrum of fluctuating pressures and \( A \) the correlation area of fluctuating pressures, in the case of flat head terminal, we obtain

\[ S_r(x,w) = \frac{1}{4\pi c} \int \Theta(w) a^{2r} \left[ \cos \Theta(y) - a \right] dA(y) \]

\[ + \frac{2}{(4\pi c)^2} \int \Theta(w) a^{2r} \left[ \cos \Theta(y) \cos \Theta(y) - a \right] dA(y) \]

\[ + \frac{1}{4\pi c} \int \Theta(w) a^{2r} \left[ \cos \Theta(y) \right] dA(y) \]

\[ \cdot \left[ \int \frac{a^{2r}}{\left( 8 \right)^{1/2} \cos \Theta(y) \cos \Theta(y)} dA(y) \right] \]

(2)

where \( r_i \) is the distance from the source point to the edge of flat head terminal, \( r_2 \) is the distance from source point to the central point of the flat head terminal and \( \theta \) (i=1,2) is the angle between \( r_i \) and the axisymmetric body's axis.

Formula (2) expresses sound pressure power spectrum density as a function of such parameters as Reynolds number, position of transition point, displacement of boundary layer and so on. Comparison between calculated values and measured values of one model's noise is illustrated in Fig. 4, from which we can see a better consistence of the two results. Fig. 5 is obtained as the theoretical curve of sound spectrum density at the center of head section if the impedance characteristics of axisymmetric body's surface is changed. Fig. 6 shows the comparison of noise levels of models of same dimension but different materials, i.e. Metal and glass-reinforced which is regarded as flexible. Theoretical analysis stated above is proved in this figure. We can also know from Fig. 4 that the noise level's variance with frequency acts in closed coordination with speed. These facts agree with the results of model tests perfectly.

Interaction between water medium and bodies moving in water occur in most cases, especially the revolution shell, which is different from flat plate, has a great tendency to produce flexural vibration which results in sound radiation.

Equation for radial displacement vibration of the shell is

\[ D_1 \ddot{r} + r \ddot{r} + m \ddot{r} = F \]

Wave equation of inside and outside of shell are listed below:
\begin{align}
\eta^2 p_1 - \frac{1}{c_1^2} \frac{\partial^2 p_1}{\partial t^2} &= 0 \\
\eta^2 p_2 - \frac{1}{c_2^2} \frac{\partial^2 p_2}{\partial t^2} &= 0
\end{align}

(4)

Let \( X_{\text{sh}} \) be the response function of the shell vibration, the displacement \( \mathbf{w}(\omega) \) is

\[ \mathbf{w}(\omega) = \sum_{n,m} F_{nm}(\omega) Y_{nm}(\omega) \]

(5)

With the use of saddle point method, we obtained the self sound power spectrum in far field.

\[
\Phi_{\text{sn}}(r_1, \phi, z, \omega) = \sum_{n,m} \frac{(2l+1)\pi}{\pi} \frac{H_{nm}(\omega)H_{nm}(\omega_0)}{R^2} \Phi_{\text{sn}}(\omega)
\]

(6)

Where \( \Phi_{\text{sn}}(\omega) \) is model power spectrum density of fluctuating pressure on the axisymmetric body's surface, \( R \) is the distance from source point to field point.

Making use of formula (6), we can obtain radiated sound power spectrum levels at different locations and different distances from axisymmetric body.

Sound radiation impedance can be obtained from the following presentation:

\[
\int \frac{\mathbf{S}_{\text{sn}}(k_0) + \mathbf{K}_{\text{sn}}(k_0) \mathbf{d} k}{\mathbf{K}_{\text{sn}}(k_0) \mathbf{d} k} \mathbf{X}_{\text{sn}}(\omega) \mathbf{X}_{\text{sn}}(\omega) \]

(7)

With the use of the formulae listed above, we have calculated sound power spectrum levels of sound radiation due to axisymmetric body's vibration induced by flow. The results show that the highest power spectrum level drops by 13 dB, when shell's damping factor \( \eta \) increases from 0.01 to 0.05. Numerically calculated spectrum is illustrated in Fig. 7.

CONCLUSIONS

1. To the surface of metal axisymmetric body, pressure fluctuation spectrum can be predicted theoretically, but great changes may occur in the case of bodies of different materials.

2. To the radiated sound induced by fluctuating pressure, its theoretically calculated values agree with experiment results fairly well, thereby different noise levels at head terminal can be obtained in the cases of different speeds, different surface impedance boundaries and hydrodynamic parameters.

3. The theoretically calculated result of sound radiation due to axisymmetric body's vibration induced by flow show that shell's vibration induced by flow turns out to be one of the main discrete spectrum sources at low frequencies, and this result can lead to the obtaining of sound power spectrum levels at different locations and different distance from the axisymmetric body.
PRESSURE AND TEMPERATURE DEPENDENCE OF BORIC ACID ABSORPTION IN SEA WATER

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INTRODUCTION

It has been verified by ocean and laboratory measurements that boric acid absorption in sea water depends markedly on the pH value. Although the maximum absorption per wavelength \( \alpha(\lambda) \), of the boric acid relaxation derived from the \( \alpha \)-sea measured data appeared to be independent of temperature and pressure, our resonator measurements showed that it increases with temperature [1]. On the basis of the absorption model given previously by the author and of the related oceanographic literature, we have estimated theoretically the pressure and temperature dependence of boric acid absorption in sea water.

THEORY

It has been determined that boric acid absorption in sea water is caused by an apparent single-step reaction [1]:

\[
B(OH)_3 + CO_3^{2-} \text{(total)} + H_2O = B(OH)_2 \text{(total)} + HCO_3^-(\text{total}),
\]

where the word "total" means that both the free ions and their complexes partake in the reaction. The maximum absorption per wavelength caused by Eq.(1) can be written as

\[
\alpha(\lambda) = \frac{2\pi c_0 \rho \sigma(\lambda)^2 (\Delta \nu')^2}{2RT} \left[ \frac{1+x_c}{x_c E_0} + \frac{1+x_p^2}{2x_p \Delta v} \right],
\]

where

\[
x_c = K_{B}^{-1/2} = 10^{10.38+K_B^{-1/2}}, \quad x_p = K_{AO}^{-1/2} = 10^{10.88+K_{AO}^{-1/2}},
\]

\[
K_{B} = a_B[BOH]_i/[BOH], \quad K_{AO} = a_B[C0^2_]i/[HCO_3^-],
\]

\[
A_o = [B(OH)]_o = [B(OH)]_i + [B(OH)]_k,
\]

\[
E_o = [HCO_3^-]_o = [HCO_3^-]_i + [HCO_3^-]_k
\]

\[
\Delta \nu' = -\frac{\partial \ln K_B}{\partial p} \text{RT} + \frac{3}{2} \frac{\partial \ln K_{AO}}{\partial p} \text{RT} = \Delta \nu_B - \Delta \nu_A
\]

\( \rho \) is the density of sea water, \( c_0 \) is the sound speed in sea water, \( R \) is the gas constant, \( T \) is the absolute temperature, \( p \) is the static pressure, \( K_B \) is the first dissociation constant of boric acid, \( K_{AO} \) is the second dissociation constant of carbonic acid, \( \Delta \nu' \) is the apparent molar volume change for the reaction of Eq.(1). \( \Delta \nu_B \) and \( \Delta \nu_A \) are the molar volume changes for the first dissociation reaction of boric acid and the second dissociation reaction of carbonic acid, respectively. The square brackets in Eqs. (4)-(6) denote molar concentrations. The subscripts \( t \) denote total concentrations - the sum of the free ion concentration and the complex concentrations. \( a_B \) is the hydrogen ion activity (in this paper \( pH=log_{10} a_B \)). \( A_o \) is the overall concentration of boric acid (dissociated and undissociated).

The data of \( K_B \) and \( K_{AO} \) measured at various pressures showed that both \( \ln K_B \) and \( \ln K_{AO} \) are linear functions of pressure (up to 1000 atm) [2]. Therefore, according to Eq.(7), \( \Delta \nu_B \), \( \Delta \nu_A \), and \( \Delta \nu' \) are independent of pressure. The values of \( \Delta \nu' \) derived from experimental data given by different authors are summarised in Ref.[3]. It is found that: (i) the values of \( \Delta \nu' \) obtained by resonator method by Qiu et al. and that by electro-chemistry method by Culberson-Pytkowics and Dieterich-Dieterich agree very well; (ii) the \( \Delta \nu' \) value is nearly independent of temperature, pressure, and salinity; (iii) for natural sea water, \( \Delta \nu' = -12.5 \text{ cm}^3/\text{mol} \).

The dependencies of \( pK_B \) and \( pK_{AO} \) on temperature and chlorinity at atmospheric pressure can be expressed as [4]

\[
pK_B^{(t)} = \frac{-2291.90}{T} + 0.01754T - 3.3850 - 0.23061 \text{Cl}^{1/2},
\]

\[
pK_{AO}^{(t)} = \frac{-2002.30}{T} + 0.03737T - 6.4710 - 0.4693 \text{Cl}^{1/2},
\]

where \( \text{Cl} \) is the chlorinity of sea water in parts per thousand. The chlorinity is related to salinity \( S \) by

\[
S = 1.90566 \text{Cl}.
\]

The values of \( pK_B \) and \( pK_{AO} \) at pressure of \( P \) atm can be calculated from the following relationships [2]

\[
pK_B^{(P)} = pK_B^{(t)} - \Delta pK_B, \quad pK_{AO}^{(P)} = pK_{AO}^{(t)} - \Delta pK_{AO},
\]

\[
\Delta pK_B = 1.89 \times 10^{-3}P - 4.518 \times 10^{-4}PT - 1.69 \times 10^{-7}P^2
\]

\[
+ 1.75 \times 10^{-13}P^3T^2
\]

\[
\Delta pK_{AO} = -0.015 + 8.39 \times 10^{-4}P - 1.908 \times 10^{-4}PT + 1.82 \times 10^{-7}P^2
\]

For the computations of \( \alpha(\lambda) \), we use the following approximate expressions [8]:

\[
A_o = \left(S/35\right) \times 0.43 \times 10^{-3},
\]

\[
c_0 = 1412 + 3.21I + 1.10S + 0.0167D
\]

\[
\rho_o = 1.028 + 0.45 \times 10^{-4}D
\]

where \( D \) is the depth (m), \( I \) is the temperature (°C), \( c_0 \) is
in m/s, \( \rho \) is in g/cm\(^3\). For sea water, we have \( \rho_B < 1 \) and \( \rho_I < 1 \). Then, it can be seen from Eq. (2) that \( (\alpha_l) \) will increase with temperature and pressure owing to increases of \( K_\beta, K_\theta, \) and \( \sigma_0 \).

COMPUTATION RESULTS AND CONCLUSIONS

Taking \([\text{HCO}_3^-]_{\text{p}} = 2.3 \text{ mmol/L, } \Delta V' = -12.5 \text{ cm}^3/\text{mol,} \]
\( S=35, \) and using Eqs. (2)-(16), we can obtain the dependencies of \( (\alpha_l) \) on pH, pressure, and temperature as shown in Figs. 1-4. It is seen from Figs. 1 and 2 that, in the ranges of ocean factors, \( (\alpha_l) \), not only increases with the pH value, but also increases approximately linearly with depth and temperature. In Fig. 3 the dashed line for \( D = 1000 \text{ m} \) and \( t = 5^\circ\text{C} \) coincides with the solid line for \( D = 0 \) and \( t = 10^\circ\text{C} \) over the whole sea water pH range. The partial cancellation of the pressure dependence by the temperature dependence is one of the causes for the low correlation coefficients of \( (\alpha_l) \), with temperature and depth obtained from the analysis of the at-sea measured data.

Representative results of computations show that, for the Pacific and Atlantic, the difference between the \( (\alpha_l) \), value calculated from Eq. (2) and that calculated on the basis of the Browning-Mellen equation \([8]\) is small as long as the depth is not greater than about 2000 m; however, the difference increases gradually with depth when the depth is greater than 2000 m. At 4000 m depth for the Pacific and 5000 m for the Atlantic, the values of \( (\alpha_l) \), calculated from Eq. (2) are \( 0.7 \times 10^{-4} \) and \( 1.2 \times 10^{-4} \), respectively. They are obviously greater than the respective values of \( 0.55 \times 10^{-4} \) and \( 0.8 \times 10^{-4} \) calculated on the basis of the Browning-Mellen equation. Therefore, in order to estimate the long-range attenuation of a low-frequency sound signal travelling via a path, the major portion of which is located at great ocean depths, the effect of pressure on B(OH)\(_2\) adsorption should be taken into account.

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MECHANISMS UNDERLYING TRANSITIONAL AND TURBULENT BOUNDARY LAYER (TBL) FLOW-INDUCED NOISE IN UNDERWATER ACOUSTICS

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INTRODUCTION

This research focuses on improved understanding of the mechanisms underlying turbulent boundary layer (TBL) flow-induced noise in underwater sound, through experimental investigations in the laboratory and numerical modeling. The feature parameter of the investigation is wall-pressure since finite area transducers are typically used in underwater acoustics and their sensitivity is known to increase. Transducer geometry is important in this study because it influences the pressure impulse on the transducer face. Experimentation involves measurement of the wall-pressure fluctuations by station probe mounted along the wall of an asymmetric body of revolution during its buoyant ascent from the bottom of a deep water test basin. Experimental measurements are evaluated using modeling techniques from Linear Stability Theory and from Dynamical Systems Theory (L: Chaos). Finally numerical simulation of a wall bounded shear flow at moderate Reynolds number, 0(10^4), is performed to obtain the representative wall pressure impulse on the face of an array (22 x 22) point transducers.

EXPERIMENTATION

The tests were conducted on a buoyantly-propelled, 53 cm diameter axisymmetric test vehicle during vertical ascent from the bottom of a deep, fresh water basin. The vehicle was released from an anchor point at a depth of 300 meters, and pressure sensor measurements were recorded on magnetic tape during its vertical ascent. Steady state measurements were obtained just a few seconds following release. At a depth of 45 meters, control fins were activated to maneuver the vehicle horizontally before reaching the surface. Vehicle speed could be altered by changing cylindrical weights in the ballast section. Approximate steady-state velocities varying between 12 meters/second and 22 meters/second were attainable in this way. Length Reynolds numbers is order approximately 10^6 based on length scales defined by probe position (i.e., 25 to 25 cm) downstream of the body the centerline from the leading edge. Data used for analysis were taken from the constant velocity portion of the run. Constant velocity data endured from 5 to 15 seconds, depending on vehicle speed. The pressure sensors were of the piezoelectric type measuring, 53 cm in facial diameter. An elastomeric wall cover of 0.187 cm in thickness overlying the transducers was used to provide a smooth wall surface. Both facial diameter and wall cover thickness will affect the accuracy of the measurement, the details of which are discussed elsewhere.[4]

Most of the activity in the instability growth and boundary layer development occurs between an axial displacement of 20 to 35 cm. The station probe measurements were processed through a 14-channel analog-to-digital converter using a sample rate of 10^4 seconds. Records of up to 20,000 samples were collected for analysis.

SPECTRAL ANALYSIS

Power spectral densities obtained from time-sequences of the measured wall-pressure fluctuations are provided on Figure 1 as a function of the non-dimensional Strouhal number and as a function of downstream probe position (going from top to bottom of figure). The expected frequency range of peak amplitudes is shown by the vertical lines according to the Orr-Sommerfeld Theory[4]. Note the agreement between measurement and theory is quite good, particularly in the early stages of instability growth. Also note the evident breakdown of T-S waves in the temporal power spectrum following the spatial migration into the fully developed turbulence region of the flow.

DYNAMICAL SYSTEMS PROCESSING

Using the method-of-delays technique,[4] phase portrait constructions were obtained from transitional and turbulence time series measurements. We applied a three point moving average to the turbulence time series to clean up the orbit by removing the higher mode components, outside the region of T-S frequency. Representative phase plots (Figure 2 and 3) for transition zone and turbulence zone time series measurements show striking similarity following low pass filtering of the turbulence time series. The coalescing of points at a delay interval of T = 1 x 10^{-3} is evidence of the stroboscopic orbit of the T-S wave, shown here at 1/2 wavelength. The similarity in data structure between transitional and turbulence measurements is clearly not observed using the linear data processing technique (i.e., temporal power spectrum).

Figure 1 Temporal Power Spectra Versus Downstream Displacement

Figure 2 Phase Portraits from Transitional Time Series Measurements (delay = T x 10^{-3})

Figure 3 Phase Portraits from Full Turbulence Time Series Measurements (delay = T x 10^{-3})
DIRECT NUMERICAL SIMULATION OF THE WALL-PRESSURE

We employ a direct numerical simulation code to extract the pressure fluctuations at the wall in a wall-bounded channel flow (i.e., prototype boundary layer) model. The model resolves essential turbulent scales, with no subgrid decomposition required for Rayleigh numbers up to order 10^6. The model is used for diagnostic evaluation of the wall-pressure fluctuations in wall-bounded shear flows. We assume periodic boundary conditions in the horizontal plane. The model is based on the spectral method for streamwise and spanwise modes and Chebyshev Polynomial expansion in the normal direction. In the vertical plane, we employ collocation grid points given by \( z = n \pi / N \), where \( n = 0, 1, 2, \ldots \). For the sake of brevity we pick up at a point in the simulation where the velocity and vorticity have already been obtained (i.e., they will be assumed to be given) and our interest is to compute the wall-pressure.

To evaluate pressure at the wall we obtain the following form of the momentum equation at the wall (i.e., \( z_2 = -1 \)):

\[
0 = -\frac{\partial p}{\partial t} + u \frac{\partial u_1}{\partial x_1} + v \frac{\partial u_1}{\partial x_2} + w_1 \frac{\partial u_1}{\partial x_3} + (u_1 - \text{momentum}),
\]

\[
0 = -\frac{\partial p}{\partial t} + u \frac{\partial u_2}{\partial x_1} + v \frac{\partial u_2}{\partial x_2} + w_2 \frac{\partial u_2}{\partial x_3} + (u_2 - \text{momentum}).
\]

At the wall (\( z_2 = -1 \)), we also obtain for the vorticity, with the no-slip boundary condition imposed:

\[
\nu_3 = -\frac{\partial u_3}{\partial x_1}, \quad \nu_1 = \frac{\partial u_3}{\partial x_2}, \quad \nu_2 = \frac{\partial u_3}{\partial x_3}
\]

from (3) we get:

\[
\frac{\partial \nu_3}{\partial z_2} = -\frac{\partial u_1}{\partial x_1} \frac{\partial \nu_3}{\partial x_1} + \frac{\partial u_3}{\partial x_2} \frac{\partial \nu_3}{\partial x_2} + \frac{\partial u_3}{\partial x_3} \frac{\partial \nu_3}{\partial x_3};
\]

since \( u_1 = u_2 = u_3 = 0 \) at the wall and so, from continuity, \( \frac{\partial \nu_3}{\partial z_2} = 0 \) there also for all \( x_1 \) and \( x_2 \).

Recalling that the viscous shear stress at the wall is non-zero, the first non-zero term for gradients in \( z_2 \) go as \( \frac{\partial u_3}{\partial z_2} \), and from continuity we deduce:

\[
\frac{\partial u_1}{\partial z_2} \frac{\partial u_3}{\partial x_1} + \frac{\partial u_2}{\partial z_2} \frac{\partial u_3}{\partial x_2} + \frac{\partial u_3}{\partial z_2} \frac{\partial u_3}{\partial x_3} \neq 0
\]

Combining (4) with (1) and (2), we obtain an equation for the pressure in terms of the first and third component of vorticity:

\[
\frac{\partial \nu_3}{\partial z_2} + \frac{\partial \nu_3}{\partial x_2} = \nu \left[ -\frac{\partial u_1}{\partial x_1} \frac{\partial \nu_3}{\partial x_1} + \frac{\partial u_2}{\partial x_2} \frac{\partial \nu_3}{\partial x_2} - \frac{\partial u_3}{\partial x_3} \frac{\partial \nu_3}{\partial x_3} \right]
\]

We take the Fourier Transform of the pressure, streamwise and spanwise directions to obtain time series of the wall-pressure fluctuation. The vorticity gradients in \( z_2 \) are determined by employing the collocation differential operator, so that at \( z_2 = -1 \):

\[
\frac{\partial u_3}{\partial z_2} \bigg|_{z_2=-1} = \sum_{i=1}^{N} DIF(i) \nu_{i}(0)
\]

where

\[
DIF(i) = (-1)^i (\nu_{N} - \nu_{i}) \quad i = 0 \]

\[
DIF(i) = 2(-1)^{i+1}/(\nu_{N} - \nu_{i}) \quad i = 1, 2, \ldots, N - 1
\]

\[
DIF(i) = -(2N^2 + 1)/6 \quad i = N
\]

We conclude by showing illustrations in Figures 4 and 5 of the velocity perturbation field and corresponding T-S pressure wave at the lower wall, following implementation of conversion formula. We show here a primary T-S instability mode (2-D) through a cut along the vertical plane at a snapshot in time. Equation (7) provides a benchmark for computing the wall-pressure impinging on an array of point transducers, given a 3-dimensional velocity field.

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**Figure 4** Simulated Velocity Perturbation Field

**Figure 5** Simulated Wall-Pressure (Lower Wall Shown in Center of Box)

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**SUMMARY AND CONCLUSIONS**

Time sequences of wall pressure fluctuations were obtained from measurement probes positioned along the wall of an axisymmetric body during vertical ascent from a deep (fresh) water test basin. Digitally recorded data were processed to obtain both temporal power spectra and phase portraits to study dynamical mechanisms in a developing boundary layer. The power spectra obtained from the field measurements show amplification of the Tollmien Schlichting (T-S) modes in accordance with Orr-Sommerfeld Theory predictions, while phase portraits obtained from similar measurements show evidence of persistence of the T-S wave structure into the turbulence regime, that does not appear in the power spectra. Moreover, numerical simulation of T-S instability modes in our prototype model shows that one is able to derive wall pressures that are consistent with T-S instabilities obtained from experimental measurement.

3. Predictions were performed using the Transition Analysis Program Systems (TAPS) Code, Naval Underwater Systems Center, Newport, RI, 02841, USA.
5. Original velocity simulation code was developed by K. Breuer Center for Fluid Mech., Brown University, Current address, Dept. Aeronautics and Astronautics, MIT, Cambridge MA, Visualization graphics performed by G.T. Miller, Technology Applications, Inc., New London, CT, USA.
SOUND ATTENUATION CHARACTERISTICS OF BARRIER FOR UNDERWATER SOUND PROPAGATION

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INTRODUCTION

In recent years, marine sports in the shallow water area as skin diving or scuba diving are popularized and simultaneously marine accidents arise along with it. One of the most simple and effective ways for preventing the diving accidents is to transmit an underwater acoustic signal directly to divers. The use of an audio signal is very efficient since the divers usually have no communication apparatus. In such a case, it is fundamentally important to investigate in advance how the audio signal radiated from the sea surface is attenuated in underwater.

Investigations of underwater sound attenuation characteristics have already been done. In most of these studies, however, theoretical and/or experimental considerations have been confined to the case of free sound field. When a sound insulation barrier like a protect wall exists during the sound propagation path, underwater sound signal seems to be diffracted by it. Then, it is really necessary and worthwhile to study the sound attenuation effects by barrier so as to design the barrier for screening the ship noise and to estimate the power of a signal source to send toward the divers.

In this paper, sound attenuation characteristics of barrier for underwater sound propagation is experimentally investigated. Single diffraction by a thin barrier is first considered through a cistern experiment. Furthermore, a practical equation is derived for estimating the sound attenuation in underwater and it's validity is confirmed using the data measured in an actual shallow water area. Next, we discuss a possibility for the prediction of sound attenuations in the case of a multi-diffraction process for a wide barrier.

EXPERIMENTAL SITUATION IN CISTERN

Figure 1 shows a experimental arrangement of the barrier, sound source and receivers in cistern. Measurements were carried out for two cases as follows:

1. Single diffraction: the barrier was made from a thin aluminum plate pasted both side by a polyethylene mat. Thickness of the barrier has become to be 1.5 cm. The sound source (denoted by S1) locates at the same level with top of barrier so as to avoid the influence of the thickness of barrier.
2. Double diffraction: the barrier was substituted by a square lumber shaped the corner exactly. As the thickness of barrier affects sensitively to the experimental results. Thickness of the barrier was 3 cm. The sound source (denoted by S2) locates the position entered 5 cm inside against the level of barrier edge.

In both cases, a sinusoidal wave (frequency is 70, 100, 150 kHz) with modulated pulse amplitude was adopted as a sound signal source. The receiver R was set up especially, so reference waves can be obtained from this in the absence of barrier. Accordingly, distance of S1R and S2P is equal in each other. Sound sources S1, S2, and receivers R, P are fixed 15 cm depth under the surface of water. The path length deferences were changed variously by moving the two receivers (P, R) at the same time (described by arrows in Fig.1). The path length defences were determined by measuring the distance y in Fig.1 accurately.

RESULTS AND DISCUSSION

SINGLE DIFFRACTION

In order to compare the attenuation characteristics in underwater with those in air, the experimental data was plotted in the so-called Mackawa's chart. Figure 2 indicates the summarized chart where all experimental results obtained in this study have been plotted. As is well known, the vertical axis in the Mackawa's chart shows the deference of sound level at receiver between the cases that the barrier exists or not, and the horizontal axis the Fresnel number N which is the deference of sound propagation path length divided by a half wavelength. Full line A in Fig.2 indicates the sound attenuation characteristics in the air obtained by Dr. Mackawa and dot line shows ones derived theoretically by well-known Kirchhoff's approximations. From the experimental results in the

Fig.1 The experimental situation for observing the sound attenuation effect.

Fig.2 Sound attenuation characteristics by a thin barrier for underwater sound propagation.
The values of $a$ and $b$ in eq. (1) were determined to be 13.1 and 16.3, respectively, by the least squares method. Dot and chain line in Fig. 2 is the calculated curve with the parameters obtained above in the region $N \geq 20$.

In order to confirm the validity of eq. (1), acoustic measurements were made in the actual sea area. The experimental procedure was the same as in cistern except for the frequency of sound signal. Experimental results obtained at several propagation path lengths are shown in Fig. 3. It is found that dot and chain line grasps roughly the tendency of experimental points though the number of data is not many. This fact implies that eq. (1) is also fully useful for estimating the sound attenuations in actual sea area.

**CONCLUSION**

Sound attenuation characteristics of barrier for underwater sound propagation was experimentally investigated. Results thus obtained must be helpful to create the underwater acoustic transmission system in the future. Furthermore, a useful equation was derived for the prediction of sound attenuations.

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STUDY ON SEABED TYPE DISCRIMINATION USING PULSE ECHO METHOD AND DATA PROCESSING

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1. INTRODUCTION

Recently, salvage, cable laying and dredging require seabed discrimination with the need to increase efficiency and reduce danger in the operation areas.

The authors devised a new method of seabed material discrimination by using pulse echo method and multi-layered neural network. In this method, the first echo signal returning from seabed are received by five hydrophones set in line array and structural features of the echo signals due to the changes in seabed conditions are extracted by electrical processing. Processing them by the multi-layered neural network, the seabed type discrimination are performed. Then, experiments to confirm the validity of this method were performed in the Uchiura Bay (Nagano, Shizuoka, Japan). And, satisfactory results were obtained.

The paper describes the details of the new method of seabed type discrimination and the results of the experiments.

2. MULTI-LAYERED NEURAL NETWORK

This neural network is Perceptron type which consists of three layers (input, hidden & output layers). Figure 1 shows the structure of multi-layered neural network. In input layer, five units are set in accordance with the number of hydrophones, and three units are set in output layer according to the number of seabed types to be discriminated. Units of each layer are coupled with all the units of preceding layer through synapses weight. Figure 2 shows the structure of the unit. Output value of the units of each layer are obtained by multiplying the input output value of the units of preceding layer by synapses weight and summing them up, adding offset value to it and calculating by the input-output function of the unit. As for the input-output function, Sigmoid function is adopted which is shown in Fig. 3.

In the learning process of multi-layered neural network, the offset value of synapses weight and unit are corrected to make output value of the output layer converge at the teaching signal (evaluating signal of the learned results). The correction was performed by using the error back propagation algorithm. The teaching signals given for the three seabed type are shown in Table 1. For the offset initial value of synapses weight and unit, random numbers in the range from zero to one were given. Then, the learning data in the case that the seabed is rock are input to the input unit, and (1,0,0) as for teaching signals are given. Input data is successively calculated to find the output value of each layer according to the offset value of synapses weight and unit. The calculated value in the output layer is compared with the teaching signal and the offset value of synapses weight and unit is corrected for the output value of the output layer to converge at teaching signal (1,0,0). Next, giving teaching signal (0,1,0) for the sand and (0,0,1) for the silt, similar calculation as mentioned above is done. Thus discrimination of seabed of rock, sand and silt is made. Discrimination of seabed is made by using the multi-layered neural network.
The measuring method will be described below. Firstly, the transmitter radiates an AC pulse with carrier frequency of 50kHz and pulse width 400 μsec. Figure 6 shows an example of radiated signal waveform. The echo signals from the seabed are received by five hydrophones set in line array. The echo train envelope is obtained by converting received echo signals with DC converter and recorded on a data recorder. When it is being reproduced, the echo train envelope is converted to digital signal and fed to a computer. Figure 7 shows one example waveform of echo train envelope. As is evident from Fig.7, time width of the received echo signal is confined within 1 μsec, when the seabed material is sand and silt. However, it extends over 5 μsec, when the seabed material is rock. Then, the integrated area (all of 5 μsec) under the echo train envelope representing the first echo is exploited for “intensity” of the first echo signal. As the intensity of the first echo signal provides an index of seabed “roughness” or “hardness”, it can be one of the useful informations to discriminate the seabed conditions. Measurements were made 50 times to one measuring area.

![Image](image.png)

**Fig. 1** Position of transmitter and hydrophone array

4. THE RESULT OF MEASUREMENT AND DISCRIMINATION

Figure 8 shows one example of intensity of echo signals measured at three kinds of seabed. The intensity of echo signals measured at each seabed considerably scatters and lies one upon another. It is considered that the continuous fluctuation by choppiness of sea in the array of the transmitter and the hydrophones influences radiating and receiving of the ultrasonic pulse. So, this scattering was reduced by averaging the intensity of received echo signals at similar point. One example of the averaged intensity of seven data is shown in Fig.9. Therefore, the average value of the intensity of the received echo signals was fed to the computer software (multi-layered neural network) as learning data. The learning was repeated 500 times.

Discriminating ratio was examined using the data obtained at the point where the discrimination seemed to be difficult because the difference in intensity level of echo signal among the three kinds of seabed was small. An example of distribution of discriminated data is shown in Fig.10. The number of data that could make satisfactory result by learning and discrimination using multi-layered neural network was inquired by changing the number of data to be averaged. The result is shown in Fig.11 as the relationship between the number of data to be averaged and the discriminating ratio.

This figure led to the following: in discrimination of three kinds of seabed (rock, sand & silt), satisfactory discrimination is made if calculation to discriminate is done using the average value of seven data.

5. CONCLUSION

A new method of seabed type discrimination using pulse echo method and multi-layered neural network was proposed. And the usefulness of this seabed material discrimination method was verified by experimental results.
THE SOUND SCATTERING EXPERIMENT OF THERMOCLINE IN SHALLOW SEA

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I INTRODUCTION

The temperature, salinity and pressure of sea water change with depth, which forms the sound velocity vertical distribution in the sea. The variation gradient of impedance profile in the sea can be calculated from signals scattering from the nonuniform media. Since the density of sea water does not change much from layer to layer in the shallow sea, so the sound velocity mainly depends on the temperature of the sea water, which provides a possibility by using the acoustic technique to measure the depth of thermocline.\(^\text{1}\)

In 1984 B. C. Grachev et al did experiments on sound scattering record of fine structures in deep thermocline.\(^\text{2}\) In 1986, B. Colbourne and A. E. Hay recorded the scattering signals, at 192KHz frequency and 47 M depth, from the thermocline layer.\(^\text{3}\)

In this paper, the experiments were done at two stations, depth 34 M and 44 M respectively, in the Southern Yellow Sea from July 26 to August 2 of 1988 and 1989. The center frequencies of sound sources were 120 KHz, 160 KHz and 200 KHz and impulse widths were 0.46 ms and 1.0 ms respectively. During the experiments, CTD was used to measure the temperature, salinity of different depth. Plankton samples were also hauled layer by layer to get the vertical distribution of phytoplankton and zooplankton.

The experimental results showed that the scattering signals can be received with the center frequency at 200 KHz and the impulse width at 0.46 ms when the temperature gradient of the thermocline is greater than 1.3°C/M, while those of 160 KHz were weak. The scattering signals were not significant difference between two signals at width 0.46ms and 1.0ms respectively. The analysis of the plankton samples showed its density in the thermocline is not certainly greater than the upper or lower layer, and it is however closely related to the horizontal distribution of water mass and dominate species. On the other hand, the amplitude variations of scattering signals by digital simulations match the experimental records. It therefore proves that scattering signals are caused by the nonuniformity of medium parameters such as thermocline.

II EXPERIMENTAL SETUP

The experimental system is showed in Fig. 2.1. Four transmitting-receiving transducers were used, two of them at 200 KHz with sensitivities 187 dB (directional angle of 6°) and 171 dB (directional angle of 10°). The located position of the transducers according to the interference of transmitting sources and the oceanic condition. After being formatted by BAK 4440 Gating System, signals were sent to a impulse power amplifier. The receiving signals were sent to the tuned detector-amplifier with gain 60dB and band width 2 KHz. The signal-noise ratio was greater than 5 while input signals were 200μV. Finally, the output signals of the tuned detector-amplifier were displayed by HP3562A and stored by HP9163C. While recording the scattering signals, CTD was used to record the temperature, salinity at different depth every four hours, and in the meanwhile the plankton were sampling around the thermocline.

III ANALYSIS AND CONCLUSION

The sound scattering of the thermocline in shallow sea were caused by the non-uniform structure of media or the plankton living around the thermocline. Due to the unsteadiness of thermocline, the influence of water mass circulation as well as the limitation of sampling technique, one do not

<table>
<thead>
<tr>
<th>Sample</th>
<th>Date</th>
<th>Time</th>
<th>Depth (m)</th>
<th>Temperature (°C)</th>
<th>Salinity</th>
<th>Plankton Density</th>
</tr>
</thead>
<tbody>
<tr>
<td>Sample</td>
<td>01</td>
<td>02</td>
<td>03</td>
<td>04</td>
<td>05</td>
<td>06</td>
</tr>
<tr>
<td>Sample</td>
<td>01</td>
<td>02</td>
<td>03</td>
<td>04</td>
<td>05</td>
<td>06</td>
</tr>
</tbody>
</table>

\(^*\) * means the recorded signals are not clear enough to tell the time when the scattering signals were recorded in the experiment.
have enough sufficient data to make a conclusion on the formation mechanism. To search for the formation mechanism on thermocline, besides the plankton samples collected at site to use in the analysis, according to E. Kind’s method of reflection coefficient spectrum of plane wave in randomly distributed, nonuniform layered system, we did the numerical simulation to get scattering signals for comparison to the recorded signals.

Table 2

<table>
<thead>
<tr>
<th>Name</th>
<th>Plankton</th>
<th>Total (10⁶)</th>
<th>Name</th>
<th>Plankton</th>
<th>Total (10⁶)</th>
</tr>
</thead>
<tbody>
<tr>
<td>00012</td>
<td>Phytoplankton</td>
<td>270.02 126.9</td>
<td>00012</td>
<td>Phytoplankton</td>
<td>270.02 126.9</td>
</tr>
<tr>
<td>00022</td>
<td>Zooplankton</td>
<td>130.99 87.77</td>
<td>00022</td>
<td>Zooplankton</td>
<td>130.99 87.77</td>
</tr>
<tr>
<td>00032</td>
<td>Phytoplankton</td>
<td>874.11 366.6</td>
<td>00032</td>
<td>Phytoplankton</td>
<td>874.11 366.6</td>
</tr>
<tr>
<td>00042</td>
<td>Zooplankton</td>
<td>159.47 107.4</td>
<td>00042</td>
<td>Zooplankton</td>
<td>159.47 107.4</td>
</tr>
<tr>
<td>00052</td>
<td>Phytoplankton</td>
<td>1228.99 624.0</td>
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<td>Phytoplankton</td>
<td>1228.99 624.0</td>
</tr>
<tr>
<td>00062</td>
<td>Zooplankton</td>
<td>117.71 80.1</td>
<td>00062</td>
<td>Zooplankton</td>
<td>117.71 80.1</td>
</tr>
<tr>
<td>00072</td>
<td>Phytoplankton</td>
<td>48.09 29.26</td>
<td>00072</td>
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<td>48.09 29.26</td>
</tr>
<tr>
<td>00082</td>
<td>Zooplankton</td>
<td>324.52 153.7</td>
<td>00082</td>
<td>Zooplankton</td>
<td>324.52 153.7</td>
</tr>
</tbody>
</table>

The dominant species of phytoplankton were Pyramphora, Rhizosolenia, Corenoxon, and Mixtilina; their body width was about 8–15 μm and length about 18–100 μm.

The dominant species of zooplankton were Sagitta, Coepepod and its larva; their body lengths was about 1–4 mm.

The total amounts of plankton in thermocline was more than other only in this example.

The results from digital simulations comparing to those from experiments are shown in Tab. 1. The time for most scattering signals reaching the receiver matches the time obtained by digital simulations. Table 2 shows the vertical density distribution of plankton around the thermocline. Both the total amounts of plankton and its dominant species in the thermocline were less than those in the upper.

Fig. 3.1 shows the records, in 30 times average of the scattering signals at the frequencies 200 KHz and single recorded of 200 KHz, 150 KHz and 120 KHz (Δφ = 6°, Δφ = 10° and pulse duration 0.46 ms) respectively. The pulse on the left of the recorded signals were the trigger pulse and the obvious strong pulse on the right were the echoes of sea bottom, and in the middle were the scattering signals from non-uniform layers. From Fig. 3.1 one can see that the scattering signals were clear and the amplitude variations of the scattering signals match those of the digital simulations. In Fig. 3.2, the main thermocline was at 21 meters below the sea surface, and the sound velocity gradient in the upper layer was negative. The recorded scattering signal was smaller than those from the digital simulation, which was when main thermocline was deep, the media absorption and the diffusion loss should be considered. Sometimes the signals would be with radar disturbance, while upper the second oceanic condition. Fig. 3.3 showed a typical record made by using double frequency sonar in 1988.

The results of experiments showed that the reaching time of scattering signals matched the time obtained by digital simulation, and the total amount of plankton in the thermocline was lower than those either upper or lower layer (see Tab. 2). It can be concluded that the major mechanism that causes sound scattering is the nonuniform media distribution in shallow sea. This makes a possibility to use acoustic technique for the measuring the depth of nonuniform layers, such as thermocline.

ACKNOWLEDGMENT

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REFERENCES

Consider the case of plane wave incidence on an elastic sphere lying near the interface of fluids as shown in Fig. 1. Let radius of sphere be \( a \), density and sound speed of upper and lower medium be \( \rho_1, c_1 \) and \( \rho_2, c_2 \) respectively, density and Lamé constants of elastic sphere be \( \rho, \lambda, \mu \). Use cylindrical coordinate system where plane \( z = 0 \) coincides with the interface, and \( z \)-axis passing through the center of sphere. If \( \beta_1, \beta_2 \) are the corresponding incident and reflected angle of incident wave, \( \theta_1 \) be the height of sphere's center, potential functions of sound field in upper and lower medium can be written as following with the help of multipole expansion of general spherical waves(1):

\[
\phi = -i e^{-i(kz - \omega t)} \sum_{n=0}^{\infty} \sum_{m=-n}^{n} (2\beta_{n,m}) A_{mn} \cos m\alpha (\cos \theta_1(k, z)P_{n}(\cos \theta_1) + \Psi_{n,m}(z)),
\]

where the time factor \( e^{-i\omega t} \) had been omitted, and \( n, \theta, \alpha \) are spherical coordinates with origin at center of sphere, \( V \) and \( D \) are coefficients of reflection and refraction of plane incident wave on interface functions \( \Psi_{n,m}, \Psi_{n,m}, \Psi_{n,m}, \Psi_{n,m} \) are the corresponding reflective and refractive field potentials of general spherical wave order \( n, m \). Their expressions as well as high frequency approximations are

\[
\Psi_{n,m} = \frac{i}{-i} \int \frac{d\nu}{\nu_1} \frac{1}{\nu_1} \frac{2 \nu_1 - \nu_2}{\nu_1 - \nu_2} \left( \frac{\nu_1^{1/2}}{\nu_1^{1/2}} \right)^{1/2} J_n(k_1 \nu_1) P_n(\cos \theta_1(\nu_1)) \frac{d\eta}{\eta}
\]

The analytic solution of given problem is too complicated to get any physical understanding, numerical calculation and model experiment had been done, using steel spheres with radius from 0.6cm to 1.0cm. Frequency of incident wave was chosen to be 412 kHz. The upper and lower medium were Diesel oil and salty water respectively. Following remarks may be done from the work:

1. When frequency of incident wave is sufficiently high, the behavior of elastic sphere is approximately...
the same as a fluid one with the same density and sound speed equal to longitudinal wave speed of elastic sphere which can also be seen from the high frequency approximation of $u$ and $v$ given above.

2. The interference of incident and reflected waves on interface form a partially standing wave field in the direction of $x$-axis, so it may be expected that when $x_0$ changes monotonically, intensity of scattered field receiving at some fixed position will fluctuate with the change of $x_0$. In Fig.3a the experimental results for steel sphere with radius 1.27cm were given for two different incident angle. In Fig.3b experimental and calculated values of backscattering strength of steel sphere with radius 0.9cm at incident angle 81.6° are given for different $x_0$. The discrepancy between experimental and calculated values is not only due to the approximations used in calculation and the unavoidable errors in experiment, but also due to limited impulse length used during experiment that only first circumferential wave contributes to the scattered field measurement.

3. The series expansion of incident wave in $x_0$ converges quickly with increase of $n$, therefore the values of $A_n$ (or $B_n$) determined from minimum of $W$ will decrease quickly with increase of $n$, which implies that the azimuth variation of scattered field will be dominated only by terms of small $n$. Fig.4 gives the experimental results of scattered field measurement at different azimuth angle for four values of $x_0$ which all give good agreement with the theory.

4. Fig.4 gives the change of backscattering strength with radius of sphere, it had been found both by experiment and by theoretical calculation that when $x_0$ is sufficiently large the frequency dependence of backscattering strength will be weak.

LITERATURE
THE EFFICIENT CALCULATION AND DISPLAY OF DISPERSION CURVES FOR A THIN CYLINDRICAL SHELL IMMERSED IN A FLUID

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Much of the extensive literature on shell theory centers around the computations of dispersion curves, usually thought of as the roots of a determinant $D$ which relates the angular frequency $\omega$ to the wave numbers for free waves traveling on the shell. For a cylindrical shell, the free wave paths are actually helical curves on the surface of the shell, having both a circumferential and an axial wave number. In this case, it is useful to indicate explicitly that the dispersion relation is a function of three variables, i.e., $D(\omega, m, n) = 0$, where $m$ is the number of half wavelengths of a wave traveling in the axial direction, and $n$ is the number of full wavelengths of a wave traveling around the shell. We shall refer to this function $D(\omega, m, n)$ as the dispersion volume density, since it is a function of three variables. Normally, one sees various two-dimensional displays of the dispersion relation, which are created by plotting any two of these variables against each other, while holding the third variable fixed, as illustrated in Fig. 1.

Dispersion relations for a fluid-loaded cylindrical shell have been developed previously by many authors. However, as pointed out in a lengthy article by Scott [1], some of these previously published works (e.g., [2] and [3]) made the mistake of searching for roots of the real part of the dispersion equations, which led to the computations of spurious non-physical branches of the dispersion curves. Although Scott has convincingly made this point, his actual computation procedure is cumbersome. In fact, it is difficult to determine the complex roots of such dispersion relationships by any root-following method. In the present work, the complications of calculating and following these complex roots have been simplified by using a regular computational grid for each parameter of the dispersion function, and then making the root loci evident by the appropriate use of graphics to display what we term the dispersion volume density.

The purpose of this paper is to describe and illustrate an efficient procedure for calculating and displaying the dispersion curves. We will base the results on Fliggie shell theory, which has some advantages for air-backed thin steel shells immersed in water. Only the critical steps and results in the derivation of the dispersion relationships will be given here.

The displacements and forces are expanded in a combined Fourier series and transform to allow for completely general motions due to asymmetric forcing functions. With a time-harmonic factor of $e^{i\omega t}$ suppressed, the displacements $u^r, u^f, u^z$ of the mid-surface of a straight-sided cylindrical shell of mean radius $a$ and thickness $h$ are a function of $z$ and $\phi$ only. In general, the forcing function $F$ will also have three degrees of freedom, but $F$ will be identically zero for a shell in vacuo, and $F^r = F^f = 0$ when only the force causing or resisting motion of the shell is $p_o(x, \phi, z)$, the total (incident plus scattered) pressure on the shell due to an incident acoustic wave. The shell material will be characterized by three parameters: Poisson's ratio $\nu$, Young's modulus $E$, and the density $\rho_s$, the fluid by a density $\rho_f$ and a speed of sound $c_f$. Other (non-dimensional) parameters describing the solution are

- thickness/mean radius = $h/a$
- length/mean radius = $l/a$
- thickness parameter $\beta^2 = h^2/(2a^2)$
- $s = z/a$
- $\Omega = k_o = \omega/c_f$
- $\sigma = E/\rho_s(1-\nu^2)$
- $\gamma = r/a$.

We assume that $u^r(\phi, s)$ can be expanded in a Fourier series in $\phi$, i.e.,

![Fig. 1](image-url)
\[ u'(\phi, s) = \sum_{n=-\infty}^{\infty} \psi_n(s) e^{-i\alpha n} \]

and that the axial variation of \( \psi_n(s) \) can be expressed as a Fourier transformation

\[ \psi_n(s) = \int_{-\infty}^{\infty} U_\zeta(\zeta) e^{-i\alpha \zeta} d\zeta, \]

where the transform variable \( \zeta \) is a non-dimensional axial wavenumber. Similar expansions are assumed for \( u'(\phi, s) \) and \( u'(\phi, s) \). Because of the orthogonality properties of the Fourier expansions, the application of the operators to the expansions for the displacement results in an infinite set of matrix equations for each \( \zeta \), one for each harmonic order \( n \). The differential operators all become matrix operators, resulting in a set of coupled linear equations. We also represent the acoustic pressure in a double Fourier expansion which allows us to express the forcing term as a function of \( U_\zeta(\zeta) \), the double Fourier expansion of the radial displacement. This permits us to move the forcing term in the transformed equations over to the left-hand side and write a completely homogeneous matrix equation for the Fourier components of shell displacement:

\[
Ku = \begin{bmatrix}
K_{20} & K_{21} & K_{22} \\
K_{10} & K_{11} & K_{12} \\
K_{00} & K_{01} & K_{02}
\end{bmatrix} \begin{bmatrix}
U_\zeta(\zeta) \\
\frac{\partial U_\zeta(\zeta)}{\partial \zeta} \\
\frac{\partial^2 U_\zeta(\zeta)}{\partial \zeta^2}
\end{bmatrix} = \begin{bmatrix}
0 \\
0 \\
0
\end{bmatrix}.
\]

(3)

The components of \( K \) in equation (3) are given by

\[ K_{00} = -\frac{c^2}{2} \frac{(1-\nu)}{n^2} + (k_0)^2 - \frac{\beta^2}{2} \frac{(1-\nu)}{n^2} \]

(4)

\[ K_{01} = K_{10} = -\frac{\rho_c c}{2} \frac{1+\nu}{n} \]

(5)

\[ K_{02} = K_{20} = -\alpha c \zeta - i \zeta^2 + \frac{\kappa c}{2} (1-\nu) \]

(6)

\[ K_{11} = -\frac{(1-\nu)}{2} \zeta^2 - n^2 + (k_0)^2 - \frac{3(1-\nu)}{2} \zeta^2 \]

(7)

\[ K_{12} = K_{21} = -\frac{\zeta^2}{2} \]

(8)

\[ K_{22} = 1 + \beta^2 (n^2 + \zeta^2) - (k_0)^2 + \beta^2 (1-2n^2) - \frac{\Omega^2 \kappa c}{\rho_s h} + \frac{\Omega^2 \kappa c}{\rho_s h} \]

(9)

in which the additional terms to account for the resistance and reactance components of the fluid loading are

\[ r_s(\zeta) = \omega \rho_a \text{Im} \left[ \frac{H^{(1)}(\gamma)}{y H^{(2)}(\gamma)} \right], \]

(10)

\[ m_s(\zeta) = -\rho_a \text{Re} \left[ \frac{H^{(1)}(\gamma)}{y H^{(2)}(\gamma)} \right], \]

(11)

and \( y = \sqrt{(k_0)^2 - \zeta^2} \) for frequencies such that \( k_0 \geq \zeta \). When \( k_0 < \zeta \), the so-called evanescent region,

\[ r_s(\zeta) = 0 \quad \text{and} \quad m_s(\zeta) = -\rho_a \frac{K_n(\zeta)}{\kappa K_n(\zeta)}, \]

(12)

where \( z = \sqrt{(k_0)^2 - \zeta^2} \).

Equation (2) can have a non-trivial solution whenever its determinant vanishes, which can occur for particular combinations of frequency \( \Omega \), axial wavenumber \( \zeta \), and circumferential wavenumber \( n \) when there is no fluid loading. These combinations represent the dispersion relations for the shell in vacuo. Fluid loading introduces (acoustic) damping or loss. Even in this case, the absolute value of the determinant can become very small, and the resultant forced solution can be very large for particular combinations of \( \Omega \), \( \zeta \), and \( n \). Thus minima of the absolute value of the determinant of \( K \) will define the dispersion relations for the fluid-loaded case.

The simplified computational procedure is to compute the coefficients of the matrix \( K \) and then evaluate the determinant of \( K \) at regular increments in the principal variables which are the non-dimensional frequency \( \Omega \), axial wavenumber \( \zeta \), and circumferential wavenumber \( n \). The power of the method lies in the fact that modern graphical workstations with appropriate visualization software can not only perform these calculations but can quickly display the results in a manner which allows one to understand the dispersion relations in a global sense (over the entire meaningful range of \( \Omega \), \( \zeta \), and \( n \)). The actual program code was written in an efficient high level mathematical language.

Once the dispersion volume density is calculated for an appropriate range and density of \( \Omega \), \( \zeta \), and \( n \), the results can be displayed in a variety of ways. One of the simplest and most effective means is to compute the logarithm of the reciprocal of the absolute value of \( D \), and then create a smooth interpolated color image of the results. If the color table is properly chosen, this has the effect of showing the root loci as lines whose brightness indicates the nearness to zero and whose width is indicative of width of the null and of the resolution in the appropriate parameters (\( \Omega \), \( \zeta \), and \( n \)). If calculated in this way, it is also possible to animate the display of these color images, which enhances the understanding of how the location of the modes change with changes in various parameters. All of these ideas will be illustrated in the presentation of this paper, but cannot be adequately shown in this brief black and white hardcopy.

In summary, the present work confirms the earlier work of Scott as to the correct dispersion loci, but provides a simpler and more easily interpreted method of calculating and displaying the dispersion volume density over the entire relevant range of frequency and wavenumbers.

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A METHOD FOR CALCULATION OF ACOUSTIC FIELD OF A CYLINDRICAL PROJECTOR WITH ONE END CONICALLY BAPPELED

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INTRODUCTION

In the underwater acoustics equipment used in transponders or moving bodies, a projector is mounted on one end of a cylindrical tube. From this configuration, the reflective wave from the tube end is induced, and the directivity pattern of the projector is degenerated. To minimize this degeneration, a correctly shaped reflector (baffle) changing the direction of the reflected wave is usually introduced between the projector and the tube. However, since there are few methods for calculating the acoustic field of such a projector, how to design the reflector shape to obtain a desirable directivity pattern is not yet completely obvious. Hence, in this paper, we suggest a useful method of calculating the acoustic field so that we can design the vertical directivity pattern of the cylindrical projector with a conical shape reflector. At first, we explain how to derive the formula for the proposed method, then compare the calculated results of the proposed method with the actual experimental results, and prove that this proposed method is correct. Finally, we discuss about vertical directivity patterns of the projector with the conical reflector's slope varied.

TEORETICAL ANALYSIS

Calculation of Acoustic Field by the Least Squares Method

Let us consider the axisymmetrical vibration of a cylindrical projector mounted on a usual axisymmetrical rigid structure in the infinite fluid region (as shown in Fig. 1). The center of the cylindrical projector is placed at the origin of an orthogonal coordinate system \((x, y, z)\) as shown in Fig. 1, and an arbitrary point \(Q\) in the fluid region is represented by a cylindrical coordinate \((r, \theta, z)\). It is assumed that the projector should vibrate radially with a constant amplitude \(F\) and with an angular frequency \(\omega\), and the other structure except for the projector itself should be completely rigid. The position vector connecting with the origin and an arbitrary point \(P\) on the external surface of either the structure or the projector is specified as \(S\) and the normal vibrating velocity of the projector is denoted as \(N(S)\). Hence, the velocity of the projector \(V(S)\) is expressed as \(\omega F S\). If the velocity potential generated in the infinite fluid region is named as \(\psi(r)\), where \(r\) is the position vector directed to arbitrary point \(Q\) in the fluid region, \(\phi(r)\) is given by the solution of the wave equation which satisfies the radiation condition in the infinite region as follows

\[
\psi(r) = \sum_{n=0}^{\infty} a_n h_{in}^0(\kappa r) P_l(\cos \theta) \tag{1}
\]

where, \(h_{in}^0(\kappa r)\) is the \(n\)th order spherical Hankel function of the second kind, \(P_l(\cos \theta)\) is the \(l\)th order Legendre function of the first kind, \(\kappa\) is the wave number, and \(a_n\) is the unknown constant to be determined by the boundary conditions. Moreover, if the normal particle velocity on the surface of either the structure or the projector is called \(V(S)\), then by using Eq. (1) based on the theory of velocity potential, we can obtain

\[
V(S) = -\frac{\partial \psi(r)}{\partial r} \bigg|_{r=S} = \sum_{n=0}^{\infty} a_n \phi_n(S) \tag{2}
\]

where,

\[
\phi_n(S) = \frac{\partial}{\partial r} \left( h_{in}^0(\kappa r) P_l(\cos \theta) \right) \bigg|_{r=S} \tag{3}
\]

where, \(\phi_n(S)\) means the unit vector outward normal direction. In general, the unknown constants \(a_n\) cannot be strictly obtained, unless the projector has a simple geometry, such as an infinite cylinder or a sphere. On the contrary, for the projector with a shape of finite cylinder or with a part of a sphere, the method to obtain these constants by using the least squares method was proposed. That is the method to obtain \(a_n\) from the minimization of the square error between \(N(S)\) and \(\phi_n(S)\) overall the surfaces. The square error has the form

\[
E = \sum_{i=0}^{\infty} \left( |N(S_i)| - |V(S_i)| \right)^2 ds \tag{4}
\]

where \(ds\) is the differential surface element. The minimization condition \(\partial E / \partial a_i = 0\) leads to the following simultaneous equations with respect to \(a_i\) (\(i = 0, 1, \ldots, L\), \(L\) is the truncated term number)

\[
\sum_{n=0}^{\infty} \phi_n(S_i) = (\phi_n(N(S))) \tag{5}
\]

where \(\phi_n(N(S)) = \int \phi_n(S) N(S) ds \tag{6}\)

where, \(i = 0, 1, \ldots, L\) and \(\phi_n^*\) indicates a complex conjugate. But, this method is not applicable to the acoustic field calculation for more complicated projector.

In this paper, in order to carry out the integration of Eq. (1) and (4) for an arbitrary boundary shape, we proposed to divide the surface into several conical elements, and derive the formula for \(\psi(r)\) and \(V(S)\) in each divided element region.

Deriving the formula for 1 and 4 to 6

As shown in Fig. 3, the region is divided into 1 element regions, and the border points of the regions are defined as \(P_1, \ldots, P_k\). The symbols \(e_1, e_2, e_3\) are defined respectively as unit vectors towards axial, radial, and circumferential in the cylindrical coordinate system. Because we discuss the axisymmetrical vibration, \(\partial \phi \partial \theta \partial r\) in Eq. (2) is given in every region by using only unitvectors \(e_1, e_2, e_3\) as follows

\[
\frac{\partial \phi}{\partial m} = \left( \frac{\partial \phi}{\partial r}, \frac{\partial \phi}{\partial \theta}, \frac{\partial \phi}{\partial z} \right) \cdot n \tag{7}
\]

After obtaining \(\partial \phi / \partial r, \partial \phi / \partial \theta, \partial \phi / \partial z\) from Eq. (1), and by substituting these into Eq. (4), we can obtain relative terms of coefficients \(a_i\), then \(\psi(r)\) and \(V(S)\) can be obtained as follows

\[
\begin{align*}
\psi(S) &= \left( [\frac{\pi}{4} \omega h_{in}^0(\kappa r) P_l(\cos \theta) + \frac{\partial}{\partial r} h_{in}^0(\kappa r) P_l(\cos \theta)] e_1 + \frac{\partial}{\partial r} h_{in}^0(\kappa r) P_l(\cos \theta) \right) \cdot n \tag{8}
\end{align*}
\]

Fig.1 Configuration of the cylindrical projector mounted on an axisymmetrical rigid structure

Fig.2 Diagram of the structural regions (cross section with z axis included)
where, the primes indicate differentiation with respect to the arguments of the functions. In this proposal, we represent \( \phi_i (x) \) for the regions 1, i.e., 1 = 1, 2, \ldots, 11 as \( \phi_i (x) \) and \( \phi_i (y) \) is formulated as follows.

\[
\phi_i (x, y) = \phi_i (x, y) + \phi_i (x, y) P_i (x, y) E_i
\]

where,

\[
\begin{align*}
\phi_i (x) &= \cos \left[ \tan^{-1} \left( \frac{r \cos \phi_i - r_i \cos \phi_{i-1}}{r \sin \phi_i - r_i \sin \phi_{i-1}} \right) \right] \\
\beta_i (x) &= -\sin \left( \tan^{-1} \left( \frac{r \cos \phi_i - r_i \cos \phi_{i-1}}{r \sin \phi_i - r_i \sin \phi_{i-1}} \right) \right) \\
\gamma_i (x) &= \frac{r \cos \phi_i - r_i \cos \phi_{i-1}}{r \sin \phi_i - r_i \sin \phi_{i-1}} \\
E_i (x) &= \cos \left[ \tan^{-1} \left( \frac{r \cos \phi_i - r_i \cos \phi_{i-1}}{r \sin \phi_i - r_i \sin \phi_{i-1}} \right) \right] \\
+ \sin \left[ \tan^{-1} \left( \frac{r \cos \phi_i - r_i \cos \phi_{i-1}}{r \sin \phi_i - r_i \sin \phi_{i-1}} \right) \right] \xi
\end{align*}
\]

\( r, \theta \) are the spherical coordinates of border point \( i \) \( (i = 0, 1, \ldots, 11) \). And the differential surface element for the divided region \( (i = 1, 2, \ldots, 11) \) is shown as follows.

\[
dS = \sqrt{\frac{r_i^2 \sin \phi_i (\theta - \theta_i) \sin \phi_i (\theta - \theta) + r_i \cos \phi_i (\theta - \theta) + r \cos \phi (\theta - \theta)}{r_i \sin \phi_i - r \sin \phi_i}} \, \sin \phi \, \cos \phi \, d\theta
\]

According to Eqs. (10) and (15), it is possible to carry out the surface integral at the region closed by arbitrary polygonal lines. Hence, we substitute Eqs. (10) and (15) in Eqs. (6) and (7), and carry on integration for \( \theta \) and \( \phi \). If we solve Eq. (5), then we can obtain the unknown constants \( \alpha_i \). And the acoustic pressure \( p(x) \) generated by the projector at an arbitrary point \( Q \) in the field can be obtained by the following equation.

\[
p(x, y) = \rho_s \sum_{i=0}^{11} \phi_i (x, y) P_i (x, y) E_i (x, y)
\]

where, \( \rho_s \) is the density of fluid.

**NUMERICAL CALCULATION EXAMPLES AND DISCUSSION**

To evaluate this method, we compared the experimental and calculated results of directivity for the cylindrical projector with a conically shaped reflector. As for the frequency, we selected 31 kHz available for the equipped apparatus. The dimensional shape of the projector used for this experiment is shown in Fig. 3. The calculated result and the experimental result are shown in Fig. 4. The difference between the experimental value and the calculated value is less than 5 % at about 90° direction. One of the reasons is that despite of being made of hard polyvinyl chloride, the conical baffle is treated as a rigid body in the calculation. In other words, the condition that particle velocity on the conical surface always remains zero is not satisfied. In addition, there is still some difference in the experimental and calculated results, as shown at 180° area. We think this is caused by the different treatment in analysis regarding the boundary condition at the underneath of the projector. That is, a rubber-molded cable terminal happens to protrude from the underneath of the experimental projector. However, the characteristics such as the sharpness of directivity as shown in the experimental result are well reproduced in the calculated result. So this proves that this analytical calculation method is correct. Moreover, we carried out calculations to investigate the influence of the conical part slope on the directivity. Here, as shown in Fig. 5, we indicate a calculation example with the slope angle only varied while the length of conical side line \( \omega \) remains constant. For the cases of different slopes an \( \omega, \theta, \phi \) in Fig. 5 a, the calculated results of vertical directivity at f = 31 kHz are shown in Fig. 5 b. As the slope angle from z-axis becomes larger, the directivity level becomes lower at about 180° direction. On the other hand, as the slope angle becomes smaller, the directivity becomes omni directional. Also, within an angle less than 90°, the change in directivity level by the slope angle is very insignificant. Directivity pattern in case of changing the conical baffle side slope like this way has not been able to be estimated without formerly making a sample model. But our method can calculate the directivity of a projector not only combined with a conical baffle for variable conical slope angles, but also combined with an arbitrarily shaped baffle expressing the surface approximately by composition of conical elements with different slope angles.

**CONCLUSIONS**

An analytical method for the calculation of acoustic field generated in the infinite fluid region from the cylindrical projector mounted on the axi-symmetrical structure was proposed. The cylindrical projector with a conically shaped reflector at one end was taken, and its vertical directivity was calculated and compared with experimental result. As a result, we obtained accurate directivity patterns which agreed with the experimental characteristics very well, proving the correctness of this method.

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NEARFIELD ACOUSTIC HOLOGRAPHY (NAH) ON UNDERWATER SOUND SOURCES

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INTRODUCTION

The technique of nearfield acoustic holography (NAH) was developed in previous decades. The basic principles underlying NAH have been given a detailed description in the works of Williams et al.[1][2][3]. The major advantage of NAH is that there is no limitation in its resolution of the sound field and it can be used to radiation problems at low frequency. This makes it especially suitable for the study on vibration and sound radiation of the subaerated structures because of the higher sound speed and correspondingly larger wavelength. But how to verify the accuracy of the predicted sound fields? At present, the reconstructed fields via NAH are mostly compared with the measurement results. This will correspondingly cause some problems. First, measurement will unavoidably induce systematic errors and measurement deviation. Second, in most cases, it is very difficult to measure the surface velocity and the surface pressure of an underwater sound source. Some attempts have been made to compute the vibration and sound radiation of the underwater sources to compare with the predicted results via NAH, but because the fields of a subaerated un baffled plate are impossible to determine analytically, the results are not satisfactory.

The present paper addresses the use of modal expansion technique (MET) to the calculation of the vibration and sound radiation of a subaerated un baffled plate which acted as a source for NAH. The agreement between the theoretical results calculated by MET and the predicted values provided by NAH demonstrates that the technique of NAH is valid for underwater sound sources.

REVIEW OF THE PRINCIPLES FOR NAH

Assume a monochromatic plane source of finite extent from Green's theorem, the sound pressure amplitude and phase at source plate represented by complex quantity \( P(x', y') \) can be used to calculate the sound field at any point \( (x, y) \) in a source-free half space.

\[
P(x, y, z) = \int P(x', y') G(x-x', y-y', z) dx' dy'
\]

(1)

\[
G(x, y, z) = \exp(ikR) / (2\pi R^3) \left( 1 + \frac{(x-x')^2 + (y-y')^2 + z^2}{R^2} \right)
\]

(2)

Obviously, this is a convolution computation. The convolution theorem is applied to Eq. (1), yields:

\[
P'(k_x, k_y, z) = \mathcal{F}^{-1} \mathcal{F} \{ P(x', y') \} G'(k_x, k_y, z)
\]

(3)

Where, \( P'(k_x, k_y, z) \) and \( P(x', y') \) represent the 2D Fourier transform of \( P(x, y, z) \), respectively; and

\[
\begin{align*}
\mathcal{F}^{-1} \mathcal{F} \{ P(x', y') \} &= \exp \left( \frac{i \omega k_x x' - k_y y'}{2} \right) \quad (k_x^2 + k_y^2 < k^2) \\
G'(k_x, k_y, z) &= \exp \left( \frac{i \omega k_x x + k_y y - k z}{2} \right) \quad (k_x^2 + k_y^2 > k^2)
\end{align*}
\]

(4)

Assuming that \( P(x, y, z) \) at plane \( H (z=2a) \) is known, by Eq. (3), we have:

\[
P'(k_x, k_y) = P'(k_x, k_y, 2a) G'(k_x, k_y, 2a)^{-1}
\]

(5)

Eq. (5) shows the backward projection (reconstruction of the field at the source plane from the pressure measurement at plane \( H \)). Similarly, the Eq. (3) may show the forward projection. Eq. (3) and Eq. (5) are the basic formula for NAH.

With \( P(x, y, z) \), one can also determine the particle velocity \( v(x, y, z) \) of acoustic medium by considering Euler equation:

\[
V = \frac{1}{(\omega \rho)} \nabla P
\]

(6)

Furthermore, sound intensity can be obtained from:

\[
I(x, y, z) = \text{Re} \{ P(x, y, z) \times \overline{P^*(x, y, z)} \} / 2
\]

(7)

EXPERIMENTAL WORK OF NAH

The experimental setup of NAH in anechoic water tank is shown in Fig. 1. A square steel plate of dimensions 40 cm x 40 cm x 2 mm is hung under water with the entire edges simply supported. The plate is driven by a small piezoelectric exciter at the center with driving frequency 10 KHz. A hydrophone controlled by a microcomputer scans on a prescribed 20 plane (with a distance of 2 cm from the plate source). The data of sound pressure and phase recorded at 3 cm plane are fed to the input of signal processor 7718S through the GPIB interface. After a series of processing procedure, including data sampling, 2D DFT, spatial filtering and convolution computation, the recorded data can be used to predict the pressure and normal vibration velocity on the plate.\(^{[3]}\) (Fig. 2)

THE THEORETICAL SOLUTION FOR THE VIBRATION AND SOUND RADIATION

Consider the plane source described above: a plate of dimensions \( L_x \times L_y \times h \), with edges simply supported and excited by a force \( F_0 \) at its center. The equation of motion for the plate is as follows:

\[
(\alpha^2/m_0^2 + \omega^2/g^2) V(x, y) = F_0 \delta(x-a) \delta(y-b)
\]

(8)

Where, \( V(x, y) \) is the displacement at any point \( (x, y) \) on the plate, \( p(x, y, 0) \) is the acoustic pressure, \( \omega \) is the frequency. Using the modal expansion technique [MET], suppose:

\[
V(x, y) = \sum_{n=0}^{\infty} \sum_{m=0}^{\infty} \cos(k_x n \pi/L_x) \cos(k_y m \pi/L_y)
\]

(9)

Where, \( k_x, k_y, k \) are the eigenvalues corresponding to the boundary conditions for simply supported as \( k_n^2 = (n/2L_x)^2 \) and \( k_m^2 = (n/2L_y)^2 \).

From Eq. (8) and Eq. (9), we obtain:

\[
L_x \omega \rho \left( \omega^2 - \omega^2_0 \right) V_{\text{mean}} \omega \text{mean}
\]

(10)

where \( \omega = \omega (\rho \alpha^2 - \gamma) \) and \( \alpha = (D/\rho \alpha^2)^{1/2} (k_x^2 + k_y^2) \).
From Eq. (8), the pressure \( p(x, y, z) \) is as follows:

\[
p(x, y, z) = \frac{ip}{(2\pi)^2} \int_{-L}^{L} \int_{-W}^{W} p(x, y, 0) \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

From Eq. (12), the pressure \( p(x, y, z) \) is as follows:

\[
p(x, y, z) = \frac{ip}{(2\pi)^2} \int_{-L}^{L} \int_{-W}^{W} \left[ \frac{1}{i17} \frac{1}{i17 + 1} \right] \exp\left[ i(k_x x + k_y y) + i\frac{(k_x^2 - y^2 - \gamma^2)}{2z} \right] \, dx \, dy
\]

\[
= \frac{-1}{i17} \frac{1}{i17 + 1} \int_{-L}^{L} \int_{-W}^{W} \frac{1}{i(k_x^2 - y^2 - \gamma^2)} \, dx \, dy
\]

(13)

With \( z = 0 \), after a lot of simplification, the \( F_{in} \) in Eq. (12) can now be determined:

\[
F_{in} = F + ip = \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} W_n' \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

\[
= \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

\[
= \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

(14)

Where:

\[
W_n' \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

For the higher-order modes of a plate, \( k_x = \frac{\pi x}{L} \) and \( k_y = \frac{\pi y}{W} \), Eq. (15) simplifies to:

\[
W_n' = \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

\[
= \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

\[
= \frac{w}{(4\pi^2 v_m)} \int_{-L}^{L} \int_{-W}^{W} \cos(k_x x) \cos(k_y y) \, dx \, dy
\]

(16)

So \( F_{in} \) can reduce the displacement \( V_{in}(x, y) \), velocity \( V_{in}(x, y) \) can be obtained as follows:

\[
F_{in} = F + ip = \frac{L_s}{k_{in}} \int_{-L}^{L} \int_{-W}^{W} \left[ \frac{1}{i(k_x^2 - k_y^2 - \gamma^2)} \right] \, dx \, dy
\]

\[
V_{in} = \frac{1}{i} \int_{0}^{1} \int_{0}^{1} \left[ \frac{1}{i(k_x^2 - k_y^2 - \gamma^2)} \right] \, dx \, dy
\]

\[
= \frac{1}{i} \int_{0}^{1} \int_{0}^{1} \left[ \frac{1}{i(k_x^2 - k_y^2 - \gamma^2)} \right] \, dx \, dy
\]

(17)

The radiating field of the plate can be determined:

\[
p(x, y, z) = \frac{ip}{(2\pi)^2} \int_{-L}^{L} \int_{-W}^{W} G(x', y') \, dx' \, dy'
\]

\[
G(x, y, z) = \frac{1 - ik_R}{(2\pi R^2)}
\]

(18)

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STONELEY AND FRANZ WAVES ON A SPHERE IMMERSED IN WATER

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Introduction

When a plane wave striking a target immersed in water two waves types appear. The first, the so called inner or elastic surface waves, propagates in the solid an in the surface neighborhood. It is constituted by the Rayleigh and whispering gallery waves or inner waves [1]. The second waves type has the acoustic energy confined in the fluid and along the solid surface. It is constituted by the Stoneley and Franz waves (or outer waves). These surface waves are very important to detect the target immersed in fluid.

We propose in this paper to study the outer surface wave circuunavigation one a sphere immersed in water. The exact solution of the scattering problem by a sphere is obtained. This solution is converted by the Watson transformation to contour integral that is evaluated numerically. An asymptotic evaluation of the pressure gives a physical interpretation of these surface waves.

Modelling

We examine the diffraction of a harmonic plane wave \(e^{i\omega t}\) \((\omega = \text{pulsation})\) by an elastic sphere of radius a (medium 2) immersed in water (medium 1). We have azimuthal symmetry and the the Hemholtz potentials \(\phi_1\) in water, \(\phi_2\) and \(\chi_2\) in sphere are :

\[
\phi_1 = \frac{\phi_0 e^{i\omega \xi}}{2}\sum_{n=0}^{\infty} (n+1/2) \left[ \frac{h_n(x_1)}{h_n(x_3)} + A_n \frac{h_n^*(x_1)}{h_n^*(x_3)} \right] P_n(\eta), \tag{1_1}
\]

\[
\phi_2 = \frac{\phi_0 e^{i\omega \xi}}{2}\sum_{n=0}^{\infty} (n+1/2) B_n \frac{h_n(x_3)}{h_n(x_1)} \left[ \frac{h_n^*(x_1)}{h_n^*(x_3)} \right] P_n(\eta), \tag{1_2}
\]

\[
\chi_2 = \frac{\phi_0 e^{i\omega \xi}}{2}\sum_{n=0}^{\infty} (n+1/2) C_n \frac{h_n(x_3)}{h_n(x_1)} \left[ \frac{h_n^*(x_1)}{h_n^*(x_3)} \right] P_n(\eta), \tag{1_3}
\]

where \(h_n^*(h_n)\) is the spherical Hankel’s function of the first (second) kind. The amplitudes \(A_n, B_n\) et \(C_n\) are given by Cramer's method [2]. Thus by this choice the radial propagating terms are equal to 1 on the sphere. We obtain \(A_n = D_n^{(1)} / D_n^{(0)}\),

where \(D_n^{(0)}\) is the determinant of the equations that affirm the continuity of fields through the boundary \(r = a\) (see annex) and \(D_n^{(1)}\) is the determinant \(D_n\) when we have substituted the first column by the right member (the excitation of the sphere immersed).

Stoneley and Franz waves scattering pressure field.

The modal expressions (1) that may be evaluated numerically to find the incident, reflected, transmitted and the surface waves (outer and inner) [1,2]. To realize an accurate study of Stoneley and Franz waves we consider the second term of (1_1). We use the Watson transformation that consists to substitute the series by the contour integral in the complex \(v\) plane [1,2,3] :

\[
\phi_0 = \frac{\phi_0}{2i} \int_C \frac{v e^{i\omega \xi} (v^{-1/2}/x_1 A(v^{-1/2}/x_1) h_n^2(x_3) P_{v^{-1/2}}(\eta))}{v x^{1/2} \cos \pi v} \, \mathrm{d}v \tag{2_1}
\]

where the path \(C\) runs long the axis \(Re(\nu) > 0\). The path \(C\) is described in the retrograde sense in the complex \(v\) plane where \(v\) is the complex extension of \(n + 1/2\) (integer \(n\)) and \(A(v^{-1/2})\) is the complex extension of \(A_n\).

We can distort the contour \(C\) into \(C^*\) closed to infinite by \(C^*\) (fig.1). We show that the integrand in (2) has poles into contour \(C^*\) and \(C^*\) (fig.2). The poles into \(C^*\) correspond to the inner waves whereas the one located into \(C^*\) near the anti-stoke line \(h^*\) contains the Franz (\(oo\)) and Stoneley (\(o\)) waves that we interest. After some calculus we show that the contribution of Stoneley and Franz waves to the scattering field in water is given by :

\[
\Pi^{(s)} = \frac{\phi_0}{2i} \int_C \frac{v e^{i\omega \xi} (v^{-1/2}/x_1 A(v^{-1/2}/x_1) h_n^2(x_3) e^{i\nu^{1/2}} P_{v^{-1/2}}(\eta))}{x^{1/2} \cos \pi v} \, \mathrm{d}v \tag{3}
\]

By a modified Newton-Raphson method we obtain numerically the poles \(v_n(x_1)\) (\(q = S, F\)) of \(A(v^{-1/2})\). \(\nu(v_n)\) is inversely proportional to the velocities of Franz \(C_F\) and Stoneley \(C_o\) surface waves. The figure 3 gives the evolution of

\[
\frac{\nu}{\nu_C} = \frac{C_S}{C_F}, \quad (q = S,F), \quad \text{Re}(\nu_C) \quad C_F
\]

when \(x_1\) increase (4 to 100). The curve \(n^2\) gives the evolution of \(C_S\) in connection with \(C_F = 1493 \text{ m/s}\) (sound velocity in water).

The numerical data for the elastic sphere are \(p_0 = 2.7 \text{ g/cm}^3, C_F = 6.420 \text{ m/s} (\text{longitudinal wave velocity}), C_T = 3.040 \text{ m/s} (\text{transversal wave velocity}). The curve \(n^2\) gives the evolution of \(C_F\) corresponding to a rigid sphere (radial displacement and stresses are zero on \(r = a\)). We verify that \(C_S\) and \(C_F\) are smaller than \(C_T\) and do not tend on the same limit because the search of the pole \(v_n(x_1)\) as far as \(x_1 = 40.000\) gives at this frequency \(C_T = 0.99950 \ C_F < C_T = 0.99965 \ C_T\). To distinguish these two waves we have drawn (fig.4) their attenuation (\(Im(v_n)\)). The Franz wave attenuation (\(Im(v_F)\), curve \(n^1\) is greater than Stoneley wave (\(Im(v_S)\)) till \(x_1 = 100\). After, \(Im(v_F)\) increases exponentially whereas \(Im(v_S)\) increases moderately, attains his maximum at \(x_1 > 1050\), what next decreases and tends uniformly to zero for \(x_1 > 10^4\). Also the attenuation allows to distinguish these two waves.

Interpretation

Now the integral (3) may be expressed as a residue series. Also, after some calculus the scattering fields associated with Stoneley wave is given by :

\[
\Pi^{(s)} = \pi \phi_0 \sum_{n=0}^{\infty} \frac{e^{i\omega \xi} \frac{D_n^{(1)} D_n^{(0)}}{D_n^{(0)}}}{x^{1/2} \cos \pi v} e^{i\nu^{1/2}} h_n^2(x_3) \frac{d}{dv} \tag{4}
\]

where \(D\) is the derivative by \(v\) of \(D\). This expression is asymptotically evaluated by using the asymptotic expression of the Hankel function (Debye approximation) and the exponential approximation of the Legendre polynomial. From that time the Franz wave pressure in water is given by :

\[
\Pi_F^{(s)} = 2 \phi_0 \rho \omega e^{i\omega \xi} \sum_{m=0}^{\infty} (-1)^m A \exp ik x_2 (r^2 - x_2^2)^{1/2} \psi_m (r, \theta) e^{i\nu^{1/2}}, \tag{5}
\]

\[
\psi_m (r, \theta) = -i \exp ik a (x_2 + 2 m \pi) + \exp -i k a (\gamma_0 - 2 m \pi), \tag{6}
\]

\[
A = \frac{1}{\delta a} \left( x_2 - x_2^2 \right)^{1/4} \left( r^2 - \delta x^2 \right)^{1/2}, \tag{7}
\]

\[
\delta_a = \theta - \theta_0 , \quad \gamma_0 = \gamma_0 + \theta_0 , \quad \theta_0 = \arccos \frac{a}{r}. \tag{8}
\]
The interpretation of this solution is easily. In the simple case $m = 0$ that corresponds to the first term in (4) using $k_2 = \omega / C_F$ the phase of these two first terms is:

$$\psi = \left( (r^2 - a^2) \right)^{1/2} + 2 \omega \delta_0 / C_F.$$ 

Also, the common factor $(r^2 - a^2)^{1/2} / C_F$ corresponds to the time to run the path $T_1P$ or $T_2P$, the phase from $a \delta_0 / C_F$ and $- \gamma_0 / C_F$ corresponds respectively to the curvilinear paths $T_1T'_1$ and $T_2T'_2$ one the sphere (Fig. 5). This, the first term $\psi_0$ in (4) corresponds to the superposition of two Franz waves; the direct one $\psi_0^+$ and the retrograde one $\psi_0^-$. These waves are excited by the tangential incident wave respectively at point $T_1$ and $T_2$ and leaving the sphere tangentially (point $T'_1$ and $T'_2$). The values of $m = 1, 2, 3, ...$ indicate that all creeping waves have encircled the sphere $m$ times.

**Conclusion.**

The numerical evaluation establishes a clear distinction between Franz (rigid scatterers) and Stoneley (elastic scatterer) waves. The velocities of these one are smaller than the velocity of sound in water. The Franz wave attenuation is greater than Stoneley one that attains its maximum about $x_1 = 1050$. After, the Franz wave attenuation exponentially increases whereas Stoneley decreases and uniformly tends to zero. The asymptotic field evaluation leaves to an interpretation of Stoneley and Franz waves; these one encircle the sphere $m$ times from either side.

**Figures**

![Fig 1. Contours for Watson transformation](image)

![Fig 2. Contours for the Franz and Stoneley waves.](image)

![Fig 3. Curve 1 (2) velocity evolution of the Franz (Stoneley) wave.](image)

![Fig 4. Curve 1 (2) attenuation of the Franz (Stoneley) wave $0 < x_1 < 10^5$.](image)

![Fig 5. Travels of Stoneley and Franz waves.](image)

**Annex**

We note:

$$\chi = \rho_2 \left( \frac{1}{r} \right)^{-1}, \quad m^2 = 2 (n+2) (n-1) - \epsilon_1 \epsilon_2 \epsilon_3 \left( \frac{h_2(x)}{h_3(x)} \right).$$

The continuity of stress vector and radial displacements across $r = n$ gives:

$$\alpha_{n+1} A_n + (\alpha_{n+2} + \alpha_{n+3}) B_n + (\alpha_{n+3} + \alpha_{n+4}) C_n = - \alpha_{n+1},$$

$$\alpha_{n+1} A_n + (\alpha_{n+2} + \alpha_{n+3}) B_n + (\alpha_{n+3} + \alpha_{n+4}) C_n = - \alpha_{n+1},$$

$$\left( \alpha_{n+2} + \alpha_{n+3} \right) B_n + (\alpha_{n+3} + \alpha_{n+4}) C_n = 0,$$

$$\alpha_{n+1} = \frac{4 h_2^2 (x_1)}, \quad \alpha_{n+2} = - 1, \quad \alpha_{n+3} = - n (n+1), \quad \alpha_{n+4} = 1,$$

$$\alpha_{n+2} = - 2 \chi \left[ 4 h_2^2 (x_2) - (4 + m^2) \right], \quad \alpha_{n+3} = 2 (n+1) \chi h_3^2 (x_3) - 1,$$

$$\alpha_{n+2} = 2 (2 + m^2) - 2 h_3^2 (x_3).$$

**References**

ACOUSTIC SCATTERING AT OBELIQUE INCIDENCE BY ELASTIC CYLINDRICAL SHELLS COATED WITH AN ABSORBENT

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INTRODUCTION

The aim of this paper is to experimentally examine the influence of an absorbent coating matrix, on the acoustic resonant scattering of elastic cylindrical shells immersed in water. In a first time, the study is carried out at normal incidence in order to evaluate the damping for each absorbent/elastoc tube structure, and to detect the resonant reemissions of surface waves propagating along the elastic tube. After word, this last study is generalized to the oblique incidence.

Let us consider an air-filled cylindrical stainless steel shell. The phase velocity of a longitudinal wave in the material is $c = 5790$ m/s, and the phase velocity of a shear wave is $c_s = 3100$ m/s. The ratio of its inner radius b to its outer radius a (a=27.1 mm) is $a/b = 0.94$. Its axis length is $L = 200$ mm. In the following, this target is labelled “reference tube”. Two other targets (respectively named “tube n°1” and “tube n°2”) of identical characteristics are considered, but each one is uniformly coated with a different 10-mm-depth absorbent polymer matrix of various efficiency.

NORMAL INCIDENCE

Evaluation of the damping

At first, the targets are experimentally investigated at normal incidence (the transmitter/receiver transducer acoustic mean beam is perpendicular to the target axis) in the 40 kHz-800 kHz frequency range. When the tube n°1 is insonified by a short incident pulse, the damping of the first acoustic echo reflected by the structure is about 2 dB/cm, with regard to the amplitude of the first echo reflected by the reference tube. The damping of the second echo with regard to the previous one reemitted by the tube n°1 is about 24 dB/cm. In the case of the tube n°2, with the same experimental configuration, it is only possible to measure the damping between the echo reflected at the absorbent/stainless steel interface with regard to the first one reflected by the reference tube (11 dB/cm).

Tube n°1

The frequential analysis of the scattering at normal incidence by tubes n°1 and n°2 is now performed. The Method of Isolation and Identification of Resonances (M.I.I.R.[1]) is carried out. The resonance spectra obtained from the scattering by the tube n°1, when compared to the results obtained from the reference tube provide the following results:
- In the 40-150 kHz frequency range, resonances related to the reemission of the Whistling Gallery wave l=2 are isolated and identified (nodes n=1 to 4);
- In the 200-800 kHz frequency domain, it is also possible to isolate and identify several Whistling Gallery wave resonances (l=2, n=11 to 18, figure 1-A). The frequency of each corresponding resonance peak is shifted towards lower frequencies with regard to the one of the peak isolated in the case of the reference tube (figure 1-B). The measured shift is about 10 kHz. This shift (about 2.5 %) is superior to the experimental error.

Tube n°2

In the particular 100 kHz-250 kHz range, many resonance peaks are pointed out. With regard to the ones isolated on the reference tube resonance spectrum, some peaks can be related to the resonant scattering of the Scholte-Stoneley waves. These waves are external ones: there main support of propagation is the liquid [2]. The other peaks are due to the finite length of the shell that induces a new condition of wave stationarity in this dimension and then additional resonances.

OBELIQUE INCIDENCE

In the case of the oblique incidence, we use the experimental configuration given on figure 2. A receiver transducer is set such that the mean beam of the wave reflected by the target with respect to the Snell-Descartes law under the incidence angle $\alpha$ coincides with the transducer axis.

The frequential analysis of the reemitted signal is obtained by means of the bistatic Method of Isolation of Resonances (forced signal and resonance spectra).

Reference tube

We focus our study in the particular 100-300 kHz frequency domain: indeed, a complete study of the resonances and the associated waves is previously performed in the case of the reference tube at oblique incidence. The conclusions of this investigation are given in the following (One should notice that these results are similar in the case of the reference tube of acoustically “infinite” length [3]).
- The forced signal spectra plotted from 0° to 30° show an identical general shape: a succession of regularly spaced wide minima are pointed out in the 100-250 kHz frequency range (figure 3).
- Wide resonance peaks are isolated in the same frequency range on a given resonance spectrum (figure 4): each frequency location of a minimum pointed out on the forced signal spectrum coincides with the one of a resonance peak observed on the corresponding resonance spectrum.
- The comparison between resonance spectra plotted for two close values of the incidence angle shows that there is a very small frequency shift of the resonance peaks towards higher frequencies when the incidence angle increases. It is obvious that this observation is also valid for the minima frequency locations in the case of the forced signal spectra.
- The isolated resonances have been experimentally identified at $\alpha=0°$, and a comparison with theoretical works[4] has allowed to correlate them with the reemission of Scholte-Stoneley waves. A follow-up of the Scholte-Stoneley resonance frequencies for $\alpha$ varying from 0° to 30° is given on figure 5.

Tubes n°1 and n°2

The same observation are done in the case of tubes n°1 and n°2 (one forced signal at $\alpha=18°$ and one resonance spectrum at $\alpha=20°$, obtained from the tube n°1 are respectively given on figure 6 and 7).

Then, the follow-up of the Scholte-Stoneley
resonances is also possible in the case of the tube n°1 (figure 6-A) and the tube n°2 (figure 6-B).

CONCLUSION

We have shown that, although the absorbent coatings damp the amplitude of the surface waves reemitted from the elastic target, it is possible to isolate Scholte-Stoneley wave resonances and to follow the evolution of their respective frequencies when the incidence angle increases.

NUMERICAL STUDIES OF SEABED QUALITIES USING REFLECTED DIVERGENT ACOUSTIC BEAMS

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INTRODUCTION

The study of the seabed is of great importance in relation to exploration of seabed minerals, to off-shore constructions, to positioning of power and communication lines and to underwater sound propagation modelling. In spite of the fact that many authors[1,2,3] have paid considerable attention to the reflection of sound beams from seabed interfaces from the theoretical, numerical and experimental sides, several aspects of beam reflection still need to be studied. In the present paper, a brief theoretical analysis of the interaction of sound beams with a multi-layered seabed is given based on a versatile numerical simulation program permitting computation of signal reflection from multi-layered viscoelastic materials below a water column. Emphasis is in particular put on divergent acoustic beams in order to simulate real beams in underwater acoustic studies of seabed qualities.

The paper aims at an illumination of sound beam interaction mechanism at a layered seabed in order to be able to develop an “inverse procedure” for remote monitoring of physical qualities of seabed materials.

The paper deals with some theoretical and numerical studies performed as a part of task 2 of the EC financed research project A.C.I.D. under the Marine Science and Technology (MAST) programme.

I. THE GENERATION OF THE SOUND BEAM

Sound beams can be considered as a composition of a family of plane waves. In the case of 2D beams, the sound field can be expressed by the following Fourier integral transform pair:

$$\nabla(x, z) = \int_{-\infty}^{\infty} \Phi(t) e^{i(kz - \omega t)} dt, \quad \omega = \sqrt{k^2 - \xi^2}, \quad \text{Im} \xi > 0, \quad (1)$$

$$\Phi(t) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \nabla(x, \xi) e^{i\xi t} d\xi, \quad (2)$$

where \(\nabla(x, \xi)\) is the field of the incident sound beam at the interface which is located at \(z = 0\), and \(\Phi(t)\) is the wave-number spectrum of the incident beam at the same interface. Each plane wave component propagates in the medium and is reflected and refracted at the interface independently. Different plane wave components are incident on the interface at different angles, \(\theta = \arccos(1/k_0)\), and different reflection and transmission coefficients are related to each individual plane wave component. As to the reflected sound beam, it is a composition of all the reflected plane wave components:

$$\nabla_r(x, z) = \int_{-\infty}^{\infty} R(\xi) \Phi(t) e^{i(kz - \omega t)} dt, \quad (3)$$

where \(R(\xi)\) denotes a plane wave reflection coefficient.

A divergent sound beam can be produced from a focused sound beam after its focus (see Fig.1). A focused sound beam is established by using the Green’s function in the numerical calculations:

$$\nabla(t) = \frac{1}{2\pi} \int_{-\infty}^{\infty} V_c(t') e^{i\omega t} \omega d\omega, \quad (4)$$

where \(V_c(t)\) is the surface velocity of the line source, which is assumed to be described by a Gaussian distribution function (see Fig.2) given by:

$$V_c(t) = V_0 \exp(-N\tau^2 a^2), \quad (5)$$

and the corresponding field amplitude distribution is now:

$$\Phi(t) = V_0 a (2\pi)^{1/2} \exp(-a^2 \xi^2 / 4N), \quad (6)$$

where \(N\) is a numerical constant, equal to 0.345 when \(a\) is the 3-dB radius of the source, and where \(V_0\) is the sound pressure amplitude at the center of the source. By changing the amplitude of the \(V_c\) function, the curvature, and the geometrical focal length of the line source, different kinds of sound beams can be generated.

II. THE NUMERICAL PROGRAM AND QUALITIES OF THE MEDIUM USED IN THE SIMULATION OF BEAM REFLECTIONS

The numerical calculations are based (1) on the Thomson-Haskell matrix method for reflection of monochromatic waves from a layered medium and (2) on the FFT technique. After a sound beam is established using given parameters, it is expanded in the wave-number domain. Next, the reflection coefficient for each plane wave component is calculated. Complex number calculation is used for the reflection coefficients since these coefficients are complex for the angle of incidence greater than the critical angle. A reflected sound beam is then formed by taking an inverse FFT in the wave-number domain to comprise all the plane wave components reflected from the interface.

From the practical point of view, qualities of a real seabed with a sediment layer overlying a semi-infinite rocky half-space are chosen for evaluation. Parameters for the seabed and for the water column used in the numerical calculations are: For the water, \(c_w = 1500 \text{ m/s}, \quad \rho_w = 1000 \text{ kg/m}^3, \quad \alpha = 0\); for the sand, \(c_s = 1677 \text{ m/s}, \quad c_s = 477 \text{ m/s}, \quad \rho_s = 1830 \text{ kg/m}^3, \quad \alpha = 0.0007 \text{ dB/} \lambda\); for the rock, \(c_r = 3500 \text{ m/s}, \quad c_r = 1600 \text{ m/s}, \quad \rho_r = 3000 \text{ kg/m}^3, \quad \alpha = 0.0007 \text{ dB/} \lambda\). Critical angles are for the water/sand interface: \(\theta_c = 63.44^\circ, \theta > 90^\circ\); for the sand/rock interface: \(\theta_c = 25.38^\circ, \theta_s = 69.64^\circ\). Thickness of the sediment layer is assumed to be 1.5 meter.

III. NUMERICAL BEAM REFLECTION RESULTS

The study of the seabed is based on the fundamental qualities of water and the seabed materials described above. A variation of some of these qualities is made in order to illuminate the individual parameter influence on the beam reflection. In addition, the carrier frequency of the sound beam, the thickness of sediment and the angle of beam incidence are varied in order to investigate the influence of these parameters on the reflected beam patterns. 50 kHz, 1.5 meter and 25.38° are used as basic values for the frequency, the thickness and the angle of beam incidence, respectively.

In describing the beam patterns, a relative measure of 15 degrees of darkness is used to represent different intensities in the dB scale. The reference intensity is always that of the largest field value which is denoted by degree 1.

1. Influence of the sediment thickness

Fig.3 shows a group of beam reflection patterns for the basic parameters given above, but with a variation in the sediment layer thickness from 1.0 to 1.5 and 2.0 meters represented by...
Fig. 3a, 3b and 3c, respectively. Fig. 3 shows in particular the influence of the layer thickness on the geometrical displacement of the second reflected beam pattern. In Fig. 3a the two reflected beams are mixed and some interference appears due to the small sediment layer thickness. In Fig. 3c the two reflected beams separate from each other with the increasing layer thickness. The intensity of the second reflected beam is also affected by the thickness. Fig. 3a shows the largest intensity and Fig. 3c the lowest.

2. Influence of the attenuation in the sediment

Fig. 4 gives a series of beam reflection patterns for the basic parameters given above with variation of the attenuation in the sediment from 0.0 dB/\( \lambda \) to 0.0007 dB/\( \lambda \) and to 0.007 dB/\( \lambda \), respectively. Since different plane wave components propagate different path lengths in the sediment, a different degree of attenuation is connected with each component. The final result is the reduction in the intensity of the reflected beams.

3. Influence of the carrier frequency of the incident beam

Fig. 5 shows a comparison between reflected beam patterns produced by different carrier frequencies in the incident sound beam. Two kinds of effects are shown in the Figs. 5a to 5c. Due to the fact that the figure is plotted in a scale of wavelength, and the higher the carrier frequency, the smaller the wavelength, the first effect is that the separation between the two reflected beams gets larger as the frequency becomes higher, and the second effect is that the attenuation increases while the frequency becomes higher since the attenuation is related to the wavelength. In Fig. 5a the two reflected beams seem to reduce to one and to three degrees of darkness. In Fig. 5b the two reflected beams separate from each other and two darkness layers appear, while in Fig. 5c the second reflected beam disappears.

4. Influence of the angle of incidence, \( \theta_s \)

Fig. 6 compares the reflected beam patterns under different angles of incidence. Full reflection is expected to happen when the angle of incidence is larger than the critical angle for shear waves. Thus, in the case of reflection from a real seabed, no full reflection can be expected from the first interface assuming the sediment to be able to support shear waves. Even so, the angle of incidence is still an important factor affecting the beam reflection. In Fig. 6a, \( \theta_s = 20^\circ > \theta_c \), the two reflected beams join to some extent and only one degree of darkness exists in the second reflected beam. In Fig. 6c, \( \theta_s = 30^\circ > \theta_c \), the two reflected beams separate and two degrees of darkness exist in the second reflected beam. In Fig. 6b, \( \theta_s = 25.38^\circ = \theta_c \), the second reflected beam is the strongest among the three pictures.

CONCLUSIONS

Several numerical computations involving the reflection of divergent sound beams by a 2-layer seabed have been made. The influence of the sediment parameters on the reflected sound beam patterns are analyzed. The thickness of the sediment, the attenuation in the sediment and the carrier frequency of the incident sound beam can highly affect the intensity of the reflected sound beams. Also, the angle of beam incidence is a factor with a considerable influence on the reflected sound beams. The numerical studies reported are being used in an "inverse procedure" to evaluate seabed qualities and to measure the attenuation in the sediment, thus forming a basis for a characterization of the sediments.
ASPECTS REGARDING THE ACHIEVEMENT OF PIEZOELECTRIC TRANSDUCERS FOR UNDERWATER COMMUNICATIONS

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The ultrasounds are widely used in different activities related to marine medium, such as measurement of ships noise, underwater communications, location of underwater objects, marking of objectives and so on.

To achieve ultrasound transducers, piezoceramic materials of PZT type (lead titanate zirconate) have been used, being preferable for their special properties: large electromecanic coupling coefficient, good chemical stability, increased operating temperature and others.

The laboratory studies allowed us to realize ceramic masses of Pb(Zn_{0.53}Ti_{0.47})O_3 type with mineralizing additives (bismuth, niobium or manganese oxides).

These materials were used to construct cylindrical piezoceramic elements with variable heights and wall thicknesses, for which physical and electrical measurements in air were done, according to IRE standard, with the purpose to obtain optimum parameters.

Our cylindrical piezoceramic elements were used to build a piezoceramic transducer as a component of an underwater communication device.

The piezoceramic transducer (Fig. 1) is made of a metallic body (1), on which the cylindrical piezoceramic element of PZT type (3) is fixed by sticking. The electrodes of the cylinder are connected to the electronic monitoring system. This assembly is isolated from the working medium through a rubber cap (2), with a specific acoustic impedance close to that of the seawater.

Between the piezoceramic element and the rubber cap is placed the coupling fluid 4 (castor-oil), also with an acoustic impedance close to that of the working medium. All this assembly is protected against blows by a protection element (5). The transducer is tightly fixed to the device by an O-ring.

Fig. 1
1 - metallic body; 2 - rubber cap; 3 - cylindrical piezoceramic element; 4 - coupling fluid (castor-oil); 5 - protection element.

To obtain the item (2), several compositions of natural rubber with additives were studied and acoustic measurements were made on them.

The piezoceramic transducer operates both in emission and reception regimes.

To determine the operating parameters, the transducer was the subject of a measurement program in a phonosorbent pool. The following parameters were determined: the directivity characteristic, the frequency dependence of the sensitivity and the impedance, the electroacoustic efficiency and the acoustic output power.

Fig. 2
The directivity characteristic of the transducer in vertical plane is presented in Fig.2, while the same characteristic in horizontal plane is shown in Fig.3.

![Fig.3](image)

In tab.1 are presented the results of the measurements regarding the answer in frequency at reception around the working frequency (21 kHz).

<table>
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<th>f(kHz)</th>
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<th>$\gamma (\mu V/deg)$</th>
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The frequency dependence of the sensitivity is presented in Fig.4, noting the high values of the sensitivity around the frequency of 21 kHz.

The comparison method was used to calibrate the transducer. The measurement pool is provided with Brüel & Kjaer measurement and control equipment.

Our devices were also tested in real environment conditions (marine medium), obtaining a reliable link between divers at distances up to 1200 m.

The devices have been realized in three ways, using piezoceramic elements with walls of various thickness; the maximum receiver sensitivity was observed for transducers having ceramic elements with less thick walls.

![Fig.4](image)

The same devices were used with good results for a reflector type underwater location equipment.

The devices own a good mechanical strength against the action of outside factors and also an adequate insulation.

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THE EQUIVALENT TARGET STRENGTH OF BUBBLE SCREEN IN CONSIDERATION OF MULTIPLE SCATTERING AMONG THE BUBBLES.

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INTRODUCTION

Bubbles have been studied in underwater acoustics for several decades. Many researchers studied mechanics and acoustic properties of a single bubble [1-2], however acoustic properties of bubble screen are complicated due to multi-scattering interaction among bubbles in water [3-4].

It is known, when a bubble in water excited to volume pulsation by sound wave dissipates greater part of its energy in thermal conduction process, smaller part in sound radiation. Sound attenuation of bubble screen is much more complicated than that of a single bubble. E.L.Carstensen and L.L.Foldy’s measurement [5] showed that damping constant of every bubble in the bubble screen is several times that of calculated on the basis of single bubble theory. This difference is mainly due to interaction among bubbles.

In this paper, a theoretical formula of estimating target strength of bubble screen on the basis of multi-scattering theory among bubbles in water is derived. The theoretical values calculated by the formula is quite close to the ones measured.

EQUIVALENT TARGET STRENGTH OF BUBBLE SCREEN IN CONSIDERATION OF MULTI-SCATTERING AMONG BUBBLES.

We define a target strength of unit area as equivalent one of bubble screen.

\[ TS_0 = 10 \lg \left( \frac{I_s}{I_o} \right) \cdot \left( \frac{1}{S} \right) \]  

(1)

where

- \( I_s \) is incident sound intensity. \( I_s / I_o \) is backscattering wave’s sound intensity at 1m from the bubble screen.
- Practical target strength of the screen with side surface area \( S \) can be given by \( T_s \).

\[ T_s = TS_0 + 10 \lg S \]  

(2)

The equivalent target strength formula of the screen with uniform distribution bubbles at resonance is firstly derived on the basis of classical theory on bubble. Consider that sound wave with finite beam width is normally incident upon the bubble screen of thickness \( W \). On assumption that the screen locates in the far field of the transducer; incident sound wave is plane one, and bubble distribution is uniform. The total echo intensity is obtained by means of energy superposition.

\[ I_s = \int_0^\infty \left( S_s Sdx / 4\pi^2 \right) I_o e^{-1/2} \frac{S_s / S}{1 - e^{-2S / W}} \]  

(3)

where \( S_s \) and \( S_e \) are total bubble scattering and absorption cross section in unit volume respectively.

In general, \( ScW > 1 \) and let \( r = 1m, \) thus

\[ TS_0 \approx 10 \lg (S_s / 8\pi S_e) \]  

(4)

finally, we get

\[ TS_0 = 10 \lg (R / f_s / 4\pi c_0 S_e) \]  

(5)

where, \( R \) is the bubble radius at resonance frequency, \( S_e \) is a damping constant of bubble, \( c_0 \) is sound speed in liquid, \( f_s \) is resonant frequency of bubble.

To evaluate accurately equivalent target strength of bubble screen, equivalent target strength formula considering multi-scattering interaction among bubbles is obtained by correcting the formula (5) on the basis of theory of sound interaction among bubbles.

The linear second-order differential equation of motion for a bubble in sound field can be written in bubble volume as [7]

\[ m_\delta \ddot{\delta} + b_\delta \dot{\delta} + K_\delta \delta = -P_\delta e^{i\omega t} \]  

(6)

where \( P_\delta e^{i\omega t} \) is incident plane wave, \( m_\delta = p / 4\pi R \) is oscillating mass of liquid around bubble, \( V \) is an instantaneous volume variation of bubble, \( \dot{V} \) and \( \ddot{V} \) are the first and the second derivative of \( V \) with respect to time, respectively, \( b_\delta \) is dissipation coefficient of a bubble, \( R \) is equilibrium bubble radius, \( K_\delta \) is stiffness constant.

As the multi-scatterings reach a final equilibrium state, the equivalent incident sound pressure at any point is [6]

\[ P_{sph} = P_\delta e^{i\omega t} / \left| 1 - \int_0^\infty n(R) A_0^{(1)} dR / k_0^3 \right| \]  

(7)

where \( A_0^{(1)} \) is a amplitude of the first scattering wave, \( k_0 \) is incident wave number, \( n(R) \) is concentration of bubble with radius \( R \) in unit volume.

In general, \( \left| \int_0^\infty n(R) A_0^{(1)} dR / k_0^3 \right| < 1 \). equivalent equation of bubble pulsation in first approximation is

\[ m \ddot{\delta} + b \dot{\delta} + K \delta = -P_\delta e^{i\omega t} \]  

(8)

This equation means that the multi-scattering interaction makes an increase of damping constant which reaches maximum value at resonance. The resonant damping constant with multi-scattering interaction can be written as

\[ \delta_r = \omega b / k = (\omega b_\delta / K_\delta) + (\pi N / k_0^3) \]  

(9)

where \( N \) is bubble number in unit volume, \( k_0 \) is the wave number at resonance, \( \delta_r \) is determined by the theoretical curve shown in Fig-28.2 of reference [1], and was compared with the damping constant \( \delta \) and one measured. It is found that \( \delta \) is close to measured one.
The dispersion of sound speed in consideration of multi-scattering interaction is relevant to bubble concentration, damping constant, and in general is close to dispersion of sound speed given by the classical theory of a single bubble [8], therefore we can take \( C \approx C_0 \) for resonant bubble.

After consideration of effect of multi-scattering interaction among bubbles on the damping constant, the formula (5) can be written as

\[
T S_e = \frac{10}{\rho} \left( \frac{\rho R_p}{4 C_0^2 \delta} \right)
\]

(10)

The curves of equivalent target strength with and without considering the multi-scattering among bubbles are shown in Fig.1, and are compared with measured one.

![Fig.1 Equivalent target strength of bubble screen versus frequency.](image)

**MEASUREMENT OF EQUIVALENT TARGET STRENGTH**

The experiment was performed in a tank (L = 12m, W = 5m, D = 4m) at our institute. The block diagram of the measurement system is shown in Fig. 2.

![Fig.2 Block diagram of the measurement.](image)

1. source (B / K 1402) 2. Gate (B / K 4440) 3. Power Amp. (B / K 2713) 4. PreAmp. (B / K 2636) 5. Oscilloscope 6. Record (B / K 7004) 7. Signal Analyzer

T – Transmitting transducer  H – Hydrophone

Transmitting transducer with band of 20kHz~40kHz, was used. The transducer and the hydrophone are 1.7m and 0.7m away from the center of the screen respectively, and 2m high from the bottom. the transmitted pulse width of the transducer is 0.64ms, and pulse period is 16Hz.

Since target strength of bubble screen is relevant to radiated area of the screen, and beam width of piston transducer determines the radiated area of the screen, we measured the frequency response of main beam width of transmitting transducer.

A bubble generator supplied with compressed air was made. The equivalent target strength spectrum of bubble screen by signal analyzing is shown in Fig.3.

![Fig.3 Measured equivalent target strength spectrum.](image)

**CONCLUSION**

This paper presents a theoretical formula of evaluating equivalent target strength of bubble screen on the basis of multi-scattering interaction among bubbles. Since distribution of bubble concentration is random, the multi-scattering interaction is quite complicated.

Therefore we take account of the multi-scattering interaction on the basis of the formula derived by classical bubble theory. So long as we correct mechanics parameters (m, k, b) of bubble. Volume pulsation, we can reduce 'multi-body' scattering into 'single-body' one, and obtain a bubble screen's equivalent target strength formula of considering the multi-scattering by replacement of C, δ in the formula (5) with C_0, δ_0 respectively.

The measurement results show that values calculated by corrected equivalent target strength formula are quite close to ones measured.

**REFERENCES**

PROGRAMMABLE ARRAY SIGNAL SIMULATION
SYSTEM FOR DIGITAL SONAR

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1. INTRODUCTION

The array signal simulator is a very important tool in sonar performance testing. Now, with the development of high-speed DSP chips and the application of signal processing technology, it becomes possible to design a general-purpose programmable multi-channel array signal simulation system for digital signal processing.

The algorithm presented in this paper is to simulate the array signal utilizing a standard signal table, and the noise is generated by a long shift register with feedback of module 1/2 summation. With these algorithms, we can provide multi-channel array signal efficiently.

The platform by which the algorithms are supported is designed to be a TMS320C25-based parallel pipeline architecture. This multiprocessor platform is task-oriented, that is, each processor is loaded with special tasks. These processors work concurrently to yield the array signal of digital sonar.

2. SIGNAL SIMULATION

The single frequency array signal is expressed as

\[ G(t) = \cos 2\pi f(t - t_0) \text{ where } 0 \leq t < 1/f \]

After sampled with frequency \( f_s \), k-th sample is

\[ g(k) = G(k/f_s) = \cos(2\pi k/f_s - 2\pi f_k) \]

For a specific value of \( f \), we can consider \( g(k) \) as a standard data base. By delaying it with an interval \( p < N \). The new signal with frequency \( f_p \) is expressed as

\[ f(k) = g(kp) = \cos(2\pi kp/f_s - 2\pi f_{kp}) \]

In order to generate array signal by using the standard data base \( g(k) \), it is necessary to calculate the phase term \( 2\pi f_{kp} \).

For a specific array geometry and the incident signal angle \( \theta \), it is easy to calculate the phase delay value \( \varphi(\theta) \) e.g., for a equi-spaced line array with space \( d \), the \( i \)-th element time delay can be expressed as

\[ \varphi(\theta_i) = \left( i - 1 \right) d \sin \theta \]

where \( c \) is the sound velocity.

Therefore, the different phase term \( 2\pi f_{kp} \) is easy to compute. By varying the angle \( \theta \), we can get the whole array signal in 360\(^\circ\) range.

3. NOISE SIMULATION

In this section, an algorithm for generating a multi-channel pseudo-random noise is expressed.

The algorithm of generating pseudo-random numbers looks like that of a linear feedback shift register except that each stage of the former is 16-bit long instead of 1-bit as in the latter.

The first step of generating the white noise is to select \( k \) seeds on \([-1/2,1/2]\), and put them in \( k \) shift registers. The random number is supposed to be uniformly distributed in the interval \([-1/2,1/2]\). The next step is to calculate the sum modulo 1/2 of the 1st sample and the \( k \)th. The result is the \( k+1 \)th sample of the noise sequence. With this result, we have \( k+1 \) samples of the noise, and at last we can output the 1st sample of the noise and save the rest in registers accordingly with a shift of one stage. At this time, we can make another round of calculation in order to yield the \( k+2 \)th sample and output the 2nd sample. By making this procedure on, we will get a white noise sequence that meets our request.

The 1/2 modulo summation can be expressed as

\[ x(n) = (x(n-1) + x(n-k)) \mod 1/2 \]

where \( \mod \) denotes

\[ x(n) = x(n-1) + x(n-k) \]

if \( x(n-1) + x(n-k) < 1/2 \);

\[ x(n) = x(n-1) + x(n-k) - 1 \]

if \( x(n-1) + x(n-k) > 1/2 \);

Decimating the generated noise with interval \( p \), we can get \( p \) channels of noise. Both the decimation theory and the simulation of the algorithm on computer prove that the \( p \) channels of noise is white, independent, and have a uniform probability distribution.

The correlation function and cross-correlation function of the \( p \) channels of noise are shown below

\[ r_{ij} = \sum (x(n) - \bar{x})(x(n-i) - \bar{x}) \]

\[ r_{ij} = \sum (x(n) - \bar{x})(x(n-j) - \bar{x}) \]

The mean value and variances of the generated noise also prove the rightness of the algorithm.

When we have simulated \( p \) channels of
unit-amplitude point signal that each has a different time delay and p channels of white noise ranging from -1/2 to 1/2, we can yield p channels of sensor signal by summing the weighted sum. The weight is decided by SBS requested.

In the case of multi-source, we have to specify the different weight value for each source.

4. PLATFORM OVERVIEW

The platform by which the algorithms are supported consists of two TMS320C25 chips. These two DSP chips form a two-stage pipeline system. The data stream flows from the first stage to the second, while each C25 chip makes special processes on the data stream simultaneously.

The two C25 chips is controlled by a control module, and the technique by which the data stream flows between two processors takes advantage of a RAM architecture which will work efficiently, cooperating the control module.

Besides those modules above, there are one module for parameter input, one module for analog result output, and one module for result display.

5. COMMON RAM

One technique that the simulation system uses to exchange data while each processor works concurrently is the common RAM. Which is the RAM that is used by either the first processor or the second. After processes the data, the first C25 will save the results in this RAM area. Later, in some time, the second C25 will read the data from this area and processes the data. We name this RAM area shared by two C25 chips as the global data area.

6. SAME ADDRESSES

To make two C25 chips work efficiently, we design that while the first C25 is saving the processed data, the second one reads the data from the global data area.

To prevent two C25 from colliding with each other, we organize the global RAM using another technique, which is called same addreeses. That is the global data area is composed of two groups of RAM chips that have same addresses. In fact, under the administration of the control module, while the two C25 are working concurrently, they use different RAM chips although the addresses they use to access the RAM chips are same.

7. CONCLUSION

This system is a real time simulation system. It can simulate multichannel array signal of sonar system, and is fit for checking sonar systems and training sonar operators without doing much work to convert the signal from analog to digital as before. Moreover, the array signal it simulates can be conveniently changed, while the addition of new work that needs us to do is to input parameters into the system from the input module.

8. REFERENCES

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AN EFFICIENT ESTIMATE METHODS FOR ACOUSTIC SOURCE LOCATING (QUASI MAXIMUM LIKELIHOOD METHOD)

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ABSTRACT

This paper presents a quasi maximum likelihood method (QML) for passive ranging problems in underwater environment. Performance analysis and simulation results are presented. Compared to a recently developed spherical interpolation (SI) type method, it shows that QML have improving localization performance significantly, while computation effort still retain simpler.

1. INTRODUCTION

Range estimate in passive location system is very difficult. For traditional culture based ranging method with three sonar linear array, although in theory it has range observability, but in many practical situations, this observability is too weak to work. There are three ways to solving this problem, i.e., spatial diversity by multi-sensor (SI type method); synthetic diversity by successive measurements (various filtering); and channel diversity by multipath measurement.

This paper considered a practical case: smaller array size, larger range to baseline ratio, faster target motion, and locating in all direction. It is essential to utilize temporal and spatial diversity in full. Thus suggest a maximum likelihood method, for its adaption to our situation and its unbiased and minimum variance estimate property in asymptotic sense.

2. PROBLEM FORMULATION

Consider a plane locating system having five sensors, see fig.1. orthogonal type array ensure all direction tracking. The parameter vector $\Theta$, describing target motion is

$$\Theta = (r, b, v, c)$$

where $r, b$ denote the range and bearing of target at current time, $v, c$ denote target velocity and heading. Taking sensor $\Theta = 0$ as reference, and making $N$ times range difference measurements (RD). Thus having RD's set

$$Z = \{z_{1}, z_{2}, \ldots, z_{N}\}$$

where $z_{i}$ denote range difference about true parameter, $n, i, j$ denote sensor number and sampling time, $v$ denote measurement noise. Assuming measurement noise to be i.i.d. normal with $\mathcal{N}(0, \sigma^2)$, the likelihood function of parameter (joint PDF of the observations $z$ conditioned on $\Theta$) can be written as

$$L(\Theta) = \prod_{i=1}^{N} \mathcal{N}(z_{i})$$

where $c$ denote a constant, $\Theta$ denote hypothetical value. Thus to obtain the estimate of parameter, one must solve

$$\hat{\Theta} = \arg\max L(\Theta)$$

This is a complicated multiparameter optimization problem, tend to be compute complex and time consuming, SI like method [1], [2], [3] by introducing equation error to simplify computation drastically, but temporal diversity are limited, thus limited its performance improve.

On the other hand, to perform ML estimate [4], apart from compute problem, for passive ranging system due to stronger non-linearity, insufficient information and correlation among parameters, result numerical stability problem, thus not only time consuming but also necessitate specific stability algorithms. So motivate the development of quasi ML method.

3. QUASI MAXIMUM LIKELIHOOD

Table 1 list a typical fisher information matrix (FIM) and its eigenvalue. It shows that range's observation are very weak, while bearings are very strong. The ratio of maximum to minimum eigenvalues is very large, show that not only the uncertainty ellipsoid is very elongated and the variance in estimated target parameter will be large, but also that the FIM is ill-conditioned, and numerical difficulty may arise in the actual implementation of an algorithm. Furthermore as range baseline ratio increases, theory and practice shows that due to model nonlinearity and insufficient information (the larger range baseline ratio, larger RD noise, the more little information), target estimate tend to be increasingly based in the practical finite sample statistics.

With these factors having been mentioned, a implementation scheme for QML are now present. see fig.2. It possess two feature.

1. Estimate target parameter seperately according to their observability, i.e., first estimate bearing then estimate range parameter, because bearing estimate accurately it is all most equivalent to joint estimate in practice.
2. In optimization, targets velocity and course are acting on function value only through range and bearing, because bearing having estimate accurately, and range's relative change are small in a short time interval. Thus allow us to simplify problem as follow. If range change less then range resolution, it may be seen no varied in a short interval. If range change great then resolution, it may be seen varied linearly.

In so doing yet no much influence on estimate in practice. Thus multiparameter optimization reduce to one (or two)dimension problem, simplify compute and improve stability. Numerical test shows its performance approach CRIB. and even the most sophisticated signal processing techniques including full ML estimate can yield only marginal gains.

4. NUMERICAL TEST

A monte-carlo simulation was conducted to demonstrate and compare the performance of
the SI and OML algorithm. See Fig 1 for simulated scenario. Target true parameter are: bearing (b=45°), course (c=45°), normalized speed (c=1) and range varies from 200 to 600 meters. The measurement were generated synthetically and the noise was additive, zero-mean, independent Gaussian with HD noise level (σ=1 mm). In all test, parameter (b, a, c) and noise level (σ) are fixed, while range and measurement number (n) are selected. Results for each estimate property were obtained by averaging the solutions from 100 trials, in which different noise sequences were used.

Table 2 and Table 3 lists the statistical property of OML and SI estimator separately. It shows that for OML estimator bias is reasonably small, and its standard deviation a small approach to CMB. While for SI estimate the bias is too large, there is no sense to speak of a meaningful bias at low range to baseline ratio or low HD noise or the product of both i.e. the so called effective noise level. The two methods approximate the ideal. While the OML method provides an improved estimate that exhibits a more graceful departure from the CMB. Table 4 shows the effect of temporal diversity. When measurement number (n) increase, the standard deviation decreases with a factor (AR), agrees well with theory result.

5. CONCLUSION

Based on the simulation results reported in the previous section, it is concluded that compared to SI like method, OML provides an effective scheme for locating and tracking maneuvering target in larger effective noise. In large effective noise, using temporal diversity is important, it can extend the usefulness and effectiveness of an existing array by increasing temporal processing gain.

Due to their applicability to a class of source-geometries, including three dimensions scenario, and it need only one or two dimension optimization, its compute are rather simplicity and stability. The OML are recommended as alternative method for passive locating algorithm.

REFERENCES


![Fig 1. Notation for target trajectory and sensor geometry](image1)

![Fig 2. OML location algorithm block diagram](image2)

![Table 1. Fisher information matrix (FIM) for a typical case](image3)

![Table 2. Statistical property of OML estimator (measurement number n=21)](image4)

![Table 3. Statistical property of SI estimator (measurement number n=21)](image5)

![Table 4. Effect of temporal diversity for OML estimator (Taking measurement number (n) from 5 to 21, range=60 m)](image6)
CHI–SQUARE TEST IN ADAPTIVE KALMAN FILTER OF PASSIVE TRACKING

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In passive localization the two time delay measurements are imperfect in most cases due to the short observation array and the complexities of ocean acoustics environment. In order to overcome these difficulties, much effort has been spent in the development of tracking system for many years. In this paper the test in Kalman filter for detection of target maneuver is performed to improve the state estimates which includes target speed and course estimates, the filter structure accepts fundamentally two stage filtering.

TWO STAGE KALMAN STRUCTURE

PREFILTER OF TIME DELAYS $t_1$ AND $t_2$

In the first stage filter the tracking filter models are kept as simple as possible. The dynamics model and the observation model are following

$$
\begin{bmatrix}
    \tau_1(k+1) \\
    \tau_2(k+1)
\end{bmatrix} =
\begin{bmatrix}
    1 & T \\
    0 & 1
\end{bmatrix}
\begin{bmatrix}
    \tau_1(k) \\
    \tau_2(k)
\end{bmatrix} + \begin{bmatrix}
    W_1(k) \\
    W_2(k)
\end{bmatrix}
$$

(1)

where $V(k)$—observation noise, $W(k)$—plant noise. $\tau_1$ is the time delay of $t_1$, or $\tau_2$; $\tau_2$ is the rate of time delay $t_2$; $T$ is interval between two samples.

SECOND STAGE FILTER

The cartesian coordinate system is used here for course and speed estimates. The results of first stage filter give less noisy estimates of each of time delays, then the intermediary range and bearing estimates are obtained. Range and bearing having been changed into cartesian coordinate, the state vector elements $(R_x, R_y)$ consisting of X–Y components of source range are yielded. The second stage filter have following dynamic model and observation model

$$
\begin{bmatrix}
    V_x(k+1) \\
    V_y(k+1)
\end{bmatrix} =
\begin{bmatrix}
    1 & 0 & 0 & 0 \\
    0 & 1 & 0 & 0 \\
    T & 0 & 1 & 0 \\
    0 & T & 0 & 1
\end{bmatrix}
\begin{bmatrix}
    V_x(k) \\
    V_y(k) \\
    R_x(k) \\
    R_y(k)
\end{bmatrix} +
\begin{bmatrix}
    W_{v_x}(k) \\
    W_{v_y}(k) \\
    W_{r_x}(k) \\
    W_{r_y}(k)
\end{bmatrix}
$$

(3)

where $W_{v_x}$, $W_{v_y}$, $W_{r_x}$, $W_{r_y}$ are plant noise. $V_{v_x}$, $V_{v_y}$ are observation noise. If the range estimate is only required, the resulting range of first stage filter will be filtered.

CHI–SQUARE TEST FOR MANEUVER DETECTION

In order to raise the performance of Kalman filter in passive tracking, it is required that source maneuver should be detected, then at the time of maneuver occurring the Kalman filter can be reinitialized. In this paper a chi–square test in Kalman filter is applied to detect the maneuver.

TWO ELLIPSOID TEST[1]

Before introduction of $\chi^2$ test, the two–ellipsoid test is shown here. The problem of maneuver–detection is modeled as that of detecting a signal of unknown magnitude that occurs at an unknown time. Assume the following model for the system state $x(k)$ and the observation $z(k)$:

$$
\begin{cases}
x(k+1) = \Theta(k+1)x(k) + B(k+1)q(k+1) + \phi(k, \theta)
\end{cases}
$$

where $\phi(k) = H(k)x(k) + r(k)$

(5)

The event of maneuver is represented by the Kronecker delta $\delta(k, \theta)$ which is unity for $k = \theta$, where $\theta$ is the time at which the maneuver occurs, and zero otherwise. The two–ellipsoid test uses two estimates. The first is $\hat{x}_1(k)$ which is the estimate obtained using the online measurements $Z(k)$ via a Kalman filter, and the second is the estimate $\hat{x}_2(k)$ which is computed from the a priori information only. The first estimate $\hat{x}_1(k)$ is obtained from the following equation

$$
\hat{x}_1(k+1) = [1 - K(k+1)H(k+1)]\hat{x}(k+1) + K(k+1)Z(k+1)
$$

(6)

The second estimate $\hat{x}_2(k)$ is yielded from the following equation

$$
\hat{x}_2(k+1) = \hat{x}_2(k) + \delta(k+1)\hat{\epsilon}_2(k)
$$

(7)

These are two types of estimate of the same point. The two state estimates and associated covariances define two Gaussian probability density functions (PDF), one is the PDF of $x(k)$ using on line measurements $Z(k)$. The other is the PDF of the state conditioned estimate on the hypothesis of no maneuver in which the online measurements are not used. The first moments $\hat{x}_1(k)$ and $\hat{x}_2(k)$ of these two PDF’s may be considered as point estimates of $x(k)$ with uncertainty $P_1(k)$ and $P_2(k)$ respectively. Confidence regions $R_1(k)$ and $R_2(k)$ can be placed about each estimate.

For one dimension, if the confidence regions are intervals and the test is direct. For two or more dimension, the boundaries $c_1$ and $c_2$ are ellipsoids. The estimate $\hat{x}_1(k)$ reflects the on line measurements $Z(k)$ which indicate the actual situation of $H_q$ (no maneuver) or $H_f$ (maneuver) processed by a Kalman filter that assumes $H_q$. The estimate $\hat{x}_2(k)$ reflects only the a priori information (no on line data) and assumes $H_q$. If the two confidence region’s $R_1(k)$ and $R_2(k)$ overlap, the true state may be in both confidence regions, and it
is reasonable to conclude that no maneuvers have occurred. If the two confidence regions do not overlap, the true state cannot be in both regions simultaneously and a maneuver is declared. This is the two-ellipsoid overlap test.

\[ \chi^2 \text{ TEST} \]

The two-ellipsoid overlap test can be viewed as a geometric method. An alternate approach is $\chi^2$ statistics test, which yields a closed form for two and higher dimensions, and computationally straightforward. Define the estimation errors.

\[ \beta(k) = e_1(k) - e_2(k) = \hat{x}_1(k) - \hat{x}_2(k) \]  
\[ B(k) = E(\beta(k)\beta^T(k)) \]

The covariance of $\beta(k)$

$\beta(k)$ is Gaussian with zero mean and covariance $B(k)$. Since it is the linear combination of two Gaussian random variable $e_1(k)$ and $e_2(k)$, define a scalar test statistic

\[ \lambda(k) = \beta^T(k)B^{-1}(k)\beta(k) \]  

the test statistic $\lambda(k)$ is chi-square distributed with $n$ degrees of freedom for $n$-dimension. The test for maneuver detection is

\[ \lambda(k) > K \quad \text{maneuver} \]
\[ \lambda(k) < K \quad \text{no maneuver} \]

where $K$ is a threshold.

SIMULATION RESULTS

The particular source—observer geometry selected for the simulation, an initial contact at a range of 20 Kiloyards and a true bearing of 180 °, the source moves on a course of 90 ° at a speed of 20 knots, while the observer maintains a constant course of 120 ° with a speed of 10 Knots. All angles are referenced to north. The data are available at equal intervals time delays are corrupted by additive, zero—mean, white—Gaussian noise with a standard deviation of 5 μs. After 10 minutes, source turns the course with initial value 0 ° is changed to 90 °, the speed 20Kn is up to 30 Kn, the corrupted noise still keeps 5μs, the turning rate of course and speed are 3 degree per second, and 0.5 Knot / per second respectively. Fig.1(a) (b) (c) (d) show their mean error and standard deviation of range, speed and course and the test statistic respectively.

RESULTS OF REAL DATE

The sea experiment has been conducted in a shallow water with the depth of 35 m, the velocity gradient of sea—water is −0.3 1/s. The sound velocity near the surface is 1481 m/s, the receiver with-constant speed 5 Knots and constant course is under sea 8~9m, the spacing of each sensor is 22 m, the source keeps linear motion, the signal passed 2KHz−10KHz band-pass filter has been recorded by data recorder. Real data recording in sea experiment are reproduced in lab, the sampling interval is 30 micro second for each channel, which is satisfied with the sampling theo-

rem but the precision of time delay estimates is too lower without interpolation. The method of interpolation used in this paper is published in Ref(3). Fig.2(a) shows the range estimate by correlation method, in the Fig, the dotted line shows the range estimate without post—processor, the black points show true location of the target which is measured by radar, the solid line shows the range estimate of Kalman tracking. Fig.2(b)(c) show the estimates of speed and course respectively.

REFERENCES

BEAMFORMING OF AN UNKNOWN-SHAPED LINE ARRAY

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Introduction

The uncertainty of element location in a towed line array is produced due to influences of a towing vessel's motion, internal wave and turbulence, etc., so that the performance of the array may undergo distorted loss. In order to reduce influence of the uncertainty, the location vectors of the sensor and the signal bearing are jointly estimated by series model according to the criterion of mainlobe sharpness, in terms of the signal spectrum component in a plane wave sound field, so that dynamic compensation can be performed. The comparison of results between the dynamic compensation and straight-line approximation processing shows that the signal to noise ratio (SNR) gain improvement of about 4dB would be obtained by dynamic compensation for a small towed array when the distortion is medium and the tow speed is low.

Analysis

A line array with N omni-directional elements is uniformly placed along the Z axis with spacing d. It is in a sound field with plane wave signal pulse uncoherent noise. The array is towed by a towing vessel with slow speed and straight heading. Under the condition satisfying the sample theorem, the output of every element is uniformly sampled at a time interval \( \Delta t = 1 / (a \lambda) \) where \( a \) is the sample rate, \( f_c \) is the upper limit frequency of a signal. Therefore, the complex time-space sample matrix

\[
X(n,i) = U(n,i) + jV(n,i)
\]  
(1)

can be obtained, and it consists of both signal and noise. If the position vector of \( n \)-th element is \( (y_n, z_n) \), then the beam control delay of \( n \)-th element in \( m \)-th beam may be quantized processed, and the quantized deviation becomes

\[
I_{mn} = \text{INT}_C \frac{d}{c} \left[ y_n \sin \theta_m + z_n \cos \theta_m \right]
\]  
(2)

where \( \text{INT} \) is an integrated function, \( c \) is the sound speed.

The signal and noise of an output can be seen as random processes throughout all states, so we can use a time average instead of an assemble average.

Suppose that the length of one signal processing period is \( \Delta t \), and the distortion of array shape is small, it is easy to evaluate quantized deviation \( I_{mn} \), when the array is considered as ideally straight without the random disturbance part of location vector. Then the space response of average beam output for a randomized array with noise can be written as:

\[
B(m) = \frac{1}{K} \sum_{k=0}^{K-1} \left( \sum_{n=0}^{N-1} W_n U(n, k \Delta + I_{mn}) \right)^2 + \sum_{n=0}^{N-1} W_n V(n, k \Delta + I_{mn})^2
\]

(3)

where \( W_n \) is the weighting function, \( k \) is number of signal processing period. However, even for a sound source with constant intensity, the maximum of average beam output is variant with time, so that the output SNR is also variant and the performance bears significant losses because of the influence of the uncertainty of element location.

In order to decrease the losses of array shape distortion, it is possible to estimate the location vector of elements in time domain in terms of a plane wave in a sound field, so that dynamic compensation can be performed. The time series \( x(n) \) has been taken Fourier transform and separated out the component of average signal spectrum \( X(n,k) \) of intensity plane wave to increase the input SNR. As the period of shape variation is much larger than that of signal processing, the array shape can be considered to be time-invariant in one signal processing period, and the expression of it can be expanded according to a series. After coordinate of the reference point \( (y_k, z_k) \) is given, the optimum coefficients \( A_k \) of the array shape can be derived by iterative operation according to the criterion of maximum mainlobe sharpness for \( k \)-th component \( X(n,k) \) of intensive plane wave signal. Therefore, the estimated \( (\tilde{y}_k, \tilde{z}_k) \) of the location vector of elements can be derived as:

\[
\begin{cases}
\tilde{y}_k = \sum A_f f, \tilde{y}_k \in C \\
\tilde{z}_{k+1} = \tilde{z}_k + \frac{d}{\sqrt{1 + (dy_k / dz_k)^2}}
\end{cases}
\]

(4)

The beam output \( B(m,k) \) is a space Fourier transform of the signal spectrum component \( X(m,k) \) and the location vector \( (\tilde{y}_k, \tilde{z}_k) \) of elements. Then the output space response for a narrow band beam can be represented as:

\[
B(m,k) = \sum_{n=0}^{N-1} W_n X(n,k) \times \exp\left[ j \frac{d}{c} \left( y_n \sin \theta_m + z_n \cos \theta_m \right) \right]^2
\]

(5)

where \( \phi_n \) is the angular frequency of the \( k \)-th signal component.

---

Fig.1 The variation of SNR and source bearing with time for the 32-element line array when towed speed \( v = 4.2 \text{Nm} / \text{R} \)

Fig.2 The variation of the estimated maximum normal displacement with time for the 32-element line array when towed speed \( v = 4.2 \text{Nm} / \text{R} \)
Experimental Result

Based on a computer simulation, above mentioned theory analysis has been verified by an experiment. In the towing experiment, it is shown in Fig.1 that the output SNR and signal source bearing versus relative time when the experimental array system with 32-element is considered as a straight line approximately. As shown in Fig.1, the variation of source bearing with time is an approximately and periodically oscillatory curve and the amplitude fluctuation of SNR may exceed 10dB due to the influences of the element-location uncertainty. It shows that the performance deteriorates and the stability of the array decreases if the array is treated as a straight line. This is one of the major factors that limits the performance of the towed line array sonar.

The rate of the maximum value of the estimated normal displacement to wavelength, versus relative time is shown in Fig.2 for the above mentioned array system, the maximum has exceed tolerance limit 0.1. Comparing Fig.1 to Fig.2, we find that $\Delta y / \lambda$ is large when SNR is small at the same time, while $\Delta y / \lambda$ is small when SNR is large. This variational trend is quite reasonable, which shows that the performance deterioration is caused by the distortion of the array shape, it also proves that the estimating method is practical.

It is shown in Fig.3 that typical beam output of dynamic compensation according to estimated location vector versus a space bearing at some moment, Fig.3a is a wide band beam in time domain, Fig.3b is a narrow band beam in frequency domain. For convenient comparison, the results of the approximate processing when the array is considered as a straight line are represented in the figures with dotted line.

The variable $R_0$ on the top of Fig.3a is the maximum value of dynamic compensational beam, while $R_0$ corresponds to the same value of the approximate straight line processing, they are the utilized factor respectively, while Fig.3b are utilized by the maximum value of dynamic beam.

In all experimental processes, the dynamic compensation may eliminate space obscurity caused by the uncertainty of element locations in most cases, so that the SNR gain improvement on variant extent would be achieved, the typical value is about 4dB when the distortion of array shape is medium. Comparing to the wide band processing in time domain, the narrow band frequency domain processing may not be more advantageous according to our up-to-date experience. Since the criterion of mainlobe sharpness is adopted the computational time for an iterative operation is considerably reduced and the high speed processor develops very rapidly, it is possible to perform real-time operation on signal processors in the near future.

Fig.3 Beam output of dynamic compensation
(a) wide band beam  (b) narrow band beam

References:
AN ALGORITHM OF ADAPTIVE STEP SIZE FACTOR IN ADAPTIVE FILTERING

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Correction Match of Sound Propagation Channel

The ocean sound propagation channel is generally a time-space variant channel. The underwater sound signals as a body carrying the information, often distort strongly after passed through the channel. The distorted signals will reduce the processing gain of a coherent receiver. In this paper a model of transversal filter (moving average model) is developed to describe the channel. An adaptive match method for the time-variant channel is presented and the least mean square error filtering method (LMS method) is used to solve the weighting coefficients of the filter. In the LMS filtering, an adaptive iteration equation is used both in the steepest descent method and in the LMS algorithm. Let \( X_n(n) \) be a set of sampled signal and \( d(n) \) be the desired signal, then their cross-correlation function and auto-correlation function are

\[
P_N = E\{d(n) \cdot X_N(n)\} \quad (1)
\]

\[
R_{NN} = E\{X_N(n) \cdot X_N^*(n)\}. \quad (2)
\]

Using the least mean square error filtering a iteration equation for weighting coefficients can be obtained

\[
W_{n}(n + 1) = W_{n}(n) + 2\mu [P_N - R_{NN} W_{n}(n)] \quad (3)
\]

\[
\alpha = 2\mu
\]

\[
0 < \alpha < 2 / \lambda_{\text{max}}. \quad (4)
\]

where \( \lambda_{\text{max}} \) is the maximum eigenvalue of the auto-correlation matrix \( R_{NN} \).

Equation (3) shows that if the steepest descent method is used to solve the iteration equation, the known 2nd-order statistic of the time series \( d(n) \) and \( X(n) \) are required. It means that \( d(n) \) and \( X(n) \) should be the stationary time series. On the other hand, if \( d(n) \) and \( X(n) \) are the non-stationary time series, then a random adaptive algorithm called the least mean square error (LMS) can be used. The LMS method is different from the steepest descent method in the use of auto-correlation function and cross-correlation function of \( d(n) \) and \( X(n) \). In the case of LMS method, the estimated auto-correlation and cross-correlation function should be used and there will be fluctuation in the obtained weighting coefficients.

Suppose \( \hat{V} \) denotes the true value of the gradient while \( \hat{V} \) denotes its estimated value

\[
\hat{V}_{\hat{w}}[d(n)] = -2\hat{P}_{\hat{w}} + 2\hat{R}_{\hat{w},w} W_{\hat{w}}(n) \quad (5)
\]

where \( \hat{R}_{\hat{w},w} \) and \( \hat{P}_{\hat{w}} \) are the estimated values of auto-correlation and cross-correlation functions

\[
\hat{V}_{\hat{w}}[e(n)] = \frac{2}{\hat{W}_{\hat{w}}} E\{e^2(n)\} = 2E\{e(n)\} \frac{\partial e(n)}{\partial W_{\hat{w}}} \quad (6)
\]

where \( e(n) = d(n) - W_{\hat{w}}(n) X_{\hat{w}}(n) \) is the error function

\[
\frac{\partial e(n)}{\partial \hat{W}_{\hat{w}}} = X_{\hat{w}}(n) \quad (7)
\]

As showed in (1) the iteration equation becomes

\[
W_{n}(n + 1) = W_{n}(n) + 2\hat{E}(e(n)) X_{\hat{w}}(n) \quad (8)
\]

where \( \hat{E} \) is the estimated expected value.

To reduce the calculation, an approximate method must be utilized to derive the LMS algorithm. Instead of the expected value of the function, the instantaneous value is used:

\[
\hat{E}(e(n)) = e(n) X_{\hat{w}}(n) \quad (9)
\]

or finally

\[
W_{n}(n + 1) = W_{n}(n) + 2\mu e(n) X_{\hat{w}}(n) \quad (10)
\]

Modeling and Matching of the Multipath Sound Channel

It is well known that the ocean sound propagation channel often appears to be multipath channel. If a parameter model is to be used to describe the channel, then a transversal filter which is a moving average model is a suitable one. In sonar signal processing a copy-correlator is often used to increase the processing gain. Due to the influence of the channel on the sonar signal, the cross-correlation coefficient of the received signal and the reference signal is often decreased in a great deal amount. To compensate the influence of the channel, the adaptive correction must be carried out. The principle scheme of the adaptive correction match is shown in Fig.1. The weighting coefficients can be calculated from the iteration equation (10), where \( e^2(n) = |y(n) - X(n)|^2 \) is the error energy. From the scheme shown in Fig.1 and the iteration equation (10) we know that, to realize the correction match of the time-variant channel, the weighting coefficients of the adaptive filter must be calculated as rapidly as possible. For this purpose, two measures have been taken: 1. Substituting into (10) the instantaneous value instead of the expected value; 2. Using an adaptive iteration step size factor \( \mu_{a} \) instead of a fixed \( \mu_{f} \).

The adaptive \( \mu_{a} \) is derived as follows:

\[
\mu_{a} = \frac{e(n + 1)}{2e(n) \cdot R_{xx}} \quad (11)
\]

It can be seen from (11) that the iteration step size factor must be adapted to the change of the error, and a suitable adaptive iteration step size will increase the convergence speed. This result is shown by computer simulation.

Analysis of the Computer Simulation Results

In this section the behaviors of the fixed \( \mu_{f} \) and the adaptive \( \mu_{a} \) at different signal to noise ratio and the quantitative comparison of the fixed \( \mu_{f} \) with the adaptive \( \mu_{a} \) under certain conditions are shown by the computer simulation. For this purpose a tapped delay line with 9 equal interval taps is used as a transversal filter. Each tap has a weight and the weighting coefficients are supposed to simulate the paths.
of the multipath channel. To simplify the calculation without losing common significance, suppose that only the fourth weighing coefficient has an amplitude equal to 1 and all the other weighing coefficients are equal to zero.

The computer simulation results are shown in Figs. 2–7. Some conclusion can be drawn as follows:

1. The difference of using the fixed iteration step size factor and the adaptive one in the absence of noise is found in comparison of Fig. 2 and Fig. 3. In the case of adaptive $\mu_a$, the error energy reaches to zero when the number of iteration approaches to 300. While in the case of fixed $\mu_f$, in order to keep the same convergence speed a quite large fixed $\mu_f$ should be taken ($\mu_f = 0.01$) such a large $\mu_f$ results in the oscillation of the error energy function.

2. Fig. 4 and Fig. 5 show that when the stationary convergence is required for both cases of fixed $\mu_f$ and adaptive $\mu_a$ and the signal to noise ratio is equal to $-6$ db, the case of adaptive $\mu_a$ has two times greater convergence speed than the case of fixed $\mu_f$.

3. The signal was delayed on the fourth tap of the tapped delay line (the transversal filter). Fig. 7 shows that when the signal to noise ratio is reduced to $-26$ db, the calculated weighting coefficient is still a little larger than the other weighting coefficients. However, in the case of signal detection where certain detection probability is required, the signal to noise ratio of $-20$ db is more reliable (see Fig. 6).

Conclusion

An adaptive iteration step size factor $\mu_a$ derived in this paper can be used to speed up the convergence of the adaptive iteration equation. This result has been proved by the computer simulation. Once a transversal filter is used to model the time-variant multipath sound channel, a rapid solution of the channel response function must be found to realize in real time the correction match of the sound channel influence. We know from Reference [2] that the time-variant sound propagation channel mainly includes 3 parts: the rapid random-variant part, the regular slow-variant part, and the stationary part. If an iteration algorithm of rapid convergence is used to find the weighting coefficients of the sound channel response, then the second and third parts of the time-variant channel can be corrected in real time.

References

THE EFFECT OF SOURCE MOTION ON MATCHED-FIELD PROCESSING IN OCEAN ACOUSTICS

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Matched-field localization schemes often show a high sensitivity to acoustic variabilities due to mismatch between assumed and actual environments. In this paper, we focus on the effect of source motion or doppler on matched-field processing (MFP). To accomplish this, MFP is extended to treat a moving source problem with normal model description of the sound field. The extension involves both the temporally nonstationary and spatially inhomogeneous nature of the sound field generated by a time-harmonic point source moving uniformly in a stratified oceanic waveguide. It is demonstrated that the impact of source motion, though the velocity of a moving source is much smaller than the sound velocity of the waveguide, can be significant to MFP. In addition, the criteria for minimizing the effect of doppler on MFP is discussed.

INTRODUCTION

Matched-field processing is a generalization of conventional planar wave beamforming which exploits complex multipath propagation conditions to facilitate three-dimensional localization of underwater acoustic sources [1]. Typically, matched-field source localization involves measured narrowband vertical array outputs with replicas of the expected field, predicted by an acoustic propagation model, for a set of hypothesized source positions. As discussed in several recent articles [2,3,4], major difficulties facing this approach is the localization process is sensitive to errors in the propagation model such as sound speed, water depth, and geoaoustic bottom parameters. In this paper we focus on the effect of source motion or doppler on MFP, which did not attract much attention before. The objective of this paper is to analyze a moving source problem and demonstrate how significant the source motion can be to MFP.

Acoustic Field Due to A Moving Source

A formal expression of the acoustic field generated by a time-harmonic moving point source has been derived by hawke-b7 in terms of normal modes of horizontally stratified ocean with the assumption that the source motion is uniform and linear at constant depth.

(a) Side View  (b) Top View

Figure 1. Source and receiver geometry and ocean environment: (a) side view, (b) top view.

For simplicity, consider the case of a source moving radially away with respect to a stationary receiver as in Fig. 1 (\( R_0 = V/2 \)).

\[ p(z,t) = e \sum_{p} \phi_p(z) \frac{\phi_p(Z_0)e^{-ikpR_0}}{\sqrt{k_p R_0}}e^{i \eta_p V_{R} t}, \quad t \in [T/2, T] \]

where \( \phi_p \) is the \( p^{th} \) mode function and \( \phi_p \) is the horizontal wavenumber associated with the \( p^{th} \) mode. \( R_0 \) denotes the source and receiver separation at time \( t=0 \), the center of the observation interval \( T \).

Optimum Receiver

Using complex signal representation, the received signal along a vertical array of dimension \( L \), demodulated at the source frequency \( \omega_s \), can be written as

\[ \tilde{r}(t,z) = \sqrt{E} \tilde{b}_s(t,z,A) + \tilde{w}(t,z), \quad t \in [-T/2, T/2], \quad z \in [0, L] \]

where \( \tilde{w}(t,z) \) is an additive complex white noise process, and \( \tilde{b}_s(t,z,A) \) is a complex Gaussian random variable [6]. \( \tilde{b}_s(t,z,A) \) is a signal normalized in time \( T \) and space \( L \) as

\[ \int_{-T/2}^{T/2} dt \int_{0}^{L} dz \tilde{b}_s(t,z,A) \tilde{b}_s^*(t,z,A) = 1 \]

\( A = (R_{0}, Z_0, V) \) is the parameter vector we want to estimate. Then the localization problem is formulated in the context of non-linear multiparameter estimation problem in the presence of additive white noise. Assuming the parameters are unknown and nonrandom, an optimum receiver based on Maximum Likelihood (ML) estimate is shown in Fig. 2 [7].

Figure 2. Optimum receiver for maximum likelihood (ML) algorithm

In the process of development, we obtain a function called Generalized Ambiguity Function (GAF) in analogy with radar and active sonar theory. Using the orthonormality conditions of mode functions with the assumption that the vertical array is sufficiently long and densely populated, GAF becomes simplified as [8]

\[ |G(A)| = \left| \sum_{p} \phi_p(Z_0)e^{-ikpR_0} \sin[k_0(V-V_0)/2] \right|^2 \]

and

\[ G(A) = G(Z_0, R_0) = \left( \sum_{p} \frac{\phi_p^2(Z_0)}{k_p} \right)^{-1/2} \]

The ambiguity function shows the response of the estimator to all possible values of hypothesized parameters A specified. We note that if there is no mismatch of source motion, i.e. \( V = V_0 \), GAF reduces to the range-depth function defined in modal beamforming except the normalization constant \( G(A) \). The peak of this function yields source range and depth estimates simultaneously.

Simulation Results

We pick horizontally stratified Arctic ocean environment as shown in Fig. 3 and choose 50 Hz as a
source frequency.

Figure 3. Arctic ocean environment for simulation

Then the waveguide supports a total of 19 propagating modes. For convenience, introduce uncompensated doppler distant $X=(X-D)T$. For a source at 30 km range and 1000 m depth, Fig. 4 shows the range-depth ambiguity function with and without mismatch of doppler distance $X$, respectively. The two plots look almost similar except the change of peak contour level. However, reduction of about 30 dB in the mismatched case indicates decrease of signal to noise ratio, which degrades the localization process dramatically. For a source speed of 1 knot, $X=30$ m corresponds to only 1 minute of observation time. The reason is as follows. The effect of source motion is confined in the argument of sinc function with discrete horizontal wavenumber $\theta$, and the observation time window $T$, where $X=30$ m corresponds to the first null of the sinc function. Secondly, there is a slight difference in $\theta$ between the modes, making the sinc function term just like a scale factor. In other words, the doppler spreading of the modes is narrow enough to be considered as a single doppler shift.

In this paper, the effect of source motion on MFP has been investigated. With the models constraining the geometry, the dynamics, and the signals, we formulated the moving source localization problem in the context of multiparameter estimation problem. The optimum receiver based on the maximum likelihood estimates produces GAP which illustrates the global properties of the estimator. The principal result is that uncompensated source motion could significantly degrade the performance of MFP. In addition, a moving source problem can be treated as a stationary source problem if the uncompensated doppler distance $X$ is less than half the wavelength of the discrete modes (3 dB reduction).

Reference


Figure 4. Range-depth ambiguity function: (a) without mismatch, (b) with mismatch ($X=30$ m).

Conclusions
ANALYSIS OF DIRECTIVITY OF RANDOMIZED LINE ARRAY

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Because of existential random error of the position of installation of element and of existential random change of array shape, many utility acoustical array is randomized array. The element position is synthetic result of definite designing position and random position. For convenient research, we are interested in randomized array of random change of array shape.

FUNDAMENTAL ANALYSIS OF THE ARRAY SHAPE OF RANDOMIZED LINE ARRAY

Compared with definite line array, spatial change of randomized line array is respectively ΔX, ΔY, ΔZ. See figure 1.

(figure 1)

According to spatial change of randomized array, we can see local geometry of randomized array on figure 2:

(figure 2)

On figure 2, each section makes an angle of θi to Z axis and each section makes an angle of ψi to Y axis, i = 1, 2, ..., n. n is number of element.

MATHMATICAL DESCRIPTION OF DIRECTIVITY OF RANDOMIZED LINE ARRAY

On figure 2, vector of each element is shown by \( \mathbf{r}_i \) = (x_i, y_i, z_i). Set a bearing of acoustical entry, make an angle of θ to Z axis and make an angle of ψ to Y axis, \( \mathbf{e}_r = \text{e}_x \text{e}_y \text{e}_z \), \( e_x = \sin \theta \sin \psi \), \( e_y = \sin \theta \cos \psi \), \( e_z = \cos \theta \). Make reference to origin of the figure.

2. phase difference ψi of i element is then:

\[
\Delta \psi_i = (2 \pi / \lambda) (x_i \sin \theta \sin \psi + y_i \sin \theta \cos \psi + z_i \cos \theta)
\]

(1)

Set \( (\theta_0, \psi_0) \) bearing of the peak, then phase difference between bearing of the peak and origin is \( \psi_{00} \):

\[
\Delta \psi_{10} = (2 \pi / \lambda) (x_0 \sin \theta_0 \sin \psi_0 + y_0 \sin \theta_0 \cos \psi_0 + z_0 \cos \theta_0)
\]

(2)

Therefore directivity factor of randomized line array of n element shown by \( R(\theta, \psi) \):

\[
R(\theta, \psi) = (1/n) \left( \prod_{i=1}^{n} \left( \sum_{i=1}^{n} \psi_{i0} \right)^{-1} \left( \sum_{i=1}^{n} \psi_{i0} \right)^{-1/n} \right)
\]

(3)

If \( y_i = z_i = 0 \), then line array overlaid on the X axis. \( R(\theta, \psi) \) shown for following equation:

\[
R(\theta, \psi) = n \sin \left( 2 \pi d (\sin \theta \sin \psi - \sin \theta \sin \psi_0) \right) / n \sin \left( 2 \pi d (\sin \theta \sin \psi - \sin \theta \sin \psi_0) \right)
\]

(4)

In (4) n is number of element, d is interval of element. In the case, we may think towed line array as stable. Equation (4) is directivity factor of even line array of n element with equal interval d, the fact is familiar to us.
ROBUST ARRAY SIGNAL PROCESSING

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INTRODUCTION

Array processing is to process the signal carried by propagation waves using the array data which are obtained by an array of sensors with some specific geometry to detect signals and estimate their parameters. Obviously, the processing conditions have a great deal of influences on array processing. In general, the processing model can be idealized approximated as the planewave processing in a limitless homogeneous medium. Based on this planewave model, the direct and simple method of array processing is called conventional beamforming by means of time delay (phase shift) and sum. In effective the conventional beamforming is the spatially Fourier transform. The wavenumber analysis which the transform gives out is at least a little square fit to planewave model. Hence, in the condition of a planewave signal and homogenous noise background (specifically spatial white noise) this method is an optimal array processing. Present of the plane wave interference, using adaptive beamforming can improve the performance of sensor array potentially. A variety of formulations of the optimum array processing problems arise to the same spatial processor. The basic concept is to use measured background spatial correlation characteristics to reject noise and interference, thereby improving the array gain. In order to reject noise without rejecting signals, it is necessary to make some assumptions about the signal characteristics. The problem of mismatch arises when the true signal characteristics differ from the assumed ones. In this case the degradation of the performance of adaptive beamforming can be much faster than that of conventional beamforming. A robust processor should be relatively insensitive to small errors in the assumed signal characteristics. The important sources of error which occur in beamformers are approximately uncorrelated from sensor to sensor and degrade their performance in a way which is similar to adding an amount of spatially white noise to each sensor. The array gain against spatially white noise (white noise gain) is a measure of robustness and its reciprocal is a measure of sensitivity to tolerable errors.

This paper is to present an improved adaptive beamforming algorithm which includes both a linear equality constraint and a quadratic inequality constraint. The quadratic inequality constraint may be used to ensure that the beamformer is robust, not highly sensitive to small errors or perturbations. It limits signal suppression effects and limits the increasing of the adaptive weights which is important in digital implementations. Thus, it controls sensitivity to tolerable errors.

ARRAY GAIN SENSITIVITY TO ERRORS

For simplicity, the following derivation assumes that the signal and noise are uncorrelated zero-mean complex narrow-band series respectively. The array of interest consists of N sensors of known but arbitrary geometry. The sensor outputs can be arranged into a column vector x(t). The covariance of the array data is represented by R. The power output of the beamformer is

\[ s = w^* R w \]  

Where w is the row vector of weights and * represents the conjugate transpose. When a plane wave signal of strength \( \alpha \) from direction \( \theta \) impinges on the array, \( \mathbf{R} \) will include a term \( \alpha^2 \delta(\theta) \delta^*(\theta) \).

where \( \delta(\theta) \) is a row vector of a signal from direction \( \theta \) for the specific array geometry under consideration. Decomposing \( \mathbf{R} \) into signal and noise components as follows:

\[ \mathbf{R} = \mathbf{w}^* \delta(\theta) \delta^*(\theta) + \mathbf{w}^* \mathbf{Q} \mathbf{w} \]  

The noise covariance matrix \( \mathbf{Q} \) normalized to have its trace equal to the number of sensors \( N \) so that the input signal-to-noise ratio is \( \alpha^2 / \sigma^2 \). The array gain \( G \) is the improvement in signal-to-noise ratio due to beamforming, that is,

\[ G = \frac{|w^* \delta(\theta)|^2}{w^* \mathbf{Q} w} \]  

When the background is spatially white noise, \( \mathbf{Q} \) becomes the identity matrix \( \mathbf{I} \) and the array gain becomes the "white noise gain" that is,

\[ G_w = \frac{|w^* \delta(\theta)|^2}{w^* w} \]  

The sensitivity of array gain to errors and perturbations can be examined by assuming the signal to be perturbed by small zero mean random errors with normalized covariance matrix \( \mathbf{A} \), so that the signal covariance matrix becomes \( \mathbf{A} + \mathbf{E} \), where \( \mathbf{E} \) is a strength parameter. The fractional sensitivity \( s \) of array gain to these random errors is

\[ s = \frac{dG}{dE} = \frac{w^* \mathbf{A} w}{G} - \frac{1}{|w^* \delta(\theta)|^2} \mathbf{w}^* \mathbf{w} \]  

The fractional sensitivity is equal to the reciprocal of the array gain against noise with the covariance \( \mathbf{A} \) of the random errors. When the errors and perturbations are uncorrelated from sensor to sensor, the sensitivity is equal to the reciprocal of the white noise gain \( (\mathbf{Q} = \mathbf{A} = 0) \). When the sensitivity to tolerable errors. The white noise gain is a useful and convenient measure of robustness.

In reference [4], we have derived the array gain and the white noise gain of the conventional and the adaptive beamforming respectively. The adaptive beamforming is to minimize the total output power subject to a constraint of unity undistorted signal response from the desired look direction which is called MDR beamforming. It has been proved that the performance of the adaptive beamforming is better than that of the conventional beamforming at the cost of degradation of the white noise gain and that the conventional beamforming can make the white noise gain reach to the maximum value, i.e., the array elements number \( N \). Thereby, in the presence of errors and perturbations the array gain of the adaptive beamforming degrades faster than that of the conventional beamforming does. To solve the mismatch problem existing in adaptive beamforming, we introduce a new adaptive beamforming algorithm that can maintain better performance in the present of small errors and perturbations. The new adaptive beamforming is called robust adaptive beamforming.

ROBUST ADAPTIVE BEAMFORMING

Based on the MDR beamforming algorithm, we introduce an inequality constraint on the white noise gain, that is,

\[ \min \{G_w \} \text{ subject to } w^* \delta(\theta) = 1, G_{w} > \delta^2 / \sigma^2 \]  

The constraint value \( \delta^2 / \sigma^2 \) must be chosen less than or equal to the maximum possible white noise gain \( N \). The solution of (6) is

\[ w = (R + \mathbf{E})^{-1} \delta(\theta) \]  

where \( \mathbf{E} = \mathbf{A} - \mathbf{I} \).
where $s$ is adjusted to satisfy the white noise gain constraint. The Lagrange multiplier $s$ provides a continuous nonlinear parameter between the WNR beamforming $(a = 0)$ and the conventional beamforming. Unfortunately, the relationship between $s$ and the constraint value $\delta$ is not a simple function.

We wish to work with $\delta$ directly and avoid using the intermediate parameter $s$. Decomposing $w$ into its orthogonal components $w_\perp$ and $w_\parallel$ and using the constraint $w^H\delta w = 1$, (4) becomes

$$
G = w^H w = 1 - \delta
$$

(8)

where $\delta = 4\gamma^2 / \gamma$. The constraint (8) may be written as

$$
\delta = 1 / N - 1 / \gamma = \lambda^2
$$

(9)

Thus, the white noise gain constraint can be replaced by a constraint on $\delta$ which is the projection of $w$ onto the space orthogonal to the signal direction. Since $w_\perp$ is independent of the array data, the constraint may be written in $\delta$.

The weights vector of the robust adaptive beamforming is

$$
w = w_\perp + \frac{1}{\delta - 1} w_\parallel
$$

(10)

SIMULATIONS AND EXPERIMENTS

Since the effects of the perturbations on plane wave front and the elements position errors are similar to those of steering vector, we can assume that the entries of the steering vector are added by the zero mean uncorrelated random errors with a variance $\gamma^2 = 1 / 2$. Considering seven-element array whose elements are spaced uniformly at $d/4 = 0.3$ in a spherically isotropic noise field. The noise power is 0 dB, and the plane wave is coming from broadside signal power in 0 dB. The array is steered to broadside based on the perturbation steering vector. In simulations setting $\delta = 0.034$.

For comparison, we compute the array gain and white noise gain of the conventional, WNR and robust adaptive beamforming in ideal and perturbation conditions respectively. In robust adaptive beamforming the constant value of white noise gain is $3\delta / 8468$. The array gain and white noise gain of three beamformers in two conditions are shown in Table 1.

<table>
<thead>
<tr>
<th>Method</th>
<th>Array Gain (dB)</th>
<th>White Noise Gain (dB)</th>
</tr>
</thead>
<tbody>
<tr>
<td>conventional beamforming</td>
<td>6.40</td>
<td>8.26</td>
</tr>
<tr>
<td>WNR beamforming</td>
<td>7.34</td>
<td>1.96</td>
</tr>
<tr>
<td>robust adaptive beamforming</td>
<td>8.58</td>
<td>4.84</td>
</tr>
</tbody>
</table>

Fig. 1. conventional beamforming output

Fig. 2. WNR beamforming output

Fig. 3. robust adaptive beamforming output

A new adaptive beamforming algorithm has been presented. Its performance has been proved in simulations and experiment researches. In the face of the inevitable small errors, the algorithm is robust. In the aspects of signal power protection and interference rejection it has better performance synthetically.

REFERENCES

A FAST AND ACCURATE ESTIMATOR OF DOA AND FREQUENCY

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INTRODUCTION

Estimating DOA (Direction of Arrival) and frequency is common practice for many detection equipments such as sonars. It is well known that these kinds of estimation problems can be solved via spectrum analysis methods.

The spectrum estimation algorithms that are based on eigen-analysis of data correlation matrix have been widely studied since the 70s, which feature a very high resolution for detecting harmonic components embedded in noise. In order to improve computation efficiency for eigen-analysis and the ability of tracking the estimated parameters, the recursive algorithm[1] and adaptive algorithm[2] which can estimate eigenvalues and eigenvectors directly from array data have been proposed. Further, for the special case of a single monocular plane wave, the estimation problem can be much simplified. In this paper, we proposed a new algorithm for DOA and frequency estimation which incorporated the adaptive eigen-analysis and maximum likelihood estimation. Computer simulation and implement in hardware have proved its perfect performance.

PRINCIPLES

There are some conventional algorithms which can be used for eigen-decomposition of the correlation matrix, but they usually demand a lot of computations. And the computation burden for estimating correlation matrix is also heavy. So it is often difficult for them to be handled in real time. In contrast, the computation burden that the adaptive eigen-decomposition algorithm [2] (AEDA) needs is much smaller and it can obtain the eigen-vectors containing the information of DOA or frequency directly from array data. This advantage is particular outstanding for large array. And this algorithm also provides a perfect tracking ability for non-stationary situation.

The AEDA may be implemented with a constrained adaptive filter structure.

Define object function J:

\[ J = E \{ y y^H \} = E \{ \tau \{ y y^H \} \} = \tau \{ W^H RW \} \]

and minimize J subject to the orthonormality constraint

\[ W^H W = I \]

where \( y \) and \( W \) are the output vector and weight matrix of the filter respectively, and superscripts \( T \) and \( H \) denote transposition and Hermitean transposition of matrices, respectively. \( R \) is observation data covariance matrix.

A constrained-gradient search procedure is used to adjust weights. It is easy to know the minimum of the object function is the smallest eigen-value of covariance matrix \( R \), while the weight-vector that maximizes the object function is the eigen-vector corresponding to the largest eigenvalue.

Consider a uniform linear array with \( L \) identical sensors, a monocular plane wave, centered at frequency \( \omega \), impinging on the array from direction \( \theta \). Using complex signal representation, the temporal sample sequence (for \( i \)th sensor) and the spatial sample sequence (for \( m \)th snapshot) can be respectively expressed as:

\[ X_i(m) = a \exp(j \omega m \Delta) + n_i(m) \]
\[ m = 0, 1, \ldots, M-1 \]

\[ X_i(i) = a \exp \left( j \omega d \sin(\theta) / c \right) \Delta i + n_i(i) \]
\[ i = 0, 1, \ldots, L-1 \]

where, \( a \) is the complex amplitude of the signal, \( d \) is the spacing between the sensors, \( c \) is the propagation speed of wave front, \( n_i(m) \) and \( n_i(i) \) are the additive noise at the \( m \)th moment and at the \( i \)th sensor, respectively. \( L \) and \( M \) are the spatial and temporal dimension of the data, respectively.

Then, spatial covariance matrix and temporal covariance matrix are:

\[ R_x = E \{ X_i \} X_i^H \]
\[ R_y = E \{ X_i \} X_i^H \]

where \( E \{ \cdot \} \) denotes expectation operator.

Obviously, \( \hat{\varepsilon} \) and \( \hat{\varphi} \) are the largest eigen-vectors of covariance matrix \( R_x \) and \( R_y \) respectively. Thus for the monocular plane wave case, we can directly obtain the estimations of DOA and frequency simply through the largest eigen-vector rather than computing pseudo-spectrum. This can reduce the computational demanding with a great extent, and thus lay a foundation for the rapidity of the algorithm.

In addition, the maximum likelihood estimator (MLE)[3] may be employed to reduce the estimation variance. But it is effective only for high signal to noise ratio (SNR) situation. When SNR is low, the modelling error will ruin its performance.

Now we propose a new algorithm for simultaneously estimating DOA and frequency.

\[ \hat{\theta} = \arcsin \left( |c \hat{\varphi}| / (2 \pi \hat{\varepsilon}) \right) \]

Assume the estimates of \( \hat{\varepsilon} \) and \( \hat{\varphi} \) are \( \hat{\varepsilon} \) and \( \hat{\varphi} \), respectively. Then,

\[ \hat{\omega} = \frac{1}{2 \pi \Delta} \sum_{m=0}^{M-2} \arg \left( \sum_{m=1}^{M} \epsilon(m) \epsilon(m+1) \right) \]

![Diagram](image-url)
\[ \theta = \arcsin \left( \frac{1}{2} \sum_{\alpha} w(\alpha) \text{arg} \left[ \hat{s}(\alpha) \hat{s}^*(\alpha+1) \right] \right) \]

where \( w(m), w(i) \) are ML weight coefficients[3].

SIMULATION RESULTS AND HARDWARE IMPLEMENTATION

In this section, we present the results of computer simulation and hardware test. Consider an uniformly spacing linear array consisting of 12 elements with half wavelength space in between and 5 points of time sequence for each element, four kinds of DOA estimation algorithms are used to compare there performance.

![Graph showing estimation variances versus SNR for four algorithms.](image)

Fig. 2 The comparison of the estimation variances for four algorithms. \( \hat{\theta} = \sin(\theta) \)

Fig. 2 presents the estimation variance versus SNR. Each curve in Fig. 2 is an average for 100 estimates. The curve 1 corresponds out new algorithm (algorithm I) and the curve 2 is the same as the curve 1 except that ML estimator is not included (algorithm II). The algorithms III and IV corresponding to the curve 3 and 4 are ML estimator only and direct phase estimator respectively. The later uses the data vector to estimate signal phase directly. The 5th curve represents Cramer-Rao (CR) bound.

Obviously, the variance of algorithm I attains CR bound in appropriate SNR (above a threshold of 0dB). The algorithm III has similar property, but with a threshold of 6dB. The algorithms II and IV will never attain CR bound no matter how big SNR is. It is easy to see that the adaptive eigen-decomposition takes an effect of reducing the threshold of the estimation variance by means of increasing time-bandwidth product.

The same conclusions can be reached for the frequency estimation.

Our simulation has also proved that the adaptive algorithm is effective for varying estimated parameters. For example, when SNR=6dB and \( \theta \) changes from 0° to 45° in 600 samples, the DOA estimates can track input proces perfectly.

We also performed a hardware test with a TMS320-C25 signal processor. For the same model above, it takes only 5.7ms to get an estimate with the algorithm I and the estimation variance is in the same order as that of the computer simulation results. Obviously, it is not difficult to implement the new algorithm in real time with high speed signal processor such as TMS320-C25.

CONCLUSION

For single monocular plane wave case, the proposed algorithm features simple structure, small computation burden, deduced estimation variance and tracking ability. It is also easy to implement in hardware. These conclusions have been proved through computer simulation and a hardware implementation research.

REFERENCES

AN ADAPTIVE KALMAN FILTER FOR NOISE TRACKING

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The noise which a ship produces can be received by sensors of sonar and processed to obtain the information of the target track such as range and bearing. The authors present a new algorithm of filtering, called Fractionally Linearized Adaptive Kalman Filtering (FLAKF), which is more suitable for processing the ship noise received by the sensors or the information which is seriously disturbed. The major advantages of the algorithm are that it depends little on initialization, has a high processing gain, and never becomes divergence when the target makes maneuvers. Both range and bearing signals can be processed by the filter.

I. THE BASIS OF FLAKF

Usually the target range which is derived from the received noise varies nonlinearly with time. So does the bearing. Assume the state components as

\[ X_k = \begin{bmatrix} X_k(1) \\ X_k(2) \end{bmatrix} \]

where \( X_k(1) \) is target range component (in metres) at time \( k \) and \( X_k(2) \) is range rate component (in m/s). Then the state equation should be linearized,

\[ X_{k+1} = \Phi X_k + W_k \]

where \( \Phi = \begin{bmatrix} 1 & T \\ 0 & 1 \end{bmatrix} \), \( T \) is sampling period, \( W_k \) consists of state noise and linearization error.

The measurement equation is

\[ Y_k = H X_k + V_k \]

where \( H = \begin{bmatrix} 1 & 0 \end{bmatrix} \) and \( V_k \) is the measurement noise.

According to literature [1], although the linearization error of state equation can be accumulated and becomes larger and larger with time, it affects little the filtering accuracy at the early stage of filtering. If the whole target track is divided into many fractions and the standard Kalman filtering is performed in each fraction respectively, the linearization error and its accumulation in each fraction might be so small that it can be ignored. Meanwhile because of the great inertia of ship, the state noise appears very small and can also be ignored. Therefore, approximately we have

\[ W_k \approx 0, \quad E[W_k] \approx 0, \quad \text{var}[W_k] \approx 0. \]

The adaptive estimation of variance of \( V_k \) is presented in literature [2].

II. THE ALGORITHM OF FLAKF

According to the simplification made above, the equations of FLAKF are rather simple:

\[ X_{k+1} = X_{k+1}/X + K_{k+1} s_{k+1} \]

\[ X_{k+1}/X = \Phi X_k \]

\[ s_{k+1} = Y_{k+1} - H X_{k+1}/X \]

\[ K_{k+1} = P_{k+1}/X (H P_{k+1}/X H^T + \hat{R}_k)^{-1} \]

\[ P_{k+1}/X = \Phi P_k \Phi^T \]

\[ P_{k+1} = (1 - K_{k+1} H) P_{k+1}/X (1 - K_{k+1} H)^T + K_{k+1} \hat{R}_k K_{k+1}^T \]

\[ \hat{R}_k = (1 - c) \hat{R}_k + c (s_{k+1}/X)^2 (1 - H P_{k+1}/X H^T) \]

\[ c_k = (1 - b)/(1 - b^{k+1}), \quad 0 < b < 1 \]

In each fraction the filter works iteratively by using these equations. The initial values of the first fraction can be obtained by simple ways. Before filtering, a number (say, 12) of measurement data are recorded. They are averaged to serve as \( X_{0,1}(1) \). Here the first subscript "0" represents zero time and the second subscript "1" the first fraction and so forth. Since we do not know if the range rate is positive or negative, let it be zero. The square of the standard deviation of these recorded data serves as \( \hat{R}_0 \). As to \( \hat{P}_0 \), let

\[ \hat{P}_{0,1} = \begin{bmatrix} a \hat{R}_{0,1} & 0 \\ 0 & 0 \end{bmatrix} \]

where parameter "a" is positive and smaller than 1, "c" is selected according to the signal being processed. If the range signal (in metres) is being processed, \( c = 1 \sim 10 \) and if the bearing signal (in rad.), \( c = 0.0001 \sim 0.00001 \).

The initial values \( X_{0,1} \) and \( \hat{P}_{0,1} \) have been selected for the first fraction. Thereafter the filtering value \( X_{i,i} \) of the ith fraction can serve as the initial value \( \hat{X}_{i+1}/X \) of the \( (i+1) \)th fraction. The initial value \( \hat{P}_{i+1} \) of the \( (i+1) \)th fraction is determined by the value \( \hat{R}_{i+1} \) in the same way as the first fraction.

There is a transient time at earlier period of each fraction in which the filtering accuracy is not high enough. Therefore the filtering values during the transient time should not be output. The filter of the \( (i+1) \)th fraction can start to work before the filter of the ith fraction ends work. In this way the filter of the \( (i+1) \)th fraction will have been steady by the time when the filter of the ith fraction terminates and the output can be obtained without being broken.

During each fraction the filter works only for a short time. Thus there will be no filtering divergence due to \( \hat{P}_{k,i} \), attenuating to zero. Actually the filter of the present fraction has terminated before \( \hat{P}_{k,i} \), attenuates to zero. Therefore the decrease of filtering gain is limited and it will not occur that the filter refuses to accept new measurement data due to the filtering gain decreasing to zero. That is the idea of GAIN LIMITA-
TION.

It can be seen that the new filter can work with the state equation unknown as long as the signal changes approximately linearly in each fraction. Consequently the FLAKF can be utilized widely. If we define the first component of the state as target bearing and the second as bearing rate, and consider the bearing as observation, the new filter can process bearing signal with only a few parameters modified.

III. SIMULATION AND REAL DATA PROCESSING

We will use the new filter (called filter A) to process simulation data and real data, compare it with the filter (called filter B) presented by Deng Zhi-li and Guo Yi-xin, which is an improvement of the filter presented by Sage and Husa. [1]

The performance of each filter is indicated by the parameter called PROCESSING GAIN as defined below:

\[ G_A \text{ or } G_B = 10 \log (\sigma^2_{a}/\sigma^2_{aw}) \]

where

\[ \sigma^2_a = \frac{1}{N-1} \sum_{i=12}^{N} (y_i - x_i)^2 \]

\[ \sigma^2_{aw} = \frac{1}{N-1} \sum_{i=12}^{N} (\hat{x}_i - x_i)^2 \]

\[ i = 0, 1, 2, \ldots, N \]

\( x_i \) is the real target range; \( y_i \) is the measurement; \( \hat{x}_i \) represents the estimation of range; \( N \) is the total number of sampling data. The first twelve data are ignored because of the transient period at the beginning of filtering. Apparently the greater the processing gain is, the better the performance becomes.

The two filters work on the same condition. All of the initial values are obtained by the method described above. The factor \( b = 0.995 \). In filter A, there are 60 steps in each fraction and 25 steps coinciding in the adjacent fractions. As to filter B, assume that:

\[ \hat{q}_a = 0, \quad \hat{q}_b = \begin{bmatrix} 0 & 0 \\ 0 & 0.1 \end{bmatrix}, \quad \hat{r}_a = 0. \]

1) Simulation Results

Example 1 Assume that a ship moves at the speed of 20 Kn (10.3 m/s) in a straight path. At the beginning of filtering, the target is at the distance of 40 Cab (7412 m) and the relative bearing from the target to the sonar is 90°. Let the sampling period \( T = 2s \) and the noise (disturbance) be white Gaussian with zero mean and variance of \((20/\sqrt{2}) X_i^2\) which, of course, is unknown to the filters. The total number of sampled data \( N = 600 \). Then \( G_a = 36 \text{ dB} \) and \( G_b = 13 \text{ dB} \).

Example 2 To examine the performance of the two filters when the target makes maneuver, let the ship move directly to the sonar for 600 seconds and turn back immediately. The other assumptions remain unchanged. Then \( G_a = 24 \text{ dB} \) and \( G_b = 3.9 \text{ dB} \). It seems that the tracking performance of filter A is much better than that of filter B.

2) Real Signal Processing Results

In practice, noise which disturbs the measurements is not white Gaussian. In this case we still use the two filters to process the data of sonars.

Example 3 A target moves to the sensors of DUXX—5 sonar. Sampling period \( T = 3s \). Then \( G_a = 37.9 \text{ dB} \) and \( G_b = 4 \text{ dB} \). \( G_a \) is much smaller than \( G_a \) because the initial values have great errors. If a more accurate set of initial values is given by some way, \( G_b \) will increase to 23dB, which is much greater than before. But it is hardly possible to have such a set of initial values.

Example 4 Another series of sampled data from a type of sonar is processed. The sampling period \( T = 2s \). Then \( G_a = 18 \text{ dB} \) and \( G_b = 2.6 \text{ dB} \).

IV. SUMMARY

According to the results above, we conclude that:

1) The new filter presented in this paper has a better performance in filtering and tracking noisy target. The approximation for the noise (disturbance) \( W_i \approx 0 \) will not result in divergence.

2) The new filter depends little on the initial values. The initial values can be obtained by simple ways. But the filter presented in literature [2] depends much on the accuracy of initial values.

3) The new filter, when its parameters are determined, can work very well no matter what kind of path the target follows. But the parameters of the old filter should be modified when given a different target path in order to obtain a better performance.

4) The new filter can also be used to process bearing data with a few adjustments of parameters.

REFERENCES


A METHOD FOR IDENTIFYING OUTLIERS IN DATA OBSERVED FROM SONARS

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I. Introduction

Measurement of the direction of arrival to a source observed by a linear array sonar is made from differences of arrival times of signal at different sensors. When the signals of source, the noise of the ocean and the thermal noise of the sensors are all stationary stochastic processes, "good" direction-of-arrival (DOA) data can be observed. If a strong noise disturbance is suddenly added or the linearity of propagation of acoustic wave is severely destroyed, outliers occur in the DOA data.

The often-used least-square estimation is not a robust estimation, for the residuals of observations from computing the least-square estimates using the data containing outliers do not necessarily indicate the locations of the outliers. In other words, observations with large residuals may not be outliers, and those with small residuals may not be "good" observations. This is the Mask phenomenon. In this case, deleting the outliers by use of residuals from least-square estimates and then computing the least-square estimates may finally lead to a convergence to a false solution.

The method to be presented in this paper is based on two features of source DOA measurement by a array sonar. First, the source motion within a short time interval (say two minutes) can be taken approximately as a uniform motion along a straight line. Second, the data observed by the array sonar may contain a high proportion (say, 30% or higher) of outliers. Therefore, the method to be presented has a breakdown point of 50%. If the probability of outlier occurrence reaches , the method fails. So long as the proportion of the outliers contained in the observed data is less than 50%, the method to be presented in this paper can work effectively.

II. Robust Estimation

Suppose that a source moves uniformly along a straight line within a short time interval. At time , the source DOA with respect to the linear array is , and the change rate of DOA is . The observation at time is denoted as . Then, there is the linear model:

\[ y_i = \theta_0 + \delta_0 (t_i - t_0) + e_i, \quad 1 \leq i \leq N \]  

where is the observation error, and is assumed to be a sequence of independent random variables with identical distributions . A typical form of distribution is

\[ F = (1 - \epsilon) \Theta + \epsilon H \]  

where , a normal distribution function with mean zero and variance , denotes the distribution of "good" observations, an arbitrary symmetric distribution function, denotes the distribution of outliers, and stands for the probability that an outlier occurs.

The problem now is to determine which of observations are "good" and which are outliers. To this end, it is necessary to find out two "good" observation from first and then test the others starting from these two "good" ones. Consequently, the first step is to find out, without any prior knowledge of , two "good" observations from first which contain outliers.

One random drawing method can serve our purpose. That is, each time a subsample of two points is randomly taken from and from these two points an estimate of is established. Denote the estimate established by the th sampling as . Suppose the number of sub samples (say, ) is sufficiently large such that the probability of drawing at least one subsample consisting of two "good" points should be close to one. Then, one of (, ), , must be "good" (in the sense that it is obtained from "good" observations; however, it may not be accurate.)

When we have a number of estimates (, ), , of , and know at least one of them is "good", the second step is to find out this "good" one. Several objective functions for robust estimation can be employed for this purpose.

The commonly used robust estimations can be classified into two types, [1], [2]: M-estimates (maximum likelihood type estimates) and L-estimates (linear combinations of order statistics).

Now substitute the m estimates (, ), , of into the objective functions of the above-mentioned M-estimations or L-estimations that have high breakdown points. Those that minimize the objective function will be "good" estimates.

That kind of random drawing method, however, is hard to apply. To ensure that one subsample consisting of two "good" observations is taken out in probability of nearly one requires a great number of sub samples. We thus will propose an alternative in the next section to find two "good" observations from .

III. A New Method

Suppose is a multiple of 4, and let . Divide into groups with being the first group and the th group. When the number of outliers in is smaller than , among the groups there should exist at least one group that contains no more than one outlier. Apply a robust method to delete two points from each group in such a way that, if a group contains only one outlier, the outlier can be deleted in probability of nearly one. Therefore, after the operation of the groups there must be one group in which the remaining two points are "good".

Now we investigate this robust method. Suppose there are observations , , satisfying the linear model Eq.(1) and one of which is an outlier. We consider a robust method to delete the outlier. For a linear model of form Eq.(1), we can find a linear estimate of ,

\[ \hat{\theta}_n = \sum_{i=1}^{n} \beta_i y_i \]  

We require that is an unbiased estimate of , i.e., . Thus, must satisfy the following constraint
conditions:

\[ \sum_{i=1}^{4} \beta_i = 0 \]  
\[ \sum_{i=1}^{4} \beta_i r_{a_i} = 1 \]  

where \( r_a \triangleq t_{a_i} - t_{b_i} \). Write

\[ r_{b_i} = y_{b_i} - \hat{\theta}_{b_i} r_{a_i} \quad (1 \leq j \leq 4) \]  

and it is easy to verify

\[ r_{a_i} = \hat{\theta}_{a_i} + e_{a_i} - r_{b_i} \frac{1}{4} \sum_{i=1}^{4} \beta_i e_{b_i} \]  

where \( \hat{\theta}_{a_i} \) is unknown, and will not be estimated because it is a constant and only displaces the whole sequence, without changing the order of the elements, when we rearrange \{r_{a_i}\}, 1 \leq j \leq 4, in a descending order.

We will choose the coefficients, \{\beta_i\}, which satisfies both constraint conditions of Eq.(4) and some robust conditions that if a certain \( y_{a_i} \) is an outlier, the corresponding \( r_{a_i} \) is, in probability of nearly one

\[ r_{a_i} = \max_{i \leq 5} r_{a_i} \quad \text{or} \quad r_{a_i} = \min_{i \leq 5} r_{a_i} \]  

If we can make it, deletion of the maximum and minimum from \{r_{a_i}\}, 1 \leq l \leq 4, leaves two "good" points.

Now we turn to the question what robust conditions the coefficients satisfy. By Eq.(6), assuming \( y_{b_i} \) is an outlier we have

\[ r_{a_i} = \hat{\theta}_{a_i} + (1 - \beta_i) r_{a_i} - r_{b_i} \sum_{i \neq j} \beta_i e_{b_i} \]  

where \( 0 < 1 - \beta_i r_{a_i} < 1 \). Since \( e_{b_i} \) is the error of the outlier and the other \( e_{b_i}, i \neq j, \) are error of "good" observations, the second terms on the right-hand side of the above formula play a key role in rearranging. From Eq.(6) it is easily seen that to validate Eq.(7) in probability of nearly one, we demand

\[ \mu_j \triangleq 1 - \max_{i \leq 5} \beta_i (t_{a_i} - t_{b_i}) \]  

sufficiently large for every \( j, 1 \leq j \leq 4 \). In the ordinary least-square estimation, there exist the so-called leverage points, or in the present case \( t_{a_i} \) and \( t_{b_i} \), at which the value of Eq.(9) is near zero (in the non-equi-interval sampling case) or even less than zero (in the non-equi-interval sampling case).

To make all \( \mu_1, \mu_2, \mu_3, \mu_4 \) sufficiently large, every \( |\beta_i| \) must be sufficiently small. However, since \( \beta_i \) should satisfy the unbiased condition of Eq.(4), not all \( \mu_i \) can be made large enough simultaneously. A plausible way is to decrease \( \mu_2, \mu_3 \) and increase \( \mu_1, \mu_4 \) so that \( \mu_2 \) and \( \mu_3 \) are equal to \( \mu_4 \) or \( \mu_3 \). This is the robust condition \{\beta_i\} should satisfy. To achieve such coefficients, we require that, under condition of Eq.(4), \{\beta_i\} minimize the following sum of weighted squares:

\[ \sum_{i=1}^{4} \frac{1}{4} |\beta_i|^2 \]  

where \( \delta_1 = \delta_4 = 1 \)

\[ \delta_1 = \frac{T_1}{T_1 + T_2} \]  
\[ \delta_4 = \frac{T_2}{T_1 + T_2} \]  

with \( c \) a suitably chosen constant.

The solution to that constrained extremum problem is

\[ \beta_i = \delta_i t_{a_i} \quad (l = 1, 2, 3, 4) \]

where

\[ \lambda = 1 / \sum_{i=1}^{4} \delta_i r_{a_i}^2 \]
\[ r_{b_i} = t_{b_i} - t_{a_i} \]

To make \( \mu_1, \mu_2 \) equal to \( \mu_3 \) or \( \mu_4 \), we take

\[ C = \begin{cases} T_2 + T_3 & \text{if } T_2 \geq T_1 \\ T_1 + T_2 + T_3 & \text{if } T_1 \geq T_3 \end{cases} \]

IV. Identification of Outliers

The "good" estimate \((\hat{\theta}_0(m), \hat{\theta}_2(m))\) we have found could not be used directly for identification of outliers in \( \{y_i\}, 1 \leq l \leq N \), because it is obtained by only two points and therefore is not accurate. To acquire an accurate estimate of \((\hat{\theta}_0, \hat{\theta}_2), (\hat{\theta}_0(m), \hat{\theta}_2(m))\) must be substituted into M-estimate for iteration. When the iteration converges, the solution obtained is the accurate estimate \((\hat{\theta}_0, \hat{\theta}_2)\) of \((\theta_0, \theta_2)\). From that estimate we then compute the residual of each observation:

\[ r_i = y_i - \hat{\theta}_0 - \hat{\theta}_2 (t_i - t_0) \]

and make decision. If

\[ |r_i| > C \sigma_0 \]

holds, we decide that \( y_i \) is an outlier; otherwise it is a "good" observation. In Eq.(16), the threshold \( C \) depends on \( N \). If \( N = 30 \), the probability of the corresponding false alarm is 1%, and \( C=3.236 \). For details, the reader may refer to the table of fractiles of normally distributed ordered statistics.

We now write out the iterative formula of M-estimation.

\[ \left( \begin{array}{c} \hat{\theta}_0 \\ \hat{\theta}_2 \end{array} \right) = \left( \begin{array}{cc} \sum_{i=1}^{N} & 1 \\ \sum_{i=1}^{N} (t_i - t_0) & \end{array} \right) \left( \begin{array}{c} 1 \\ \sum_{i=1}^{N} (t_i - t_0) \end{array} \right) \left( \begin{array}{c} w_1 y_i \\ w_2 y_i \end{array} \right) \]

where

\[ w(r) = \frac{1}{\lambda_0 |r|} \quad |r| \leq \lambda_0 \sigma_0 \]
\[ \lambda_0 \sigma_0 < |r| \leq \lambda_1 \sigma_0 \]
\[ 0 \quad |r| > \lambda_2 \sigma_0 \]

where we may take \( \lambda_0 = 2, \lambda_1 = 3, \lambda_2 = 8 \).
DIVER NAVIGATION AND TRACKING USING A PROGRAMMABLE DIVE COMPUTER AND AN INTELLIGENT TRANSPONDER ARRAY

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A decompression computer, the ACE PRO-FIILE, has been reconfigured as a navigation aid which allows a diver to determine his own position. The wrist-mounted computer calculates and displays the diver’s coordinates by measuring the arrival times of acoustic pulses from a fixed array of transponders. These coordinates are then encoded and transmitted by the computer so that the diver’s position can be tracked remotely.

INTRODUCTION

The central feature of our system is a commercial decompression computer, the ACE PRO-FIILE (Quatek Ltd), which is programmed to allow a diver to monitor his position at any time with respect to some arbitrary datum point on the sea bed. The position can also be tracked by a dive supervisor on board a ship or in an underwater habitat. The computer is used with three acoustic transponders which are dropped onto the sea bed to form an arbitrary-shaped long base line triangular array. One of the transponders is an "intelligent" device (the master) and the other two are "semi-intelligent" devices (the slaves). With this arrangement, any position can be found by triangulation.

An acoustic command triggered by the dive supervisor starts the transponders operating in their calibration mode. The master interrogates the two slaves, which in turn interrogate each other; the master then emits three pulses on a common frequency which enable the computer to calculate the side lengths of the array. The transponders then change to a navigation mode in which each emits a pulse, also on a common frequency, at regular intervals under the control of the master.

The dive computer is programmed with an algorithm that calculates the coordinates of any position from measurements of the arrival times of the three pulses from the transponders. Positions are then displayed in real time in both alphanumeric characters and graphically. The algorithm allows only two-dimensional position-fixing because for the majority of conceivable applications the diver is interested only in his position on the sea bed rather than in mid-water. The computer’s built-in depth gauge provides a readout of depth.

It was a necessary condition in the development of the navigation system that the ACE PRO-FIILE was not changed significantly. In practice, we made only the following changes:
- reprogrammed the memory space with an accurate position-fixing algorithm;
- added an omnidirectional acoustic transducer, a narrow band filter and decoding circuitry at the input port to receive acoustic pulses from the transponders;
- added a transmitter at the output port to transmit coded data for tracking purposes (using the same transducer);
- changed the original screen to a dot matrix liquid crystal display module in order to achieve acceptable spatial resolution in the graphical display of the diver’s position.

TRANSPONDERS

In principle, three "intelligent" transponders can be used for navigation but in this system only one truly intelligent transponder, the master, is necessary. The design of this system is based on a modified PRO-FIILE using a 65C102 processor running at a clock frequency of 13.1072 MHz. When the master is used with two slave transponders, the array as a whole behaves in an intelligent manner, but with the attendant advantage of having much less hardware. In the development of the transponder layouts, we have designed and simulated the necessary frequency generators, decoding logic and fifth-order filters by computer-aided engineering.

The bandwidth is dictated by the resonant frequency (69.75 kHz) and Q-factor (5.5) of the type of acoustic transducers used, which are D/170 electrostrictive spheres (Universal Sonar Ltd, U.K.). We designed the filters to have the following frequencies:

- f0 = 71.23 kHz;
- f1 = 74.470 kHz;
- f2 = 68.267 kHz.

The common reply frequency f0 was chosen to be as close as possible to the resonant frequency so that maximum power could be transmitted; the other frequencies were chosen to be as far as possible apart yet within the transducer’s -3 dB bandwidth of 63.5 to 76.5 kHz. The four frequencies are derived from a 3.2768 MHz crystal oscillator which was chosen because, when divided, this frequency provides digital pulse trains with equal mark-space ratio. The frequency spectrum of each train has a single peak at the centre frequency of one of the filters, whereas the spectrum for trains of uneven mark-space ratios has two components that lie above and below the centre frequency and outside the bandwidth of the filter.

Before the navigation system can be used, the transponders are dropped onto the sea bed in an arbitrary triangle. Typically, they would be about 100 m apart but their exact locations are not critical because the algorithm can be applied for any triangular array. To explain the mode of operation, we suppose that the three transponders are A, B and C and that their separations are AB = d1, AC = d2 and BC = d3. Initially the transponders lie inert until the master is triggered by a pulse transmitted from a "dry" system operated by the dive supervisor; this starts the calibration mode, which is illustrated by Figs 14 and 2. Meanwhile, we designed the transponder Figs 3 and 4.

On receiving and decoding this pulse, transponder A (the master) replies on a common reply frequency f0 and passes a "token" pulse at frequency f1 to transponder B (the slave). On receiving this token, B responds by transmitting a new token at frequency f2 to both A and C (the second slave). On receiving this, A transmits to the diver’s computer at f3, thus providing a measure of d2, as shown in Fig. 2. Meanwhile, having received the token from B, C responds by transmitting a new token at f5 and blocks any further calibration transmissions from B; this token is then received by both A and B. Transponder A transmits another pulse at f6, which provides a measure of d1 + d3, while transponder B transmits another token at f7 which when received by A results in the transmission of a third pulse at f1 to give a measure of 2(d1 + d3). When the diver’s computer has received all three pulses on f3, it calculates the dimensions of the triangular base and is then ready to receive further pulses from the transponders when they switch automatically to their navigation mode.

In the navigation mode, all three transponders transmit at frequency f0. The mode of operation, illustrated by Figs 3 and 4, is as follows: A "pings" at f0 and transmits a token at f1; B receives this token, pings at f0 and transmits a new token at f2; C receives the token, pings at f0 and transmits a new token on f3; A receives this token then goes into a "wait" state (to allow time for the diver’s computer to complete the navigation algorithm) before starting the procedure again.
Additional modes of operation, such as direct interrogation from the surface or transponder recalibration can be achieved by adjusting the length of the "start" pulse to the master. In its intelligent role, the master can detect a breakdown in communication, for example due to an obstruction between any pair of transponders or between one of the transponders and the computer. In this case, the master pings five times at its normal rate (to allow for the possibility of a temporary obstruction) then transmits a "panic" signal to the "dry" system.

SOFTWARE

The software for the diver's navigation computer consists of two quasi-concurrent processes, "Position Fix" and "Contact Check". Position Fix runs continuously, whereas Contact Check is interrupt driven. The frequency of the interrupt depends on whether the computer is submerged.

Contact Check monitors the input from four metal contacts on the exterior of the computer. On the surface, this routine is called every 0.5 s and accesses different user menus depending which pair of metal contacts is bridged with moist fingers. One of the necessary initialization operations is to enter the velocity of sound in water, which is used in the navigation algorithm. When the computer is immersed, all five contacts are deactivated by contact with the water and the interrupt frequency is reduced so that Contact Check simply tests every 30 s to determine if the dive has ended.

Position Fix is initiated by the input from the hydrophone's receiver circuit. Depending on the length of the input pulse, it can determine whether the transponders are in calibration or navigation mode. The time differences between the subsequent pulses are measured using software counters, and calculations are made to determine the transponder locations or the diver's position.

The navigation algorithm is essentially a quadratic equation and as such gives two mathematically correct solutions. In practical terms, the solutions yield two points that are often widely spaced apart in the general area around the transponders so further calculations are necessary in order to select the correct one. Position Fix determines the true diver position through a strict series of checks. In essence, these track the diver during the course of the dive and reject positions that the diver could not possibly reach in the given time.

The software was prototyped in "C", which is an industrial standard. This allowed rigorous testing of the navigation algorithm through the use of a simulation program built around the basic program. Once the software was proven it was cross-compiled to run on the embedded 65C102 microprocessor of the PRO-FIFILE. The input and screen driver routines were written directly in assembly language.

FUTURE DEVELOPMENT

The ultimate aim is to use the system in the vicinity of a habitat in which a team of scientific divers (breathing nitrox) will live for about two weeks. It is anticipated that eventually an array of transponders will be deployed so that up to six divers, each carrying their own navigation computer, can find their own positions simultaneously and independently and navigate themselves back to the habitat. At the same time, they will be tracked from inside the habitat. To achieve a system of this complexity may require time-sharing of the available frequencies, which would imply slower updating of the position fix for any one diver. This is the theme of our future research.

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SHORT RANGE UNDERWATER ACOUSTIC COMMUNICATION SYSTEMS

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INTRODUCTION

Along with the advance of ocean development, the underwater acoustic communication systems (ACS) are more and more important. It can transmit and receive commands, information and images under the water. Underwater acoustic channel is a multipath, time varying and dispersion channel. Acoustic transmissions are complex considerably. Its bandwidth is less than that of cable more than two magnitude. In order to design a reliable underwater acoustic communication system, integrating techniques will be developed. In general, there are three kinds of underwater acoustic communication systems, i.e. short range systems (<4km), medium range systems (4-20km) and long range systems (>2000km). In this paper, a short range system is introduced. Three techniques are used herein. The first is the signal design, the signal is modulated by MFSK and detected by FFT; the second is the error correction codes, i.e. convolutional codes and Viterbi algorithm (VA); the third is the narrow beam sonar array. This system is tested in water tank systematically, where multipath effects are in existence. Based on the experimental results, we know these techniques are useful. Recently, the adaptive equalizers and new type synchronizers are being studied.

STATISTICAL CHARACTERISTIC OF CONVOLUTIONAL CODES IN UNDERWATER ACOUSTIC CHANNEL

In some cases, the underwater acoustic channel may be considered as a Rice or a Raleigh channel (1), the symbol error probability \( P_e \) of MFSK codes is the error probability \( P_b \) of FSK convolutional codes is obtained when four frequencies are used to transmit data and the signal and noise ratio are the same. The \( P_b \) of convolutional code (constraint length \( k = 3 \)), code rate \( r = 1/2 \) is the same as the \( P_b \) of MFSK linear block code \((N = 6)\). We obtain that, at the same signal to noise ratio, the frequency band width required by linear block code is larger about 6 times than that of convolutional code.

In some cases, the direct wave and the first reflected wave of sea surface are the dominated waves in the multipath signals. If the fluctuation of signal is not considered, bit error probability \( P_b \) of convolutional code \((k = 3, r = 1/2)\) is obtained:

\[
P_b = \frac{1}{2} \left( \frac{1 - D_b D_r}{D_b D_r} \right)^{1/2}
\]

where

\[
D_b = \exp(-h_b/N_0), \quad D_r = \exp(-h_r/N_0)
\]

where \( h_b, h_r \) are called intersymbol interference coefficients (3), \( N \) is the one-side noise power spectral density, and they are shown in reference (3). In detail, the values are shown in Fig.1, where \( T \) is the ratio of code element impulse space with impulse width and \( \beta \) is the reflection coefficient of sea surface. Pb at \( \beta = 0.3 \) is better than that at \( \beta = 0.1 \) obviously, this is the one of superior characteristics of convolutional codes. In the case of Fig.1, the space of code element is large enough, so that the first reflected wave of sea surface can be resolved by decoding due to the deconvolutional operation and can be aligned and superposed on the direct wave, then the \( P_b \) is reduced. The value of curve 2 is better than that of Rayleigh channel (1,3). In practice, since the random component exists in direct wave and reflected waves of sea surface due to the fluctuation of medium, and \( T \) is large in some cases, the \( P_b \) of convolutional codes is inferior than that of curve 2 in Fig.1. The \( P_b \) of curve 2 in Fig.1 is the characteristic of convolutional codes in the ideal conditions.

Based on the above results, we know the convolutional codes are good codes against the multipath in underwater channel.

SHORT RANGE UNDERWATER ACS

The performances of short range underwater ACS are:

- Modem function: transmit and receive 4-8 frequencies or more.
- Data frequency band: 480-550 KHz
- Syn. frequency: 470 KHz
- Codes: convolutional codes \( k = 3, r = 1/2 \) and others
- Decoder: Viterbi algorithm
- Rate: 1200-4800 bauds
- Max. Range: 1000m

In this system, four data frequencies and \( k = 3, r = 1/2 \) convolutional code are used usually, but more complex modem function and algorithm are designed (4). The characteristics of this ACS are tested in water tank whose length, width and the maximum water depth are 3m, 1m and 0.8m respectively. The diagram of this system is shown in Fig.2.

There are two kinds of transmit transducer. 1. The beam width in horizontal plane and in vertical plane are \( \theta_a = 1.8^\circ \) and \( \theta_r = 3^\circ \), respectively. 2. \( \theta_a = 1.8^\circ \) and \( \theta_r = 60^\circ \). There are two kinds of hydrophone also: (1) \( \theta_a = 3^\circ \), \( \theta_r = 60^\circ \); (2) Omnidirectional.

The received signal through receiver is fed into data acquisition board which is plugged in PC-286 computer. A speedup board made at DSP Tms32020 is plugged in computer also. The sampling data are analyzed with computer or with the help of speedup board. Since the multipath and fluctuation characteristics in underwater acoustic channel, an effective nonlinear quantizer is used as a soft quantizer to decide the results obtained by FFT. Then the results are analyzed with Viterbi algorithm with computer or with the help of speedup board.

The experiment results are:

1. The transmit transducer 1 and the
hydrophone 1 are settled at two ends in a water tank. In this case, the conditions generated the multipath are in existence due to the upper and lower surface, the ahead, behind, right and left wall of the tank. Since the beam width of the transmit and the receive transducer are narrow, the multipath signal never been observed in the experiment. The experiment results shown that, even when the error correct codes are not used, the phenomena of error bits are not observed. Since the alignment between beams is difficult and the narrow beam array is large, the narrow beam array as a method to overcome the multipath is used only in few cases.

2. The transmit transducer 2 and the hydrophone 2 are settled at the two ends in the water tank. The multipath are observed and its frequency spectrum are shown in Fig.3. If the encodes are not used and only the WFSK signals are used, the Pb is between $10^0$ and $10^4$. If the $K=3$, $r=1/2$ convolutional code implemented by four frequencies and VA with hard quantizer are used, the Pb is between $10^0$ and $10^2$. If all are the same as above, but a soft quantizer is used, the Pb is lower than $10^2$. The Pb in a Rayleigh channel is the largest, the Pb obtained in this experiment is the medium, and the theoretical value considered only the first reflected wave of sea surface is the lowest. The reasons are obvious.

The integrating techniques which consist of WFSK, convolutional codes and VA with soft quantizer offer a good method to overcome the multipath in underwater acoustic channel.

References


![Fig.1 Error probability of convolutional code, K=3, r=1/2.](image1)

![Fig.2 Diagram of short range underwater acoustic communication system](image2)

![Fig.3 Spectrum of multipath in underwater channel](image3)
REAL TIME UNDERWATER IMAGING SYSTEMS

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After the great development of real time imaging in medical echography last years, this technique is more and more often used in other domains such as underwater imaging. We present here two systems developed in our Laboratory : a frontal imaging system for remote operated vehicles and a multibeam echosounder dedicated to fishery applications.

FRONTAL IMAGING SYSTEM

Principle

The device insonifies the sea floor at low site angles and builds the image from backscattering such as in side scan sonar (fig 1). The emitted acoustical beam is widely open in the vertical plane and strongly focused in the horizontal direction at the interception of the floor. Information is achieved through the modifications of reverberation or through the shadows created by obstacles laid on the floor rather than specular reflection. As in other acoustical systems, the distance R of obstacles is related to the time of flight of ultrasound.

The scanning is realized by an electronic rotation of the beam in the $\Theta$-azimutal direction to obtain a sectorial R-$\Theta$ image of the sea floor. This rotation is achieved through the scanning of a focused aperture on a circular array of transducers.

The specifications retained are a 32° sectorial image and a range varying from 2 to 100 m with an azimuthal resolution of 0.7°.

![Fig. 1: frontal imaging principle: objects are detected by their projected shadow.](image)

The antenna

We have developed a transducer technology [1] based on composite ceramic (PZT rods embedded in a polymer matrix) which offers high coupling coefficient, wide frequency band, good electrical and mechanical matching and low losses. In this application, they can work from 600 to 900 kHz which is convenient for our ranges (<100 m). We have chosen a circular geometry because the azimuthal scanning is obtained by a simple translation of the focused aperture along the antenna and the delays needed are small enough.

The antenna is made of 64 transducers regularly spaced with a pitch of 0.5° (7 mm) on a circle of 0.8 m radius. The aperture is realized by 32 adjacent transducers; its width is about 220 mm compatible with the 0.7° resolution required even for the lower frequency. In the vertical plane, we need a wide aperture (40°) to insonify a large enough area of the seafloor. This is obtained by using cylindrical transducers instead of linear ones. Though, the final geometry of the antenna is toroidal.

![Fig. 2: Theoretical directivity pattern of a focusing aperture made of 32 transducers at 7 mm pitch on a circle (radius: 0.8 m) working at 600 kHz; focus: 24 m; parabolic shading.](image)

Beam forming

Among the possible shading patterns [2], we use a parabolic weighting for the reception (fig. 2); this pattern reduces the sidelobe level down to -18 dB but does not increase too much the 3 dB width of the directivity pattern: 0.65° instead of 0.58° without shading for a frequency of 600 kHz.

We have to take into account that using a pitch (7 mm) greater than the wavelength (1.7 to 2.5 mm) leads to the presence of grating lobes (fig. 3). We can avoid artefacts induced by these lobes by using an adequate pattern at emission.

![Fig. 3: Presence of grating lobes; same aperture than on fig. 2 except no shading.](image)

We use the dynamic control of focalisation versus distance that we have developed for medical imaging systems [3], i.e. the signals coming from the different elements of the aperture are injected according the parabolic law used in Fresnel approximation in an analogic delay line made of inductance/capacity cells where the capacity is controlled through a DC polarisation during reception.

Image formation

For each range (6, 12, 24, 48 or 96 m), the sectorial image is built from 64 rays spaced with a pitch of
0.5°. For each ray, the signal coming out the delay line is amplified to compensate attenuation and diffraction effects, digitized on 8 bits, numerically detected and compressed to give a value on 4 bits for each one of the 512 pixels in R. Then the 64 x 512 R-R image is written on a 256 x 512 X-Y memory in a way ensuring a good interpolation in θ. The image is then displayed on a classical video monitor according a standard video rate.

The main problem in getting real time underwater images is due to the time of propagation. A pure sequential treatment would lead to long acquisition times (2s at 24 m). To increase the image rate, two solutions:

- Processing in parallel 4 beams by using parallel delay lines. This increases the complexity of electronics but reduces the acquisition time down to 0.5 s.
- Processing 4 pixels according 4 angles instead of one by rotating the focusing aperture during the reception time. This technique combined with the previous one increases the image rate to 8 images/second but deteriorates the reception conditions and furnishes images of lower quality which can be used when the vehicle moves rapidly.

In both cases, the emission beam has to be wide enough to cover the 4 or 16 beams built at reception. Though, the shading pattern is carefully chosen to minimize the side lobe level to avoid artefacts due to the grating lobes.

We present images of cylindrical and spherical concrete objects lying on a sediment bed under a few meters of water which show the good behavior of the system.

**MULTIBEAM ECHOSOUNDER**

**Principle**

We have chosen to display the information through a topographic sectorial (32° wide) image in the vertical plane orthogonal with respect to the ship axis. The lateral resolution is about 2° within the range which can vary from 50 to 800 m. We present here the numerical version of the sonar [4].

**The antenna**

We use the same kind of transducers made of composite ceramic than previously. The working frequency (70 to 130 kHz) is adapted to the range.

For the same reasons, we also retain a circular geometry for the antenna. The focusing aperture must be wide enough (> 30 λ) to achieve the 2° resolution. As we shall see later, the image is built with only one shot, so we cannot accept the presence of grating lobes at reception; this limits the pitch of transducers (p < 1.5 λ).

The retained antenna is made of 48 transducers at a pitch of 1° (18 mm) on a circle (radius : 1.1 m). In the orthogonal direction, their length is 225 mm which leads to a directivity of about 5° which is sufficient to ensure a good recovering between emission and reception beams even with a pitch due to heavy swells.

**Beam forming**

To simplify the system, beams are collimated at infinity. The deterioration is significant only in the area near the sounder (less than 10 m) which is out the area we want to observe. On the other hand, it remains important to reduce side-lobe level with a shading pattern. It is also important to take into account a temporal shading : out of the steering direction, the signals coming from the transducers do not completely overlap and the edge effects may be important; we can reduce their influence with a modulation of the signal. We can appreciate the influence of this modulation on the directivity pattern on fig. 4.

**Image formation**

The image is built with 16 simultaneously processed beams. The emission beam must be wide enough to cover the whole imaged area. The system provides 5 ranges : 50, 100, 200, 400 and 800 m with an emission duration which corresponds to the length of 2 pixels of the images.

The first step is an heterodyn detection which lowers the frequency down to about 8 kHz. In each beam, the signals coming from the transducers are digitalized, weighted with the shading coefficients, delayed and added. The detection is done by measuring the maximum peak-to-peak value within the pixel. A coefficient compensates the effects of attenuation and diffraction and the value of the pixel is given on 4 bits after a logarithmic compression. The pixel is adjusted to obtain 512 pixels in the R-direction for any range.

The 16 beams are selected within 96 preformed beams spread over 48° with a 0.5° pitch. The choice is controlled at each depth R with the information given by an inclinometer to compensate the effects of rolling during reception. As in the previous system, the R-th image is converted and interpolated in a X-Y image which is displayed on a monitor using false color encoding. Since all the beams are parallelly processed, the refreshing rate is only limited by the range and varies from 1 to 10 images per second.

The multibeam echosounder have been tested on the N.O. Thalassa of IFREMER during several sea trials. The images obtained show its capabilities : the detection of shoals or fishes (at lower ranges) is rather good even close to the sea bottom; furthermore, the width of the image permits to locate targets out of the axis of the ship and gives some information on the size of shoals. We get also informations from seabottom : its topography even for large depths; detection and localisation of obstacles which may be dangerous for trawls.

**CONCLUSION**

The two systems presented show that is possible to achieve real time operating in underwater imaging. The main difficulty is due to the time of propagation of ultrasound on wide ranges and pure sequential techniques have to be given up. A partial or complete parallel process may reduce the acquisition time in a ratio up to 16 and offer a rate from 10 down to 0.5 images per second.

**REFERENCES**

**TIME-FREQUENCY RELATIONS BETWEEN THE ACTIVE SONAR CHANNEL, WAVEFORM AND RECEIVER**

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1. INTRODUCTION

The classical method of the sonar design is based on the sonar equation which is a logic relation between the energy parameters of the sonar. In fact, the processes of sonar target detection, signal classification or the parametric estimation etc. is the one of the information extraction from the received signal or the echo.

Active sonar process is simplified as follow

\[ u(t) \rightarrow h(t, t) \rightarrow v(t) \rightarrow \text{receiver} \rightarrow y(t) \]

\[ (1) \]

where \( u(t) \) is the input signal, \( h(t, t) \) is the response function of the sonar propagation process, \( v(t) \) is the received signal of the sonar, \( y(t) \) is the output of the receiver.

\[ (2) \]

Owing to the ocean medium is random and non-homogeneous and the target is movable and range-spread, \( h(t, t) \) and \( v(t) \) are all random and time-variant.

From the point of view of theory of information and communication, to design a high-class active sonar system, an important thing is to select and design the sonar waveform and sonar receiver to match well the sonar channel [2].

Two key tools for the analysis of time-variant signal or channel, the one is ambiguity function (AMF) which represent the feature of the delay-Doppler resolution and the other is Wigner-Ville spectrum (WVS) which is the Wigner-Ville spectrum of the time-frequency distribution.

In this paper, the relations between the waveform channel and the receiver are discussed based on the two-dimensional WVS (TDWVS) of linear time-variant channel.

2. SONAR WAVEFORM

Generally, sonar waveform \( u(t) \) is narrow-band determined and the spectrum is \( U(f) \). The TDCF (Two Dimension Correlation Function) is defined as

\[ X_{c}(\Delta t, \Delta \varphi) = \int u(t + \Delta t) u^{*}(t) \exp(-j2\pi f t) df \]

The AMF is defined as \( \mathbb{W}_{w}(t, \varphi) = |X_{c}(t, \varphi)|^{2} \), and represents the time-frequency resolution characteristics. The WVS is defined as

\[ \mathbb{W}_{w}(f, \varphi) = \int u(t) \exp(-j2\pi f t) \exp(-j2\pi \varphi t) dt \]

and it represents the time-frequency distribution characteristics. Easily show the Fourier transform between \( \mathbb{W}_{w}(f, \varphi) \) and \( X_{c}(t, \varphi) \).

3. SONAR CHANNEL

We call all of the transverse channel, target scattering and the reverberation process as sonar channel. Assume that is a linear time-variant system. The time-variant response function and transfer function is \( h(t, t) \) and \( h(f, t) \) respectively, and other two system functions are spread function \( S(t, \varphi) \) and bi-frequency function \( B(f, \varphi) \). Usually, the sonar channel is partial coherent, i.e., \( H(f, t) \cong H(f, t) \), where \( H(f, t) \cong 1 \) \( - \Delta f = 0 \)

The correlation function of the channel is defined as

\[ R_{x}(t, t; \Delta t, \Delta \varphi) = \mathbb{W}_{w}(t, \varphi) \mathbb{W}^{*}(t + \Delta t, \varphi + \Delta \varphi) \]

\[ (3) \]

may be used to describe the fading or spread statistic characteristics of the channel. In order to analyze the influence of the channel on the signal, introduce the TDWVS (Two Dimension Wigner-Ville Spectrum) of the time-variant channel as

\[ \mathbb{W}_{x}(f, t; t, \varphi) = \int \int \mathbb{W}_{w}(f + \Delta f / 2, t + \Delta t / 2, f - \Delta f / 2, t - \Delta t / 2) \exp(-j2\pi f t) \exp(-j2\pi \varphi t) df dt \]

\[ (4) \]

\[ \mathbb{W}_{w}(f, \varphi) \]

\[ (5) \]

\[ \text{etc}[3] \]. The TDWVS of the time-variant channel represents the time-frequency fading and spread statistic characteristics.

Similarly, we introduce

\[ X_{c}(\Delta t, \Delta \varphi) = \int \mathbb{W}_{w}(f + \Delta f / 2, t + \Delta t / 2, f - \Delta f / 2, t - \Delta t / 2) df dt \]

\[ (6) \]

it is the 4-dimensional PDF of \( W_{y}(f, t; t, \varphi) \).

For the random channel, the ensemble average

\[ \langle X_{c}(\Delta t, \Delta \varphi) \rangle = \int \mathbb{W}_{w}(f, t; t, \varphi) \exp(-j2\pi f t) \exp(-j2\pi \varphi t) df dt \]

\[ (7) \]

where \( R_{x}(f, t; t, \varphi) = R_{W_{y}}(f, t; t, \varphi) \) in Eq. (4) -- symmetric correlation of the channel.

Some special channels such as fluctuate mono-path or point-target, and time-invariant multi-path or range-spread target etc., the TDWVS may be simplified. Some properties of TDWVS may be used to simplify the analysis of the complex channel [4]. For example, the echo channel \( E_{y}(f, t) \) is a series one by the transient channel \( H_{a}(f, t) \) and target scattering channel \( H_{t}(f, t) \) \( - H_{a}(f, t) = W_{1}(f, t; t, \varphi) \), and \( W_{a}(f, f; t, \varphi) = W_{w}(f, t; t, \varphi) \).

\[ \langle X_{c}(\Delta t, \Delta \varphi) \rangle \rightarrow P_{a}(\Delta f, \Delta t) \]

\[ (8) \]

where \( P_{a}(\Delta f, \Delta t) \) and \( R_{a}(\Delta f, \Delta t) \) is the coherence function and the scattering function of the channel, respectively.

For bi-temporal (B-TWVS) channel, we have

\[ R_{x}(f; t; \Delta t, \Delta \varphi) = R_{a}(\Delta f, \Delta t) \]

\[ (9) \]
and 

\[ <\psi_c(f, \tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]
\[ <\chi_c(\Delta t, \Delta t; \Delta t, \Delta t) > = R_{s}(\Delta t, \Delta t) \]

4. HYDROACoustIC SIGNAL (HAS)

HAS \(v(t)\) is the output of the sonar channel, i.e., the acoustic signal received in the ocean, e.g., the transient signal, target echo \(u(t)\), reverberation \(r(t)\) and ocean noise \(n(t)\). The main function of the active sonar is to extract the informations of the channel from the HAS. HAS is with duration \(\Delta t\) and bandwidth \(B\) and is usually random time-varying and non-stationary. Generally, \(v(t) = v(t) + v(t)\), \(\overline{v(t)} = 0\), and the correlation function is

\[ R_v(\Delta t, t) = \mathcal{L}(\Delta t/2 > v^*(t- \Delta t/2) \]

and may prove that the WVS of \(v(t)\) is

\[ W_v(f, t) = [\delta(h, f, t') \mathcal{F}_v(f, t' - t, -f) \delta(t' - t)] \exp(-j2\pi f \Delta t) dt' \]

and the ensemble average is

\[ <W_v(f, t) > = \mathcal{L}(\Delta t) \overline{W_v(f, t)} \exp(-j2\pi f \Delta t) \Delta t \]

Similarly, the TCF of \(v(t)\) is

\[ \chi_v(\tau, \varphi) = [\chi_c(\varphi, \varphi) \mathcal{F}_v(f, f + \varphi) \delta(t - \varphi)] \exp(-j2\pi f \Delta t) \Delta t \]

and the AMF of \(v(t)\) is

\[ \chi_v(\tau, \varphi) = [\chi_c(\varphi, \varphi) \delta(t - \varphi)] \exp(-j2\pi f \Delta t) \Delta t \]

and the fluctuation covariance of the energy of \(v(t)\)

\[ D_v(t, \varphi) = [\chi_c(\varphi, \varphi) \delta(t - \varphi)] \exp(-j2\pi f \Delta t) \Delta t \]

For WSSUS channel,

\[ R_v(\Delta t, t) = \mathcal{L}(\Delta t) \overline{R_v(f, t)} \exp(-j2\pi f \Delta t) \Delta t \]

For BTWSS channel,

\[ R_v(\Delta t, t) = \mathcal{L}(\Delta t) \overline{R_v(f, t)} \exp(-j2\pi f \Delta t) \Delta t \]

where \( \overline{E_s(f, t)} \) is the energy spectrum of \(u(t)\), \( \mathcal{P}_s(f, \varphi) \) is the bi-frequency correlation function of \(u(t)\). From Eq. (15), we can see that for BTWSS channel, the signal \(v(t)\) is WSS and \(W_v(f, t)\) is just power spectrum \(G_s(f, \varphi)\), and, \(f\) is not for WSSUS channel.

5. SONAR RECEIVER

By Eq. (10), we can obtain the output WVS of the time-varying receiver, only \(v(t) = v(t)\). There are different receivers for different purposes of the sonar. We discuss only the case for non-time-varying receiver \(Z(f, t) = \delta(f)\), the output is the cross-correlation function of \(v(t)\) and \(z(t)\)

\[ <\psi_v(\tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]

or

\[ <\chi_v(\tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]

the cross-AMF is

\[ <\psi_v(\tau, \varphi) > = <\chi_v(\tau, \varphi) > \mathcal{P}_s(f, \varphi) \]

In sonar detection, the output maximum value is

\[ \mathcal{L} = \gamma(0)^2 = <\psi_v > 0 \]

By Moyal's formula,

\[ \mathcal{L}^2 = \int W_v(f, t) W_v(f, t) df dt \]

For WSSUS channel

\[ <\psi_v(\tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]

This Eq. represents the decorrelation of the channel on the input signal \(u(t)\).

\[ \mathcal{L} = \int [\mathcal{F}_v(f, \varphi) \mathcal{F}_v(f, \varphi)] df \]

and

\[ <\psi_v(\tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]

From Eq. (16), we can see the matched filter is also a time-frequency correlator.

2. Receiver matched by channel: \((Z(f, t) = \delta(t)\) is the spectrum of target echo \(u(t)\), the output

\[ y(t) = \mathcal{F}_v(f, t) \]

and similarly, we have

\[ \mathcal{L} = \int \mathcal{F}_v(f, t) \mathcal{F}_v(f, t) df \]

\[ <\psi_v(\tau, \varphi) > = \mathcal{P}_s(f, \varphi) \]

where \(y(t) = v(t)\) in \(t\), and the reverberation is ignored.

We have discussed the time-frequency relations of the narrowband sonar processes. For broadband signal and for the echo channel with high-speed moving target, the time-frequency relations are more complex, and wideband AMF and WVS must be used.

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PREMIERS RESULTATS EN DISCRIMINATION SPECTRALE DES ECHOS SONAR LARGE BANDE (60-140 kHz) DE 6 ESPECES DE POISSONS

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1. INTRODUCTION

Nous rapportons ici les résultats d'une étude préliminaire de discrimination des espèces ichthyologiques à partir des réponses spectrales obtenues dans la bande 60-140 kHz, sur 45 individus de tailles diverses appartenant à 6 espèces différentes. Les animaux immobilisés par anesthésie ont été testés séparément, en vue dorsale, avec un système sonar expérimental que nous avons développé dans ce but. Ces résultats montrent que si les méthodes de "classification" donnent un taux de reconnaissance assez moyen celles de "discrimination" paraissent mieux adaptées au problème.

2. MATERIEL ET METHODE

2.1. Le système sonar utilisé

Les principales caractéristiques du système sonar expérimental utilisé sont les suivantes:
-bande passante utile de 20 à 140 kHz
-directivité constante en émission et réception (θ/2 = 11 à -3 dB)
-quasi absence de lobes secondaires
-niveau d'émission maximum : 190 dB/µPa (en valeur efficace)*
-signaux utilisables : toute forme d'onde courant la bande 20-140 kHz et dont la durée totale ne dépasse pas 1024 échantillons temporels (taille de la mémoire du générateur de signaux arbitraires).
-acquisition sur oscilloscope numérique (séquence ≤ 1024 échantillons temporels).
-pilotage du système par micro-ordinateur (BUS 13E 488).

2.2. Protocole expérimental

Les mesures ont été réalisées en cuve cylindrique (Ø = 2.5 m, hauteur = 3.1 m) sur individus isolés anesthésiés au phén oxyethanol juste avant l'expérience. L'animal était maintenu horizontalement à 1.5 m sous le transducteur par 2 attaches de nylon très fins manipulables de la surface, l'une étant fixée autour de la nageoire caudale, l'autre autour de la tête au niveau des ouies. Un lest horizontal (tige métallique de 30 cm) qui venait affleurer le fond de la cuve, fixé de la même manière au sujet, assurait sa stabilité en position qui était contrôlée visuellement au travers des hublots d'observation de la cuve.

Signal d'émission utilisé :
-durée : 0.8 ms ;
-bande spectrale : 20-140 kHz ;
-amplitude constante.
En réception, la bande a été limitée par filtrage passe haut à 60 kHz.
Un cycle sonar de 2 s. a été choisi pour assurer un évanouissement complet des échos sucsessifs fond-surface de chaque émission et la fenêtre d'acquisition a été réglée de manière à ne prendre en compte que le seul écho renvoyé par l'individu testé.
Un exemple d'écho renvoyé par un brochet est donné sur la fig. 1.

Fig. 1 : Echo sonar d'un brochet en vue dorsale. Signal d'émission de 0.8 ms et d'amplitude constante modulé linéairement de 20 à 140 kHz. Filtré passe haut 60 kHz.

Sur chaque individus 50 acquisitions successives ont été réalisées, mais seules les 10 premières d'entre elles ont été utilisées ici.

3. ANALYSE DES DONNEES

L'étude statistique réalisée a porté sur les index spectraux linéaires de réflexion des échos Ind(f_i) définis par :

\[ \text{Ind}(f_i) = \exp I(f_i)/(10 \log(10)) \]

avec \( I(f_i) = \text{Le}(f_i) - \text{Lr}(f_i) - \text{Sv}(f_i) - \text{Sh}(f_i) - G + K \log(10) + 20 \log(a(f_i)) \)

où :
\( f_i \) : fréquence discrète d'indice i
\( \text{Le}(f_i) \) : densité spectrale du signal d'émission
\( \text{Lr}(f_i) \) : densité spectrale du signal reçu
\( \text{Sv}(f_i) \) : réponse en émission du transducteur
\( \text{Sh}(f_i) \) : réponse en réception du transducteur
\( G \) : gain de réception
\( K \) : affaiblissement du signal en émission (régence)
\( a(f_i) \) : coefficient d'amortissement du milieu
\( r \) : distance émetteur-cible
Le pas d'échantillonnage fréquentiel des résultats \( f_i = f_1 \cdot i \) de 50 à 7 Hz

La fig. 2 représente l'index spectral de réflexion correspondant à l'écho de la fig. 1.

* Un transducteur CBT (Van Buren, 1983) de même type mais ne couvrant que la bande 27-54 kHz a été utilisé dans une application similaire par Simmonds et Copland (1987).
Les 6 espèces étudiées sont les suivantes : Brochet (Br) Esox lucius ; Carpe Cyprinus carpio (une distinction a été faite entre carpes cuir (Ca) et carpes écailles (Ce)) ; Omble (Om) Salvelinus alpinus ; Perche (Pe) Perca fluviatilis ; Tanche (Ta) Tinca tinca ; Truite (Tr) Salmo trutta. Le nombre d’individus utilisés, la fourchette des tailles et le nombre d’acquisitions faites dans chaque cas sont rapportés dans le tableau ci-dessous.

<table>
<thead>
<tr>
<th>Espèce</th>
<th>Br</th>
<th>Ca</th>
<th>Ce</th>
<th>Om</th>
<th>Pa</th>
<th>Ta</th>
<th>Tr</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mère</td>
<td>9</td>
<td>5</td>
<td>2</td>
<td>9</td>
<td>5</td>
<td>10</td>
<td>5</td>
</tr>
<tr>
<td>Tailles (cm)</td>
<td>40-60 39,5-42</td>
<td>35-45 29-37</td>
<td>38-49 31,4-48,1</td>
<td></td>
<td></td>
<td></td>
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</tr>
<tr>
<td>Mère acquis.</td>
<td>90 80 59 76</td>
<td>90 189 250</td>
<td></td>
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</tbody>
</table>

4. RESULTATS

834 index spectraux ont été obtenus puis normalisés par rapport à leurs maxima, chacun d’eux comportant 156 valeurs associées à la variable f_j. Diverses méthodes de classement et de discrimination ont ensuite été appliquées au tableau ainsi constitué.

4.1. Méthodes de classification

La méthode de classification dite "hiérarchique ascendante" - ou d’agrégation- ne donne pas de résultat intéressant, quel que soit le critère de classification choisi (distance inter-spectrale, euclidienne ou de "Manhattan"). Celle dite "des centres mobiles" - ou kmeans - (qui minimise, par repositionnement successif des individus dans les différents groupes, la variance intra-groupe) donne pour sa part des résultats assez satisfaisants comme le montre le tableau d’affectation ci-dessous.

<table>
<thead>
<tr>
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<tr>
<td>Br</td>
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<tr>
<td>Br 30</td>
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<tr>
<td>Ca 7</td>
</tr>
<tr>
<td>Om 0</td>
</tr>
<tr>
<td>Pa 17</td>
</tr>
<tr>
<td>Ta 26</td>
</tr>
<tr>
<td>Tr 0</td>
</tr>
</tbody>
</table>

Le taux de bon classement est en effet de 50 % alors qu’il ne serait que de 14 % par une affectation au hasard des individus.

4.2. Méthodes de discrimination

Les diverses méthodes discriminantes utilisées : discrimination linéaire ou quadratique, avec ou non sélection pas à pas des variables ou lissage des index spectraux par pondération (0.25 ; 0.50 ; 0.25) et échantillonnage- donnent dans l’ensemble des résultats apparemment satisfaisants qu’une validation croisée vient toutefois mettre en évidence qu’un rapport plus favorable entre le nombre des individus (actuellement 45) et celui des variables (156) conduirait à de meilleurs résultats.

Ainsi, une réduction à 16 du nombre de variables par sélection pas à pas, suivie d’une analyse discriminante linéaire donne 96 % de bon classement, mais la validation croisée, par une seule acquisition par individu, ramène ce résultat à 66 %. Les indices des fréquences les plus discriminantes étant alors : 1, 4, 17, 20, 21, 24, 28, 31, 32, 33, 35, 86, 120, 151, 152, 155.

5. CONCLUSION

Ces résultats confirment les conclusions d’autres travaux (Kjaergaard et al., Lebourges, Zakaria et al.) sur la potentialité d’une discrimination objective des espèces ichthyologiques par sonar large bande (qui doit toutefois être validée sur animaux actifs).

Les méthodes de traitement utilisées mettent par ailleurs en évidence la nécessité d’améliorer le rapport individu/variables ce qui peut s’obtenir par une représentation plus condensée du signal (modélisation, cepstrum...) et l’on peut penser pouvoir améliorer ces résultats par un choix d’une méthode de discrimination plus appropriée qui serait appliquée non pas au seul spectre énergétique des échos rétrodiffusés mais aussi à leurs caractéristiques temporelles.

BIBLIOGRAPHIE


ACOUSTIC DOPPLER CURRENT PROFILER
AND ITS SIGNAL PROCESSING

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INTRODUCTION

Acoustic Doppler Current Profiler (ADCP) is an advanced instrument measuring the vertical profiles of the ocean current. It can measure the tridimensional velocity components $(V_x, V_y, V_z)$ and their absolute direction in 1 to 64 cells of a vertical profile without flow disturbance. The ADCP described in this paper was developed by the Institute of Acoustics, Academia Sinica. It can be towed by a ship to measure the ocean current in a wide range, or mounted on the sea bottom for long time survey. So it is useful in scientific research and the ocean resources exploration.

The main technical specification:

High Accuracy (for velocity): $(1-3)$cm/s.
(for direction): $(2-5)$ degrees.

Remote Current Profiling:

---- Spatial Coverage : up to 150 m.
---- Measured Cells : up to 64 cells.
---- High Depth Resolution : 1 m.

Acoustic Frequency : 300 KHz.

Acoustic Beams: 4-beam JANUS, oriented 30 degrees off vertical in 90 degrees azimuth increments.

Maximum towed speed : 15 knots.

Maximum velocity measured range: +/-10 m/s.

The current data is processed in real time and displayed in numerical & graphical forms.

THE TOWED ADCP DESCRIPTION:

The Towed ADCP consists of a towfish, a electronic box, a PC-286 computer and a winch. The block diagram of the ADCP is shown in Fig.1.

Within the towfish, there are a 4-beam transducer array in JANUS configuration, a 4-channel high-power sonar transmitter (each channel can transmit one kilowatt pulse electric power), a 2-channel low noise sonar receiver, a control unit employing a single chip computer MCS-51, a liquid pressure sensor, and a thermal sensor to correct acoustic speed near the transducer array surface, and a compass to indicate the absolute direction.

The towfish is designed to be towed at the speed of less than or equal to 15 knots, and can be towed by various ships. By our experience, if the towfish is towed at the depth which is deeper than twice the draught of ship, it will not only avoid the noise of the ship, but also has a small pitch and roll angles, which are less than 2 degrees. Since the frequency of pitch and roll are generally much lower than the ping rate, the variance of the velocity estimate caused by the pitch or roll is:

$$\sigma^2 = V_o^2 \cdot \frac{\phi^2}{150}$$

where $\phi$ is a maximum pitch or roll angle in radian and $V_0$ is the velocity without pitch or roll. If $\phi = 2$ degrees, the measurement error is 0.044m/s. So it is needed to revise the pitch and roll angles in the towed ADCP.

In our ADCP, each beam width is designed to be 4 degrees. Because a wide beam will induce a large measurement error of velocity, while a narrow beam will require a large transducer array, there should be a trade-off between them. On the other hand, in a observation cell, move with the same velocity, the relative measurement error of velocity due to the beamwidth is:

$$\frac{V_o - V_\theta}{V_o} \approx \frac{\phi}{2\theta}$$

where $\theta$ is the beamwidth in radian.

In order to reduce the number of wires in the toving cable, frequency division multiplex modulator and demodulator are used to enable several channel signals to be transmitted by a coaxial line between the towfish and the electronic box.

During operation, an acoustic pulse is transmitted first, then the volume reverberation signals are received by transducers, amplified by preamplifiers, compensated by adaptive time-gain circuits, filtered by narrow-band filters and modulated into frequency division multiplex signal that can be transferred to the electronic box by a coaxial cable.

Within the electronic box, the modulated signal coming from the towfish is demodulated into several channel signals. Then the analogue signals are digitized by a multichannel A/D converter in PC-286 computer. The Doppler drift or the velocity for each cell is calculated by a two-word DSP32020/225 digital signal processor in real time. The results can be displayed on the monitor and stored in disk.

ADAPTIVE TIME-GAIN COMPENSATION:

We know that if the distribution of the scatterers illuminated by an acoustic beam is uniform, the volume-reverberation level is mainly affected by spherical spreading and acoustic absorption. Actually, the distribution of the scatterers, a large number of suspended particles and plankton moving along with the ocean current, is nonuniform and complicated, and varies with season, sunshine and region. So the attenuation of volume reverberation signal usually deviates from the pattern of the spherical spreading. In experiment, we found that a fixed time gain compensation curve used to compensate the attenuation will cause the signal distortion and large measurement error. In our ADCP, an adaptive time-gain compensation circuit controlled by a single-chip computer is used, which has 8048 dynamic range and can provide smooth optimum compensation curves by studying the envelope of reverberation signal. This adaptive time-gain compensation method tracks the variation of volume-reverberation level rapidly and precisely, and enables the ADCP to work at optimal state in various conditions. The volume-reverberation strength can be calculated through a inverse transformation of the compensation curve.
Fig. 2 and Fig. 3 are two typical time-gain compensation curves and the relevant output signals of the receiver. It is obvious that the output signal in Fig. 2B is not well compensated by Fig. 2A, that is derived from the spherical spreading and the acoustic absorption. It is the reason that the distributed density of plankton is greater near the bottom than in upper water. Fig. 3A is a better compensation curve provided by the computer after studying the envelope of signal showed in Fig. 2B, and Fig. 3B is a well compensated signal wave.

THE SIGNAL PROCESSING FOR THE ADCP

In the ADCP, it is very important to estimate the Doppler drift. Many methods, such as zero-cross, phase-locked loop, and coherence, etc., have been developed. Among these, the covariance method is considered the best one and often adopted, since it can give good estimations for mean frequency and its variance through small computation. But in some conditions, the variances are large. We consider that the large variance of mean frequency is caused by turbulent flow or eddy [1], and only a mean frequency cannot give us the information about it.

Fig. 4 is the power spectral density of the ocean current data, which is obtained by Complex Hanning Data-Adaptive weighted Burg algorithm (CHDAB) [2]. There are two or more peaks in spectra, and every spectral peak represents a Doppler drift corresponding to current velocity. We think that in some conditions the power spectrum, which can show us fine structure of the ocean current and some helpful information for turbulence study, is more useful than only a mean frequency. Comparing with the covariance and FFT methods, the CHDAB is a good algorithm for the spectral estimate in the ADCP, which has a high resolution and precision for frequency estimate. The covariance method cannot give the power spectrum and the FFT method has a low resolution of frequency in short data.

**REFERENCES**


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![Fig. 1. The block diagram of the towed ADCP.](image)

![Gain Fig. 2(A).](image)

![Gain Fig. 3(A).](image)

![Amplitude Fig. 2(B).](image)

![Amplitude Fig. 3(B).](image)

Two typical TGC curves and the output signal waves

![Fig. 4. Range evolving the incremental Doppler spectra for AB, model of 4ms, 300 kHz data](image)
SIGNAL PROCESSING IN GEOSONAR

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1. INTRODUCTION

Geosonar which has been recognized as a high frequency seismometer by geophysicists is in fact a sonar of low frequency. The spectrum of a typical geosonar with high resolution may extend from 200Hz to beyond 10kHz. As a portable equipment, the dimension of its sound array is usually limited to about half a meter. Only thinking of the fact that for a military sonar operating on the frequency of 3kHz, the dimension of sound array may be over 3 meters, it is easy to imagine how bad the S/N ratio of a geosonar will be! Especially in case of surveying in very shallow water, multiples in water layer, residual oscillations of tides, etc., produce a lot of interferences that may cause a measurement result useless. During the past decade, how to reduce this kind of interferences and improve the measuring technique of geosonar has been the aim of our lab all along. This article will introduce the signal processing technique of geosonar both in time and space domain and some application result as well.

2. SIGNAL PROCESSING IN SPACE DOMAIN

Small Planar Array with High Gain and Short Residuals

Because a great part of noise comes from the ship's propelling gear, it isn't homogeneous in space. If this feature can be utilized in the design of sound array, the S/N can be enhanced, though the directivity factor depends on geometry dimensions only. e.g., the array gain of a towed streamer may be much higher than the directivity factor. We have developed a small planar array (GPV) with multi-layer air baffle. It is of high array gain and with very low residual oscillations. Fig.1a and Fig.1b are the directionality diagrams of GPV and the common array respectively. The low sensitivity on the horizontal direction (-20dB) of GPV leads nearly to the same result just like a short streamer against the propelling noise. And because the backward sensitivity is -26dB down compared with the axis, so the "ghost" (due to air-water interface reflection) could never been seen on our record. It is the advantage that a streamer couldn't offer. It can work well in extra shallow water (1-2 meters).

Multiplicative Array

Although the multiplicative array is well known by the sonar designers, however, to the author's knowledge, the linear additive array is used in current geosonars without exception. In case of measuring in the area with sandy sediment, the operation frequency needs to be much lower than 1000Hz. If the frequency is 300Hz, then the wavelength will be 5m, there isn't any directivity for a small planar array. In that case, a multiplicative array only with two hydrophones (Fig.2) could be utilized. Fig.3 shows the output of two hydrophones, and after TVG and rectification. As seen, the merit of multiplicative array is obvious. For a additive array with two elements, the theoretical processing gain is AG=10lg(n), where n=2, AG=3db, it is quite low. But in the case of multiplicative array, a high processing gain can be obtained due to the time-arriving difference of the short-duration pulse. For example, the direct arrival interference which is frequently seen on the profiling chart working at low frequency, almost disappear thoroughly. Also, water layer multiples and residual oscillations are restrained by this method. However, in case of detection under strong noise background, the S/N ratio may be a problem, but it can be improved by other way of signal processing, e.g., stacking, etc.

3. SIGNAL PROCESSING IN TIME DOMAIN

Time-varying Filtering

To satisfy the demands of resolution and penetration simultaneously, the spectrum of a typical high quality geosonar may extend from several hundreds to beyond 10kHz. However, the high frequency components decrease much more rapidly with the penetration than the low frequency ones do. Hence, for keeping the output S/N ratio maximum all the time, a time-varying filter should be used. The central frequency of this filter is

\[ \omega_{c,\alpha} = \frac{2\pi}{T} \left( \frac{E_1 \cdot A_1 \cdot \beta_1}{\sum_{i=2}^{n} E_i A_i \cdot \beta_i} \right) \]

where
- \( E_1 \) -- thickness of 1st layer
- \( A_1 \) -- absorption of 1st layer
- \( t \) -- average sound velocity in sediment
- \( \beta_1 \) -- average absorption of sediment
- \( A_1 \) -- reflectivity of sea bottom

A switched capacitor filter adapt to this purpose. It not only can give simultaneously the highest resolution and maximum penetration being able to offer by the system, but also suppress the multiple interferences, i.e., it can bring all potentialities of the system into full play. The output reverberation to signal ratio is

\[ \text{Signal to Noise Ratio} = \frac{S}{N} = \frac{A_0}{A_1} \]

where
- \( N \) -- peak value of the echo arriving at the same time with mth multiple
- \( A_0 \) -- reflectivity of sea bottom
- \( A_1 \) -- reflectivity of sea bottom

Signal Stacking

While the noise is random process, the signal from sediment is a fully constant one, so the successive echoes are stacked, the S/N must increase. Under ideal circumstances, after stacking n times, the gain of S/N is 20log n. When the surveying route is along a linear line, every time of transmitting and receiving conduct at different spot, the effect of stacking just like to form a virtual streamer. The synthetic arraylength will be \( l=(n-1)\cdot T \), where \( V \) is the surveying velocity; \( T \) is the shot period.

Synthetic Planar Array

If we combine the preceding ideas about multiplicative array and virtual streamer, a planar array suitable for low frequency use can be synthesized by two hydrophones mounted on both boards of the ship and with a microcomputer.

Seabottom Tracking and Dual TVG

Seabottom tracking can effectively prevent from strong interference in water layer. Dual time-varying gain control (compensating the spherical spreading losses and sediment absorption separately) can minimize the dynamic range of the output. Both are beneficial to the diagram quality and automatic
4. EXPERIMENTAL RESULTS

Since 1980s above-mentioned processing techniques have been successfully used in GPY and PGS subbottom profilers. The typical results are as follows:

Fig. 4 illustrates the function of TVF. In front of Harbor Beilun there was a air-filled layer, a few meters beneath seabed. The strong multiple interference, formed by that layer and seabed, diminished due to TVF.

Fig. 5 was a record of harbor Amoy. Due to all above processing techniques, the stratifications above bedrocks are shown very clearly. We almost couldn't find any interference and false events.

Fig. 6 was taken in Hangzhou Bay. Some interesting geological phenomena near an underwater trench reveals vividly.

Fig. 7 was the profile of Port Helsinki, where acceptable penetration and resolution has been difficult to acquire by common profilers.

Through these tests, it seems to be possible to develop a system fulfilling the needs of most marine geologists.

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1. Introduction

Echo-sounder or fish-finder is usually used to observe the underwater area and sea floor. It is difficult to take optical photographs of wide range underwater area because of great transmission loss of light in the water. When we search the wide range underwater area for a certain object, we infer the image of the target object from multi data produced by the echo-sounder. A frontal image or a contour map is displayed on the echo-sounder through cross fan-beam or the acoustical imaging system. However, it is extremely difficult to understand three dimensional shapes of the object and situation from the display. Especially when the object exists under sea floor or covered with other objects and material, the countour map cannot display the object exists under sea floor. To solve these problem, in our experiment, we set up a special prototype system to get three-dimensional bird eye image of sea floor and buried object in the marine sediments using personal computer to process various data obtained by echo-sounder equipped mono pulse cross fan beam.

2. Outline of the prototype system

Fig.1 shows block diagram of the prototype system used in the experiment. The echo-sounder which consists of cross array transducers is used to this system. Three disk type inductance transducers are arranged on line and the mono pulse sound is transmitted 5 times per second with 14°×48° fan beam. The reflected sound is received by 8ch linear array hydrophones with 20°×4° fan beam and recorded on the magnetic data tape by each channel on boat. The recorded data are played back and beamformed through the phase synthesizer and stored on the floppy disk of personal computer on shore. Three dimensional bird eye images correspond to reflected acoustic level and displayed with suitable angle. The area of display is limited by the capacity of the floppy disk. In our experiment, 32 horizontal data by each 1° for athwart-ships direction(x-axis), 256 data sequence by each 0.2 second along the ship course (y-axis) and 32 vertical data by each 15cm (z-axis) are stored by beamforming on the memory. Accordingly, the area of display will be: 2×(depth)×tan20° for X-direction, (velocity of ship)×(0.2 sec.)×256 for Y-direction and 480cm for Z-direction.

3. Experiments of three dimensional display

In the experiments, we applied this system to display acoustical bird eye images of objects and a school of fishes over the sea floor. The sediment layers as well as buried objects in the sediments in the shallow water.

3.1 Display of the objects and a school of fishes over the sea floor

Fig.2 shows example displays of the object and a school of fishes in about 12m depth sea. Fig.2(A) shows the analog record by commercial echo-sounder with 200kHz. Fig.2(B) shows the dot line area of Fig.2(A) color displayed by color fish-finder. Fig.2(C) shows three dimensional bird eye image of a school of fishes and object. We notice on the display the part of a school of fishes swarms over the far side of the object. The threshold level for this display is set at low level to detect the school of fishes because the reflected wave level from the school of fishes is weaker than from the sea floor and the object. On the other hand, Fig.2(D) shows the case in which the threshold level is set at high level as the sea floor surface and the object are only displayed and the school of fishes excluded. Over view angle of these bird eye images is 30°.

3.2 Display of the buried object in the sediments

Fig.3 shows examples of three dimensional bird eye images of the objects which are assumed to be buried in the sediment about 2m under the sea floor. The all original data are displayed in Fig.3(A). The sea floor is generally flat and the buried objects are found under the sea floor. However, it does not clearly show the sea floor image. Fig.3(B) is the thin out display of the same data to find the buried objects under the sea floor. Further in Fig.3(C), the image of only these objects buried about 2m under sea floor is displayed with 60° over view angle. We can find the distribution of large and small buried objects clearly from this image.

4. Conclusion

Three dimensional bird eye images of objects and the school of fishes and the buried objects in the sediments can be displayed by personal computer in processing data from echo-sounder equipped with mono pulse cross fan-beam. These images serve to understand three dimensional shapes and situations more easily than the contour map or the cross section of the echo-sounder. This system is useful especially to detect sunken obstacles on the sea floor or buried in the marine sediments.
Fig. 1 Block diagram of the prototype system.

Fig. 2(A) Analog record of the object and a school of fishes by the echo-sounder.

Fig. 2(B) Display by the color fish-finder.

Fig. 2(C) Three dimensional bird eye image of a school of fishes and the object.

Fig. 2(D) Three dimensional bird eye image of the object and sea floor surface.

Fig. 3(A) Three dimensional bird eye image of sea floor and buried objects.

Fig. 3(B) Thin out display of Fig. 3(A).

Fig. 3(C) Three dimensional bird eye image of only buried objects.
ACCURACY ANALYSES OF BISTATIC SONAR LOCALIZATION

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1. INTRODUCTION

With the development of underwater countermeasure and counter-countermeasure, from 1950's, a separated source and receiver sonar system was studied, which is called bistatic sonar [1]. If the system consists of more than one source or receiver, it is called multistatic sonar. In the late 1960's and early 1970's, Hazeltine Corp. made many experiments on bistatic-multistatic sonars in north Atlantic and Mediterranean[2].

The studies were put aside due to limited technology. Recently, the development of bistatic sonar comes into a new phase as the result of the successful application of bistatic radar. In 1989, Henry Cox expended bistatic sonar comprehensively and systematically[3].

This paper starts from the bistatic geometry, and discusses three ways and their accuracy. In order to obtain the concept directly on the condition of presumed parameters, the localizing errors are drawn with curves and are compared. The engineering realizability is taken into consideration.

2. LOCALIZATION METHODS AND THEIR ACCURACY

The bistatic geometry is a triangle of source, target and receiver. The equi-time of arrival ellipse is shown in Fig.1.

It is necessary to consider the notion system from the point-of-view of the receiver. It is useful to define an equivalent monostatic range, $R_m[3]$, as follows,

$$R_m = \frac{(R_1 + R_2 - R_3)}{2}$$  \hspace{1cm} (1)

2.1 Method 1: with $\gamma \rightarrow R_3 - R_m$
By[3],

$$R_2 = \frac{R_m + R_3}{1 + \left( \frac{R_3}{R_m} \right) (1 - \cos \gamma)}$$  \hspace{1cm} (2)

Generally, mean square error (M.S.E.) is used in error estimation:

$$\sigma^2 = \left( \frac{\partial R_2}{\partial R_m} \right)^2 \Delta R_m^2 + \left( \frac{\partial R_2}{\partial R_3} \right)^2 \Delta R_3^2 + \left( \frac{\partial R_2}{\partial \gamma} \right)^2 \Delta \gamma^2$$  \hspace{1cm} (3)

Where $\Delta \gamma$ is composed of two terms: $\Delta \gamma_1$(measurement error related to the beamwidth of Receiver), $\Delta \gamma_2$(error caused by Receiver position imprecision).

Two curves are drawn with two groups of data in Fig.2. The presumed parameters are,

$$R_m = 15km, R_3 = 35km, \Delta R_m = 50m, \Delta \gamma_1 = 0.1^\circ$$

2.2 Method 2: with $\alpha \rightarrow R_3 - R_m$

The same as Method 1,

$$R_2 = \frac{2R_m^2 + (1 - \cos \alpha)(R_3 + 2R_m R_3)}{2R_m + R_3 (1 - \cos \alpha)}$$  \hspace{1cm} (4)

Now, $\gamma$ is unknown. In order to implement localization, $\gamma$ should be evaluated,

$$\gamma = \text{arcSIN} \left[ \frac{\text{SIN} \alpha}{R_2} \left( 2R_m + R_3 \right) - \text{SIN} \alpha \right]$$  \hspace{1cm} (5)

2.3 Method 3: with $\alpha \rightarrow R_3 - \gamma$

By the law of sines, the $R_2$ is given out,

$$R_2 = \frac{R_3 \text{SIN} \alpha}{\text{SIN}(180^\circ - \alpha - \gamma)}$$  \hspace{1cm} (6)

To localize, it must be ensured that $\gamma + \alpha < 180^\circ$

3. DISCUSSION AND COMPARISON

With Method 1, when $\gamma = 180^\circ$, the error of localization is equal to that of monostatic case. When $\gamma$ is of other degree, the error is larger than monostatic case. When $\gamma = 45^\circ$, the error reaches its maximum. When $\gamma > 120^\circ$, the difference between bistatic and monostatic range error is small, and the influence of $\Delta R_3$ is minor.

For implementation of Method 1, the receiver directivity needs to be sharp, but the source directivity doesn’t. Localization needs no more information of source, which makes it convenient for the receiver to work independently.

To compare Method 1 with Method 2, convert the variation of $\alpha$ to that of $\gamma$, and draw curves in Fig.2. Obviously, when $\gamma$ is small ($\gamma < 60^\circ$), Method 2 leads to less range error than Method 1 does.

Evaluating $\gamma$ is inevitable to be obtained the bearing with Method 2. Within the range of measurement, bearing error is rather small. The maximum error of bearing is about $0.6^\circ$
For implementation of Method 2, the
directivity of transmitter is expected to be sharp, and
there is no requirement on receiver directivity. \(\alpha\)
must be known, therefore, it requires an one-way
communication from source to receiver.

Method 3 uses two angles of the triangle. It is
very unfavorable if one of the inner angles is
large(close to 180°) or small(close to 0°). To com-
pare with the former, Method 3 is discussed with the
equi-time arrival ellipse limit, as the former. Choose
\(\alpha\) as a variable, delete some large value points and
Fig.3 is drawn.

The localization accuracy of Method 3 is worse
than the former. Comparatively, when 30°
\(<\gamma<70°\) ,the error is small and varies slowly, and
this method can be used for coarse localization.

Directivity of both source and receiver need to
be sharp for Method 3 implementation, and beams
should be synchronized strictly. \(\alpha\) should be known,
and hence, there must be an one-way communi-
cation from source to receiver and the data rate is high.

4. CONCLUSION

Method 1 does not need synchronous beams of
source and receiver. It is convenient to implement.
The receiver can work independently. When \(\gamma\)
is large, the range error is close to monostatic case. Its
accuracy is high. It is a better way for bistatic realiza-
tion.

Method 2 has the same advantage which does
not need synchronous beams. Its error is minimum
among the three ways. Communication is
required. It is a good way.

For Method 3, when the angles are medium,
the error is comparatively small. The method re-
quires communication and the data rate is high. It
can be used as a coarse measuring way.

![Fig.1 Equi-time of arrival ellipse](image)

S:Source; R:Receiver; T:Target; \(\alpha\):TSR; \(\beta\):STR;
\(\gamma\):TRS; \(R_1\):range,Source to Target; \(R_2\):Target to
Source; \(R_3\):Source to Receiver.

![Fig.2 Range M.S.E of Method 1 and Method 2](image)

1:Method 1,\(\Delta R_1=50\)m; 2:Method 1,\(\Delta R_1=10\)m;
3:Method 2,\(\Delta R_1=50\)m;4:Method 2,\(\Delta R_3=10\)m;
5:Monostatic case with the same parameters

![Fig.3 Range M.S.E of Method 3](image)

1:\(\Delta R_3=50\)m; 2:\(\Delta R_1=10\)m.

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PERFORMANCE ANALYSIS OF THE MUSIC ALGORITHM FOR DIRECTIONS-OF-ARRIVAL ESTIMATION OF COHERENT SIGNALS

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I. INTRODUCTION

The MUSIC method can estimate the directions of arrival (DOAs) of incoherent and partially coherent signals and yield high resolution. For coherent signals, this method encounters serious difficulties. In order to overcome these difficulties, Evans et al. developed the spatial smoothing technique [1]. In order to improve the performance of solving the coherent signals, we also presented the spatial averaging scheme [2]. However, their performances are affected by many factors. In this paper, we analyse the performance of the MUSIC algorithm with the spatial averaging scheme.

II. MUSIC WITH SPATIAL AVERAGING SCHEME

Consider a uniform linear array consisting of \( M \) identical sensors and receiving signals from \( q \) narrow-band coherent sources which arrive at the array from directions \( \theta_1, \theta_2, \ldots, \theta_q \). The received signals at \( i \)-th sensor can be expressed as

\[
x_i(t) = \sum_{i=1}^{M} S_k[t - (i-1)\tau_k] + n_i(t) \tag{1}
\]

where \( S_k(t) \) represents the \( k \)-th source signal received at the first sensor, \( \tau_k = \frac{d \sin \theta_k}{c} \), \( d \) represents the separation between sensors, \( c \) denotes the propagation velocity of signals and \( n_i(t) \) denotes the additive noise at the \( i \)-th sensor. It is assumed that the signals and noises are stationary, zero mean uncorrelated random processes, and further the noises are assumed to be uncorrelated and identical between themselves. The array output covariance matrix is defined as

\[
R_M = E[X(t) \cdot X^H(t)] \tag{2}
\]

where \( X(t) = [x_1(t), x_2(t), \ldots, x_M(t)]^T \) \( E[\cdot] \) denotes the expectation operator, \( T \) denotes transpose, and \( H \) denotes complex conjugate transpose. Let \( r_y = E[x_i(t)x_i^*(t)] \). Since \( r_y = r_y^T \), \( R_M \) is a Hermite matrix. When the received signals are uncorrelated between themselves, \( R_M \) also is a Toeplitz matrix. The MUSIC with spatial averaging scheme is carried out in two steps. First, we evaluate the average covariance matrices of each interferer, and construct a \( N \times N(\leq M) \)-dimensional Toeplitz matrix \( R_M \), with all the elements on the \( k \)-th diagonal being

\[
r_k = \frac{1}{M-k} \sum_{i=k}^{M} r_{ii+k} \quad k = 0, \ldots, N-1 \tag{4}
\]

Then, through the eigendecomposition of matrix \( R_M \), we obtain the MUSIC algorithm with spatial averaging scheme.

III. PERFORMANCE ANALYSIS

For convenience, we only consider two coherent signals \( S_1(t) \) and \( S_2(t) \) from directions \( \theta_2, \theta_2, \ldots, \theta_q \), i.e., \( S_2(t) = \alpha S_1(t - \tau_{12}) \), where \( \alpha \) and \( \tau_{12} \) denote the relative amplitude attenuation and the time delay between the two coherent signals received at the first sensor. Hence, the received signals at the \( i \)-th sensor can be expressed as

\[
x_i(t) = S_1[t + (i-1)\tau_1] + S_2[t + (i-1)\tau_2] + n_i(t) = S_1[t + (i-1)\tau_1] + \alpha S_1[t + (i-1)\tau_2 - \tau_{12}] + n_i(t), \quad i = 1, \ldots, M \tag{5}
\]

where \( \tau_1 = \frac{d \sin \theta_1}{c} \) and \( \tau_2 = \frac{d \sin \theta_2}{c} \). Further, \( r_y = E[x_i(t)x_i^*(t)] \)

\[
= E[S_1[t + (i-1)\tau_1] \cdot S_1^*[t + (i-1)\tau_1]] + E[S_2[t + (i-1)\tau_2] \cdot S_2^*[t + (i-1)\tau_2]] + E[S_1[t + (i-1)\tau_1] \cdot S_2^*[t + (i-1)\tau_2]] + E[n_i(t) \cdot n_i^*(t)]
= r_{11}[\tau_{12} - (i-1)\tau_1] + r_{22}[\tau_{12} - (i-1)\tau_2]
+ r_{12}[\tau_{12} - (i-1)\tau_2] + \sigma_n^2 \delta(i-j) \tag{6}
\]

When \( S_1(t) \) and \( S_2(t) \) are uncorrelated, the third and fourth terms in (6) are zero. If \( S_1(t) \) and \( S_2(t) \) are correlated, these terms are no longer equal to zero, i.e., \( r_{11}[\tau_{12} - (i-1)\tau_1] = E[S_1[t + (i-1)\tau_1] \cdot S_1^*[t + (i-1)\tau_1]] = E[S_1[t + (i-1)\tau_1] \cdot \alpha S_1^*[t + (i-1)\tau_2 - \tau_{12}]] = ar_{11}[(i-1)\tau_{12} - (i-1)\tau_1 - \tau_{12}] \)

and \( r_{12}[\tau_{12} - (i-1)\tau_2] = ar_{11}[\tau_{12} - (i-1)\tau_1 - \tau_{12} + \tau_{12}] \)

Since the values of \( r_{11}[\tau_{12} - (i-1)\tau_1] \) and \( r_{12}[\tau_{12} - (i-1)\tau_2] \) are only related to the difference \( (j-i) \), but not to the absolute \( i \) or \( j \), they are not changed after averaging the elements on each diagonal of matrix \( R_y \) arithmetically. It can be seen from (7) and (8) that the values of \( r_{11}[\tau_{12} - (i-1)\tau_1] \) and \( r_{12}[\tau_{12} - (i-1)\tau_2] \) not only are related to the difference \( (j-i) \), but also to the absolute \( i \) and to the time delay \( \tau_{12} \). Because \( S_1(t) \) and \( S_2(t) \) are narrowband signals, their autocorrelation functions are approximately periodic.
Hence, in order to make the values of $r_{ij}((j-l)\tau_1-(i-l)\tau_2)$ and $r_{mi}((j-l)\tau_1-(i-l)\tau_2)$ be smaller after averaging arithmetically on the condition that spatial average times is definite, it is required in the statistical sense that these values vary largely with the different i. Therefore, from (7) and (8), it is necessary that $\Delta r$ takes larger value to decorrelate the coherence well, while

$$\Delta r = |\tau_1 - \tau_2| = \frac{d\sin\theta_1}{c} - \frac{d\sin\theta_2}{c}$$

$$= \frac{2d}{c} \left( \frac{\theta_1 + \theta_2}{2} \right) \left( \frac{\theta_1 - \theta_2}{2} \right) \sin \frac{\theta_1 - \theta_2}{2}$$

(9)

From (9), we give the discussion in detail as follows:

1) On the condition that $\theta_1 - \theta_2$ is constant, $\Delta r$ takes maximum when $\theta_1 + \theta_2 = 0$, and $\Delta r$ tends to zero when $\theta_1 + \theta_2$ tends to $\pm \pi$. That is to say, in the case that two coherent sources are located at the symmetry positions with respect to the array broadside, spatial averaging can decorrelate the coherence more effectively.

2) When $\theta_1 - \theta_2 = \pm \pi$, $\Delta r$ takes maximum, and $\Delta r$ tends to zero when $\theta_1 - \theta_2$ tends to zero. Hence, the larger the direction angle difference between coherent signals is, the better the performance of decorrelation is.

3) The larger the value of $d$ is, the larger the value of $\Delta r$ is, and the better the decorrelation of coherent signals is. But the condition $d \leq \frac{\lambda}{2}$ must be satisfied, with $\lambda$ being the wavelength of the incident signals. If $d > \frac{\lambda}{2}$, the spurious peaks will appear in the spatial spectrum.

Because $r_{m1}[(j-l)\tau_1-(i-l)\tau_2]$ and $r_{m1}[(j-l)\tau_1-(i-l)\tau_2]$ are related to $\tau_1$, the spatial averages of $r_{m1}[(j-l)\tau_1-(i-l)\tau_2]$ and $r_{m1}[(j-l)\tau_1-(i-l)\tau_2]$ are still related to $\tau_1$ for the finite average times. For narrowband signals, $\tau_1$ can be represented by phase $\varphi_{11}$.

Through above analysis, we come to the conclusion that the performance of the MUSIC algorithm with spatial averaging scheme is related to the DOAs of coherent signals, the separate angle between the sources, the relative phase between coherent signals, the separation between sensors, and the spatial averaging times.

V. REFERENCES


![Fig.1(a)](image1a.png) ![Fig.1(b)](image1b.png)

![Fig.1(c)](image1c.png) ![Fig.1(d)](image1d.png)

![Fig.2](image2.png) ![Fig.3](image3.png)

![Fig.4](image4.png) ![Fig.5](image5.png)
RECONNAISSANCE DE FORMES, TRAITEMENTS D'IMAGES ET INTELLIGENCE ARTIFICIELLE POUR LE TRAITEMENT DE BRUTS SOUS-MARINS

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Introduction
La reconnaissance acoustique en milieu marin a pour but de détecter des cibles ou des obstacles ou d'en estimer certains paramètres comme leur vitesse, leur distance, leur immersion, etc. Cependant, de l'avis des experts du domaine, l'importante évolution technologique des hydrophones laisse entrevoir une augmentation conséquente de la quantité d'informations à traiter et il est donc nécessaire de développer rapidement des systèmes d'aide à l'identification de cibles en milieu marin. Nous proposons un environnement permettant l'acquisition et la paramétrisation du signal, la détection ou la description de phénomènes acoustiques, l'intégration et la validation de savoir existant sur le domaine.

Nous décrirons tout d'abord les fonctionnalités de notre environnement pour la paramétrisation du signal, le traitement des images, la reconnaissance de formes et le traitement des connaissances. Nous aborderons ensuite l'étude et la symbolisation de bruts sous-marins tels que les émissions sonars ou les bruits d'hélice.

La paramétrisation du signal
Le module de paramétrisation associé à un signal audio une représentation numérique en temps et fréquences. L'onde acoustique est décrite par une suite continue de vecteurs temporels, dont les composantes ou paramètres standard sont obtenues à partir des spectres d'une Ftt calculée sur 1024 points. Afin de rendre compte de tous les sons rayonnés par un navire, les conditions de calcul des ces paramètres sont dynamiquement modifiables pendant une session de travail.

Cependant, l'extraction automatique des bruts sous-marins nécessite l'étude de variations simultanées de plusieurs paramètres standard. Aussi, des paramètres temporaires sont définis par combinaison de paramètres standard, de scalaires ou d'autres attributs temporaires. Référencés par leur nom dans des prédicats ou des fonctions, ils sont calculés seulement lors de l'activation de la règle qui les utilise et respectent ainsi le processus de backtracking induit par PROLOG [1]. Par exemple, la règle suivante

```prolog
esm(R, F, G) :-
  max(esm(R, F, G, moyenne([F,G]))) = 0;
```
détermine pour chaque tranche de la zone temporelle z la valeur maximale entre l'entier zéro et la densité d'énergie de la bande de fréquences F/G à laquelle on a soustrait sa valeur moyenne. Comme le laisse entrevoy cette règle, nous utiliserons généralement ces outils pour décrire les représentants génériques de toute une classe de paramètres temporaires.

La reconnaissance de formes
Les outils proposés dans le cadre de la reconnaissance des formes[2] constituent les éléments de base pour la localisation, la description et la classification de phénomènes acoustiques pertinents. Sur nos paramètres de travail, nous pouvons, soit effectuer une mesure particulière dans un intervalle de temps ou de fréquences donné (recherche de périodicité, calcul des extrêms, etc.), soit modéliser et/ou symboliser les évolutions temporelles ou fréquentielles de ces paramètres. Par exemple, la naissance et la disparition d'un phénomène sont modélisés par le prédicat colline([I], [P, AR, LGM, ICROW, G, C, GB, D, D, D, D, D, D]) énumérant par backtracking toutes les collines du segment <I> dont le plateau, d'une largeur au moins égale à l'arg, ne varie pas au dessus du seuil fixé par m ICROW et dont la pente gauche (resp droite) est contrainte par un seuil minimum c, g (resp c, d).

Le problème du bruit en milieu marin
Appelée le monde du silence, la mer est en fait un milieu extrêmement bruité [3][4][5][6]. La conséquence première de ce bruit est le masquage par lui total de certains des phénomènes étudiés. La nature du bruit marin et des phénomènes nous étant inconnue, nous ne pouvons appliquer les techniques classiques de filtrage du signal en amont de notre paramétrisation. Aussi, en vue d'une meilleure extraction et d'une meilleure identification des formes, nous avons définis des outils améliorant localement le rapport information/signal de notre représentation paramétrique et/ou de nos paramètres provisoires.

Des opérateurs de lissage permettent d'élimer les valeurs erratiques présentes sur nos paramètres de travail. Un gestionnaire d'image intégrant les techniques classiques de traitements d'images (seuillage, filtrage, recherche de contour, etc [7][8][9][10]) nous permet, en considérant notre représentation temps-fréquence comme une image (temps, fréquence, énergie spectrale), de concevoir une chaîne de traitements en situation de reconnaissance ou une série d'essais de traitements en situation d'assistance à l'opérateur.

Outils pour le traitement des connaissances
Exprimée sous une forme déclarative, les connaissances de notre système sont regroupées d'une part dans un ensemble fixe de règles, et d'autre part dans un treillis de résultats évoluant au cours de l'analyse. Une définition naturelle et concise de cette connaissance est rendue possible grâce à un ensemble de prédicats permettant une écriture symbolique, non seulement de l'information, mais aussi des opérations arithmétiques ou logiques qui lui sont appliquées. Enfin, un contrôle fin du processus d'analyse et la définition de stratégies de reconnaissance propres aux bruts sous-marins sont facilités par un ensemble d'outils spécialisés dans la gestion du choix de l'efficacité des prédicats de recherche.

Le traitement de la cavitatio d'hélice
La cavitatio d'hélice est particulièrement génante dans la mesure où elle s'accompagne généralement d'une érosion de l'hélice. Ne pouvant plonger la tête sous l'eau pour examiner le caractère destructif de ce phénomène, nous allons tenter de l'étudier au travers du bruit qu'il rayonne.

Suite à nos travaux, il apparaît que le bruit de cavitatio d'hélice se présente comme une succession de transitoires de tailles diverses (de 300 Hz à 3500 Hz) et de formes multiples (virgule, apostrophe). En outre, non seulement le temps séparant deux transitoires est constant, mais en plus il est fonction de la vitesse de rotation de l'hélice.

La localisation spatiale des transitoires repose sur cette périodicité du phénomène. Tout d'abord, grâce au prédicat

```prolog
decouper_spectrogramme_BDF([z]) ->
tous les (enfermer intervalle([25f]), periodicite([esm(z), vort_min, vort_max]), creer_classe([z, BDF, periodicite()]));
```
nous décompons le spectrogramme en une série de bandes de fréquences dont on détermine la périodicité la plus probable (ces bandes de fréquences sont stockées en mémoire sous forme d'un treillis de classes). Nous déterminons alors l'étendue spectrale des transitoires en associant entre elles toutes les bandes de fréquences adjacentes dont les périodicités sont voisines ou multiples.

La localisation temporelle des transitoires de cavitatio s'effectue par une stratégie de recherche de plateaux de cavitatio sur les paramètres temporels esup5, esup10, esup15 et esup20 définis sur la bande de fréquences correspondant à l'étendue spectrale des transitoires de cavitatio.

```prolog
plateau_cavitation([p, ar]) ->
plateau([p, ar, 10]),
esup20([z]) ->
max(esm(moyenne([5]) + 20)];
```

B10-10
A cause du bruit ambiant, la caractérisation des transitoires de cavitation ne peut se faire directement à partir de l'étude de paramètres temporels ou fréquentiels. Aussi, connaissant les localisations fréquentielle et temporelle des transitoires, nous créons pour chaque transitoire une image le contenant. Puis, sur chacune de ces images, nous appliquons un opérateur graphique de type Nagao spécialisé dans la détection et le renforcement des segments de droites.

Le contraste entre le transitoire et le reste de l'image ayant été augmenté, nous évaluons alors localement le paramètre fréquence-énergie maximum. La forme du transitoire est déterminée grâce à la position du centre de gravité de la colline fréquentielle présente localement, tandis que la pente du transitoire est évaluée par une régression linéaire sur les valeurs de ce paramètre.

Conclusion

Nous avons réalisé un environnement sous PROLOG facilitant l'acquisition, la représentation et le traitement des connaissances nécessaires à la reconnaissance automatique de bruits sous-marins. Cet environnement utilise des techniques appartenant à des domaines scientifiques aussi différents que le traitement du signal, la reconnaissance de formes, le traitement d'images et l'intelligence artificielle. En outre, la souplesse des outils, le formalisme et les concepts de programmation introduits autorisent une interaction naturelle des connaissances sans discontinuités depuis le signal sonore jusqu'aux niveaux symboliques les plus éloignés du phénomène physique.

Sur des bruits enregistrés dans des conditions réelles de navigation, cet environnement nous a permis d'étudier des émissions sonars (100% de succès dans l'extraction et la caractérisation de ce phénomène[11]), des bruits mécaniques de rotation, certains bruits hydrodynamiques, et de mettre en évidence une nouvelle composante du bruit de cavitation d'hélice. En situation de reconnaissance, plusieurs heuristiques ont été mises en œuvre pour la caractérisation automatique de ces bruits sous-marins.

Bibliographie


Résultats et commentaire

Les résultats ont été obtenus sur une base de sons constituée de seize bruits enregistrés durant 20 à 40 s dans des conditions réelles de navigation (présence d'un très fort bruit ambiant) et classés par l'expert suivant un critère d'écoute "facile", "moyen" et "difficile".

Le pourcentage de réussite pour la localisation temporelle et fréquentielle des transitoires de cavitation, ainsi que pour la détermination du nombre de tours de l'hélice avaroisi les 90%. Les causes principales des erreurs rencontrées sont d'une part la superposition dans le signal de plusieurs bruits sous-marins de nature périodique, et d'autre part l'extraction par le système d'un
INTRODUCTION

The application of the complex autocorrelation technique in the measurement of ADCP

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INTRODUCTION

The estimation of mean Doppler shift is one of the main tasks of the Acoustic Doppler Current Profiler (ADCP). Many methods have been developed for this estimation, but the complex autocorrelation technique is the one used most currently.

In the measurement of current ranging in large scale depth of the ocean, using this technique we only need to deal with less signal sampling and mathematical operations, and better accuracy can also be obtained even though under the condition of lower signal to noise ratio.

METHOD OF THE COMPLEX AUTOCORRELATION TECHNIQUE

Scattered echo signal \( X(t) \) received by the ADCP with fluctuating amplitude and phase can be written as a complex signal

\[
X(t) = I(t) + jQ(t),
\]

where \( I(t) \) is the in-phase component of the signal and \( Q(t) \) is the quadrature component.

According to the references [1][2], it is demonstrated that the mean Doppler shift of the return echo is

\[
\hat{f}_d = \frac{1}{2\pi} \text{Arg} [R(\tau)],
\]

where \( R(\tau) \) is the autocovariance function evaluated at time \( \tau \). The average estimated of \( R(\tau) \) at lag \( \tau \) is

\[
R(\tau) = \frac{1}{N} \sum_{m=1}^{N} X_m \overline{X}_{m-\tau},
\]

The asterisk stands for the complex conjugate. \( N \) is the number of data pairs. The estimate of the mean Doppler shift of the echo is given by

\[
\hat{f}_d = \frac{1}{2\pi} \text{Arctan} \left( \frac{I_m Q_{m+1} - I_{m+1} Q_m}{I_m I_{m+1} + Q_m Q_{m+1}} \right).
\]

On the basis of this method, the quadrature and in-phase signals are sampled and digitalized, then passed to the micro-processor which calculates \( \hat{f}_d \). Figure 1 is the block diagram of the electronic hardware.

FIELD TEST RESULTS

A field test was performed in the sea near Qingdao of China in August, 1990. The transducer of the ADCP was put at 2.7m depth under the sea surface by a rod fixed on one side of a moored ship. At the same time a Direct Read Current Meter (DRCM) was hung at 10m depth in the field in order to make a comparison between the ADCP and the DRCM. The water depth is about 35m. Only the currents of two bins were measured by the ADCP, that is, of 10m and 18m bins.

The test lasted for 9 hours and a set of current data was recorded in every 15 minutes. Obviously, there are 39 sets of data obtained. There are two curves in Fig.2 to present the currents of these two bins respectively. It is clear that the two currents are different, but with some trend, and are in accordance with current distribution of this field.

In Fig.3 there is the contrast results of the ADCP and DRCM at the 10m depth.

![Fig.1](image1)

![Fig.2](image2)

![Fig.3](image3)
By comparison of the two curves from the ADCP and DRCM it is obvious that they fit each other very well in the measurement period, except for one set of data which is appeared at time 02:45. By the method of regression analysis, we derive that the correlation coefficient of the curves in 0.998, and the rms difference between these two data group is 3cm/sec. These results are similary to the results of Pettigrew and Irish[3].

CONCLUSION

According to explanation above it may be considered that the complex autocorrelation technique used in the ADCP has notable benefits in the measurement of current profile, such as less arithmetic and better accuracy. It is believed that this technique will be applied more widely.

REFERENCES

APPRECIATIVE RELATIONSHIP OF RESONANT FREQUENCIES OF SIMILAR SHELLS SUBMERGED IN WATER

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INTRODUCTION

This paper deals with the prediction of the resonant frequencies of thin shells submerged in water. In cases of large shells, vibration analysis have to turn to the use of models. The accepted way of doing this is to scale shell and model faithfully in every respect, shape, thickness and all, and to employ the classical scaling law. But the scaling law, however, sometimes proves to be too restrictive, since it is not always convenient, for instance, to scale the thickness of the shell in proportion with typical surface dimensions. Thereby, there remains the problem of thickness distortion between the models and the prototype. How to predict the behavior of the shell from the model's has been put forward in engineering [1].

In this paper, by the natural mode expansion method and a few of reasonable assumptions, an approximate relationship was derived for predicting the resonant frequencies of similar thin shells with dissimilar thickness submerged in water. As an example, experiment was made for cylindrical shells with one end fixed and another free. The results show that the formula discussed in this paper is feasible and effective.

THEORETICAL DERIVATION

It is well known that the relationship between two completely similar shells is very simple with the use of the scaling law. So, in the following derivation, we need only to seek the relationship of resonant frequencies between two shells with same surface dimensions and boundary conditions but with different thickness.

Equations of Motion

The equations of motion of an uniform thin shell can be written as the following general forms [2]:

\[ L_i \left( u_i, u_j, u_k \right) = \rho h \ddot{u}_i - q_i \]  \hspace{1cm} (1)

\[ L_2 \left( u_i, u_j, u_k \right) = \rho h \ddot{u}_i - q_2 \]  \hspace{1cm} (2)

\[ L_3 \left( u_i, u_j, u_k \right) = \rho h \ddot{u}_i - q_3 \]  \hspace{1cm} (3)

where \( u_i, i=1,2,3 \), are displacements in the curvilinear orthogonal coordinates on the neutral surface of the shell where 3 indicates the direction perpendicular to the surface; \( q_i, i=1,2,3 \), are three forcing terms; \( L_i \), \( i=1,2,3 \), are linear differential operators; \( \rho \) is the shell mass density and \( h \) is the thickness.

Natural Mode Expansion

According to the orthogonality property of natural modes, the general solution of Eq.(1)-(3) can be expanded in an infinite series as:

\[ u_i(x,y,z) = \sum_{k=1}^{\infty} \eta_k U_k(x,y) \]  \hspace{1cm} (4)

where \( \eta_k(t) \) are unknown factors needed to be determined, and \( U_k \) are the natural mode components in three principal directions when the shell is vibrating with mode \( k \) in air.

Through some substitutions and by the using of the properties of natural modes, one can at last get:

\[ \eta_k + \omega_k^2 \eta_k = -F_k \]  \hspace{1cm} (5)

where \( \omega_k \) is the natural frequency of mode \( k \) and

\[ F_k = \frac{1}{\rho h N_k} \int \left( q_1 U_1 + q_2 U_2 + q_3 U_3 \right) dS \]  \hspace{1cm} (6)

\[ N_k = \int \left( U_1^2 + U_2^2 + U_3^2 \right) dS \]  \hspace{1cm} (7)

where the integration is throughout the surface of the shell.

Since the shell is submerged in water, the resonant vibration will cause the variances in pressure in water. So the forcing terms on the surface are

\[ q_1 = q_2 = q_3 = 0 \]  \hspace{1cm} (8)

\[ q_2 = F_0 + F_1 = F_2 \]  \hspace{1cm} (9)

where \( p_0 \) and \( p_2 \) are pressure distributions on each side of the shell.

If the shell is excited vertically, only modes that are vertical are dominant. And then

\[ F_2 = \frac{1}{\rho h N_k} \int \left( p_0 U_2 \right) dS \]  \hspace{1cm} (10)

Assumption and Predicting Formula

In Eq.(5), since \( p \) is caused by the vibration of the shell and it also affects the vibration in turn, the whole problem comes to be very complex. Some reasonable assumptions should be adopted to get a simple and feasible formula of prediction:

1. \( \text{water be considered to be uncompressible} \)
2. \( \text{the mode shapes be the same when a shell vibrates with the same mode in air and in water, because there would be little difference between them.} \)
3. \( \text{the mode shapes be also the same between similar shells with same surface dimensions but different thicknesses if their modes be the same.} \)

According to assumption 2, addition with the continuity of vertical velocity on the surface of the shell, the pressure on the shell could be written as:

\[ p = p_0 \hat{n} \]  \hspace{1cm} (11)

where \( p_0 \) do not vary with time and its distribution is determined only by mode shape.

Assuming that the resonant frequency is \( \Omega_k \) if the shell vibrates in water with mode \( k \), then Eq.(5) becomes

\[ \omega_k^2 = -1 - \frac{1}{\rho h} C \]  \hspace{1cm} (12)
where $C_r$ is only determined by vibration distribution on the shell and has no business with $U_{4r}$ and frequency.

Considering two similar shells with only different thicknesses, one is denoted by superscripts "1" and the other is not, and using assumption 3 gives the predicting formula as:

\[
\frac{\omega_1^2}{\omega_0^2} = \frac{m}{A^2} \frac{1}{1 - \frac{m}{A^2}}
\]  
(13)

which appears very simple.

**Discussion**

The exactness of the formula is mainly based on the three assumptions. According to assumption 1, the near-field pressure acts as an inertial force, which is true especially in low frequency range. Assumption 2 is often adopted in the calculations in structural dynamics [3]. It should be noticed that the thinner the shells are, the more different the mode shapes are. The same with assumption 3 if the distortion of thickness is too large. The assumptions define the application limits of the formula which should be examined by experiments.

**EXPERIMENTAL VERIFICATION**

A number of experiments were conducted on shells with different thicknesses and shapes to examine the exactness of the relationship derived above. Here, because of limited space, only the results of a group of cylindrical shells with one end fixed and another free are presented. The cylindrical shells were 305mm long and with diameters of 150mm. Their thicknesses varied from 1mm to 4mm.

Tests were conducted on these shells to measure the natural frequencies and mode shapes in air and in a large tank filled with water. Each shell was excited radially at a point near its free end. The experiment scheme is shown in fig 1.

The results are shown in Table 1.

**Table 1. Resonant Frequency and Mode**

<table>
<thead>
<tr>
<th>Thickness (mm)</th>
<th>1</th>
<th>2</th>
<th>3</th>
<th>4</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mode</td>
<td>$\omega_0^2$ (Hz)</td>
<td>262.0</td>
<td>290.0</td>
<td>323.0</td>
</tr>
<tr>
<td>m=1</td>
<td>$\omega_0^2$ (Hz)</td>
<td>262.0</td>
<td>290.0</td>
<td>323.0</td>
</tr>
<tr>
<td>n=2</td>
<td>$\omega_0^2$ (Hz)</td>
<td>115.0</td>
<td>162.1</td>
<td>214.0</td>
</tr>
<tr>
<td>Mode</td>
<td>$\omega_0^2$ (Hz)</td>
<td>383.0</td>
<td>660.5</td>
<td>936.0</td>
</tr>
<tr>
<td>m=1</td>
<td>$\omega_0^2$ (Hz)</td>
<td>146.0</td>
<td>342.5</td>
<td>525.1</td>
</tr>
</tbody>
</table>

In Table 1, m is the mode number in axial direction and n the mode number in circumferential direction. Mode "K" indicates the mode of "m".

Fig 2 and Fig 3 show two typical frequency response curves of a cylindrical shell, one in air and another in water. It can be seen from the figures that the frequency curve become smooth when the shell vibrate in water. The resonant frequencies with high modes can not be distinguished. That is the reason why only low modes are listed in Table 1.

To examine the exactness of Eq.(13), the resonant frequencies in water of the shell 1mm thick were calculated using the resonant frequencies of other shells with thicknesses of 2, 3 and 4mm. The predicting frequencies were 112.7Hz, 119.2Hz and 114.3Hz for m=1 and n=2 mode. For the mode of m=1 and n=3, the predicting frequencies were 150.9Hz and 139.6Hz. Comparing with the frequencies measured (115.0Hz for m=1 and n=2 mode, 146.0Hz for m=1 and n=3 mode), the maximum error was less than 5%.

**CONCLUSION**

In this paper, it was shown that one can easily predict the resonant frequencies of submerged shells from model experiments even when the scaling law was not applicable. The agreement between the predicted and measured resonant frequencies was very well, which verified the exactness of the predicting formula.

**REFERENCES**


**Fig 1. Diagram of Apparatus and Instrumentation**


**Fig 2. Frequency Response Curve in Air**

**Fig 3. Frequency Response Curve in Water**
THE WKBZ MODE APPROACH TO SOUND PROPAGATION IN HORIZONTALLY STRATIFIED OCEANS

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INTRODUCTION

Though the normal-mode approach can give an exact solution to the acoustic field in a stratified medium, the conventional mode approach is awkward to apply to higher frequency and broadband propagation due to complication of evaluating a great number of mode eigenvalues and eigenfunctions. So, some approximate methods were proposed. In this paper, a new mode approach based on the WKBZ approximation to sound propagation in ocean channels is proposed and numerically applied to sound propagation in range-independent channels.

WKBZ MODE THEORY

In order to overcome the shortcoming of singularity at turning depths of the conventional WKB approximation, the generalized phase-integral approximations (WKBZ approximations) were proposed. The eigenfunction \( \Psi_i(z) \) of the turning modes from the WKBZ approximation with two parameters \( E \) and \( D \) is given by

\[
\Psi_i(z) = \sqrt{\frac{2}{S_i}} \times \frac{\exp\left(-\int_0^z \sqrt{\frac{1}{\omega_i} - k_i^2(x)} \, dx\right)}{\sqrt{2 \left\{ b_i \frac{k_i^2}{E} - D \left[ k_i^2(x) - \nu_i^2 \right] + 16k_i^2(x) \right\}^{1/4}}} \frac{\sinh\left(\sqrt{\frac{1}{\omega_i} - k_i^2(x)} \, dy + \pi / 4\right)}{\sqrt{2 \left\{ b_i \frac{k_i^2}{D} - b_i \frac{k_i^2}{E} \left[ k_i^2(x) - \nu_i^2 \right] + 16k_i^2(x) \right\}^{1/4}}} \frac{(-1)^i \exp\left(-\int_0^z \sqrt{\frac{1}{\omega_i} - k_i^2(x)} \, dx\right)}{\sqrt{2 \left\{ b_i \frac{k_i^2}{E} - D \left[ k_i^2(x) - \nu_i^2 \right] + 16k_i^2(x) \right\}^{1/4}}} \left(1\right)
\]

where \( E = 2.152 \) and \( D = 1.619, \nu_i \) is the eigenvalue, \( k(x) \) is the distribution of the mode, \( \eta_i \) and \( \zeta_i \) are respectively the turning depths above and below the channel axis.

To illustrate the accuracy of the WKBZ approximation, we discuss the model of a bilinear channel. Fig.1(a) is the velocity profile, Fig.1(b) the integrated eigenfunction for \( f = 10 \), Fig.1(c) the enlarged part of the eigenfunction near the turning depth \( \eta_i \). It can be seen from Fig.1 that the WKB approximation diverges at \( \eta_i \), while the WKBZ approximation has certain accuracy everywhere.

When the upper turning-depth is near the surface, it is necessary to take account of the surface phase-shift correction to Eqs.(1). The surface phase-shift \( \varphi_s \) can be expressed as

\[
\varphi_s = \begin{cases} \pi + 2\arctan \left( \frac{v(x)}{w(x)} \right) & \text{for } \beta > 0 \\ \pi \left( \frac{2}{2 + 2\arctan \left( \frac{v(x)}{w(x)} \right) - 2\omega_0} \right) & \text{for } \beta < 0, \end{cases}
\]

where \( t_0 = \left( \nu_i^2 - k_i^2(x) \right)/b(x) \), \( b(0) = \left| \frac{dk^2(z)}{dz} \right|_{z=0} \), \( \omega_0 = \frac{2}{3} \left| t_0 \right|^{1/2} \), \( u(0) \) and \( v(0) \) are Airy functions. The phase-shift \( \varphi_s \) varies fast with \( t_0 \) when \( -1 < t_0 < 1 \), and it approaches to \( \pi / 2 \) and \( \pi \) when \( t_0 > 1 \) and \( t_0 < -1 \), respectively. Incidentally note that the surface phase-shift correction is important for evaluating the eigenvalues and the WKBZ eigenfunctions of the near surface turning modes.

APPLICATION TO THE NORTH PACIFIC

For illustrative purposes, the WKBZ mode approach is numerically applied to sound propagation in the North Pacific. The sound velocity profile is shown in Fig.2.

For a frequency of 100 Hz, a source depth of 25 m, and a receiver depth of 200 m, the curves of transmission loss versus range computed by WKBZ approximation and NM (convolutional mode approach) are shown in Fig.3, whereas WKBZ-NM denotes the difference of the transmission loss less than 100 dB between WKBZ and NM. It can be seen from the figure that the convergence-zone structures computed by the two mode approaches are coincident very well.

At last, the transmission loss versus range and depth calculated by the WKBZ mode is shown in Fig.4, where the range is up to 100 km and the vertical depth is up to 4000 m. This figure shows the interesting acoustic field distribution in the three-dimension space.

REFERENCES

Fig. 1 The WKBZ approximation versus the exact solution and WKB approximation

Fig. 2 A North Pacific velocity profile (Ref. 8)

Fig. 3 Transmission loss versus range for a source depth of 25m, a receiver depth of 200m, and a frequency of 100Hz.

Fig. 4 Three-dimension acoustic field distribution for a source depth of 200m and a frequency of 100m.
COMPARISON OF DETECTION PERFORMANCE OF LINE-COMPONENT SPECTRUM AND DEMODULATED LINE-COMPONENT SPECTRUM IN BACKGROUND-NOISE

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INTRODUCTION

The research work on radiated-noise of vessels for recognition purpose have been done for many years. Except for extracting the propeller features from demodulation spectrum, the propeller features can be also obtained directly from line components of the spectrum of radiated-noise in low frequency range.

ANALYSIS OF REAL DATA

Fig.1(a) shows the spectrum of radiated noise of ship A. The maximum frequency of the spectrum is F_max, which is scores of Hz. The averaging number is 8. We have known the nominal rotation speed of propeller shaft is M_1, rev. per minute, M_1 = 60F_1. The frequency of first line component is F_1 Hz which corresponds just with the nominal rotation speed M_1. The level of line component of F_1 Hz is quite strong, which is more than 20 dB over the continuous spectrum. Their harmonic components are plenty. Fig.1(b) illustrates the demodulation spectrum of ship A, whose line components corresponding with propeller are weak.

DETECTION OF DEMODULATED LINE COMPONENTS

As modeling the periodically locally stationary random process, the radiated noise can be given in following expression:

\[ S(t) = [1 + m(t)]x(t) \]

(1)

where \( x(t) \) is the narrow band stationary Gaussian white random process, \( m(t) \) is called the modulation function which varies at slowly periodical rate, and its frequency is much lower than that of \( x(t) \). \( S(t) \) is the periodically locally stationary Gaussian process, \( m(t) \) is the periodical function. That is, \( m(t) = m_0 \cos 2\pi F_p t \), where \( m_0 \) represents the peak of the modulation depth, \( \sigma^2_m \) is the variance of \( m(t) \). Assume the interfering noise \( n(t) \) to be the white Gaussian noise with the variance \( \sigma^2_n \), it may be verified that the locally stationary Gaussian process \( S(t) = [1 + m(t)]x(t) \) added to the stationary Gaussian noise \( n(t) \) makes another locally stationary Gaussian process \( Y(t) \), namely

\[ Y(t) = [1 + am(t)][R(t)] \]

(2)

where \( a \) in the expression is the factor of the proportion without dimension, that is

\[ a = \sigma^2_x / (\sigma^2_m + \sigma^2_n) \]

and \( R(t) \) is narrow band stationary white Gaussian process, its variance is given by

\[ \sigma^2_R = \sigma^2_m + \sigma^2_n \]

(4)

The demodulation power spectrum estimation of radiated noise in background-noise can be expressed in Fig.2. The intensity of the power spectrum density estimation at 0Hz is given by

\[ S(0) = \frac{2}{\pi} (\sigma^2_m + \sigma^2_n) \]

(5)

the mean of the continue spectrum is

\[ \mu_e = \frac{B}{2B} \frac{1}{\pi} (\sigma^2_x \text{ and } \sigma^2_m) \]

(6)

the standard deviation follows that

\[ \sigma_e = \mu_e \sqrt{\frac{1}{2B}} \]

(7)

The intensity of line component is higher than that of the continue spectrum, the intensity difference of them is equal to

\[ S(F_n) = \frac{2\sigma^2_m}{\pi} \frac{2}{\pi} (\sigma^2_x \text{ and } \sigma^2_m) \]

(8)

In order to illustrate the detection performance, the probability density distribution functions with the line component (hypothesis H_1) and without line components (hypothesis H_0) are shown in Fig.3. The horizontal axis \( x_1 \) corresponds to the situation of \( S(0) \) normalized. For convenience, \( x_1 \) axis is normalized with \( \sigma^2_0 \), further, and get \( x_2 \) axis. \( \frac{\sigma^2_m B}{2B} \) is defined to be the "signal level", it is the level of line component over the continue spectrum. The higher it is, the better detection performance is. \( \sigma^2_0 \) as 1/\( \sqrt{B_1} \) indicates the normalized fluctuation of the continue spectrum estimation, here it plays a role of background-noise for detection of the line component, it causes the probability of false alarm. When the interfering noise \( \sigma^2_m \) get higher or signals \( \sigma^2_x \) get lower (the target is far away from the reciver), then \( x \) tends small, the signal level descends and the detection performance gets worse with \( x \) getting small. In terms of the signal level \( \sigma^2_m B / 2B \), it is seen that \( \sigma^2_0 \) The larger the modulation depth \( m_0 \), the more easily the detection is achieved, as \( m_0 \) is small, it will be little value of application, consequently the features extraction of the propeller's speed bases on that how large enough the modulation depth \( m_0 \) is. As mentioned above, \( x(t) \) and \( n(t) \) are assumed to be stationary Gaussian signals, the line level and \( B / 2B \) is the bandwidth of the band-pass filter in pre-processor; \( B \) is the frequency resolution of the post-processor vary in direct proportion, therefore, \( B / 2B \) may play an important role, for instance, when \( B_2 \) and \( B \) are 0.125Hz. 125Hz respectively, \( B / 2B \) reaches 1000, it means that the line level can be enlarged to be 1000 times.

DETECTION OF DIRECT LOW FREQUENCY LINE
COMPONENT

The detection performance of direct line component in background-noise concerns the problem of detecting sinusoidal signal from the stationary Gaussian noise ROC (Receiver Operating Characteristic) can represent the performance. Suppose sinusoidal signal power is $S$, then the mean of probability density distribution with line component (hypothesis $H_2$) is $\sigma^2_{01}S$, $\sigma^2_{02}$ is the variance of background-noise, its normalized deviation is $1 / \sqrt{B_1T_1}$, where $B_2$ is the bandwidth, $T_2$ is the integrare time of low-pass filter. For the case of direct low frequency line component the probability density distributions under hypotheses $H_2$ and $H_1$ are also shown in Fig.4, but it corresponds to $x_3$ axis. $x_3$ axis can be normalized with $\sigma^2_{02}$ to be $x_4$ axis. We can see, curves of probability density distribution are basically same (Suppose to take same value of $B_2$ and $T_2$). For convenience, the difference of the mean of probability density distribution in frequency $F_0$ under hypotheses $H_1$ and $H_0$ is defined as "signal level" $H$. The signal level of demodulated line component is $a^2m^2B / B_2$ and the signal level of direct line component is $S / \sigma^2_{02}$.

COMPARISON OF DETECTION PERFORMANCE BETWEEN THE DIRECT WAY AND THE DEMODULATION WAY

From Fig.3. We can know that the signal level is getting low as signal-to-noise ratio goes down, but they decline at different rate. Fig.4 shows the relationship between the signal level and signal-to-noise ratio. The curve a in Fig.4 indicates relationship in direct way, the curve b shows the relationship between the signal level and signal-to-noise ratio in demodulation way while $S / N = 30dB$ and $m^2B / B_2 = 30dB$. Corresponding with right vertical axis $a^2$ the curve b shows the relationship between the proportional factor $a^2$and signal-to-noise ratio. Table 1 shows the relationship between $a^2$and signal-to-noise ratio $\beta$. In the case of the direct way the signal level is proportional to signal-to-noise ratio. In the case of demodulation way, if signal-to-noise is high, the signal level drops slowly with $S / N$ decrease; if $S / N$ is low, it drop rapidly with $S / N$ decrease. When signal-to-noise ratio is high, i.e. $\sigma^2 / \sigma^2_{02} > 1$, and $a$ is approximately equal to 1, so the signal level $a^2m^2B / B_2$ is getting to a constan. Refer to the right part of the curve b which does not vary with $S / N$. When signal-to-noise is low, $a$ is close to $\sigma^2 / \sigma^2_{02}$, the signal level is proportional to the square of $S / N$ ratio, refer to the left part of the curve b which is a straight line, if signal-to-noise drops by 10dB, the signal level will descent by 20dB.

CONCLUSION

The analysis of the experiment at sea shows that there are small strong superposed line components in direct spectrum of radiated-noise occurring at discrete frequencies which correspond with the rotation speed of propeller shaft, or propeller blade frequency, or their harmonics frequency. Behaviour of these line components depends on the propeller. So the line components can provide the features of target recognition. It will be attractive to use these line components for the towed-line array sonar or the seashore sonar, etc. It will be also more valuable for long range detection and recognition due to its small transmission loss. Hence the further research work on these line components in low frequency range will be necessary.

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PERFORMANCE ANALYSIS OF A PARTIALLY ADAPTIVE SIDELOBE CANCELLER APPLIED TO A TOWED LINE ARRAY

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FRACTIONAL DISPERSION AND ITS COMPENSATION

Performance of adaptive sidelobe cancellers in suppressing strong interferences is related to the array aperture, the incident direction and the bandwidth. A parameter called "fractional dispersion" can be used to describe these factors[1,2]. For a line array of equally spaced elements, the fractional dispersion caused by an interference may be written as

\[
FD = \Delta \tau \cdot BW = \frac{1}{2\pi} \left[ \frac{(N-1)d\sin \theta}{C} \right] \cdot 2\pi \cdot BW,
\]

(1)

where \(\Delta \tau\) denotes the delay time for an interference to run across the whole array, \(BW\) is the interference bandwidth, \(N\) is the number of elements, \(d\) is the distance between any two neighbouring elements which is generally taken as a half wavelength \((\lambda/2)\) of the center frequency \((f_0)\), and \(\theta\) is the angle of arrival. The above equation reveals that FD describes the amount of phase shifting over the whole array for a broadband interference. The bigger FD means the need for more phase compensation in the adaptive sidelobe canceller. The commonly-used device for compensating the phase shifting in time domain is tapped delay line (TDL). If its unit delay time \(T\) is taken as \(1/BW\) [3], the number of taps in the TDL

\[
N_t \geq \Delta \tau_{\text{max}} / T = FD_{\text{max}}.
\]

(2)

Equation (2) shows that the value of the fractional dispersion can be used for determining the least length of a TDL.

PERFORMANCE DESCRIPTION FOR A PARTIALLY ADAPTIVE SIDELOBE CANCELLER

Fig.1 shows a typical partially adaptive sidelobe canceller (PASC) applied to a line array of \(N\) equal spaced elements. We assume the system to be operating on an algorithm of the Howells–Applebaum type. In other words, the desired signal on the main beam is absent and the adaptive weights are adjusted for minimum array output which includes only interference and noise. A cosine squared on a pedestal window is introduced into the forming of the main beam.

Let \(V\) be the column vector of signals and \(W\) be the row vector of complex weights.

\[
V = (V_0, V_1, \ldots, V_M)\end{aligned}\]

\[
W = (W_0, W_1, \ldots, W_M),
\]

(3)

(4)

where \(W_0 = -1\) is needed for holding the main beam signal to be constant. The system output power is

\[
P = W \cdot \Sigma V^* \cdot W^*.
\]

(5)

Let the matrix \(A\) be \(\Sigma V\) which is the signal covariance matrix. W.D White has proved that the minimum residual power \(P_0\) is equal to the reciprocal of the up-left element (denoted by \(A_{\infty}\)) in the matrix \(B = A^{-1}\). We may define the cancellation ratio as

\[
CR = \frac{E[V_0^2V^*_0]}{P_0} = A_{\infty} B_{\infty},
\]

(6)

where \(A_{\infty}\) is the up-left element of the matrix \(A\) and represents the main beam power.

RESULTS OF NUMERICAL CALCULATION

Numerical calculation gives the dependence between CR and FD with a set of parameters such as the number of elements \((N)\), the angle of arrival \((\theta)\) and the relative bandwidth \((FBW = BW / f_0)\). In the computation, we let \(N = 32\) and \(d = \lambda/2 = 1.5\) meters. The selection of auxiliary channels is changeable, whereas the conventional auxiliary channels composed of element 2,6,11,18,26,29 are often selected. 45 dB of the interference/noise ratio in each receiving channel is also assumed.

Performance with no TDL in each auxiliary channel

Fig.2 illustrates that the residual power \(1/B_{\infty}\) grows steadily and the cancellation effect becomes poor with the increasing of the relative bandwidth and the angle of arrival. \(A_{\infty}\) decreases with the increasing of \(\theta\) because of the window on the main beam. These curves in the figure show that the big fractional dispersion strongly influences the performance of the adaptive system.

Performance improvement with TDL in each auxiliary channel

Fig.3 illustrates the reduction of residual power output of the PASC with TDL which length has met the requirement of the equation (2), for a set of 4 bandwidths of the interference. When \(FBW = 60\%\), which is approximate to one octave, CR is larger than 17dB even if the interference arrives from the endfire direction.
Influence of the length of TDL

Fig. 4 gives some calculated residual powers for different lengths of TDL. The number of taps $N_t$ is between 2 and $FD_{\text{max}}$. Computation results indicate that few variation occurs in the cancellation performance for the case of the small angle of arrival and the narrow bandwidth, because the PASC has enough degrees of freedom which come from all auxiliary channels while FD is relatively small.

Influence of the location of auxiliary channels

Fig. 5 shows different calculated residual powers which are corresponding to three kinds of locations listed in the below table in which "" denotes the used auxiliary element and "number" denotes element numbers between two used auxiliary elements. From Fig. 5 we know that "two ends cluster" type location is the best to suppress interferences. It gets extra 10 db cancellation ratio by comparison with "uniform" type when $FBW=40\%$ and $\theta=90^\circ$.

<table>
<thead>
<tr>
<th>No</th>
<th>type</th>
<th>locations</th>
</tr>
</thead>
<tbody>
<tr>
<td>I</td>
<td>central density</td>
<td>$10 \times 3 \times 2 \times 1 \times 2 \times 3 \times 10$</td>
</tr>
<tr>
<td>II</td>
<td>uniform</td>
<td>$3 \times 5 \times 5 \times 5 \times 5 \times 5 \times 3$</td>
</tr>
<tr>
<td>III</td>
<td>two ends cluster</td>
<td>$1 \times 3 \times 6 \times 7 \times 8 \times 3 \times 3$</td>
</tr>
</tbody>
</table>

CONCLUSION

Based on the fractional dispersion and the adaptive cancellation ratio, the wideband interference cancellation performance of a PASC with TDL has been examined. Numerical results show that a typical PASC can get extra 30 db of CR than a conventional beamformer for the interference with approximately one octave bandwidth near the endfire direction. The location of the auxiliary channels has more influence on the cancellation performance than the length of TDL in all channels.

REFERENCES

A FAST ADAPTIVE FILTER TRACKING ALGORITHM WITH NORMALIZED VARIABLE STEPSIZE

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I. NVS ALGORITHM AND ITS PERFORMANCE

As we have known that the iterative formula of weight coefficients for LMS algorithm is [1][2]

\[ W_{n+1} = W_n + 2\mu \epsilon_n X_n \]  

(1)

where \( \mu \) is the stepsize, \( \epsilon_n \) the error, \( X_n \) input vector at time \( n \). Theoretical analysis has indicated that when \( \mu \) is fixed, the global convergence speed of the output of ADF mainly depends on the smallest eigenvalue \( \lambda_{\min} \) of the autocorrelation matrix \( R \) of input sequences and the global missadjustment mainly depends on the largest eigenvalue \( \lambda_{\max} \). When matrix \( R \) has a large eigenvalue spread, the convergence speed of the output will be slow and the missadjustment will increase. Besides, the fact that eigenvalues of \( R \) are changed as the intensity of the input signal changes will also influence convergence speed and missadjustment, and even destroy the convergence condition. However, LMS algorithm has advantages of less computational cost, stabilization and easy realization, so it has got wide applications. When one improves LMS algorithm, the advantage should be maintained and the above shortcoming must be overcome. The iterative rule of weight vector of the normalized variable stepsize (NVS) adaptive filter algorithm is

\[ W_{n+1} = W_n + 2\mu \left( \frac{\| \epsilon_n \|^2}{\hat{P}} \right) \epsilon_n X_n \]  

(2)

where \( \hat{P} \) is the estimate of the input power.

\[ \hat{P} = \hat{P}_{\text{init}} + \theta \left( \epsilon_n^2 - P_{\text{avg}} \right), \quad 0 < \theta < 1 \]  

(3)

and

\[ \| \epsilon \| = \begin{cases} C, & \text{when } |\epsilon| > C \text{with } C > 1 \\ |\epsilon|, & \text{when } |\epsilon| < C \end{cases} \]  

(4)

Here the action of \( \hat{P} \) is to make the increment of the weight vector relatively steady.

Let \( \mu = \mu_j \), it can be seen that \( \mu_j \) can react to the error quickly. Because \( |\epsilon| \) is a random variable, it may cause ADF to be deconvergent, so limiting \( |\epsilon| \) is needed.

Following the derivation of LMS algorithm, it is easy to show that the convergence condition of NVS algorithm is

\[ 0 < \mu C < \frac{1}{\left( \lambda_{\min} / \hat{P} \right)} \]  

(5)

The convergence time constant of the 1-th weight and the mean square error are

\[ \tau_{11} = T_s / \left( 2\mu \| \epsilon_n \|^2 / \hat{P} \right) \]  

(6)

and

\[ \tau_{\text{MSE}} = T_s / \left( 4\mu \| \epsilon_n \|^2 / \hat{P} \right) \]  

(7)

respectively, and where \( T_s \) is sampling frequency. And the static missadjustment is

\[ M < \frac{(\mu C / \hat{P})}{T_s} \]  

(8)

Only three extra multiplications are needed compared with LMS algorithm.

II. COMPUTER SIMULATIONS

To compare the tracking ability of NVS algorithm with LMS algorithm, a system identifier is constructed as in Fig.1. Here the weights of FIR filter is variable. The input to the system is a white noise with zero mean value and the variance of 1. For convenience to compare, let the number of weights of both ADF and FIR system be \( L = 8 \). The sampling frequency is \( f = 1000 \text{Hz} \). The tracking performances of two algorithms for the following situations have been calculated. In the first one, the weight \( W_{LS} \) of FIR system is changed as sinusoidal function with a frequency of 100Hz. The result is shown as in Fig.2. It can be seen that the tracking of weights for LMS has an obvious delay, and the tracking performance for weights will be improved as \( \mu \) increases, but LMS algorithm is still not so good as NVS algorithm. Next, let the weight \( W_{LS} \) of FIR system be a step function, and the result is shown as in Fig.3, where the mean square error in Fig.3(b) is obtained from the average of 50 samples, that is

\[ J = \frac{\sum_{i=1}^{50} e_i^2(i)}{50} \]  

(9)

Several other situations have also been investigated. The results are briefly described as follows:

(a). The weights \( W_{LS} \) of FIR system are divided into 2 and 4 groups and the weights of each group are changed as different rules. Computer simulations indicate that the weights of ADF are tracking accordingly well, but since there is some coupling among weights of various groups, the tracking error will increase.

(b). Let the numbers of weights of ADF and FIR system are not the same. The results indicate that when the number of ADF is more than that of FIR system, the tracking error does not increase obviously. However, when the number of ADF is less than that of FIR system, the tracking error will increase clearly.

(c). In NVS algorithm, when increasing \( \mu \) and \( C \), the tracking ability of ADF will be improved further, but the tracking noise will increase and hence \( \mu \) and \( C \) should be selected carefully.

(d). Approximate Newton algorithm[3] has been investigated in this paper. By replacing KLT with DCT, take the main diagonal elements of the eigenvalue matrix of \( R \) as an approximate eigenvalue estimate of \( R \). To a time variable system, the calculation for DCT is needed in each iteration, which will cause the computational cost to be increased obviously, and hence it has not been adopted in this paper.
The smoothing gradient algorithm has been studied as well. To a non-time-variable system, smoothing gradient will reduce the static missadjustment. However in time variable systems, since the gradient is time variable, it will simply introduce the bias for the gradient estimate and cause the tracking error of weights to be increased. The method has not been employed in this paper.

Both theoretical analysis and computer simulation have indicated that the tracking ability of the new NVS algorithm is advanced over LMS algorithm in all cases.

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Fig. 1: Adaptive System Identifier

Fig. 2: Weight Tracking When \( w_t \) Changes as Sinusoidal Function
(a) \( \mu = 0.001 \), \( c = 3 \), \( \theta = 0.01 \)
(b) \( \mu = 0.005 \), \( c = 5 \), \( \theta = 0.01 \)

Fig. 3: Weight Tracking and Mean Square Error of the Output
--- The Parameters of ADF are
\( \mu = 0.001 \), \( c = 5 \), \( \theta = 0.01 \)
(a) Weight Tracking
(b) Mean Square Error of the Output
A NEW SLIDING CZT WITHIN–PULSE BEAMSCANNER

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I. INTRODUCTION

The within–pulse beamscanning technique had been suggested and realized in 1960's[1][2]. But for the reason about electronic components, this kind of system was more complicated and did not obtain wide applications. Therefore, the single beam sonar with mechanical scanning has been used as short range high resolution sonar in underwater vehicles for a long time. Since the data rate of mechanical scanning sonar is not high enough to meet the requirement by the developing moving vehicles in nowadays, the within–pulse beamscanning technique has again become an attractive scheme[3][4]. A within–pulse beamscanning method using sliding chirp–Z transform (CZT) is presented in this paper, in which the multi–channel processing with a large dynamic range is simplified to a single channel processing just after pre–amplifiers, and the weighting for array elements together with the time delay compensation required for the near field focus is also performed in the channel, so that the system is very compact. An experiment system has been built in laboratory and some experimental results have obtained in a test tank.

II. THE PRINCIPLE OF SLIDING CZT BEAMSCANNER

The r–th beam output for a uniform spacing linear array with N elements can be represented as DFT of spatial sampled \( X_k \)

\[
B_r = \left| \sum_{k=0}^{N-1} X_k e^{-j \frac{2\pi}{N} rk} \right|, \quad r = 0,1,\ldots,N-1
\]  

(1)

where

\[
X_k = A_k e^{-j \phi_k} \delta_{k, \text{beam}}
\]  

(2)

is the complex demodulated signal from the k–th element, and \( \phi \) is the direction of incoming plane wave, \( A_k \) the amplitude of the k–th signal, \( d \) the spacing of elements, \( \lambda \) the sound wavelength in water. Eq.(1) can be written as the known chirp–Z transform:

\[
B_r = \left| \sum_{k=0}^{N-1} X_k e^{-j \frac{2\pi}{N} r k} e^{j \frac{2\pi}{N} (k-c)^2} \right|
\]  

(3)

where factor \( e^{-j \frac{2\pi}{N} r k} \) before the summation has been omitted, which is nothing to beamforming.

When Eq.(3) is used for beamforming, it is required that \( X_k \) should be a "snap shot" from elements at various instant. Therefore, multichannel preprocessing, including normalization and compression of the dynamic range as well as sampling and holding, must be carried on before CZT processor. If the signals of array elements are sampled serially at sampling frequency \( f_s = 2f_o \), the received signal from each element has a time delay, as a result, each beam will only predefected by an angle

\[
\theta_p = \sin^{-1} \left( \frac{C}{2f_o d} \right) = \frac{\lambda}{2d}
\]  

(4)

For a CZT of N points, the required stage number of transversal filter must be 2N–1 and can not work continuously. If Eq.(1) is changed to a sliding CZT:

\[
B_r = \left| \sum_{s=0}^{N-1} X_{r+s} e^{-j \frac{2\pi}{N} s k} e^{j \frac{2\pi}{N} s k^2} \right|
\]  

(5)

and the serially sampling is adopted, each time \( |B_r| \) is calculated, the sample \( X_k \) and the premultiplying factor \( e^{-j \frac{2\pi}{N} s k} \) all slide r samples, therefore we can use a N–stage transversal filter to calculate the value of \( |B_r| \) continuously. This algorithm is particularly suited to high resolution within–pulse beamscanning. As a result, each time a new sample is fed into CZT processor, a new beam output will be obtained at the output of CZT processor, so that a spatial continuous beamscanning is performed in this way.

III. PREPROCESSING FOR SIGNALS AND FOCUS COMPENSATION

Since signals from elements are real sinusoidal, the in–phase and quadrature component have to be produced.

Let

\[
y_k = X_k e^{-j \frac{\pi}{N} k} = R + jI
\]  

(6)

where

\[
R = X_k \cos \frac{\pi k^2}{N} + X_k \sin \frac{\pi k^2}{N}
\]  

(7a)

\[
I = X_k \cos \left( \frac{\pi k^2}{N} \right) - X_k \sin \left( \frac{\pi k^2}{N} \right)
\]  

(7b)

where \( X_k \) and \( X_k \) are the in–phase and quadrature components from the k–th element respectively. A suggested method in the early time[5] is to use the phase shift circuit, and the pre–multiplication is performed by adjusting the gain of multi–preamplifiers. \( N \times 4 \) analog switches with high speed and two normalization circuits are needed. Clearly this method is short of flexibility. Sampling outputs of elements orthogonally can simplify preprocessing, which reduces the number of analog switches to \( N \), and normalization circuits to one. This scheme is shown as in Fig.1, where \( A_k \cos \frac{\pi k^2}{N} \) and \( A_k \sin \frac{\pi k^2}{N} \) are stored in EPROMS beforehand respectively and 4 MDAC are used to produce 4 components in Eq.(7). By changing the read address the required weighting values or the ratio of the main lobe and sidelobe can be changed conveniently.

Consider the detection in the near field, the focus compensation is needed. Generally speaking, the focus to each point in a detecting sector is impossible by using hardware[6]. However, the focus to sound sources in some range and directions can improve the effect for detection in
a short range. Assume that the signal received from a
source at range \( R \) in front of the array is

\[
X_k = A_k e^{j \frac{2\pi f}{C} \Delta t_k} \tag{8}
\]

where \( \frac{\Delta t_k}{C} \) is the signal delay relative to a reference ele-
ment due to near field effect. After weighting and time-de-
lay focus compensation, it becomes

\[
X'_k = A_k a_k e^{j \frac{2\pi f}{C} (\Delta t_k - \Delta t'_k)} \tag{9}
\]

where \( \frac{\Delta t'_k}{C} \) is compensated delay. Therefore the chirp
pre–multiplying term becomes

\[
y' = X'_k a_k e^{-j \frac{\pi}{N} k^2} = X_k a_k e^{-j \pi \frac{k^2}{N} + \frac{2\pi f}{C} \Delta t'_k} \tag{10}
\]

It can be seen from above equation that if we store the
weights:

\[
W_c(k) = a_k \cos \left( \frac{\pi k^2}{N} + \frac{2\pi f}{C} \Delta t'_k \right) \tag{11}
\]

\[
W_s(k) = a_k \sin \left( \frac{\pi k^2}{N} + \frac{2\pi f}{C} \Delta t'_k \right) \tag{12}
\]

into EPROMs in advance, the near field focusing and
amplitude weighting for elements can be performed.

IV. EXPERIMENTAL RESULTS

An experimental system has been built in laboratory.
The working frequency is 330kHz, the receiving sector is
23° and the scanning rate is 20.6kHz. The system is com-
posed of receiving array of 32 elements with a beamwidth of
0.7°, transmitting array, receiver, transmitter and the
display monitor. The CZT processor consists of 4 CCD
tapped delay lines. A single chip TMS320C25 is employed
like the calculation required from rectangular to po-
lar coordinate for display, so that the update rate of picture
on the screen reaches 10 frames per second in the range of
40m.

Some experiment has been done in a test tank. The
targets are two plates with the size of 20cm x 20cm sepa-
rated by 20cm in both direction and range, and they are
about 6.5m far away the transmitting and receiving arrays.
Fig.2 shows the obtained picture of targets. It can be seen
that two targets can be resolved in direction and range.

V. CONCLUSION

A clear advantage of this scheme is that the single
channel processing can be done just after preamplifiers, so
that the system is compact. This scheme can be expected to
use in some vehicles where the space may be limited.

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Improvements of Hydrodynamic Noise Measurement in Water Tunnel

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1. Introduction

Water tunnels are the principal facility used in hydrodynamic noise research. The main purpose of noise measurements in water tunnel is to obtain the equivalent free field acoustic characteristics of models such as wings, propellers etc. Method of hydrodynamic noise measurements in water tunnels concern with the placing of hydrophones and the conversion of measured pressures in facility to the expected values in free space. Usually there were several ways for hydrophone to mounted in close-jet water tunnel: (1) hydrophone is placed in the flow, (2) hydrophone is mounted flat in the tunnel's wall, (3) hydrophone is mounted up in a small box, which mounted flat in the wall and filled with water. In open jet facility the hydrophone is generally placed in the annular part of the chamber surrounding the water jet where may be considered a dead water place [1]. Because of the complexity of the boundary condition, it is impossible to obtain theoretically the exact free field acoustic characteristics of models from the measured pressure values in tunnel, and usually experimental calibration method are used. i.e. mount a known acoustic volume source at the place where the model will be located, and measure the transfer function $H(\omega)$ of sound field between source and hydrophone, which can be defined as:

$$H(\omega) = \frac{P_T(\omega)}{P_S(\omega)}$$

where $P_T(\omega)$ is the RMS pressure 1 meter distance from the source in free field and $P_S(\omega)$ is the RMS pressure measured in the facility for known source. If $P_T(\omega)$ is the RMS pressure value measured in facility for some model, then the equivalent free field pressure at 1 meter distance from the acoustic centre of this model can be get as:

$$P_T(\omega) = \frac{P_T(\omega)}{H(\omega)}$$

In early time, methods used for noise measurement in water tunnels were widely different, so the experimental results of hydrodynamic noise obtained from different tunnels were incomparable. In 1976, 15th ITTC Cavitation Committee recommended a procedure for facility calibration, noise measurement techniques and data documentation [2], and 51 organizations were requested to comment on the usefulness of the committee's recommendations to indicate procedures followed in their facilities. Many facilities were interested in some form of comparative noise measurement among facilities. Therefore, the cavitation committee decided to initiate a comparative noise measurements and use "Sydney Express" propeller model as test body. Several facilities cooperated in performing comparative test with this model and main results were presented in ITTC81[3].

The recommended procedures requested that hydrophones mounted in a water-filled box flush mounted on the tunnel's wall and a calibration of water tunnel to obtain the transfer function from the cavitation volume to the hydrophone is done by replacing the cavitation source with a known omnidirectional source.

Fig.1 gives noise spectra of "Sydney Express" measured in facilities of Hamburg Ship Model Basin (RSTVA) of Germany, Earlslands Mekaniska Verkstad (KAMV) of Sweden, Kryloff Shipbuilding Research Institute (KRYLOFF) of USSR respectively.

![Fig.1 Cavitation noise of propeller.](image)

We can see from Fig.1 that noise spectra measured in different facilities differ largely from each other. Above 8kHz, the trend of spectra decreasing are similar, but maximum difference in level is nearly 20 dB. Below 8kHz, maximum difference in level between spectra is nearly 50 dB. In spite of that there are many incomparable factors in their experiments and data processing, defects of noise measurement method may be the principal cause.

It is known that due to reflections at boundary, the sound field in any limited space may have complicated spatial distribution, the pressure value measured at some single point could not reflect the real characteristics of the source. The noise spectrum measured at this point therefore generally will not consistent with free field measurement. Noise propagation in duct and characteristics of sound field in box flat mounted on elastic tube had been discussed in Ref.[4]. The results indicated that noise spectra measured in duct or in box flat mounted on duct can be converted to equivalent free field measurement by experimental calibration only when the noise source is a point one. But hydrodynamic noise source in water tunnel experiments always have random spatial distribution, therefore, using transfer function of a point volume source to calibrate the results of real spatial distributed source will introduce large error.

In order to obtain the free field acoustic characteristics of the noise source in facility, we propose a reverberation chamber method to measure hydrodynamic noise radiated by tested model in water tunnel, i.e. build up a water filled tank surrounding the test section of close-jet water tunnel, measure the spatial averaging pressure spectrum of model radiated noise in the tank, then convert the measured results to equivalent free field level by sound field calibration.

II. Sound field in reverberation tank

The dimension of water tank in our experiment is about 1.5x1.3x1.3m$^3$, the wall are made of plastic plate, the bottom is concrete and the upper surface is free. Assume a monopole source at $r_0$ in the tank, the average square pressure at point $r$ can be represented as

$$\int |P(r, r_0)|^2 dV = \frac{4\pi p_0^2 r_0^2}{K^4} \sum_{n=0}^{\infty} \frac{K-n}{K+n} \left( \frac{K-n}{K+n} \right)^2 \sum_{m=0}^{\infty} \left( \frac{||r||}{K+n} \right)^2 \left( \frac{||r||}{K-n} \right)^2$$

(3)
where \( \rho \) is the density of water, \( V \) is the volume velocity of the source, \( V \) in volume of the tank, \( \phi_n(r) \) is the \( n \)th eigenfunction, \( \Lambda \) is the orthogonal normalization factor of eigenfunction, \( k_n = \sqrt{k - k_0}, k_0 \) is wave number, \( k_0 \) is the \( n \)th eigenvalue, \( c_1 = \text{arg}(k - k_0) \). \( k^2 \) can be expressed as:

\[
k^2 = r^2 - \frac{1}{4V}(q - a_0^2)
\]  

(4)

where \( r \) is the eigenvalue of homogeneous wave equation under rigid boundary condition, \( q \) and \( a_0 \) determined by impedance of the real boundary. From Eq.3 we can see that the average square pressure at arbitrary position consist of two series, the first one corresponds to the contribution of adding each normal mode independently, the second one corresponds to the contribution of interference between normal modes. The value of interference term depends on position of survey point and frequency of sound wave. Because of the interference of normal modes, noise spectrum measured in tank is not similar to the spectrum of the source measured in free field. Fig.2 gives (1) the spectrum measured in tank when the source transmit white noise; (2) transmitting frequency response of the projector in free field.

![Fig.2](image)

From Fig.2 we can see that there are large difference between spectrum measured in tank and transmitted frequency response of the projector in free field. Fluctuation of spectrum measured in tank is about 20 dB, details of spectra measured at different position in tank are different. Therefore spectrum measured at single point in tank could not reflect the real spectral properties of the source in free field. Taking spatial average of Eq.3 and utilizing orthogonality of normal modes we have

\[
\bar{P}_1 = \frac{1}{V} \int \int |P(\mathbf{r}, r_0)|^2 dV = \frac{1}{V} \sum_{n} \sum_{m} \frac{|\phi_n(r_0)|^2}{k_n^2} \Lambda_n
\]  

(5)

The square pressure of spatial averaging is only determined by excitation of the source, the interference of normal modes have been eliminated. When frequency is high enough, i.e. there are a lot of normal modes be excited in a defined frequency band, and substituting \( |\phi_n(r_0)|^2 \) by its spatial averaging value will not induce large error, convert sum to integral we can obtain the result from Eq.5(6):

\[
\bar{P}_1 = \frac{16\pi \rho q^2 V}{a_0^2}
\]  

(6)

where \( a \) is the area of inner surface of the tank, \( a_0 \) is the average coefficient of sound absorption of the tank wall. If volume velocity is represented by radiated sound power, Eq.6 may be written as

\[
\bar{P}_1 = \frac{4pcw}{a_0}
\]  

(7)

where \( W \) is the sound power radiated by the projector, \( c \) is sound speed in water. Eq.7 tells us that the spatial averaged square pressure is proportional to radiated sound power of the source and inversely proportional to the sound absorption in tank. Fig.3 gives the spatial averaged pressure spectrum measured in tank while the projector excited by white noise.

![Fig.3](image)

The spatial averaging is performed by moving the hydrophone slowly and simultaneously perform sampling analysis by spectra analyzer. Comparing Fig.3 with Fig.2 we can see that the spatial averaged pressure spectrum in tank and pressure spectrum of projector in free field resemble each other. The interference of normal modes are flattened. If the volume velocity in Eq.6 expressed by the average pressure at 1 meter distance from the source in free field, the ratio of spatial averaged square pressure measured in tank and square pressure at one meter distance from the acoustic centre of the source in free field could be written as

\[
\bar{P}_1 = \frac{16\pi \rho q^2}{a_0}
\]  

(8)

\( R \) may be defined as constant of "room", and is only determined by the physical parameters of the tank and fluid filled in it. \( R \) can be measured by a calibrated source or calculated by Eq.8 using Sabine formula and measured reverberation time of the tank.

III. Conclusion

Build a water tank around the test section of close-jet water tunnel and perform noise measurement in the tank, then radiated sound power and equivalent free field spectrum of hydrodynamic noise sources could be obtained by Eq.7 and Eq.8. Which will give results comparable for hydrodynamic noise measurements of models in different facilities.

References:

AN ANALYSIS OF THE SOUND SPEED CHARACTERISTICS IN THE SEA NEAR THE ESTUARY OF THE YANGTZE RIVER AND IN THE NORTHW ESTERN PACIFIC OCEAN

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In this paper, the sound speed, which varied with depth, season, time and geographic location, at the estuary of the Yangtze River and in the sea area near Jishou Island, was calculated and analyzed on the basis of observations. The results show that the space-time variation of sound speed in the sea area under investigation was relatively considerable, and that there existed considerable seasonal variations and regional features for both thermocline and sound channel.

INTRODUCTION

It is well known that there are many research subjects in underwater acoustics and oceanography relating to the sound speed in sea water. It is therefore of great significance to study the distribution of sound speed in the ocean and its variations. In our study the following equation of Mill [1] was adopted:

\[ C_{STP} = C_0 + C_1 + C_2 + C_3 + C_4 \]

where

\[ C_0 = 1402.392 \]
\[ C_1 = -0.501109368673 \times 10^{-3} + 0.55048471272 \times 10^{-6} \]
\[ C_2 = 0.221535969240 \times 10^{-9} \]
\[ C_3 = -0.135923290781 \times 10^{-12} + 0.128955756644 \times 10^{-15} \]
\[ C_4 = 0.265484716608 \times 10^{-18} + 0.15934974045 \times 10^{-21} \]

In this survey, 6 cross-sections were set up, 13 stations, of which 2 were continuous stations, were established. The measuring range was 122°30' - 128°30' E and 30°45' - 35°00' N, with depth ranging from 15m to 115m.

DISTRIBUTION AND VARIATION OF SOUND SPEED

Plane Distribution

The plane distribution of sound speed in the surface layer can be seen in Figs. 1, 2, 3 and 4. It is clear in these figures that in the sea area under investigation the surface layer sound speed was low near the coast and high in the offshore sea, and at the same time, it was low in the area under the influence of the fresh water of the Yangtze River, and high in the area under the influence of the Eurobetic current. Nevertheless, the distribution also varied with the seasons. In winter, the surface layer sound speed increased considerably from the west to the
east; its minimum, 1480.7 m/s, lay near the mouth of
the Yangtze River, and its maximum, 1527.4 m/s, lay
in the region of the main stream of the Kuroshio.
The longitudinal gradient was large. In summer, it
increased universally, with the augmented range in-
creasing from the west to the east; all the sound
speeds were above 1520 m/s, and the longitudinal
gradient was less than that in winter. Spring and au-
tumn were the transitional seasons for the hydrolo-
gical elements. In autumn, the north wind began inten-
sifying and during this period the coastal current
came to appear, the isovels were gradually tending to
parallel to the coastline. This was more obvious in
spring.
Thus it can be seen that although radiation
and evaporation cause seasonal variations of the
sound speed, the current system of high temperature
and salinity in the offlying sea is also one of the
important factors affecting the status of sound
speed distribution.

Vertical Distribution of Sound Speed

The influence of meteorological conditions on
the sea area under investigation was very consider-
able. Because the huge differences in meteorological
conditions between winter and summer, the vertical
distribution of sound speed basically existed in two
types: the winter type and the summer type (Fig. 5).

![Fig. 5. Vertical distribution of sound speed at Station A3 in different seasons.]

1. Positive Gradient Type

In this case, the sound speed increased with
depth.

2. Thermocline Type

The sound speed, beginning with a certain depth
decreased abruptly with depth, with its vertical
gradient \(|\Delta C/\Delta z| \geq 0.5\) m/s.
The vertical distribution of sound speed of the
thermocline type generally can be modeled by a
three-layer structure. The first layer was the upper
positive gradient layer; the second layer was the
negative thermocline; and the third layer was the
lower positive gradient layer. In such a case, the
first layer was mainly affected by temperature and
salinity; the vertical distribution of sound speed in
the second layer was largely consistent with that of
temperature. Thus it can be seen that the varia-
tion in sound speed mainly resulted from tempera-
ture. The third layer was mainly affected by pres-
sure.

Analysis shows that the Yellow Sea Cold Water
Mass played an important role in the formation and
distribution of the thermocline of the Yellow Sea.

Time Variation of Sound Speed

As a whole, the highest sound speed in the sur-
face layer often appeared in early autumn and the
lowest in early spring. The diurnal variation of sound speed was in accordance with the diurnal vari-
ation in the temperature of seawater. The diurnal
variation in the bottom layer was considered to be
related to tides.

**SOUND CHANNEL**

**Surface Sound Channel**

The above-mentioned upper positive gradient
layer was the surface sound channel; its axis was at
the sea surface. This channel existed in all seasons.
The thickness of this channel was only 8.4 m in sum-
mer and 66.8 m in winter.

**Underwater Sound Channel**

The depth of the axis of this channel undulated
sharply in summer (Fig. 6), the maximum depth being
at 100 m, and the minimum being at 20 m. The fluctu-
ation of sound speed at the axis was quite large.
After summer this channel gradually vanished.

![Fig. 6. Depth variation of sound channel axis.]

**Double Sound Channel System**

This channel formed only in summer, the depth of
the axis for the upper channel being 22 m and 15
m, respectively, and that for the lower channel be-
ing 75 m and 76 m, respectively.

**CONCLUDING REMARKS**

The sound speed in the area varied mainly with
temperature except for the Yangtze River Estuary,
where salinity was also important. It was lower near
the coast than the open sea, and lower near the Yang-
tze River Estuary than the Kuroshio area. Its time
variation was basically of two types: surface and win-
ter; and its profile was of three types: surface chan-
nel (all year round), underwater channel (in summer),
and double channel (in summer).

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MFSK-DTCCA AND ITS APPLICATIONS
TO UNDERWATER ACOUSTIC TELEMETRY

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1. INTRODUCTION

There exist temporal, spatial and frequency-dependent fluctuations, multipath propagation and high noise in the underwater acoustic channel. So it is much more difficult for information to be transmitted in underwater acoustic channel than in air radio channel. Especially there are severe multipaths including macromultipath (reflected paths from surface and bottom) and micromultipath (a single path) in underwater acoustic channel. These multipaths are the main obstacle to underwater telemetry.

Much efforts have been made to underwater telemetry, the successful implementations include MFSK [1], DTCCA[2,3] and etc. MFSK is of higher signal bit rate, but it cannot be effectively implemented on the condition of deep fade; Though DTCCA has low signal bit rate, it can be successfully applied to strong signal scintillation situations.

We adopt the merits of both MFSK and DTCCA and developed a unique signal processing method—MFSK-DTCCA (Multi-Frequency Shift Keying and Digital Time Cross Correlative accumulation). Theoretical estimations and in situ experiments showed its performance is excellent. It can reject the intersymbol caused by macromultipath and suppress signal fluctuation caused by micromultipath. The system is also of high probability of detection and very low probability of false alarm on the condition of low signal to noise ratio.

2. THE MFSK-DTCCA SYSTEM

2.1. Transmitter system

The functional block diagram of transmitter system is illustrated in Figure 1. Data is transferred to a microprocessor, the microprocessor encodes the digital data, modulates a signal and transmits this signal, as well as signal for synchronization, into the acoustical medium. In order to verify that the MFSK-DTCCA method is valid for acoustic telemetry, only five frequencies are used in our prototype. (up to two hundred and fifty six frequencies are employed in the newly-developed system). Signaling structure is shown in Figure 2. Frequency fs acts as a synchronization pulse; The other four frequencies act as coded pulses. The number of combination obtainable from these four frequencies is equal to twenty four.

2.2. Receiver system

The functional block diagram of the receiver system is illustrated in Figure 3. PLL (Phase-locked loop) is employed to demodulate MFSK signal, the demodulated signal is transferred to a microprocessor. Digital time cross correlative accumulation and decoding are accomplished in the microprocessor. In order to reject the intersymbol caused by macromultipath, a specific signal processing method was developed in the receiver system; each frequency is transmitted only once in one block of coded pulses. Once a pulse of frequency fn is received, the receiver system automatically rejects this frequency fn after this time. So the system can successfully reject intersymbol caused by macromultipath in the duration of one block of pulses. In addition, in order to overcome the fluctuation caused by micromultipath, appropriate threshold Mn of cross correlative accumulation can be selected. So satisfactory probability of detection and very low probability of false alarm could be attained with the MFSK-DTCCA system.

3. THEORETICAL ESTIMATES

3.1. Probability of detection (4)

Because the signal envelopes follow Rician distribution, the probability of detection of the MFSK-DTCCA system is

\[ P_d = P_{s} \cdot (P_{s})^M \cdot N^4 \]

in which

\[ P_{s} = Q(\mu^2/2, \nu/\xi) = Q(3.5, \mu^2/2) = 0.95 \]

(ranges: 2.5 km, \mu^2/2 = 3.5 \xi e/2)

\[ N^4 = \sum_{n=0}^{N-1} P_{s} \gamma^2 \]

Table 1 lists Ps and Ps at different threshold Mn.

3.2. Probability of false alarm (4)

The probability of false alarm of the MFSK-DTCCA system is

\[ P_f = P_r \cdot (P_{r})^M \cdot N^4 \]

in which

\[ P_r = 10^{-4} \]

\[ N^4 = \sum_{n=0}^{N-1} P_{r} \gamma^2 \]

Table 1 lists Pr and Pr at different threshold Mn.

4. TEST PROGRAMME

4.1. Test in a tank

The received signal waveform in the reverberation tank is shown in Figure 4. Though there exists very severe macromultipath in the tank, the probability of detection of the system was 100 percent, and no error was detected.

4.2. Field test

The field test was conducted in Xiamen harbor, Fujian, China. One transducer was deployed in a fixed location in 20 meter of water; The other was
placed on a nonmoored vessel which was drifted about one knot. The sea state was 4-5. Table 2 lists probability of detection at the range of 2.5 nautical miles with different threshold $N_t$. The entire transmission lasted approximately one hour, no error was detected.

5. CONCLUSIONS

In situ test demonstrated that a 2.5-nautical mile transmission range with a greater than 98 percent probability of correct message detection could be achieved with the MPSK-DTOCA system.

The following general conclusions can be made from the test results and theoretical estimations.

1. The MPSK-DTOCA system can reject inter-symbol caused by macromultipath and suppress signal fluctuations caused by micromultipath.

2. The system is of high probability of detection and very low probability of false alarm on the condition of low signal to noise ratio.

3. The signaling rate of the MPSK-DTOCA system can be increased if the bandwidth is available and signal processing can be implemented within economic constraints.

6. REFERENCES


Tab. 1. Estimated probability of detection & probability of false alarm at different threshold

<table>
<thead>
<tr>
<th>Threshold $N_t$</th>
<th>$P_d$</th>
<th>$P_{fa}$</th>
<th>$P_{d2}$</th>
<th>$P_{fa2}$</th>
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<td>0.997</td>
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<td>6×10^{-17}</td>
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<td>0.946</td>
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Tab. 2. Measured probability of detection at different threshold (range: 2.5 km)

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<td>0.98</td>
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<tr>
<td>2</td>
<td>1.0</td>
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Fig. 1. Transmitter system

Fig. 2. Signaling waveform

Fig. 3. Receiver system

Fig. 4. Received signal waveform
A NEW BEAM PATTERN OPTIMIZATION METHOD FOR COMPLICATED ARRAYS

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ABSTRACT

A new beam pattern optimization method for complicated arrays is discussed. It was proposed in [6] firstly. The method was based on simulated annealing algorithm that is a powerful technique for solving multidimensional combinatorial optimization problems. In this paper, a more complete analysis about the method is given, the “energy function” and “system temperature” selection, associated with the corresponding procedure, are emphasized. A practical design example is presented. It is shown that the new method is effective generally and can be used to Sonar, Radar and other fields.

INTRODUCTION

Sensor array designing plays an important role in Sonar, Radar and other applications. One of the key tasks in array designing is beam pattern optimization. Although adaptive array processing has been developed from years, fixed beamformers are still used in many cases, especially in practical applications. The reason is that traditional beamformers requires smaller computations and simpler hardware. For a linear array, there are many methods for optimizing the beam pattern, such as Dolph-Chebyshev and Hann methods. But in the case the the configuration of array become very complex, these methods are failed to do optimization. It is same with an array of arbitrary configuration. So that a general optimization method is needed. "Notch Noise Field Technique" proposed by Y. L. Ma and J. W. R. Griffiths[4] is a good approach. Because of the complicated relation between the pattern and the configuration of array, mathematic method is not applicable. "Notch Noise Field Techniques"-NNFT consider the beam pattern optimization from a systematic view. This is the method’s key point. We follow the same idea, but take a different approach. We introduce simulated annealing algorithm to beam pattern optimization. In this paper, we give a more complete description of the new approach. The emphasis is put on the selection of the “energy function” and the “system temperature”. A Modified Simulated Annealing (MSA) algorithm, associated with its procedure, is described fully. The robust of MSA, in terms of the initial data, system temperature and energy function, is pointed out. Some useful conclusions are obtained.

SIMULATED ANNEALING ALGORITHM

Simulated annealing is a computational heuristic for obtaining approximate solution to combinatorial optimization problems. It imitates the cooling stages of annealing that is usually done very slowly to reduce the number of defects in the crystal structure and to minimize the potential energy stored in the molecular configuration. Seeking the minimum of a given objective function of many variables is a typical combinatorial optimization problem in statistics. The variables are subject to intertwining constraints, and they interact with each other in complicated ways not unlike the molecules in a physical annealing. By appropriately defining an effective temperature for the multivariable system and imitating the physical annealing process, researchers have sought to solve complicated optimization problems. Kirkpatrick (1983) first investigated the use of simulated annealing in connection with the physical design of computer, such as VLSI layout and partitioning problems. S. Geman and others (1984) used an annealing algorithm to the restoration of noisy blurred images; EL Gamal and co-workers (1987) to the designing of communication codes; Sharrman (1987) to the high resolution DOA (Direction Of Arrival ) parameter estimation. The successes of simulated annealing have resulted in a surge of interest in the methods. There are intensive researches on it. More recently, a fast simulated annealing (FSA) algorithm is proposed by H. Szu (1986). It is adopted in this paper.

The simulated annealing algorithm involves two quantities "energy function" $E(\vec{W})$ and "system temperature" $T$. $E(\vec{W})$ is a cost function that is to be minimized and $\vec{W}$ is a column vector of N parameters to be found. Our object is to find a set of parameters $\vec{W}_{opt}$ that yield a global minimum cost function.

BEAM PATTERN OPTIMIZATION

One of the most important advantage of an array is its directional properties that enable it to discriminate between signals arriving from different directions and improve signal-to-noise ratio. This property is highly desirable in most practical applications. Array directionality designing or beam pattern designing implies the appropriate selection or estimation of the optimal weights. A basic model of array processing is shown in Fig. 1. By choosing a set of weights correctly the desired directional pattern can be obtained, i.e. optimization of directional pattern. We introduce simulated annealing algorithm to the optimization problem.

Let signal vector be

$$\vec{X}(\theta) = [x_1(\theta), x_2(\theta), \ldots, x_n(\theta)]^T$$

where $\theta$ is the direction of incident signal. Let weight vector be

$$\vec{W} = [w_1, w_2, \ldots, w_n]^T,$$

so the directional pattern is given

$$G(\theta) = \vec{W}^T \vec{X}(\theta).$$

Let the desired directional pattern be $G_0(\theta)$, we define the "ener-
gy function" as
\[ E(\tilde{W}, \theta) = E_0 [G(\theta) - G(\theta)]^2. \]

In the simulation, \( \theta \) is replaced with limited-length discrete data. The optimization problem becomes to minimize the "energy function" \( E(\tilde{W}, \theta) \), i.e., to seek a set of weights \( \tilde{W}_{opt} \) in the meaning of minimizing "energy function".

Our simulated annealing algorithm is similar to FSA. It employs a Cauchy/Lorentzian generating probability density function:
\[ d_k(\theta) = \frac{a_{i_k}}{(1 + a_{i_k}^2 \theta^2)^{(N+1)/2}}, \]
a Boltzmann acceptance probability:
\[ \text{prob}(\text{acceptance}) = \left( \frac{1}{\text{exp}(-\Delta E/T)} + 1 \right)^{-1} \]
and the annealing schedule:
\[ T = C/(1+1) \]
\( C, a \) are constant.

We use \( E(\tilde{W}, \theta) \) (defined above) as the "energy function" to be minimized. \( \tilde{W} \) is varied during the schedule. The procedure is described as follows:

1. Set the initial parameters \( \tilde{W}_0, T_0 \).
2. Choose a N-dimensional random vector according to the temperature dependent Cauchy/Lorentzian density.
3. Give a perturbation to the current weights using the random vector. \( \tilde{W}_i \) is got.
4. Calculate the "energy" \( E_i(\tilde{W}_i, \theta) \).
5. If \( E_i < E_0 \), accept the new parameters with a probability according to the Boltzmann distribution.
6. If \( E_i > E_0 \), accept the new parameters with a probability according to the Boltzmann distribution.
7. Reduce the temperature \( T_i \)
\[ T_{i+1} = T_i/(1+1) \]
and go to 2.

The algorithm will continue until several energy drops or a set of stable weights are obtained.

OPTIMIZATION RESULTS AND CONCLUSIONS

To demonstrate the SA directional pattern optimization method, extensive simulations are carried out on a variety of sensor arrays. The results are generally favourable. Here we give the simulation results in part. A special sensor array is considered. The configuration of the array is shown in Fig. 2. It is a plane array and comprises 6 sensors. It is difficult for Dolph-Chebyshev and other methods to optimize the directional pattern. Fig. 3 gives the optimization results by simulated annealing algorithm. A good shaping pattern is obtained, while the mainlobe is kept very narrow (\( \psi_{30} = 30° \)), the sidelobes are suppressed under -20db. The higher the initialised temperature, the smaller the residual energy after converge. But the energy could not be less than some certain value which is determined by sensor array itself and the energy function.

We introduce simulated annealing algorithm to beam pattern optimization. Our simulation results indicated that simulated annealing appears to have the potential for the applications in sonar and radar array design. It was shown that an annealing algorithm could be used to the optimal beam pattern design problem involving in sensor array processing. The annealing solution can outperform existing methods, especially for the irregular sensor array with complex configuration.

Occasionally in some cases, the algorithm couldn't be converged to the minimum energy state. So the energy function must be selected carefully according to the different sensor array.

REFERENCES

SOUND PROPAGATION IN A MEDIUM WITH ANISOTROPIC INHOMOGENEITIES

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Recent years have seen an increasing number of papers discussing the remote measuring of sound distribution of ocean currents using acoustical scintillation technique. For coastal seas, the space-time analysis of acoustical signals not only provides the speed profile of tidal currents but also reveals the anisotropic features of these currents. The latter has been typically shown in (5), discussing the results of experiments conducted in October 1986 in Cordova Channel, B.C., Canada. Partial results of the experiments are presented in Fig. 1, where one can see various degrees of anisotropy of marked with the same letters of a tidal cycle in Fig. 1(b), where the ordinate represents the magnitude and direction of the tidal flow. Two possible explanations are suggested in Ref. 5 for this phenomenon, but none of which can reproduce this kind of distribution of arrival angle. This article investigates this phenomenon from another angle. Assume that between the source and receiver there is an ellipsoidal water mass whose location is randomly selected. The angles of arrival are calculated by ray method and their distributions are obtained, which are found to be similar to the experimental results.

MODEL AND METHOD

Assume the distance between the source and the receiver is $670m$, the background sound speed is uniform and is $1486m/s$. Between the source and receiver there is a water mass whose location is randomly selected within a cylinder whose semi-major axis connecting the source and receiver. The sound speed of the mass in

$$C = C_0 \exp \left[ -\frac{Z_0^2}{W_o^2} - \left( \frac{X - X_o}{W_x} - \frac{Y - Y_o}{W_y} \right)^2 \right]$$

where $C_0$ is the background sound speed, $\Delta$ is the strength of the sound speed perturbation, $X_o$, $Y_o$ and $Z_0$ give the water mass location, and $W_x$, $W_y$ and $W_z$ give the size of the water mass. The above expression is identical in form with that used in the ray tracing program HARP (5), but there is a significant difference. In HARP, all the position parameters are relative to the earth, therefore the orientation of the water mass or the blob is fixed. In this article, (1) is written in a rotational system relative to the earth, the angles of rotation are represented by Euler's angles $\omega$, $\gamma$ and $\phi$. Hence the orientation can be changed.

In order to increase the accuracy in eigenray tracing, we choose $\Delta = 0.0008$, i.e., the largest variation of sound speed is $1.29 m/s$, about an order of magnitude larger than the observations.

The existing methods (7) for obtaining eigenrays are not suitable for three-dimensional cases. For the sound speed distribution used in this article we developed a new method - the self-searching method to find eigenrays (9). Compared with the traditional method, this new method can reduce the number of ray calculations by over two orders of magnitude, making the calculation of a large number of eigenrays practicable. The main point of this method is as follows. A Cartesian coordinate system is set through the receiver plane, which is perpendicular to the line connecting the source and receiver and with the origin at the receiver. The coordinates $x$ and $y$, which give the position of the intersection point of the ray on the receiver plane, are calculated. The values of $x$ and $y$ are used to adjust the transmission angle and a new ray is calculated, the absolute value of the coordinate of whose intersection point on the receiver plane is to be smaller than the previous one. This process is repeated until both $|x|$ and $|y|$ are satisfied, $\xi$ being a preset small positive number (in this study we chose $0.1m$). Fig. 2 presents an example for obtaining an eigenray. In this example the position of the blob in this $(x,y)$ plane is $(31.5,2.1)$ from the source, $1.4m$ from the axis connecting the source and receiver in the horizontal direction, and $1.3m$ from the axis in the vertical direction. The rise of the blob: $W_x = 0.6m$, $W_y = 0.8m$, and $W_z = 0.6m$. The inclination $\alpha = \frac{\pi}{4}$ ($\alpha$ is the angle between the plane containing the two long axes of the blob and the horizontal plane containing the source or the receiver).

Fig. 1. (a) Distributions of arrival angle.
(b) Temporal variation of tidal flow.

Fig. 1. (a) Distributions of arrival angle.
(b) Temporal variation of tidal flow.
Fig. 2. An example of self-searching of eigen ray.

RESULTS AND DISCUSSION

Having obtained the eigenray, the horizontal and vertical arrival angles can be calculated. Fig. 3 presents the distributions of angles of arrival for different values of the water mass parameters, which are

given in Table 1, with the letters a, b, ..., f each corresponding to a plot with the same letter in Fig. 3.

Table 1. Values of water mass parameters

<table>
<thead>
<tr>
<th>Group</th>
<th>( \omega ) (deg)</th>
<th>( W_0 ) (m)</th>
<th>( W_p ) (m)</th>
<th>( W_0 ) (m)</th>
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<td>1.5</td>
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<td>b</td>
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<td>4</td>
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<td>c</td>
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<td>45</td>
<td>4</td>
<td>4</td>
<td>2</td>
<td>0.0004</td>
</tr>
</tbody>
</table>

Fig. 3. Calculated distributions of arrival angle.

Comparison between Fig. 3 and Fig. 1(e) shows that the calculated distributions of angles of arrival are similar to the experimental results when assuming there is an ellipsoidal blob between the source and receiver, with the orientation of the blob fixed and its position randomly changed within certain limits, which implies that under the circumstances of the Cordova experiments in the medium between the source and receiver there might be structures similar or equivalent to ellipsoidal water masses. The larger the difference between the sound speed in the blob and the background sound speed (\( A \) is larger), the larger the distribution range of the arrival angles; when the shape of the blob changes from ellipsoidal to being spherical (i.e., isotropical sound speed), the distribution of arrival angles changes accordingly from being approximately elliptic to approximately circular. All these are not difficult to understand.

This is only preliminary work, as there remains a series of problems to be solved. It is hard to imagine there is only one water mass between the source and receiver; what would be the case if there are more? What is the mechanism for the formation of water masses, if there are any? What are the factors dynamically determining the parameters of the water mass?

ACKNOWLEDGEMENTS

We are grateful to R.M. Jones for providing and familiarizing us with the NOAA 3-dimensional ray tracing program.

A METHOD FOR HIGH RATE TRANSMISSION OF TWO-WAY MULTIPLEXING ACOUSTIC DATA

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INTRODUCE

In some underwater exploration, it is needed to transmitt multiple channel sensor data to a remote processing centre with high rate, and, at the same time, to send control commands, DC current and power pulses downwards to the underwater data acquisition modules. When strict confinement to the sizes and weights of transmission cables exists, using a single-lead coaxial cable is superior to using a multiplex-lead cable.

On the other hand, it is a difficult technical problem to realize the high rate two-way transmission with low bit-error-rate through only a single-lead coaxial cable.

This paper proposes a method which we call Time-Domain Frequency-Guiding. It has been used in an experimental system and proved a successful method.

STRUCTURE OF MULTICHANNEL TRANSMISSION SYSTEM

This system is a time-sharing and frequency-sharing one. Different signals occupy different frequency band and transmitting time. These differences of the signals also enable us to separate them at the receiver end. The principle diagram is Fig.1. It consists of filters, timers, modulator-demodulators (Modem) and a single-lead coaxial cable. The filters at two ends of the cable have identical frequency characteristics which match the signal spectrum. The timers accomplish the time-sharing function. The Modems modulate and demodulate the 2PSK signals of acoustic data and command data. DC current and low-frequency power pulses occupy the band which is far apart the 2PSK signal carrier, so, they can be transmitted directly without modulation. For decreasing the size and weight, the signal-lead cable is chosen. It brings some difficulties to multichannel transmission. To overcome them, we design a special time-sharing and frequency-sharing technique.

METHOD OF CARRIER RESTORATION

2PSK modulation has high anticlutter ability and has found wide applications. Because of the absence of carrier component, the local-carrier restoration becomes the heart of the matter in the coherent demodulation of 2PSK signals. When we use only a single-lead cable as media in a time-sharing two-way transmission system, the demodulator at receiver end will get an intermittent input. This make trouble to the working of PLL (Phase-Locked Loop). At the beginning points of each input segments, the restored local-carrier often has not the required frequency, or shows a phase ambiguity. This will affect the bit synchronization and make bit-error-rate significantly rise.

To overcome the above difficulty, we design the Time-Domain Frequency-Guiding method. It can remove the phase ambiguity without affecting the high anticlutter ability of 2PSK modulation.

The principal points of this method can be simply described as follows. (1) At the transmitter end, a narrow pulse filled up with carrier wave is inserted at the beginning point of each 2PSK signal segment which will be sent out. We call this pulse frequency-guiding signal. It will be used in receiver for local-carrier restoration. (2) At the receiver end, when frequency-guiding signal presents, PLL turns into holding-state. In this state the control voltage of VCO (Voltage-Controlled Oscillator) keeps constant, and so the output frequency and phase do not change. The construction and working principle is shown in Fig.2.

The PLL consists of a PD (Phase Detector), a S/H (Sample/Holding circuit) and a VCO. The repetition period of the frequency-guiding signal, and equals the period of 2PSK signal segments. The width of each frequency-guiding pulse is a little more than the catching-time of the PLL. The width of the S/H control pulse is equal to or less than the width of each frequency-guiding pulse. The S/H control pulse is generated by some circuit when frequency-guiding signal presents.

The difference between such a PLL and a normal one is that a S/H is interposed between the PD and VCO of the former. This S/H plays a role of a switch in the loop. When a frequency-guiding pulse arrives, the S/H control pulse make the switch shut, then the PLL turns into tracking-state and the output of VCO will be forced to the required frequency and phase. As soon as the frequency-guiding pulse disappears, the switch is immediately opened and the output of the VCO keeps unchanged.

Applying such a method to the local-carrier restoration of intermittent 2PSK signal we can re-
move the phase ambiguity throughly.

ENGINEERING REALIZATION

Through properly selecting PD, VCO and S/H, we can construct a local-carrier restoration circuit which is able to work at high frequencies. But generally, acoustic data is of a relatively low sample rate. It's reasonable to assume that the data transmission rate is lower than 6Mbps. So the needed carrier frequency is lower than 24MHz. In such a condition, the circuit can be realized with an IC chip TA7193 and a little number of digital IC chips. TA7193 is originally designed for color TV sets. It contains many functional circuits such as amplifiers, AGC, PLL and S/H. We use only one TA7193 to finish the coherent demodulation of intermittent 2PSK signal. This makes the hardware of the receiver very simple, reliable and cheap. The transmission rate of this system reaches 6Mbps, transmission distance reaches 1000m, and the bit-error-rate is lower than 10.

Through interposing simple frequency-division network at the two ends of the cable, different kinds of signals, such as acoustic data, command data, low frequency power pulse and DC current can be two-way transmitted simultaneously. Such a system is suitable for ocean acoustic telemetry, underwater robots and sonars which need high rate data transmission.

CONCLUSION

If strict confinement to the sizes and weights of data transmission system exists, we may use a single-lead cable as media, select 2PSK as modulation, and inserting frequency-guiding signal in time-domain to realize the coherent demodulation. This has proved to be an effective method to realize the two-way high rate transmission system. This method has several advantages such as long transmission distance, low bit-error-rate and small size of cable. In addition, it can two-way transmit different kinds of signals simultaneously.

Fig.1 system diagram

Fig.2 principle of Time-domain Frequency-Guiding
A SIMPLE AND EASY EXPERT SYSTEM FOR CLASSIFYING THE SHIP RADIATED NOISE BASED ON HOMOMORPHIC FILTERING METHOD

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1. INTRODUCTION

In this paper we describe a new method to obtain the feature of ship radiated noise, based on homomorphic filter, and combine with the spectral lines among its noise power spectrum we can classify the different kinds of noise. In the classification, we use some methods of expert system and build a simple and easy expert system that include two main models, study and classification. The result of experiment shows this system can classify noise signal of underwater target efficiently, especially at low SNR. The accuracy of classifying is 90% at 0dB SNR and is 70% at -3dB SNR.

2. OBTAIN THE FEATURE OF RADIATED NOISE

Different moving underwater target can produce different noise power spectrum[1]. A ship noise spectrum is shown on figure 1. From this figure, we can describe the noise spectrum as an overlap of an envelope spectrum and spectral lines. For this reason, the feature of noise is composed of two feature quantities that reflect its envelope and spectral lines separately.

In the past investigation, we usually use Average Method to obtain the feature of spectrum envelope. That is mean we can separate an analyzed frequency scale into many bands and use average values of every bands to present its envelope. The problem can be resolved efficiently by using Homomorphic Filtering Method to obtain the envelope.

As the ship radiated noises have many similarities with voice signals, we can describe the ship radiated noise \( s(t) \) to a convolution of radiated source \( p(t) \), random noise \( n(t) \) and sound channel transfer function \( h(t) \), that is,

\[
s(t) = p(t) * n(t) * h(t)
\]

we can use Homomorphic Filtering Method to separate the different parts of noise signal[2]. As the spectrum of random noise is some small fast varied undulations that overlap on spectrum envelope, we can filter out random noise through adding cepstrum window time area. Figure 2 show a flow diagram using Homomorphic Filtering Method to obtain spectrum feature. The aim of obtaining average power spectrum from input signal is to cancel the environment noise overlapped on radiated noise. That can make us to obtain spectral lines feature more exactly.

\[
C_p(a) \text{ is cepstrum, } l(a) \text{ is cepstrum window.}
\]

\[
l(a) = \begin{cases} 
1 & |a| < a_0 \\
0 & |a| \geq a_0 
\end{cases}
\]

\( f(n), p(n) \) show the features of envelope and spectral lines separately. After the power spectrum shown on Figure 1 is smoothed by cepstrum window, it can be shown on Figure 3. In this procedure, the width of cepstrum window is 26, that is \( a_0 = 26 \).

Through Figure 3, we can see the smoothed power spectrum correctly, that make us to express a power spectrum, frequency range from 0 to 1000Hz, only by 26 cepstrum parameters, besides, this method also make the numbers that describe an envelope compressed.

3. THE REALIZATION OF EXPERT SYSTEM

Expert system includes 3 main models, study system, classification system and knowledge system. The whole system can call each model through the main menu. In the knowledge library, the knowledge is composed of spectral lines feature and cepstrum parameters. The aim of study system is to make the knowledge complemented through continuously studying. The special indication is the complement of the old knowledge. We use the method of averaging cepstrum parameters and selecting spectral lines. After synthesis, the knowledge has more representatively. In the classification, we use the cepstrum knowledge \( F(f_k) \) in the knowledge library to compare with the cepstrum parameters \( F'(f_k) \), that will be classified. Obtain Euclidean distance between them[3].
In the meantime, we should compare with spectral lines, and produce a set of weight vector $A(a_i)$ that reflects the overlap level of spectral lines. The more the spectral lines overlap, the bigger is $a_i$. Use $a_i$ to weight $d_i$ we can produce a new set of distance value.

$$d_i' = \frac{d_i}{1 + a_i}$$

From the least $d_i'$, we can certain that the classified signal is similar to the item in the knowledge library. So we can obtain the result of classification.

**Figure 1**

**Figure 3**

**REFERENCE**

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NEAR-FIELD LOCAL TARGET STRENGTH

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INTRODUCTION

The target strength TS of a body is customarily defined in far-field of the target that is for a plane wave incident of the target and for the scattered field observed an infinitely large distance away. Consequently, we have taken the target for a reradiation point source which is regarded as the acoustic center of the target. (1)

In practice, the target is illuminated by a spherical wave, and is not regarded as a "point". If for "monostatic" sonar, in a finite range $r > R_0 = L^2 / \lambda$, the amplitude and phase of the incident field are essentially uniform over the target aperture, then the definition of TS is still suitable. $R_0$ is called far-field criterion of the target, $L$ is the lateral dimension of the target, and $\lambda$ is the wavelength illuminating field over the target.

Generally, $R_0$ is quite larger in a practical measurement, for increasing the sensitivity of the system, reducing background interference or studying the reflection intensity at some parts of a large target one frequently choose a range less than $R_0$ which is called near-field of the target to make measurement. Then the target strength TS is not suitable. Measured TS is incurred large errors. (2, 3) For this reason, we advanced a new parameter NTS instead of TS.

In this paper, we give the definition and used conditions of NTS, expand the feature of NTS and relation between NTS and TS, discuss methods and results of measuring NTS of some practice target in the sea. It is showed that NTS is used for evaluating and comparing the reflection capability of different target as well as different parts of the same target in a short range from the target.

DEFINITION

When acoustic waves are incident upon a body, complicated spatial distribution of scattered radiation result near the body. In many circumstances of practical application in underwater acoustic the scattering field can be simplified. First the wave lengths of many active sonars are smaller than the dimension of the scattering body. The direct scattering component predominates and the creeping waves contributions may be neglected. Secondly, in practical measurement in the sea, a large is always located in PRAUSHOFER range of the transducers, the KIRCHHOFF approximation for scattering may be used. Scattering is still a time-resolved process. (3)

When we use a narrow pulse which length less than a range between great majority of scattering components of a target, then separated echo in the whole echo time history (ETH) can describe the main scattering contribution from every local part in the near-field target. We can use transient form of the sonar equations to define a parameter which can quantitatively express the sound reflection character of target in near-field. We called it NEAR-FIELD LOCAL TARGET STRENGTH by symbol NTS.

\[
\text{NTS} = 10 \log \frac{P_{\text{max}}}{P_{\text{max}}} \frac{2}{I - r F_0}
\]

The definition(1) differs from TS in far-field with transient form(Ref(4) eq 2.4). Here $P_{\text{max}}$ expresses a peak pressure of spherical incident wave, $r$ denotes the reference distance which lies in the direction back toward the source of sound at a distance of 1 meter from observed local surface of the target in "monostatic" sonars. $P_{\text{max}}$ is the separated echo peak pressure of the local part corresponding in the whole ETH.

Applicability of the def(1) is as follows:

High frequencies conducted K>>1, stable spherical wave field $r > S / \lambda \text{cm}$, resolvable echo pulse $t < dc/2$, $S$ is the effective used area of monostatic transducer, $r$ is the range from the nearest surface of the target, $t$ is pulse width, $d$ is the distance between scattering parts, $c$ is sound velocity in the sea.

Clearly, NTS has something to do with the type of the target and surface shape as well as relative direction to incident wave at the observed reference point, it has no relation to the work parameters and directly of monostatic sonar. It is worthy of note that NTS has relation with a measurement distance due to mutual relation between a local surface shape of a target and incident spherical waveform in near-field target. Fig.1 is NTS theoretical treatment of a fixed, rigid sphere immersed in water. It is seen that the NTS--distance relation itself contains a quantitatively described information about a local shape of the target. From fig.1, relation between NTS and TS can be showed.

When distance is short, there is a outstanding difference of them, but NTS gradually increases and is close to TS with increasing measure range $r$. For $r \rightarrow \infty$, NTS is essentially TS under situation of a far-field plane wave incident.

MEASUREMENT

It is convenient by the transient form of the active sonar equation to measure NTS of a practical target in the sea. Because an accurate knowledge of the
Transmission loss may be obtained easily.

In short pulses for monostatic, from (1):

\[
\text{NTS}(r | \xi) = \text{EL}(r | \xi) + \text{SL}(\xi) + 2\text{TL}(r | \xi)
\]  

(2)

we can obtain NTS values related to the place \( \xi \), bearing \( \Psi \) and range \( r \) of the observed local parts of the target. Where \( \text{EL} \) is peak pressure level of echo pulses, \( \text{SL} \) is source level in \( \Psi \), TL is transmission loss.

Another method is to use a reference target of an isotropic sphere (which NTS may be accurately calculated) placed at the same range of the target to be measured.

In practice measurement, owing to interference between the direct waves and the reflected waves on the sea boundaries, nonsteady of the target and the measurement system acquired echo data are always fluctuation, must take a average over data. Different from far-field, NTS is dependent on the ranges. The average may be use multi-measurement at the same range. When a precision prescribed is not high, the average can be take in a section of the ranges where transmitted loss less than the precision.

Fig.2 and Fig.3 show two examples of the variation of NTS with range, they are corresponding to two types of practice targets which observed local surface are planar and cylindrical. Table 1 presents a list of a number of NTS at the main scattering structural parts for complex target in the sea. Above results are showed that NTS can be regarded as quantitatively descriptive parameter about echo character in near-field of a target and can be used for evaluating and comparing the reflection capability of different targets as well as different parts of the same target in a short range from the target.

Table 1 NTS at the main scattering structural parts for a complex target in the sea

<table>
<thead>
<tr>
<th>Aspect angle (deg.)</th>
<th>NTS(dB) for parts</th>
<th>I</th>
<th>II</th>
<th>III</th>
<th>IV</th>
<th>V</th>
</tr>
</thead>
<tbody>
<tr>
<td>0 - 20</td>
<td>8.4 15.7 10.3</td>
<td>--</td>
<td>--</td>
<td></td>
<td></td>
<td>9.5</td>
</tr>
<tr>
<td>21 - 40</td>
<td>15.2 16.7 22.0</td>
<td>--</td>
<td>--</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>41 - 60</td>
<td>14.1 21.6 23.7 21.9</td>
<td>--</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Fig.1 The relation between the measurement distance and NTS of a fixed, rigid sphere in water

REFERENCES

VERIFICATION OF THE STABILITY CRITERION OF INVERSE
ALGORITHMS FOR LAYERED SEA BOTTOM

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INTRODUCTION

Inverse scattering problems arise in geophysics, where the properties of earth layers are to be identified from the echoes of explosions; in physics, where we wish to recover a potential from partial scattering experiments; in acoustic sounding, where material properties are to be determined from reflected sound waves and in many other fields. The inverse problem in geophysics is to determine, if it is possible, the medium properties from its response, recorded at the boundary, to some probing input signals. The crucial assumption in deriving inversion algorithms is the model of the scattering medium, and different algorithm are obtained for different elementary layer structures.

In all cases, there are two basic algorithms called layered peeling and layer-adjoining, which lead to a variety of efficient computational procedures that were arrived at in literature. Most of them did not discuss issues related to the numerical stability of algorithms. Obviously, it is quite useless to have extremely fast algorithms that provide wrong results. In this paper the DC (Downward Continuation) algorithm and the sample iteration algorithm were used. Kind's method was adopted for the acquisition of reflected signals, corresponding to layered system with the probing signals, served as standard ones for the verification of inverse programs.

The numerical simulation shows that the inversion errors are both less than 0.06%, but its capability of antisference is very poor. However, by using the stability criterion of inverse methods, the results are very satisfactory.

I. TWO INVERSE PROCEDURES

A. DC (Downward Continuation) Algorithm

B. The Sample Iteration Method

II. STABILITY CRITERION

The numerical stability of algorithms is a very important problem. In practice, noise and irrelevant scanning signal are always involved in the observed data. It has been proved that the following criteria is useful:

If \( \tilde{f}(t) = \int_{0}^{\infty} \tilde{f}(w) \cos(wt) dw \), and \( \tilde{f}(w) \) is a continuous function. The necessary and sufficient stability condition in any finite region is

\[
\inf \tilde{f}(w) > -2/\pi
\]

where \( \tilde{f}(t) \) is the recorded data, \( \tilde{f}(w) \) is the spectral function of \( \tilde{f}(t) \).

II. NUMERICAL SIMULATION

Numerical simulation is the first step in the inverse problem of the sea-bottom stratification structure, which is necessary because there is a multitude of factors which may deform the signals in the ocean, and the data collected at sea can hardly be used as standard signals for the verification of inverse methods. Kind's method was adopted in this paper for the acquisition of reflected signal corresponding to various incident signals under various circumstances which served as standard signals for the verification of inverse programs. The sound velocity profile used by D.C. Stickel is in his numerical example is given as

\[
c(z) = \begin{cases} 
c_1 & \text{if } \rho = \text{constant} \\
c_2 & \text{if } \rho = \text{variable}
\end{cases}
\]

with the origin of the x-axis pointing downwards at the sea-bottom and

\[
c_1 = 1594 \text{ (M/S)} \quad ; \quad c_2 = 1510.44 \text{ (M/S)}
\]

Fig. (1) An example of inversion for a model profile
(a) the reflected signal
(b) the model sound velocity profile
(c) the recovered profile

In Fig. (1), the curve made with (b) is obtained directly from Eq. (7). The curve made (a) is the reflected signal provided by R. Kind's method and the curve made (c) is obtained by the sample iteration.
N. MODEL EXPERIMENT RESULTS

The model experiment was conducted in a sewage pool of a paint factory. The pool is approximately 9 meters wide, 60 meters long and 2 meters deep. The deposit is dark grey, in which there are relatively solid blocks with diameters of about 4 centimeters. Sampling measurement shows that the density of the medium $\rho = 1.68 \times 10^3$ (kg/m$^3$) and sound speed $c = 1.27$ (km/s).

The source was placed at the surface to avoid the formation of a dipole sound field caused by surface reflection. The hydrophone was located about 0.7 m under the source.

Although the waveforms, before and after correction were applied, look very much the same yet the stability is much better and the final results in both inverse methods agreed well with the direct experimental observation.

ACKNOWLEDGMENT

The author would like to thank prof. Renren Zhu of Shandong University for fruitful discussions and Zhenxin Yang, Li Yin for performing the model experiment simulation.

REFERENCE


Fig. 2. The results of the model experiment (a) original recorded signal; (b) (c) recovered profiles calculated with the sample iteration and D-C algorithms; (d) the recorded signal after correction has been applied; (e) recovered profiles using the corrected signal (e).

Fig. 2(a) shows an example of waveforms received and the inversion results obtained by using DC algorithm and sample iteration algorithm are shown in Fig. 2(b), Fig. 2(c) respectively. Fig. 2(e) is the received waveform corrected by the stability criterion its inversion results are shown in Fig. 2(f) and Fig. 2(g). For the sake of comparison, the results are listed again in Table (1).

<table>
<thead>
<tr>
<th>Impedance of the sediment (x10^3Pa.s/m)</th>
<th>Error of inversion (%)</th>
<th>Impedance of the basement (x10^3Pa.s/m)</th>
</tr>
</thead>
<tbody>
<tr>
<td>the measured results of the sample</td>
<td>1.27</td>
<td>12.7 (from handbook)</td>
</tr>
<tr>
<td>sample iteration method</td>
<td></td>
<td></td>
</tr>
<tr>
<td>original signal</td>
<td>4.8</td>
<td>0.35 x 10^3</td>
</tr>
<tr>
<td>corrected</td>
<td>1.2</td>
<td>10</td>
</tr>
<tr>
<td>corrected</td>
<td>1.3</td>
<td>10</td>
</tr>
</tbody>
</table>

Table (1) The Comparison of two Inverse algorithms
Array Processing Based on Generalized Directional Signal Model with Angular Spread

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INTRODUCTION

In the early development of beamforming, two assumptions were made that the desired signal is of plane wavefronts with perfectly spatial coherence and background noise is omnidirectional (isotropic) which comes from all directions with the same power per unit solid angle. Thus the conventional Delay-Sum (DS) beamforming can provide optimum detection performance. When noise is directional, optimum beamforming first rescales the spatial noise field by the inverse of the noise covariance matrix (or by the inverse of the array data covariance matrix), then does matching processing. The noise estimator-subtractor structure has emerged as the most flexible implementation of this optimum beamformer in the sense of Minimum Variance Distortionless signal Response (MVDR). The optimum beamforming concept is a big step forward because it is freed of isotropic noise field. We are now left with a desire to drop assumption of signal with plane wavefronts as we did assumption about isotropy of noise field. Can we really do so? Regarding the assumption of plane wave signal, it is equivalent to accepting the propagation medium uni-modal and homogeneous. We are all familiar with the multipath situation (a coherent set of multipath signals in time) which is not compatible with this assumption. When the signals propagate through the practical random time and space-variant medium, their wavefronts can show the progressive losses of coherence with increasing spatial separation. These decorrelations of wavefronts result in an angular spread in the wavenumber spectrum of signal centered about the true signal direction-Of-Arrival (DOA). This signal is called Generalized Directional (GD) signal. In this paper, we present a new beamforming method for matching to the GD signal in the energy sense. Some results of the computer simulation experiments and lake-tests in Xing-An-Jiang river are given.

MATHEMATICAL MODELLING

No lossing generality, consider a linear array of \( m \) equally spaced sensors (the separation is \( d = \lambda / 2 \), \( \lambda \) is the wave-\( \text{length} \)). Thus the GD signal \( y(t) \) received at the \( i \)-th sensor is

\[ y_i(t) = s_i(t) + n_i(t) \]

where \( s \) is the mean DOA of the signal \( s(t) \) (assumed narrow band), \( n_i(t) \) is the additive isotropic noise at the \( i \)-th sensor, and \( s_i(t) \) represents the perturbed signal wavefronts at the \( i \)-th sensor. The covariance matrix of the GD signal is

\[ \mathbf{R}_y = \mathbf{R}_s - \sum_{i=1}^{m} \mathbf{R}_n \]

where

\[ \mathbf{R}_s = [s(\theta_s) \ldots s(\theta_s)] \]

\[ \mathbf{R}_n = [n_1(n_1)] \]

\[ E[\mathbf{s}_s(t)\mathbf{s}_s(t)^H] \]

\[ E[\mathbf{n}_n(t)\mathbf{n}_n(t)^H] \]

In the above expressions, \( E \) denotes a finite time average; \( , \ldots \) denotes the conjugate transpose; \( \mathbf{A} \) denotes matrix; \( \mathbf{A}^T \) denotes Schur-Rademar matrix product; \( \sigma^2 \) is the signal power; \( \sigma_n^2 \) is the noise power, and \( [ \cdot ] \) is the \( m \times m \) identity matrix. This completes the definition of our model.

ARRAY PROCESSING

For a GD signal, \( \mathbf{R}_0 \circ \mathbf{D}(\theta s) \mathbf{D}(\theta s)' \) is a full rank matrix. As the angular spread of the GD signal enlarges, the eigenvalues of \( \mathbf{R}_0 \circ \mathbf{D}(\theta s) \mathbf{D}(\theta s)' \) diverge (the maximum eigenvalue decreases). For a modest angular spread, conventional DS beamforming array gain is

\[ G_{DS} = (\lambda / d)^2 \]

Now we use matrix filtering to obtain

\[ G_{MF} = \left( \sum_{i=1}^{m} \lambda_i \right)^2 \]

where

\[ \lambda_1, \ldots, \lambda_m \]

\[ \sum_{i=1}^{m} \lambda_i \leq m \]

\[ m \geq \sum_{i=1}^{m} \lambda_i \]

\[ G_{MF} = \frac{m}{\sum_{i=1}^{m} \lambda_i} \]

Or if summing the square of each component of \( \mathbf{Z} \), we then obtain a non-coherent processing gain whose array gain \( G_{MF} \) is maximum

\[ G_{MF} = \left( \sum_{i=1}^{m} \lambda_i \right)^{1/2} \]

Obviously, we cannot use the whiteening state by the matrix filtering. Instead, we must make for spatial compensating state. This is the generalization of the time-compensation used to compensate the decorrelation which is caused
by the relative propagation time differences of the plane wave signal in the homogeneous and time and space-invariant channel. Because the spatial decorrelation is caused by the random wavefront perturbation, we can only compensate in the energy sense by the use of $\mathbf{D}(\theta)$ prior knowledge of $\mathbf{r}_0$, not compensating in the instantaneous waveform sense. We obtain the eigenvector or the signal directional vector $\mathbf{v}(\theta)$ of the compensated result of $\mathbf{r}_0$ by [2]

$$g_s = g_s^{(00)} \circ \mathbf{D}(\theta)$$

(12)

where $g_s^{(00)}$ is the eigenvector corresponding to the maximum eigenvalue of $\mathbf{r}_0$. It means that we use the compensation for seeking the wavefronts planned by the information of the decorrelation matrix. Then scan $\mathbf{D}(\theta)$ by the steering vector $\mathbf{D}(\theta)$ to obtain the energy output of the eigenbeam as shown in Fig.1. This matches to the signal model in the energy sense, whose processing gain is [3]

$$G_{\text{en}} = m.$$  

(13)

Fig. 2 gives the signal-matching performances of the eigen-beamforming:

- **Eigen-decomposition of the covariance matrix of the GD signal**
- **Spatial compensation of wavefront**
- **Energy output of eigenbeam**

Fig. 1 A general eigen-beamforming configuration

The lake-tests for practical GD signal array processing took place in Xing-An-Jian river. The water depth is 60 m. From the surface down to 20 m is a isotherm (the sound speed is 1470 m/s) and below with a small negative gradient. A linear array of 4 equally spaced sensors ($d=0.74$ m) was horizontally placed at 6.3 m depth. The source at 6.3 m depth and 18 m range from the array transmitted a 900 Hz Continuous Wave (CW) signal recorded by our experimental system, whose sampling frequency is $f_s=3$ kHz. Fig. 3 (1) gives its eigenbeam pattern.

During the experiment we collected the motor-generating signal. There are some groups of motors on the lake dam, 500 m far from the array. When working, they yield some kind of distributed noise with certain directivity which is treated as a GD signal. Fig. 3 (2) gives its eigenbeam pattern. Here, we estimate the decorrelation matrix $\mathbf{r}_0$ as following: first estimate the mean $\mathbf{D}(\theta)$ of the GD signal by DS beamforming, then we have

$$\mathbf{r}_0 = [\mathbf{D}(\phi) \circ \mathbf{D}(\theta)] \circ \mathbf{v}.$$  

(14)

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Fig. 2 The signal-matching performances of eigen-beamforming

Fig. 3 Eigens-beamformings of the practical GD signals

(a) DS beamforming
(b) eigen-beamforming
A NEW STRUCTURE OF COMPLIANT WALL FOR THE HYDRODYNAMIC NOISE REDUCTION

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ABSTRACT

A sandwich compliant wall for the hydrodynamic noise reduction is put forward by the author for the first time. Its principle of operation is explained. An experiment procedure in a water tunnel is described, and the measurement results, which indicate that the self noise is reduced remarkably, are given all over the two dimensions of inflow velocity V and cavitation index "sigma".

1. INTRODUCTION

Usually SONAR transducer is far from the propeller, the strong source of noise, and mounted in the vicinity of the stagnation of a underwater vehicle. Consequently, for a high frequency SONAR in which the mounting of the transducer is well uncoupled with the shell of the underwater vehicle, the background noise mainly comes from the hydrodynamic noise of the underwater vehicle.

In order to reduce this kind of self noise, the classical method is to select the shape of the underwater vehicle carefully, and make sure that the length of the laminate flow range is large enough, and the critical cavitation index is smaller as well.

However, since Kramer modelled the skin of dolphin in 1960's[1], people have been engaging in studying how the hydrodynamic noise can be reduced by means of the effects of a compliant wall on the hydrodynamic stability of boundary layer flow. Up to now the interaction between the compliant wall and the boundary layer hasn't been explained thoroughly, but it has been proved experimentally that the compliant wall, "Lamiflo" designed by Kramer for the aim of drag reduction, can reduce the flow noise[2]. Moreover, in many underwater vehicles, the sound transparent rubber of an acoustic window positioned in front of the transducer has been extended along the surface to the rear of the transition point, so as to allow the underwater vehicle to run at higher speed and shallower depths without cavitation.

The author recognizes that there is no absolute stability in boundary layer flow in fact. Then, a new structure of compliant wall has been put forward in this paper. Its principle of operation is different from the distributed damping one of Kramer's, and its advantage has been proved experimentally: It can reduce both the flow noise and the cavitation noise.

2. PRINCIPLE OF SANDWICH RUBBER WALL

The pressure fluctuation in boundary layer is the main source of the flow noise[3]. But it attenuates rapidly in space, and its space correlation radius is very small[4]. The pressure fluctuation applied directly on the transducer can be measured to be a strong disturbance noise, whose equivalent sound pressure is very large.

Usually, the existence of the acoustic window ensures the transducer from being affected by the pressure fluctuation. In this case, what the transducer receives is the flow noise, which is generated at the whole surface of the compliant wall by the interaction between the wall and the pressure fluctuation near the surface of the wall.

In order to reduce this interaction, the author suggests that the pressure release material in underwater acoustic engineering be used as a core between the shell of the underwater vehicle and the outer layer made of sound transparent rubber so as to form a sandwich rubber wall. An ideal pressure release core would make the pressure fluctuation not generate the momentum change on the wall, thus no noise energy would be released. Furthermore, the uncoupling effect of the pressure release core would make the pressure fluctuation not excite the shell to vibrate and to radiate noise.

Moreover, because the cavitation area appears closed to the surface of the wall, the pressure release core would weaken the compression of fluid medium, caused by the collapse of the cavitation bubble, thus reduce the cavitation noise.

3. EXPERIMENTS

3.1 Measure Noise in Water Tunnel

Two models with maximum diameter of 150 millimeters are used to measure their hydrodynamic self noise respectively in a water tunnel, whose diameter in measurement section is 800 millimeters. A hydrophone is installed in the forebody of the model, behind an acoustic window. The two models have exactly the same shape and size, but their compliant wall are different: the one is formed of sound transparent rubber, the other supplements a pressure release core, as shown in Fig.1.

3.2 Selected parameters

Because the model is complex in shape, no any local Reynolds number can be used as a parameter to express the feature of the whole flow field and the self noise data measured correspondingly, except inflow velocity V. Moreover, the flow noise is directly related to the inflow velocity also.
The hydrodynamic stability of inflow medium in its existing state is described by the cavitation index $\sigma$. That stability must be under the influence of the model and the cavitation in the flow medium must take place as well. Therefore, the cavitation index $\sigma$ can be selected as another parameter to express the data of the cavitation noise.

Thus, all the quantitative characteristics of the hydrodynamic noise in various operating condition can be expressed comprehensively in the two dimensions of $V$ and $\sigma$.

3.3 Instruments

A B&K Hydrophone type 8103 is used to receive the self noise. the output of the hydrophone is applied to the B&K Band Pass Filter type 1617 and amplified via the B&K Measuring Amplifier type 2636. The center frequencies of 1617 can be converted automatically by means of a micro-computer. The DC output of 2636 is recorded and printed via an analog–digital conversion.

4. EXPERIMENTAL RESULTS AND DISCUSSIONS

The important results obtained are the characteristics of the apparent self noise levels, as shown in Fig.2 and Fig.3. The inflow velocities in meter per second are marked at both ends of the curves.

The two figures below indicate that the characteristics of the flow noise ($V < \sigma$) don't vary with the cavitation index basically and the characteristics of cavitation noise in the vicinity of the critical cavitation index ($\sigma < 0.4$) go up rapidly with the cavitation index decreasing. There is no an absolute boundary line between the flow noise and the maximum cavitation noise. What appears is a transition condition. The noise in the transition condition is a mixture noise, which contains the flow noise and microscopic cavitation noise. When the microscopic cavitation takes place, no bubbles are visible. The existence of the transition condition expresses a process, in which the sizes of bubbles vary from small to large and the existing state of local water medium varies from separating out the contained air mainly to vaporizing mainly. When the macroscopic cavitation appears, the cavitation should be called the intense cavitation, and the local flow will more or less be destroyed.

The data in Fig.2 are measured with the model of sound transparent rubber wall, and the data in Fig.3 are measured with the model of sandwich rubber wall. It can be seen that at the same cavitation index, the levels of the two self noise are unequal: the latter is much lower than the former. This means that the hydrodynamic noise can apparently be reduced when the sound release core is used. Moreover, the intense cavitation of the latter model doesn't appear until the former model is at super cavitation. This means that the sandwich rubber wall is possessed of the advantages: It can postpone the development of the intense cavitation and enhance the critical velocity for the intense cavitation.

5. CONCLUSIONS

Main conclusions proved are as follows:

a. The sandwich rubber wall put forward by the author can reduce the hydrodynamic noise remarkably.

b. The sandwich rubber wall is effective all over the two dimensions of $V$ and $\sigma$.

c. The sandwich rubber wall can postpone the development of the intense cavitation and enhance the critical velocity for the intense cavitation.

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The wall of water tunnel

sound transparent rubber sound release material

Fig.2 Characteristics of apparent self noise level vs. cavitation index for the former model
Waveform Store Sonar Target Simulator

Lu Chuan Hou Chao-Huan
The Institute of Acoustics
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Abstract:

A novel waveform store sonar target simulator is introduced in this paper. Compared to conventional sonar target simulators, this new simulator offers much improvements in accuracy of target's characteristics, especially in simulating the moving targets, it can also simulate actual effects between underwater channels and signals.

1. Introduction

The application of using sonar target simulators to evaluate the sonar systems and equipments is a more economical and fixable way than conventional field exercises. The target simulators must have the capacity to simulate not only the characteristics of actual targets, but also the effects between underwater channels and signals, but conventional nonactive and active sonar target simulators are just functionally as a noise generator, and the inquiring-answering device respectively. The main defects of this kinds of simulators are that they are bad in simulating the actual reflectance of objects, and even worse in simulating the transmittances of the underwater channels, especially in the multi-channel characteristics.

The new sonar target simulator, we demonstrate here, is based on the principal of store waveform, it can simulate the object's reflectances and effects of the underwater channel's transmittance, it can overcomes the defects which were met in those traditional simulators.

2. Principal

In active sonar system, the reflect signal $Y(t)$ is determined by:

$$ Y(t) = A \ast B \ast X(t-t') $$

(1)

in here, the $X(t)$ is active sonar detect signal, the operator $B$ stands for the object's reflectance, operator $A$ terms the effects of underwater channels on signal, $t'$ is signal transmitting time, operator $\ast$ is convolution. Generally the interaction between underwater channels and signals is very complicated to be described clearly, especially in shallow sea area, where it is time variable.

The principles which are used in our sonar are described as following: When the searching signals by a active sonar are detected by the target simulator, they are received and stored at real time in a computer, called as "Waveform Store Computer or WSC computer". As soon as these signals are processed with the B operator, they are sent out immediately by the target simulator, what the sonar receives is those signals defined as $Y(t)$ in the equation (1). The active sonar process can be simulated as true as actually if the computer can simulate the real reflecting characteristics of objects by digital signal processing. The B operator includes the two effects of the object’s reflecting characteristics: the first is how to simulate the Doppler effect of the moving object, it is expressed here as wave contraction and extension, and can be realized by accurately controlling the speed difference of reading and writing the data in a computer; the second is how to simulate the correlation of object’s reflectance with the size, shape, and structure of the object, as well as the signal incident angles. Theoretically, the reflecting process of an object can be described as an linear system. The object reflecting signal can be stood by the convolution of the incident signal and the pulse response function of the object:

$$ Z(t) = k(t-t') \ast X(t) $$

(2)

Because the limited experiment conditions in the sea, it is usually very difficult to measure the pulse response functions of ships in practical environment. Although it can be measured by a scalable model in the experiment environment, the results are usually not satisfied yet due to lack of simulating the actual structure of ships itself. In some certain fields, the reflecting response functions can be simplified and even meet some requirements well. For example, to the far distance object, its reflecting response function can be defined as a constant number, the object reflecting density; to the near distance object for some sonars, it can be defined with the infinite numbers of reflecting densities. By the use of this simple reflecting response function models, the amount of computing (2) is dramatically reduced, and today microcomputers stand this easily. We have proved that the simplified object re-
reflecting response function model is applicable and yields good results.

3. Configuration of System Structure

The waveform store computer is the core part in the target simulator, it is a special high speed microprocessor.

As shown in Fig.2, the data output is simple, one of DMA reads data sequentially and output them to D/A stage. As shown in Fig.3, two DMA devices are used here, one reads out the wave data from high address area, after it has finished this block data, the second is toggled and the wave data in low address area are sent out. If there are more main data memories than the amount of wave form data, the ring store structure can hold them completely.

The signal identifying and controlling devices, shown in Fig.1, checks the rise and fall edges to control the WSC processor reading and writing; the Doppler generator can precisely control the output data, by using the frequency synthesis method, to accurately simulate the moving object’s Doppler effect. The center computer organizes and manages the all parts of the system. The target simulator can simulate the moving object dynamically only through changing the parameters of Doppler and in h(t) of (2).

4. Conclusion

We have discussed the basic principles and performances of the Wave Form Store Target Simulator. The practical system has been made with the Z80 CPU, and has been tested and shown very good performances in the practical under sea experiments. The second generation system with much power is underdeveloping.

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