ULTRASONICS, QUANTUM ACOUSTICS AND PHYSICAL EFFECTS OF SOUND
Negative Poisson's ratio: Isotropic Solids, Crystals.

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Materials with negative Poisson's ratio are peculiar in that they expand laterally when stretched. The stereographic projections of Poisson's ratio have been calculated for a set of cubic metals, alloys, binary compounds, acoustic crystals. From these stereographic projections the orientations of the stretch and lateral strain with extremes values of Poisson's ratio have been obtained. Angular dependence of stretch directions and corresponding to negative values of Poisson's ratio were calculated. The exotic behaviour of change of cross section of cylindrical rod during the stretching along the rod axis was revealed in cristobalite SiO₂ - the only known form of isotropic material with negative Poisson' ratio [1]. Some possible applications of materials with a negative Poisson's ratio basing on their unusual acoustic properties are discussed.

Poisson's ratio is the ratio of lateral strain in the direction to the longitudinal strain, during the stretching of a cylindrical rod along rod axis. The Poisson's ratio is a constant in an isotropic solid. In the theory of elasticity Poisson's ratio of isotropic solids is known to lie in the range of -1 and 1/2. Most isotropic solids have Poisson's ratio in the vicinity of 1/3. Materials with negative Poisson's ratio are unusual. Negative Poisson's ratio in materials is governed by the following aspects of the microstructure: the presence of rotational degrees of freedom, non-affine deformation kinematics, or anisotropic structure. Negative Poisson's ratio materials are of interest because of their unusual properties. Solids with negative Poisson's ratio easily undergo volume changes, but difficult to relation to shear changes, so G>>K , a situation opposite that seen in rubber. Negative Poisson's ratio materials expands laterally when stretched, in contrast to ordinary materials, and become fatter in cross-section.

Negative Poisson's ratio was discovered in foams, polymers, chiral structures, laminated structures, crystals. The original foam with Poisson's ratio as small as -0.7 was first developed by Rod Lakes [2]. Poisson's ratio as low as -4 have been obtained in auxetic polymers [3]. The laminted structures with negative Poisson's ratio was developed by Milton [4]. The existence of negative Poisson's ratio in some cubic metals were predicted for a stretch along a cube-face diagonal in [5]. The auxetic behaviour for a cubic materials with negative Poisson's ratio have been investigated in details in [6]. The only known isotropic negative Poisson's ratio material is an aggregate of crystals alpha-cristobalite [1]. The Poisson's ratio of an anisotropic solid depends on directions m and n and is

\[
V_{nn} = \frac{\varepsilon_n}{\varepsilon_m} = - \frac{m_n m_j S_{ijkl} n_i n_j}{m_m m_j S_{ijkl} m_i m_j},
\]

where \(S_{ijkl}\) are the elastic compliances coefficients, \(\varepsilon_m, \varepsilon_n\) - longitudinal and lateral strains, \(n, m\) - ortogonal unit vectors,

\[
\varphi(m) = \frac{\Delta S(m) / S(\bar{m})}{\varepsilon_n} = - \frac{s_{ijkl} m_j (\delta_{ij} - m_i m_j)}{2 s_{ijkl} m_j m_i m_j},
\]

where \(\Delta S(m) / S(\bar{m})\) is the cross-sectional area change of the rod.

We calculated the stereographic projections of \(V_{nm}\) for any possible directions of vectors m and n. From these stereographic projections the directions with the extreme values of \(V_{nm}\) have been obtained for a set of crystals. In our calculations we used the elastic module of crystals from ref. [7]. The results are presented in Table 1. For most metals the Poisson's ratio exhibits the negative values for the directions inside the cone around [110], corresponding the Euler angle \((45,0,90)\). In Table 1 angles \(\Delta \alpha, \Delta \beta\) are related to the directions of m and n with Euler angles \((45 \pm \alpha, 0, 90)\) and \((45, \pm \beta, 90)\) for \(V_{nm} = 0\).

A decrease of square rod cross section under the uniaxial compression was found in beta cristobalite for the directions of stretch vector m oblique to [111]. Stereographic projections of \(\varphi\) in beta
Figure. 1  Relation of Euler angles to crystallographic directions.

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Table 1.

Figure. 2  The stereographic projection of $\tilde{V}$ in beta-cristobalite, x-axis is [100], y-axis is [010], axis [001] ortogonal to this paper.

Figure 2 presented in Fig 2. The crystal exhibits cubic symmetry and elastic stiffness constants of eta cristobalite are $c_{11}=42.0$, $c_{12}=5.8$, $c_{44}=40.8$ [9].

Materials with negative Poisson's ratio materials may find a lot of applications based on the Poisson's ratio itself, or based upon it unusual physical properties. The Poisson's ratio influences the transmission and reflection of bulk acoustic waves at a free surface, the propagation of surface acoustic waves, the decay of acoustic waves, the distribution of stress field around holes and cracks. Such materials also find wide applications in sandwich panels with good sound absorbing properties, as air filters, fasteners, sandwich panels for aircraft and automobiles. Metals with negative Poisson's ratio may be used as electrodes that amplify the response of piezoelectric sensors [6].

REFERENCES

Visualisation of Elastic Heterogeneity of Diamond-Fulletrite Composites

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The interior of small hard specimens sintered from diamond-fullerite powder blend is investigated by low frequency time-of-flight acoustic microscope. B-scan mode was used for obtaining the transverse sections of the objects. B-scan images reveal the significant changes of the echo’s arrival time along the scanning line for both the longitudinal and shear wave. These time variations are related with spatial distribution of elastic properties of the composites.

INTRODUCTION

New allotropic forms of carbon created by high-pressure high-temperature (HPHT) treatment of solid C\textsubscript{60} reveal the great diversity of their mechanical structure depending on the treatment conditions from soft crystalline to superhard amorphous [1-2]. Adding of different components to the base C\textsubscript{60} can lead to the creation of the new carbon materials with interesting elastic properties. The ongoing HPHT-sintered specimens are small sized (1÷3 mm) and sometimes with the inner cracks arising due to stress relaxation during cooling. Their interior is therefore of great interest.

The acoustic microscopy is known as the unique method for study of the small opaque specimens [3-4]. It allows not only to "screen" the sample and herewith to control the presence of defects but also to follow the changing the local elastic properties [5]. Since the elastic properties are specified by the intermolecular bonding and nanostructural organisation of a matter their determination can provide the designation of specimen’s microstructure.

Our report is devoted to the study of elastic properties of the C\textsubscript{60}-diamond composites.

MATERIALS AND METHODS

The new hard carbon phases were sintered from the powder blend consisted of pure fcc fullerite powder (C\textsubscript{60}: 99.98%, C\textsubscript{70}: 0.01%) and diamond powder with diamond mass concentrations of 75 and 50%. The diamond powder included the particle with sizes from 0.003 µm to 5 µm.

Concentration of volatile admixtures was about 1% mass. The sintering was carried out at the pressure of 9.5 GPa and temperature up to 1270 K in a high-pressure tungsten-carbide chamber [6,7]. Application of such treatment to the pure fcc C\textsubscript{60} powder leads usually to the creation of amorphous 3D polymerised ultrahard sample with sound velocity comparable with ones for diamond [4].

For ultrasonic study the specimens (22 mm × 2 mm) were placed into a steel cartridge clip; the front and back faces were ground. The steel cartridge having the same thickness as the samples was then used as a reference material.

B-scan mode of acoustic microscope with operation frequency \( f = 50 \text{ MHz} \) and pulse duration 30÷70 ns was used for the investigation of the elastic heterogeneity of the specimens. The lens with the aperture \( \theta \approx 14^\circ \) created in water an ultrasonic beam with waist radius of 100÷120 µm which penetrated into the solid specimen mainly as the beam of longitudinal wave. The shear waves arise mainly due to mode conversion on the specimen’s bottom. B-scans were formed under 1-dimensional scanning and modulating the intensity of the display beam by the amplitude of returned echoes. The abscissa of the image represents the position of the point investigated, the ordinate - the propagation time of the detected echoes. So, the echo-signals formed the image of a transverse section of an object along a scanning direction. B-scans bring the information about the irregularity of the sample’s surfaces, the presence of internal interfaces and holes as well as about the elastic property variations since they all lead to the changes of the propagation time of the sound.
RESULTS AND DISCUSSION

B-scan images of sintered fullerite-diamond composites are presented on the fig. 1.

![B-scan images of HPHT-composites with diamond mass concentration of 75% (a) and 50% (b)](image)

Figure 1. B-scan images of HPHT-composites with diamond mass concentration of 75% (a) and 50% (b)

The echoes from the front face of the sample and steel give the upper straight bright line. The echoes of longitudinal and shear waves from the bottom of steel holder are depicted by down straight lines on each side of B-scans. The curved middle (M) and lower (L) bright lines in the central part of the images are formed by the echoes from the specimen.

Since the treatment condition were correspondent to the synthesis of very hard samples the M- and L-lines could be considered as the bottom echo-lines of longitudinal and shear waves, respectively. The main objections arise: how to explain the incredible big difference in velocity magnitudes for longitudinal and shear waves (the comparison with steel gives the mean values \( C_L = 20 \) km/s and \( C_S = 3.5 \) km/s, approximately) and how to specify notable distinction of their velocity gradients? These objections do not allow us to accept this explanation.

The most likely interpretation of obtained B-scans is the next: L-lines are formed by longitudinal wave echoes with mean value of \( C_L \) comparable with one for steel. M-lines on each B-scan correspond to the reflection from the internal boundary of different phases. The bottom reflections are stronger in the regions where the sound was not weakened by the internal face (fig 1,b). The bending of M-lines is mostly induced by velocity variations in the layer beyond the interface. Only a slight curvature of the internal face can be registered due to deflection of sound beam. The change of echo-signal polarity revealed the reflection from the softer phase. The bending of L-lines is caused by the velocity gradient of the whole volume. The bigger roughness of the surfaces led to the greater mode conversion, therefore the image includes the weak lines formed by shear wave echoes (fig.1b). The evaluated average value of shear wave velocities \( C_T \) are about \( 4 \) km/s. The comparison of velocity values \( C_L \) and \( C_T \) with the date presented in [4] indicates that the sintered composites consist of various 1D- and 2D-polymerised phases.

The elastic heterogeneity results from the non-uniform distribution of pressure and temperature during the synthesis. The velocity increase in the sample periphery correlates well with temperature gradient in the toroidal chamber where the heat came through the cylindrical wall of [8]. The same hardening of periphery part was observed for the samples sintered from the pure \( C_{60} \) [5].

CONCLUSION

Presented here B-scans demonstrate the possibility of checking and analysing the internal structure investigation by acoustic microscopy. The images show different kinds of heterogeneity including the coexistence the various metastable phases and a gradual variation of elastic properties over the specimens.

ACKNOWLEDGEMENTS

This work was supported by the Slovak Grant Agency VEGA (project No. 1/6072/99).

REFERENCES

Acoustic and Piezoelectric Properties of New Polar Compounds of Sn$_2$P$_2$S$_6$ Family

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The ultrasonic and piezoelectric investigations of Sn$_2$P$_2$(S,Se)$_6$ and CuInP$_2$S$_6$ ferroelectric - semiconductors single crystals have been performed. In these materials the ultrasonic anomalies at ferroelectric phase transitions (PT) were observed. Below phase transition point the comparatively strong piezoeffect was observed in these compounds. The variety of photosensitive ultrasonic phenomena such as acoustoelectric interaction, acoustoelectric voltage, relaxing ultrasonic attenuation oscillations and the photostimulated shift of phase transitions to lower temperature were observed and investigated in all these polar crystals of Sn$_2$P$_2$S$_6$ family.

INTRODUCTION

The crystals of uniaxial ferroelectric Sn$_2$P$_2$S$_6$ family are very interesting for several reasons in the context of applications and basic research [1,2]. One reason is that these crystals exhibit strong piezoelectric effect and are photosensitive semiconductors. This offers a possibility for applications in acoustoelectronic and electroacoustic devices; consequently, investigations of acoustoelectric interaction are necessary. Another reason is that the sulphur substitution by selenium in a material leads to appearance of the incommensurate phase and the Lifshitz point providing unique possibility to investigate the critical behaviour of the phase diagram near such point [3,4]. Recently it was shown, that the mobile charge carriers also play essential role in the critical behaviour of crystals [5]. The variation of the electron density in energy levels can change the Lifshitz point position in the phase diagram and strongly modifies thermodynamic characteristics, such as optic and dielectric properties. Investigations of elastic and piezoelectric properties near Lifshitz point are also interesting. Recently the new layered crystals of this family were obtained: i.e. CuInP$_2$S$_6$, which also are promising materials for microelectronics. Here we present the investigation of ultrasonic, piezoelectric and acoustoelectric properties of these Sn$_2$P$_2$S$_6$ family crystals. The ultrasonic and piezoelectric measurements were carried out by pulse-echo and resonance methods.

RESULTS AND DISCUSSION

The phase diagram of Sn$_2$P$_2$(S,Se)$_6$ crystals is complicated. Pure Sn$_2$P$_2$S$_6$ exhibits the second order ferroelectric PT near 337 K. At this PT the large critical slowing down in ultrasonic velocity is observed for all longitudinal ultrasonic modes. With substitution S by Se the PT temperature decreases and velocity anomalies increases (Fig. 1). In this figure the temperature dependencies of ultrasonic velocity measured at 10 MHz are shown for compounds with Se content $x=0$ (1), $x=0.15$ (2) and $x=0.28$ (3). The velocity anomalies at PT are accompanied by large and narrow ultrasonic attenuation peaks. For more detail discussion of anomalous ultrasonic attenuation, we refer to original papers [6,7]. With increase Se content above $x=0.28$ the PT splits to two and intermediate incommensurate (IC) phase appears. Therefore at $x_L=0.28$ and $T_L=283$ K the Lifshitz point exists in Sn$_2$P$_2$(S,Se)$_6$ system [1]. With further increase of $x$ the temperature interval of existence of IC phase increases.
and for pure Sn$_2$P$_2$Se$_6$, it reaches $\Delta T_{IC}=30$ K (Fig. 2). The low temperature transition from IC phase to commensurate ferroelectric phase is the first order PT and the thermal hysteresis of about 1.5 K appears in heating and cooling measurement cycles.

In the ferroelectric phase of Sn$_2$P$_2$(Se$_{1-x}$S$_x$)$_6$ crystals the photosensitive ultrasonic attenuation caused by acoustoelectric interaction with free charge carriers was found for piezoeactive modes. The dependencies of electronic attenuation coefficient $\alpha_{el}$ on conductivity $\sigma$ of Sn$_2$P$_2$S$_6$ and Sn$_2$P$_2$Se$_6$, crystals in ferroelectric phase. With increase of Se amount in the compound electromechanical coupling decreases and in pure Sn$_2$P$_2$Se$_6$ sample: $K_{33}=0.4$.

In CuInP$_2$S$_6$ single crystals the acoustoelectric interaction was very weak, therefore we performed piezoelectric measurements using resonance method. At room temperature for longitudinal vibrations along the c-axis the value $K_{33}=0.3$ for coefficient of electromechanical coupling was obtained. The ultrasonic velocity and attenuation anomalies were recorded in layered CuInP$_2$S$_6$ samples at $T_c=310$ K. Above $T_c$ the pretransitional effects in ultrasonic and piezoelectric properties were established. Such effects can be determined by polar clusters existing in paraelectric phase.

The other photosensitive ultrasonic phenomena such as acoustoelectric voltage, relaxing ultrasonic oscillations and the photostimulated shift of PT temperature were observed and investigated in all these polar crystals of Sn$_2$P$_2$S$_6$ family. Illumination induces the intermediate-IC phase not only in the compound Sn$_2$P$_2$(Se$_{0.28}$S$_{0.72}$)$_6$ where the Lifshitz point exists, but also in Sn$_2$P$_2$S$_6$ crystal. This follows from comparison of ultrasonic behaviour in the crystals under various illumination conditions near PT. All investigated materials have strong piezoelectric effect, nevertheless they are wide band gap semiconductors and charge carriers in deep energy levels create space charge, which modifies the elastic, dielectric and, as a result, electroacoustic parameters. This manifests itself in long time relaxation of these properties, especially near PT and in IC phase. So, for applications it is important to verify the conditions for the optimisation of electroacoustic parameters.

ACKNOWLEDGEMENTS

The support of EU under INCO-COPERNICUS-CT-97-0712 project and Lithuanian State Science and Studies Foundation is gratefully acknowledged.

REFERENCES

Evaluation of the Bindings between the Superconductor and Matrix by Ultrasonic Spectrum of AE when S-N Transition


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In the practical use of superconducting magnet, combination of the superconductor and high conductive normal metal is now indispensable. And this binding property is one of the key points whether superconductive characteristics is fully attained or its current is failed prematurely in the practical magnet. But these binding evaluations are not so easy, because even in the mechanically well-jointed conductor, it often appears poor in electric and magnetic joint in cryogenic temperature. Here we report new technique for evaluating these bindings by observing ultrasonic spectra when a normal transition is propagated along the superconductor. We have discovered a fact that good contacting conductor has an emission of this sound with sharp resonating frequency peak or peaks at around highest spectrum area; from 0.5 to 1.0MHz.

BACKGROUND

As mentioned above, even in the mechanically well-jointed conductor in room temperature, it often appears poor both in mechanic and/or electro- magnetic joint in cryogenic temperature because of so much difference of contraction between the superconductor and its matrix. And these conductor must encounter with electro-magnetic force and thermal expansion and back to re-contraction after each super-to-normal quench. So, It must be needed to examine this binding evaluation in the cryogenic in situ superconducting condition.

HOW TO EXAMINE THE BINDINGS

As like a very tight molded bell would sound a longer, higher and clear tone., the construction with superconductor and its matrix was hit with electric thermal pulse enough to initiate a super-to-normal(S-N) transition along the superconductor. Then we can obtain AE from the conductor and its ultrasonic spectrum for each transition.

CONSTRUCTION OF THE EXPERIMENT

The experimental superconductor with high purity aluminum matrix is shown in Fig.1 ,which is situated in the magnetic field of 3Tesla. Four PZT sensors were attached to the conductor for every quarter of the SC ring . We prepared above mentioned heater which made the initiation and propagation of the normal state, then each sensor sent AE signal when the propagation took place.

WHAT WE CAN GET

With the emission of above ultrasonic sound, we have discovered a fact that good contacting conductor has an emission of this sound with sharp resonating frequency peak or peaks at around highest FFT spectrum area; from 0.5 to 1.0 MHz.

WATCH AT HIGHEST PEAKS

Our discrimination of the evaluations are made by the sharpness of highest peak or peaks, more exactly saying, by the simple that we can say binding degrees as best, good or bad, according to the resonating Q-value at around highest spectrum of the sound.

EXPERIMENTAL RESULTS

Fig.2 shows the highest AE FFTspectrum ih the case of no S-N propergation happend. it means heat pulse related AE only observed within 0.5MHz. Fig.3 indicates the case of best bindings between the superconductor and matrix, showing a sharp peak at 723kHz. Fig.4 indicates the case of not the best one, showing several small peaks at around 550~750kHz. In the worst case, no peak or peaks observed at all in above area and like Fig.2.

DISCUSSION

For understanding these experiments, let us make simplified equivalent circuit like in Fig.5, where Rm is the contact resistivity between the superconductor and the matrix, and is the most important parameter here for the evaluation of the bindings, and R; matrix's
resistivity, and, Rs ; resistivity of the superconductor on and off at the critical states. With a heat pulse and S-N transition and propagation initiated at the left end in the Fig.5, Rs appear on the superconductor to a certain length, then, super current is bypassed to the normal metal through Rm, which creates another heat then proceed these process again in the next moment, vibrating electric current. In the superconducting construction, electric current vibration generates stress vibration directly because electro-magnetic Lorentz force conjugated directly to the physical force. So, in this vibration, Rm plays an important role. Sharpness of the AE spectrum is tightly represent Rm best uniformity which is reasonably explained by vibrational resonating theory.

CONCLUSION

We are going to evaluate the binding properties by the sharpness of the ultrasonic highest spectrum, because that represent for the degree of the uniformity of the bindings. In other words, we should be determined by the degree of the standard deviation of the averaged value weighted with each frequency.

**FIGURE 1.** The superconducting construction.

**FIGURE 2.** Highest AE spectrum, when no quench.

**FIGURE 3.** Highest AE spectrum, when best windings.

**FIGURE 4.** Highest AE spectrum, when not best bindings.

**FIGURE 5.** Simplified equivalent circuit of the system.
Acoustoelectric Effects in Manganites


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We report on the investigation of the acoustoelectric effect in La$_{0.67}$Ca$_{0.33}$MnO$_3$ (LCMO) films tightly bound to piezoelectric substrate. The main result is the independent confirmation of hole-like conductivity in the LCMO films and the observation of an anomalous acoustoelectric effect, which contrary to the known ordinary acoustoelectric effect is invariant under the reversal of the direction of propagation of the surface acoustic wave. We show that the anomalous effect is related to the strong dependence of film conductivity on pressure and can be expected to occur in other conducting materials with similar pressure dependence of transport properties.

Acoustoelectric technique has been proven to be useful for studying transport properties of thin films and layered structures. In this paper we present the investigations of the longitudinal acoustoelectric (AE) effect induced by surface acoustic waves (SAW) in manganite thin films, which exhibit colossal magnetoresistance effect.

The investigated La$_{0.67}$Ca$_{0.33}$MnO$_3$ (LCMO) films, 100 to 200 nm thick, were grown on +y-cut of LiNbO$_3$ (LNO) substrates by laser ablation technique and subsequently patterned into a Hall bar structure. X-ray diffraction shows that the films are single phase, epitaxial and (211) oriented with the pseudocubic lattice parameter $a_0 = 0.3853$ nm. Chemical composition of the films was checked by electron probe microanalysis.

The SAW, which was launched in $z$ direction of LNO substrate, was generated and detected by two interdigital transducers (IDT), positioned at the edges of the substrate with the LCMO film deposited centrally between them (see inset of Fig. 1). AE effect has been studied in the short-circuit geometry at 87 MHz.

Figure 1 shows the resistivity $\rho(T)$ of the investigated LCMO films. An applied magnetic field shifts the $\rho(T)$ peak towards higher temperatures and reduces the film resistivity, resulting in the magnetoresistance effect: MR $= (\rho(H) - \rho(0))/\rho(0) \sim -80\%$ at 25.5 kOe.

Results of the acoustoelectric effect measurements in the LCMO films are shown in Fig. 2. The AE current $I_{AE}$ is small ($\sim 1 \mu A$ for the SAW intensity $\Phi \sim 3 W/cm$) at high and low temperatures and approaches its maximum value of about 25 $\mu A$ near the peak of resistivity.

The acoustoelectric drag current should be odd in the SAW wave vector $q$.[1] But the AE current measured with reversed direction of SAW propagation exhibited a puzzling temperature behavior seen in Fig. 2. Such an unexpected dependence of $I_{AE}$ on the direction of $q$ can be explained by assuming two contributions to the measured AE current: $I_{AE} = I_{even} + I_{odd}$. The first term is anomalous and even in $q$, while the second one is ordinary and odd: $I_{even}(-q) = I_{even}(q)$, $I_{odd}(-q) = -I_{odd}(q)$. The results for the ordinary and anomalous contributions to the AE current for $q \parallel +z$ axis of the LNO substrate are displayed in Fig. 2. Detailed analysis shows that even AE effect dominates near the peak of resistivity and its magnitude exceeds the odd AE effect approximately twice. It should be remarked that the even $I_{AE}$ is directed along the $+z$ axis. The odd AE effect prevails at high and low temperatures, and its sign corresponds to the hole-like conductivity in the whole investigated temperature range.

In order to analyze the acoustoelectric interaction we will take into account that charge carriers in the LCMO film are affected both the strain $S_{ij}(y,z,t)$ and electric field $E(y,z,t)$ which accompany the acoustic wave in the piezoelectric substrate. The SAW induces a local modulation of the film conductivity $\sigma_{zz}$: $\sigma_{zz}(y,z,t) = \sigma_0 + \sigma_1(y,z,t) + \sigma_2(y,z,t)$, where $\sigma_0$ is the unperturbed conductivity, $\sigma_1(y,z,t) = \delta n(y,z,t) \partial \sigma_0/\partial n$ is due to the
modulation of charge carrier concentration \( n = n_0 + \delta n \) by the piezoelectric field \( \delta n \ll n_0 \), and the last term \( \sigma_2(y,z,t) \approx \sigma_0(\Pi_{3333}S_{zz} + \Pi_{3322}S_{zy}) \), where \( \Pi_{ijkl} \) denotes the modulation of \( \sigma_{zz} \) by the strain and was not considered in earlier treatments. \( E(y,z) \) generates in the film a local current density:

\[
J_z(y,z,t) = \sigma_{zz}(y,z,t)E_z(y,z,t) \equiv \sigma_z^{(0)} + J_z^{(1)} + J_z^{(2)}. \tag{1}
\]

The dc component of the local current \( J_z^{(1)} \) may be determined applying the Ingebrigtsen’s approach [2]:

\[
j_{ae}^{(1)} = q\Gamma\Phi\sigma_0(\epsilon\omega\eta_0)^{-1}, \tag{2}
\]

where \( \Gamma \) is the SAW dissipation, \( \omega \) is frequency, and \( e \) is the charge of charge carriers. \( j_{ae}^{(1)} \) is odd in \( q \) and represents the ordinary longitudinal AE current [1].

In order to calculate the dc component of the local current \( J_z^{(2)} \) due to the film deformation, it is necessary to solve Maxwell’s equations and mechanical-piezoelectric equations of state, taking into account the boundary conditions at the film surfaces. As a result we obtain

\[
j_{ae}^{(2)} = -\Gamma\Phi\epsilon_{33}(\Pi_{3333} - \nu\Pi_{3322})/p_{33}, \tag{3}
\]

where \( \epsilon_{33} \) and \( p_{33} \) are components of the dielectric and piezoelectric tensor, respectively, and \( \nu = c_{12}/c_{11} \approx 0.4 \) [3], where \( c_{11} \) and \( c_{12} \) are the components of the elastic tensor. This result clearly demonstrates that \( j_{ae}^{(2)} \) is even in \( q \), and its direction is determined by the signs of \( p_{33} \) and \( \Pi_{3333}/\nu \). The coefficients \( \Pi_{3333} \) are related to the pressure dependence of the conductivity: \( 2\Pi_{3322} + \Pi_{3333} = -3\kappa^{-1}\partial\ln\sigma_0/\partial P \), where \( \kappa \) is the film compressibility. According to the pressure experiments, in LCMO [4] the quantity \( \partial\ln\sigma_0/\partial P \) is positive, which means negative \( \Pi_{3333} \), and has a pronounced temperature dependence: it is small at high and low temperatures but peaks to about 3.5 [GPa]\(^{-1} \) at a temperature slightly lower than that of the resistivity peak. For numerical estimation we will approximate \( \Pi_{3333} \approx \Pi_{3322} \). Using \( \kappa^{-1} = 85 \) GPa [5] we find \( \max\Pi_{3333} \approx -300 \).

The above outlined theory is in a good agreement with the experimental data. As the coordinate \( z \) axis is chosen to be parallel to the \(+z\) crystallographic axis of LNO substrate, the constant \( p_{33} \) is positive [6] and, according to Eq. (3), \( j_{ae}^{(2)} \) flows along \(+z\) axis. Substituting the well known parameters of LNO and our own experimental results \( \max\Gamma \sim 2 \text{ cm}^{-1} \) into Eq. (3), we find \( \max j_{ae}^{(2)} \sim 30 \mu\text{A/cm} \) for \( \Phi \sim 3 \text{ W/cm} \). The experimental data in Fig. 2 correspond to \( \max j_{ae}^{(2)} \sim 100 \mu\text{A/cm} \). Taking into account the approximate character of our estimations of \( \Gamma \), \( \Phi \) and \( \Pi_{3333}/\nu \) we consider the agreement between the theory and experiment as satisfactory. Moreover, Eq. (3) describes very well the temperature behavior of \( j_{ae}^{(2)} \), also shown in Fig. 2. The theoretical curve in Fig. 2, normalized to the experimental maximum, was calculated from Eq. (3) using experimental dependencies \( \Gamma(T) \) (Fig. 1) and \( \partial\ln\sigma_0/\partial P \) from pressure measurements [4]. The observed maximum of \( j_{ae}^{(2)} \) is shifted towards lower temperatures by \( \sim 15 \text{ K} \) with respect to the theoretical prediction. This shift may be caused by a possible difference in temperature of the peak \( \partial\ln\sigma_0/\partial P \) between our LCMO films and bulk samples. It should be noted also that numerical estimation of Eq. (2) leads to the same order of magnitude of the odd AE current \( j_{ae}^{(1)} \) as the experiment.

**ACKNOWLEDGMENTS**

This work was supported in part by KBN Grant 2 P03B 139 18 and PBZ-KBN-013/T08/19, and the RFBR Grants 99-02-18333 and 01-02-17479.

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Hydrogen Detection with a Two-layered Surface Acoustic Wave Sensor

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Results concerning a hydrogen sensor based on a multilayer structure in a Surface Acoustic Wave dual-delay line system are presented. The sensor material consists of two layers performed in two different vapor deposition processes. The first one is a 720 nm CuPc layer, the other one – a 20 nm thin palladium film. This structure was formed in one of the dual – delay line system on a LiNbO3 Y-cut Z-propagation substrate. In such a multilayer structure can detect hydrogen in a medium concentration range (from 0.25 % to 3 % in nitrogen), even at room temperature. The sensor has a very good sensitivity, stability and is entirely reversible. The response and recovery times are very short (~200 s – 800 s), which is very important from the practical point of view.

**INTRODUCTION**

There are many methods for hydrogen detection. The first SAW devices was made by D’Amico [1]. Almost all the sensors referred above use thin palladium films as sensor material because of their well-known absorption properties relative to hydrogen. According to D’Amico [1] we have tried to use thin palladium layers (20 nm) in a SAW system to detect hydrogen. Unfortunately, the interaction was very weak, even at a relatively high (3%) hydrogen concentration. When we used only a thin copper phthalocyanine (CuPc) film (720 nm), the interaction was much better, but only at higher temperatures of about 70°C. We decided to form a multilayer structure consisting of a CuPc layer 720 nm and a 20 nm thin palladium layer in two different vacuum deposition processes - Fig. 1.

The addition of a thin palladium layer to the CuPc film greatly enhanced the sensitivity to gaseous hydrogen in a medium concentration range (from 0.25 % to 3 % in nitrogen), even at room temperature. Besides that, the sensor has a good stability, a short response, short recovery times and is entirely reversible.

**EXPERIMENTAL**

The experimental set-up is based on frequency changes in a surface acoustic wave dual delay line system, which is nowadays well known [1,2].

Response to the particular gas of the active multilayer is detected as a change in the differential frequency \( \Delta f \), i.e. the difference between the two oscillator frequencies \( f \) and \( f_0 \). A change in the mass density of the multilayer, and a change in its electrical conductivity cause significant changes in the velocity and attenuation of the SAW and consequently the changes in the frequency \( f \) and \( \Delta f \) [3].

**RESULTS**

Results for Pd and CuPc films

At first, accordingly to D’Amico [1], we used a palladium (Pd) layer as a sensitive material in experimental results. A palladium film absorbs easily hydrogen molecules and is well known material for detection of hydrogen. Hydrogen absorption cause changes in the density, elastic properties and the
conductivity of the Pd film [4]. The obtained result was unfortunately very weak - Fig. 2. At higher temperatures there were no better results.

![FIGURE 2](image1)

**FIGURE 2.** Changes of the differential frequency, $\Delta f$, versus time for two concentrations of hydrogen for 20 nm of palladium at 30 °C.

When a 720nm CuPc layer was used the sensor response was about 600 Hz, but only at higher temperatures (70 °C). The result is shown in Fig.3. The investigated CuPc layer with a thickness of about 720 nm was made by means of the vacuum sublimation method.

**FIGURE 3.** Changes of the $\Delta f$ frequency for 720 nm CuPc film for 3 % of hydrogen in nitrogen at 70 °C.

Results for a multilayer structure

For the multilayer structure (Fig.1) the obtained preliminary results are very promising. We can observe recurrent changes of the differential frequency of the measurements system in the function of the medium hydrogen concentration in nitrogen (Fig. 4).

![FIGURE 4](image2)

**FIGURE 4.** Changes of the differential frequency $\Delta f$ versus time for a multilayered sensor structure at 30 °C and six different concentrations of hydrogen in nitrogen.

The saturation level, i.e. the equilibrium between adsorption and desorption, is reached very quickly especially at lower hydrogen concentrations. Besides, it is interesting to note the great change in the interaction between 1.5 % and 2 % concentrations. This interesting interaction “jump” is completely reversible and is probably connected with surface and volume interactions and the phase transitions for the palladium film.

ACKNOWLEDGMENTS

This work was sponsored by Polish State Committee for Scientific Research within the grant Nr 8T10C 032 18.

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New Properties of SAW Gas Sensing

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An analytical expression for the SAW sensor gas response versus gas concentration (calibration curve) is deduced, at low gas concentrations. It is shown that the value of the SAW response as well as the calibration curve, the sensitivity (its slope) and the detection limit (threshold concentration) for a given gas and sorbent membrane, can be controlled by a proper choice of the SAW propagation on anisotropic piezoelectric substrate, allowing to change the partial contributions to the SAW response through mechanical displacements, coupling constant and temperature coefficient of velocity (TCV). The theoretical predictions are experimentally verified on the integrated array of SAW gas sensors implemented on ST-quartz, using Pd as a sorbent film and humid air as test gas.

INTRODUCTION

The analysis of the fundamental properties of gas sensing is of great importance in many applied areas. Among the different modern techniques used in these areas, SAW sensing is one of the most promising [1]. In this paper we present studies on the properties of the SAW response, calibration curve, sensitivity and threshold value in SAW sensors. Our measurements have been performed using integrated structure with different propagation directions along the surface of ST-cut quartz, as single crystal piezoelectric substrate; Pd, 240 nm thick, as sorbent membrane; humid air as test gas.

Adsortion of gas molecules on a substrate, coated with a sorbent film, gives rise to changes in the film properties and, thereby, in the SAW propagation velocity. The response of a SAW device is estimated to be [2]:

\[
\frac{\Delta v}{v_o} = \left(\frac{\pi h}{2\lambda}\right) \left[ \frac{\Delta \rho}{\rho} + \frac{\Delta \varepsilon_{44}}{\varepsilon_{44}} \cdot B + \frac{1 - \varepsilon_{44}}{\varepsilon_{11}} \right] - 1 \cdot D - K^2 \cdot \frac{\Delta \sigma}{\sigma} \cdot \left( \frac{\sigma^2}{\varepsilon^2 C_S^2} + 1 \right)^2 + (TCV) \cdot \Delta T
\]

where \( h, \rho, \varepsilon_{11}, \varepsilon_{44}, \sigma, T \) are the thickness, mass density, elastic constants, sheet conductivity, and temperature of the film, respectively, while \( \Delta \rho, \Delta \varepsilon_{11}, \Delta \varepsilon_{44}, \Delta \sigma, \Delta T \) represent their changes produced by the gas adsorption. \( K^2, C_s \) and (TCV) are the coupling constant, capacitance per unit length and temperature coefficient of velocity in the substrate, respectively. Coefficients A, B and D are:

\[
A = \{A_{\varepsilon_s}^2 + A_{\varepsilon_s}^1 + A_{\varepsilon_s}^3\} \cdot p \cdot v^2,
\]

\[
B = \{4A_{\varepsilon_s}^2 + A_{\varepsilon_s}^3\} \cdot c_{44},
\]

\[
D = \{4A_4^2\} \cdot c_{44}^2 / c_{11},
\]

where \( A_s, A_4, A_p \) are the normalized displacements of the acoustic wave at the surface, being \( x, y, z \) the mutually orthogonal axes along the shear, surface normal, and SAW propagation directions, respectively.

Eq. (1) shows that the response of the SAW sensor is the sum of 5 terms: the first takes into account the change in the film density \( \Delta \rho / \rho \) (mass loading), the second and third ones account for the changes in the elastic moduli \( \Delta \varepsilon_{11} / \varepsilon_{11} \) and \( \Delta \varepsilon_{44} / \varepsilon_{44} \) (elastic loading), while the fourth and fifth ones account for the changes in the film conductivity \( \Delta \sigma / \sigma \) (electric loading), and in the temperature \( \Delta T \) (temperature loading), respectively.

On limiting our attention to low gas concentrations, the relative changes in the density, elastic moduli, conductivity, and temperature of the sorbent film can be approximated, within small concentration intervals, by linear functions of the concentration \( n: \Delta \rho / \rho = a_n, \Delta \varepsilon_{11} / \varepsilon_{11} = b_n, \Delta \varepsilon_{44} / \varepsilon_{44} = c_n, \Delta \sigma / \sigma = d_n, \Delta T = e_n \), being \( a, b, c, d, e < 1 \). Under these assumptions, the relative change in the SAW phase velocity, produced by gas adsorption, is deduced from Eq. 1 as a linear function of the gas concentration \( n \):

\[
\frac{\Delta v}{v} = n \cdot \left\{ \left(\frac{\pi h}{2\lambda}\right) \cdot [-A \cdot a + B \cdot b \cdot c + D \cdot (b - 2c)] + K^2 \cdot \left(\frac{\sigma^2}{\varepsilon^2 C_S^2} + 1 \right)^2 \cdot d + (TCV)\cdot e \right\}
\]

From Eqs.(1) and (3) someone can see that the SAW response, the calibration curve, sensitivity and detection limit of the SAW sensors depend through A, B, D, K², and (TCV) on the substrate material and its orientation. This property arises from the elliptic polarization, piezoelectric activity, thermal sensitivity and anisotropy of the SAW propagation in piezoelectric single crystals.
RESULTS AND DISCUSSION

Measurements have been carried out at room temperature and atmospheric pressure on the integrated array of SAW sensors, shown in Fig. 1. This array exploits four different propagation directions along the surface of ST-cut quartz: the x direction (line d), the directions +60° and −50° off x-axis (lines b and a, respectively) and finally the direction perpendicular to the x-axis (line c). The results of the measurements are presented on Figs. 2 and 3.

The most important result following inspection of Figs. 2-3 is that the value of the SAW response as well as the shape and the slope of the calibration curves, in accordance with Eqs. (1) and (3), depend on the SAW propagation direction on anyisotropic substrate. On the other hand, in contrast with these equations, the experimental curves on Figs. 3 show a non linear behaviour. This result can be attributed to the simplified assumptions employed in our theoretical model, which does not take into account the influence of preadsorbed molecules of environmental gases, the starting conditions at the surface of the adsorbers, the mutual interaction of the adsorbed water species, etc.

CONCLUSIONS

This paper shows that the value of the SAW response as well as the calibration curve, sensitivity and threshold concentration of SAW gas sensors can be efficiently controlled without changing the sorbent material film, but simply by a particular selection of the SAW propagation directions. This property distinguish the SAW sensors favorably from all other counterparts. Further investigations on the same topic will be devoted to optimization of the material of the substrate and the propagation directions in order to enhance the anisotropy of SAW properties and, thereby, the difference in SAW sensitivity in different channels of the integrated array.

ACKNOWLEDGMENTS

This work is supported by Russian Foundation for Basic Research grant # 00-15-96652 and CNR-RAS agreement.

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Brillouin Scattering in Silica-titania Waveguides

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Silica-titania planar waveguides, produced by radio frequency sputtering have been investigated by waveguided Brillouin spectroscopy. The laser beam was coupled by a prism to different transverse electric modes of the waveguide. Brillouin spectra are mode dependent because the light propagates in the waveguide at different angles for different modes. In waveguides with a single-step profile of the refractive index, two Brillouin peaks for longitudinal phonons, and two for transverse phonons, were observed. Brillouin spectra of waveguides with a graded index profiles show a rich phenomenology of lineshapes, with additional peaks between the two main longitudinal peaks. The experimental data are reproduced by a model which considers the electric field profile in the waveguides, by using the sound velocity of the film as a parameter.

INTRODUCTION

In waveguide spectroscopy the laser light is coupled to the guide by a prism in order to excite one of the transverse magnetic (TM) or electric (TE) modes. The luminescence or the light scattering spectra, Raman or Brillouin, can be detected by collecting the emitted or scattered light from the front of the planar surface. The technique has been applied to waveguides obtained by different methods such as rf-sputtering, sol-gel dip-coating, and ion-exchange.

If the waveguide has a graded-index profile, different modes penetrate to different depths in the film, probing different structures with their typical vibrational dynamics in Raman scattering and different ion concentrations or crystal field distributions in luminescence measurements. The field profile of the excitation laser light can be calculated for each mode when the index profile is known.

Since Brillouin scattering from acoustic vibration of the glass depend on the exchanged q-vector of the scattered photon, Brillouin spectra (BS) are mode dependent.

Figure 1 shows the BS of a single-step profile silica-titania planar waveguide obtained by exciting at 514.5 nm in different TE modes by using a double monochromator with a resolution of 0.04 cm\textsuperscript{-1}. Two longitudinal and two transverse peaks are observed, with splitting which increases with the mode index. For the $m=0$ excitation no splitting is observed.

FIGURE 1. Brillouin spectra of the silica-titania planar waveguide of composition $87\text{SiO}_2-13\text{TiO}_2$ mol % and thickness 3.7 $\mu$m, obtained by excitation at 514.5 nm in different TE\textsubscript{m} modes. The solid lines are the results of a fit with the model described in the text.

A simple model can account for the measured spectra and allows to deduce the sound velocities in the film. In a ray-optic approach, which neglects the contribution to the scattering from the evanescent field in the substrate, the exciting light propagates in the $z$-direction and the plane waves have a zig-zag path in the $x$-$y$ plane (Fig. 2).
FIGURE 2. Wave propagation in a homogeneous planar waveguide. \( \vec{q}_1 \) and \( \vec{q}_2 \) are the exchanged wave vectors of the scattered light in the zig and zag paths.

Being \( \vec{k}_{\text{out}} \) determined by the geometry of detection of the scattered light, two values \( \vec{q}_1 \) and \( \vec{q}_2 \) of the exchanged wave vector will be present, as shown in Fig. 2. Therefore, the Brillouin spectrum of the waveguide will show four peaks in the Stokes and four peaks in the anti-Stokes spectrum. Two peaks are due to longitudinal phonons and the other two to transverse phonons. The corresponding longitudinal frequencies will be:

\[
\omega_L^{(1,2)} = \left( \frac{2n_2\omega_0}{c} \right) V_L \sin \left( \frac{\phi_{1,2}}{2} \right)
\]

where \( \omega_0 \) is the frequency of the laser light, \( V_L \) is the longitudinal sound velocity and the two \( \phi_{1,2} \) angles are those of Fig. 2. The transverse frequencies can be obtained by considering the transverse sound velocity \( V_T \). The appearance of a single peak in the TE\(_0\) spectrum is not due to a poor experimental resolution, but it is caused by an interference effect between the light scattered in the zig and zag paths. A detailed description of the model is reported in Ref. 5. The curves in the anti-Stokes spectra of the longitudinal phonons (Fig. 1), are obtained within this model using as free parameter only the sound velocity \( V_L \). The values of the obtained sound velocities, \( V_L = 5.75 \times 10^3 \) m/s and \( V_T = 3.46 \times 10^3 \) m/s, well compare with the available data of silica-titania bulk glasses, obtained by different techniques.

Figure 3 shows the BS of a graded-index planar waveguide in the region of the longitudinal phonons. The left column reports the refractive index profile (Fig. 3a) and the electric field patterns TE\(_m\) of some modes. The right column shows the corresponding BS, together with the results of a model (dotted line), which considers the space distribution of the exciting field in the mode. For \( m = 2, 3 \) modes, additional peaks between the two main longitudinal peaks appear in the calculated spectra. After convolution of the calculated spectra with the response function of the spectrometer (continuous line), a good fit to the experimental data is obtained, but the minor structures are lost. Figure 3b shows the BS for the mode \( m = 2 \), taken with higher resolution by using a tandem Fabry-Perot (resolution 0.15 cm\(^{-1}\)). Here, the central structure, even if not really well resolved, seems to appear.

REFERENCES

Investigation of Acoustic and Optic Confined Phonons of Tin Nanoparticles During Melting

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Confined acoustic phonons were detected in tin nanoparticles embedded in SiO\textsubscript{x} thin films, using low-frequency Raman scattering; a Sandercock interferometer was employed. The low eigenfrequencies (few cm\textsuperscript{-1}) of the modes agree with theoretical calculations for the free vibrations of an elastic sphere. Clear effects of particle melting on the acoustic modes were observed, together with the detection of damping phenomena and a coupling with a strong central peak. The observation of melting, and surface melting, through direct investigation of the nanoparticle vibrational eigenmodes represents a novel feature of this work.

Spherical nanoparticles can vibrate, according to the theory of Lamb [1-4], and selection rules [5] exist allowing only some vibrational modes to be Raman active. In this work measurements of low-frequency (LF) Raman scattering [6-9] from localized acoustic vibrations in \(\beta\)-Sn nanoparticles embedded in a SiO\textsubscript{x} film as a function of temperature are reported. The temperature evolution of these modes is shown to be a new and direct evidence of particle melting. The samples were prepared by an evaporation-condensation technique on Si substrates, as described elsewhere [10, 11]. Three films were deposited with different nanoparticle nominal size; X-ray diffraction (XRD) allowed the determination of the real particle average size: 4.5 nm\pm 2 nm (Sn\textsubscript{1}), 11 nm\pm 1 nm (Sn\textsubscript{2}), 18 nm\pm 1 nm (Sn\textsubscript{3}). The overall film thickness was 186 nm. A tin-free sample of SiO\textsubscript{x} was also deposited (Sn\textsubscript{ref}). XRD was also employed to determine the size-dependent particle melting temperature [11]: 408\pm 15 K (Sn\textsubscript{1}), 458\pm 10 K (Sn\textsubscript{2}), 478\pm 10 K (Sn\textsubscript{3}). The melting temperature of bulk tin is 505 K (232 \degree C).

Since the acoustic eigenmodes frequencies of tin particles are expected below 10 cm\textsuperscript{-1}, a high contrast Sandercock interferometer was used [11], with a free spectral range of 300 GHz, in backscattering geometry and a \(\theta \approx 50^\circ\) incidence angle. The excitation wavelength was the Ar\textsuperscript{+} laser 514.5 nm line; a power of 40 mW was incident onto the sample surface. The samples were placed in a thermocryostat with optical windows and a 10\textsuperscript{4} torr vacuum was induced. Fig.1 shows the LF spectra at room temperature. The spectrum is flat for sample Sn\textsubscript{ref}; for samples Sn\textsubscript{1} and Sn\textsubscript{2} a broad peak is detected, whose maximum is at about 55 GHz and 90 GHz respectively, with a long tail at higher frequencies. In sample Sn\textsubscript{1} only a central peak is present. Fig.2 shows the temperature evolution with temperature of the LF spectrum for sample Sn\textsubscript{3}. Two different families of modes (spherical and torsional) are possible for a sphere [1-4]. The allowed frequencies are labeled by \(n\) and \(l\) (harmonic and angular momentum number) and are proportional to the sound velocity and to the inverse nanocrystal diameter:

\[
\nu = S_{l,n}(v_t / c) \text{ [cm}^{-1}\text{]}, \tag{1}
\]

where \(d\) is the diameter in cm, \(v_t\) is the transverse sound velocity and \(c\) the velocity of light. \(S_{l,n}\) is a coefficient depending on \(l, n\), the particle composition and shape, the matrix, the boundary conditions. From selection rules [5] only \(l=0\) and \(l=2\) spheroidal modes are Raman-active at the first order. The lowest frequency mode is the \(l=2, n=0\) spheroidal (\(S_{2,0}^S\approx 0.84\)); in bulk tin \(v_t\approx 1500-1800\) m/s [12]. After subtracting a Lorentzian to account for the elastic central peak, the peaks were fitted to a damped oscillator lineshape:

\[
I(\nu) = \frac{I_0}{[1-(\nu/\nu_0)^2]^2 + [2\zeta(\nu/\nu_0)]^2}, \tag{2}
\]

where \(\zeta\) is the damping factor; we obtained \(\nu_0\approx 2\) cm\textsuperscript{-1}, \(\zeta\approx 0.25\) and \(\nu_0\approx 3.15\) cm\textsuperscript{-1}, \(\zeta\approx 0.3\) for the 18 nm and the 11 nm particles respectively. These frequencies are in agreement with the \(1/d\) dependence but correspond to \(S_{l,n}\approx 0.7\) (the small discrepancy could be ascribed to a deviation from the spherical shape [13] or to the presence of the matrix). The absence of peaks in sample Sn\textsubscript{1} is related to the fact that in the small particles the vibrational modes are overdamped (due to interaction with the matrix [14, 15]), as observed also by pump-and-probe spectroscopy measurements [16]. In sample Sn\textsubscript{3}, as the temperature increases, the intensity of the strong central peak starts to rise, reaches a maximum just before melting and then remains nearly constant (Fig. 2, right). It is still
possible to detect the vibrational peak superimposed onto the central peak as a shoulder. At 473 K the shoulder disappears, in agreement with the melting temperature of the particles. By fitting the spectra to a superposition of a Lorentzian central peak and of Eq. (2) (Fig. 2, left), a gradual softening and broadening of the vibrational mode together with an increase of the elastically scattered intensity were evidenced (see Fig. 2, right), followed by a sudden change before melting. The temperature dependence of the spectral density can be understood in terms of surface melting [17], characterized by the presence of a thin liquid layer around the solid core already below the melting temperature of the particle. The thin, viscous, liquid surface layer, coupled to the still solid vibrating sphere, damps the normal modes (ζ increases) and shifts the peak maximum to lower frequencies. At higher temperatures the melting process extends to the whole particle. So far the only evidence of surface pre-melting was static (TEM) [17]. The origin of the strong central peak [18] may be induced by the coupling of the soft modes in the single particles with relaxing degrees of freedom through a damping process (arising from dynamic inhomogeneities in the semi-liquid particles in the vicinity of the phase transition).

In conclusion, LF Raman scattering was used to study the temperature behavior of acoustic modes in Sn nanocrystals, giving new information about melting. A sudden change in the acoustic mode frequency and damping is explained in terms of surface pre-melting; no acoustic peaks can be detected after melting. A central peak, whose intensity increases with temperature, was also detected.

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Coupling between piezoelectric actuators and miniaturised thermoacoustic cavities

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Nowadays, there exists an interest for designing thermoacoustic refrigerator involving smaller dimensions than classical ones. This implies to design piezoelectric acoustic sources devoted to the thermoacoustic applications and then to investigate the behaviour of the coupling between piezoelectric loudspeakers and small thermoacoustic resonators, in order to provide the maximum acoustic energy transfer available in the resonator for the refrigeration process. The present work aims at contributing to investigate methods to describe the behaviour of this coupling, giving an accurate model which takes the form of an equivalent acoustic network.

INTRODUCTION

Nowadays, there exists an interest for designing thermoacoustic refrigerator involving smaller dimensions than classical ones, such as extreme simplicity and reliability. Thus, piezoelectric loudspeakers appear to be the best candidates to provide the expected properties. The focus here is on a modelling of the piezoelectric actuator which avoids both complex, quasi-untractable (except numerically), classical calculations and empirical modelling. A procedure for analysing source driven resonator excitation is described, which departs from conventional one ([1], [2]) in both, the inertial factors of the harmonic piezoelectric source considered and the coupling between the source and the wave propagating inside the closed waveguide (characterized by its input impedance $Z_e$) (fig. 1), are not neglected.

THE SYSTEM AND ITS MODELLING

The piezoelectric loudspeaker considered is a disk-type, employing a circular diaphragm structure formed by bonding a piezoelectric ceramic (active layer) to a metal plate. The radius of the piezoelectric ceramic and the metal plate are respectively denoted “a” and “b”, with $a < b$. The constitutive equations are written separately for two domains, that is the domain $r < a$, called domain “d” (disk), where the piezoelectric disk (radius a) and the plate are stuck together, and the domain $a < r < b$, called domain “c” (circumferential domain), which is the ring of metal plate around the piezoelectric disk. This separation into two domains is convenient more particularly because there exists a discontinuity of the neutral plane at the geometric discontinuity ($r=a$).

In place of a modelling based on a quasi-static approximation (the inertial factor being evaluated a posteriori using the corresponding quasi-static solutions), suitable complex wavenumber’s type functions are defined for individual problems, involving the backing plate and the piezoelectric plate (stuck together) and including the effect of the input impedance of the acoustic load. The wellknown equation of motion for the displacement $W$ is given by, in each domain (d and c), for an harmonic motion (the function $e^{j\omega t}$ is omitted in the following analysis):

$$\frac{1}{r} \frac{d}{dr} \left( rM_r \right) - \frac{1}{r} \frac{dM_\theta}{dr} + p\rho \omega^2 W = p_t \quad (1)$$

where the total acoustic pressure $p_t$ is the sum of the opposite of acoustic pressure rear the source diaphragm ($p_r$) and of the acoustic pressure front the source diaphragm ($p$), $\rho$ denote the effective density of the considered domain, and where $M_r$ and $M_\theta$ are the flexural moments per unit length.

Then, the basic equations for the displacement of the source plate take the form:

$$\left( \Delta \Delta - \gamma_c^d \right) W_c(r) = 0,$$

$$\left( \Delta \Delta - \gamma_c^d \right) W_d(r) = 0, \quad (2)$$

respectively for the piezoelectric domain (disc) and the passive plate domain (ring around the disc), where suitable complex wavenumber’s type functions $\gamma_c$ and $\gamma_d$ de-
pends on both the parameters of the piezoelectric plate stuck on the baking plate and the input impedance of the resonator. Explicit approximate solutions in the resonator are obtained, which provides a lumped element equivalent network (fig.2).

**FIGURE 2.** Lumped element circuit of the piezoelectric transducer loaded by the input impedance \( Z_e \) of the transducer.

**RESULTS**

Results given on figure (3) show the calculated acoustic energy transfer as a function of the driving frequency and the length of the resonator (without stack). The coupling phenomenon is emphasised when the first resonance frequency of the source (approximately 1900 Hz) and the one of the resonator (which depends on its length) are closed one to the other. The result given in figure (3) correspond to a resonator without stack; when a stack is set in the resonator the behaviour remains qualitatively the same (fig. 4) (note that in this last situation, the dissipative process in the resonator occurs essentially inside the stack region).

**CONCLUSION**

This model can be successfully used to optimise the coupling between the waveguide and the piezoelectric source in order to obtain the maximum acoustic energy transfer available in the resonator. Then, when a short stack is set at the right place over a section of a nearly half-wavelength resonator (thermoacoustic resonator, fig.1), the highest the acoustic power, the highest the heat flux from one end of the stack (cold exchanger) to the other end (hot exchanger) (for a given efficiency of the refrigerator).

**REFERENCES**

Acoustical diffraction through a Faraday phase screen

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Fink et al. for incoherent processing of ultrasonic fields have utilized random Phase Screens (RPS). These RPSs have shown interesting properties but they do not allow any modification of their randomness. It has seemed interesting to get a system where the shape of the phase screen could be controlled through a simple parameter without modifying the ultrasonic emission. We present the results we have obtained when propagating an impulsional ultrasonic wave through an instable liquid layer. The surface of the layer is vertically modulated at very low frequency (under 100 Hz). When increasing the acceleration that generates the well-known Faraday instability. The surface goes from a plane to a surface covered by waves (rolls), then to a surface where squared, hexagonal patterns or even Faraday crispation appear. This allows us to modulate the acoustical field by a phase screen of periodical, quasi-periodical or chaotic shape. When the surface is crispated or covered by hexagons, the study of the diffracted acoustical field in air shows results similar to those obtained previously for RPSs.

INTRODUCTION

Incoherent processing of pulse-echo signal has been widely used [1]. This technique has been introduced to reduce the speckle noise present in the experimental data involving estimation and spectral analysis of pulse-echo signals. It is accepted that the speckle noise is related with the spatial coherence behavior of the piezoelectric transducer. Thus, Fink et al. using a coherent transducer and a random phase screen placed in front of it obtained an excellent speckle reduction. In addition, they could optimize the relationship between directivity and signal noise rate. However, this experimental setup didn’t allow any modification of the randomness or periodicity of the phase screen.

We present the experimental results we have obtained when propagating an impulsional ultrasonic wave through a fluid layer. Upon the fluid surface patterns have been generated through the well-known Faraday instability with a very well controlled reproducibility [2,3,4]. These structures are used to modulate the ultrasonic field instead of the rigid random phase screen.

To obtain these patterns, the fluid layer is vertically modulated through at very low frequency (under 100 Hz). When increasing the acceleration the surface goes from a plane surface to a surface covered by waves (stripes), then to a surface where squared, hexagonal patterns or even Faraday crispation appear. For a sufficiently important surface the patterns generated are independent of the boundary conditions. The goal of this work is to allow to build a periodical, quasi-periodical or chaotic phase screen and to achieve the periodicity control in a simple experimental way, only using the vertical acceleration amplitude of the fluid layer as the control parameter of the shape of the screen.

When the surface is crispated or covered by hexagons, the study of the diffracted acoustical field in air shows results similar to those obtained previously for RPSs.

THE EXPERIMENTAL SET-UP

In figure 1, we show experimental setup for acoustical diffraction. The experiments are conducted using a cylindrical acrylic vessel of 280 mm diameter filled with distilled water up to a height of 3 mm.

![Figure 1. Experimental setup used to obtain the diffraction of ultrasonic field with a Faraday phase screen.](image-url)
The acoustical field has been recorded using a Bruel&Kjaer (1/8”) microphone.

**FARADAY STRUCTURES**

A fluid layer can be considered as a thin lens, whose transmittance function has a periodicity that can be used as an accurate experimental control parameter. Thus, this transmittance function can present different spatio-temporal behaviors: periodic, quasi-periodic structures, and chaotic ones.

In fig. 2, we show two patterns formation on the fluid surface. In fig. 2a, the structures developed correspond to "Stripes", there is only one unstable mode. In fig. 2b, the structures developed on the surface are "Hexagons" (three modes became unstable).

![Figure 2. Pattern formation on fluid surface.](image)

**ACOUSTICAL FIELD CHARACTERISATION**

In fig. 3, we show experimental results obtained using Faraday phase screen. We plotted the intensity of the central maximum as a function of the axis transducer distance. The acoustic field is diffracted by the four different structures developed on the surface: flat, there is not deformation on the surface; periodic structures (Stripes); quasi-periodic structures (Hexagons) and chaotic behavior (crispated surface). We have worked in Fresnel zone, at 3 mm of the fluid surface.

We want to emphasize several important results shown in fig. 3. First, the acoustical field intensity has a hard dependence with Faraday structures, especially when the surface is crispated or covered with hexagons structures. In these cases, the acoustic field intensity presents an important increase and narrowing compared with a flat surface situation. This fact evidences an increase of the directivity and efficiency of the system. Second, we can observe a displacement in the maximum position. When hexagons are developed on the surface, the maximum displacement is 0.5 mm, but when the surface is crispated, this displacement is 4 mm.

![Figure 3. Acoustic field diffracted by Faraday phase screen.](image)

Finally, in order to study the spatial correlation, we have calculated the acoustic correlation coefficient. These results are shown in fig. 4, and are similar to those obtained by Fink et al.

![Figure 4. Correlation coefficient as function the transducer axis distance.](image)

**REFERENCE**

Visualisation of Vortexes and Acoustic Sound Waves

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A new method for measuring and visualising turbulences, e.g. vortexes, and acoustic sound waves in gases is described. The measuring effect is based on the path difference $\Delta \delta$ between two interfering laser beams. This difference results from changes of the refractive index $n$, caused by pressure fluctuations $p(t)$ in a measuring gas. The method is practicable in both the audio and the ultrasonic ranges, as well as for spatially small turbulences. It is qualified for investigating and presenting numerous acoustic, fluidic and airborne sound phenomena.

MEASURING PRINCIPLE

Modified Doppler interferometers (Fig.1) can be used for measuring sound and turbulence fields in gases.

FIGURE 1. Modified Doppler interferometer

The beam splitter BS1 splits the laser beam into measuring and reference beams. The reference beam reaches BS3 via a mirror. The measuring beam hits the reflector at the end of the section of measurements, is reflected there and then guided by BS2 towards BS3. There, measuring and reference beams interfere and result in a light signal whose intensity $I$ depends on their phase difference $\Delta \phi = \phi_{M} - \phi_{R}$ and thus on the path difference $\Delta \delta$. 

$$I = \frac{I_{0}}{2} (1 + \cos \Delta \phi) = \frac{I_{0}}{2} (1 + \cos(2\pi \frac{\Delta \delta}{\lambda})) \quad (1)$$

If the reflector moves periodically with a velocity $v$, then, because of the Doppler effect, the frequency $f_{M}$ of the reflected measuring beam is different from the frequency $f_{S}$ of the emitted beam by the amount of

$$\pm \Delta f = f_{S} - f_{M} \equiv \pm f_{S} \frac{2v}{c} = \pm \frac{2v}{\lambda} \quad (2)$$

where $c$ is the velocity of light. Behind BS3, the resulting signal $R(\Delta f, t)$ is received by the detector as a frequency signal. For identifying the actual direction of the velocity $v$, the reference beam is sent through a modulator. There, a shift $f_{0}$ is superposed on the reference frequency $f_{R}$ and it follows:

$$I(f_{R}, \Delta f, t) = \frac{I_{0}}{2} (1 + \cos(2\pi (f_{R} \pm \Delta f) t)) \quad (3)$$

After a demodulation, both amount and direction of velocity $v(t)$ and shift $s(t)$ are recognisable.

For measuring temporal and spatial pressure fluctuations, as they occur in sound waves or turbulences, the reflector must be absolutely fixed. If now the pressure $p(t)$ changes within a measuring volume, the density $\rho(t)$ and also the refractive index $n(t)$ change [1][2]:

$$n(t) = \rho(t) K + 1 \quad (4)$$

$K$ is the Gladstone-Dale constant. If the pressure changes by $\Delta p(t) = p_{1}(t) - p_{0}$ with $\Delta p(t) \ll p_{0}$, under adiabatic conditions the refractive index $n$ changes by

$$\Delta n(t) \equiv \frac{n_{0} (n_{0} - 1) \Delta p(t)}{\kappa p_{0}} \quad (5)$$

At a pressure $p_{0}$, the gas in the measuring volume of the length $L$ have the refractive index $n_{0}$ and its optical length be $l_{0} = L n_{0}$. If the pressure changes by $\Delta p(t)$ then it causes proportional changes of the refractive index $n(t)$ and, consequently, changes the optical length of the measuring volume by $\Delta L(t) = L \Delta n(t)$. Now, an interferometer according to Fig.1 senses, in the case of phase demodulation, a virtual reflector shift $\Delta s(t) = L \Delta n(t)$, although the reflector itself is immovable. This shift $\Delta s(t)$ is proportional to $\Delta p(t)$

$$s(t) = \frac{L (n_{0} - 1)}{\kappa p_{0}} p_{1}(t) \quad (6)$$

In the case of frequency demodulation, a virtual reflector velocity

$$v(t) = \frac{L (n_{0} - 1)}{\kappa p_{0}} p_{1}(t) = \frac{L (n_{0} - 1) d p(t)}{dt} \quad (7)$$

is measured. It is proportional to the velocity of pressure change $p_{1}(t)$. If vortexes and sound waves are investigated interferometrically, it has to be considered that they produce time- and position-dependent pressure fluctuations $\Delta p(x,y,z,t)$ which are related to corresponding fluctuations of the refractive index $\Delta n(x,y,z,t)$. Such spatially distributed phenomena are registered appropriately with scanning Doppler interferometers, modified as mentioned above. Turbulences and sound waves that cause changes of the
refractive index $\Delta n(x,y,z,t)$ in the section of measurements are sensed as virtual changes of distance $s$ or velocity $v$ of the rigid reflector:

$$\Delta s(x,t) = \Delta n(x,y,z,t) dx ; v(x,t) = n'_{1}(x,y,z,t) dx$$

A requirement for scanning measurements is that the vortexes occur repeatedly, so that repetition conditions exist in every scanning point.

**EXAMPLES**

Fig. 2 shows the 3D-presentation of a half-spherical wave, emitted from a loudspeaker.

**FIGURE 2.** Measured sound wave of a loudspeaker

Another example are the sound fields of an ultrasonic transmitter at 71 kHz. In Fig. 3a, a momentary state of the sound field is presented. Fig. 3b shows the respective sound levels (directional characteristic).

**FIGURE 3.** Ultrasonic fields at 71 kHz

A requirement for scanning measurements is that the vortexes occur repeatedly, so that repetition conditions exist in every scanning point. Fig. 4 shows as an example a periodic air fluctuation emitted from a tube and generating a sound wave. The superposition of fluctuations and sound wave in axial direction (Fig.4b) can be described by

$$p_{1}(t,x) = \Delta p_{F} \cos(2\pi f (t - \frac{1}{a})) e^{-\alpha x} + \Delta p_{S} \cos(2\pi f (t - \frac{1}{a}))$$

with flow velocity $u$ and sound speed $a$.

In Fig. 5 a vortex field in a planar aerodynamic pipe is shown. The space on the right of the edge is a resonant cavity. The jet stream comes from below and is splitted at the edge. Clearly visible are the two parallel out-of-phase vortex paths.

**FIGURE 4.** Airborne generated sound

**FIGURE 5.** Vortexes in a planar pipe

**REFERENCES**

Theory of the proper acoustic coagulation

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Many papers of agglomerating small particles by using acoustic wave have been presented [3,4,5]. The action of the acoustic wave on small particles and theories explaining partially this phenomenon are known since long [6]. While discussing the phenomenon of transport of the small particles under the influence of acoustic wave we have not taken into account [1] the processes responsible for the coagulation of particles. We have assumed however that the appropriate growth of concentration, corresponding to the reduction of distances between particles, will cause quick coagulation in concentration zones. This paper deals with the problem of the amplitude effect, or proper acoustic coagulation and the problem of the efficiency of the acoustic coagulation. We present a theory allowing for interrelation of quantities characterizing the particles, the medium and the acoustic wave.

INTRODUCTION

Removal of particles from industrial emissions is one of most important problems in abatement of pollution. The problem of the agglomeration of suspended particles by sonic waves has been presented by several investigators [1,7,8]. In this report we assume that the coagulation occurs always in the polydispersion aerosols, as even monodispersion aerosols become polydispersional as a result of heat motions. Particles of greater dimensions vibrate in the acoustic field with smaller amplitudes, smaller particles amplitudes are greater. As a result, relative velocities occur in turn result in collision of particles. We will deal with the problem of selection of acoustic field parameters such that for a given aerosol one would obtain amplitude of vibration great enough for occurrence of the acoustic coagulation.

The problem of the efficiency of the acoustic coagulation has been discussed in this paper.

THE EFFICIENCY OF THE AMPLITUDE EFFECT OR PROPER ACOUSTIC COAGULATION

Denoting the efficiency of the acoustic coagulation by $\eta$ we may write:

$$\eta = \eta_A \eta_i$$

(1)

where: $\eta_A$ is the efficiency of the dust removal process (the efficiency of the amplitude effect) and $\eta_i$ is the efficiency of the process of inertial sedimentation of smaller particles on bigger ones.

The average distance between aerosol particles may be estimated as:

$$l_o = \frac{1}{3\sqrt{n_o}}$$

(2)

where $n_o$ is the number of particles in 1 cubic centimetre of the gas.

In reality, it is some statistical distribution in both distances and velocities of particles. In unit volume of the gas, the number of particles which are able to get in contact each other is expressed:

$$\int n(l) dl$$

(3)

where $n(l)$ is the distribution function and $A$ is the amplitude of particle vibrations.

The aerosols conform to the Gauss type distribution:

$$n(l) = n_m e^{-\left(\frac{L}{L_m}\right)^2}$$

(4)

where $n_m$, $L$ are constants.

The efficiency of the dust removal process as an result of that we call here the amplitude effect may be expressed by:

$$\eta_A = \frac{\int n(l) dl}{\int n(l) dl}$$

(5)
where: \( n_o \) is equal to:

\[
n_o = \int_0^\infty n(l) dl
\]

Making use of definition of the error function we may write \[2\]:

\[
\eta_A = \frac{1}{2} \left[ 1 + \text{Erf}[3.3(\psi - 1)] \right] \quad (6)
\]

where \( \psi \) denotes the relative amplitude of particle vibrations:

\[
\psi = \frac{A}{l_o} \quad (7)
\]

We estimate the value of the \( L \) constant and we obtain, (adopting arbitrarily the value \( \eta_A = 0.99 \)):

\[
\frac{l_o}{L} = 3.3 \quad (8)
\]

Equation (6) is an estimation formula, but one thing is for sure: the efficiency of acoustic dusting depends on relative amplitude (7).

We have obtained the relative amplitude in its final form \[2\]:

\[
\psi = \frac{\mu}{r \omega} \sqrt{\frac{2I}{\rho_0 c_o}} \sqrt[3]{\frac{3s}{4\pi \rho_p}} \quad (9)
\]

where:
- \( \mu \) drag coefficient,
- \( \omega \) angular frequency of vibrations
- \( r \) particle radius
- \( \rho_0 \) density of the medium
- \( \rho_p \) density of the particle
- \( c_o \) acoustic wave velocity
- \( s \) mass of particles contained in unit volume
- \( I \) intensity of the wave

During the process of dust removal, the above value (9) should remain constant. Parameters \( \mu \) and \( \omega \) are defined, for acoustic wave, therefore we have a condition:

\[
I \cdot s^\frac{2}{3} = \text{const}. \quad (10)
\]

Establishing all parameters except for the angular frequency \( \omega \) or frequency of vibrations \( v \) we obtain the dependence \( \eta_A (v) \). From Eqs. (6) and (9) we may calculate a maximum of the wave frequency which gives the value of \( \eta_A \) close to unity. The value of \( \eta_A \) is practically equal to unity up to several kHz and then rapidly falls to zero.

**THE EFFICIENCY OF THE PROCESS OF THE INERTIAL COAGULATION**

The quantity of the efficiency of the acoustic coagulation \( \eta_i \) is a function of the Stokes number \( n_{St} \). The problem of the \( \eta_i(n_{St}) \) dependence has been discussed in the paper \[2\], based on assumption of potential flow-around and viscous flow-around. This is the function growing from 0 to 1. In practice the problem is reduced to the value of the Stokes number:

\[
n_{St} = \frac{\tau v_w}{2R} \quad (11)
\]

where : \( \tau \) is the relaxation time of the smaller particle, \( v_w \) is the maximum amplitude of the relative velocity of particles.

Finally, we obtain \[2\] a conditional for the wave intensity in the form:

\[
I > \text{const} \cdot \frac{r^2}{\tau^2} \frac{\rho_0 c_o}{4}
\]

The calculated intensity value (12) (for average industrial aerosols) is tens thousand times weaker than the intensity required for occurrence of proper acoustical coagulation.

Some examples of this investigations and solutions for cases of special interest will be presented during the conference.

**CONCLUSIONS**

The proper acoustical coagulation is related to the amplitude effect, which is responsible for coagulating action of the acoustic field. The efficiency of the inertial coagulation may be considered as equal to unity.

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Magnetoelastic Properties of Mangnese Zinc Ferrites with Cobalt Ions

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In crystals of manganese - zinc ferrite (MZF) influence of ions of bivalent cobalt on a constant magnetoelastic interaction $B_{44}$ is investigated. Significant growth of values $B_{44}$ and the internal magnetic field is found at the increase of the doped ions content that is seen in reduction of the maximum of ultrasonic attenuation coefficient at spin-orientation phase transition in doped crystals MZF, and in change of field dependence of amplitude of the ultrasonic oscillations excited by MZF. The obtained results are discussed on the basis of theoretical models.

Introduction

On basing of magnetoelastic properties of ferrites practical devices of elements of high-frequency memory, devices for a delay and convolutions of signals, ultrasonic transducers with an opportunity of a choice of oscillations mode [1-3] are developed. Efficiency of magnetoelastic interactions can be determined by significant number of various factors, including chemical composition of used ferrite, its domain-defect structure, technology conditions, added impurities, temperature and external magnetic field.

In the present work the results on influence of ions of bivalent cobalt on magnetoelastic properties of crystals of manganese - zinc ferrite of chemical composition $\text{Mn}_0.9\text{Zn}_{0.4}\text{Fe}_2\text{O}_4$ are presented. Researches were carried out with application of two methods - an echo - pulse method of measurement of ultrasonic attenuation coefficient and the method of measurement of amplitude of ultrasonic oscillations excited on the ferromagnetic sample when placed in a combination of variable and constant magnetic fields [4]. The frequency range of researches was 5-50 MHz, temperature was varied in a range 350-100 K, the constant magnetic field changed from 0 to 200 kA/m, the concentration of cobalt ions $x$ varied from 0,02 to 0,6 mol %. For nominally pure MZF samples of such chemical composition there is the spin-orientation phase transition (SOPT) at temperature 250K investigated in details by us in an earlier paper [5]. The transition is caused by change of the direction of easy magnetization axis at change of temperature and connected with the sign inversion of the constant of magnetic anisotropy. Near SOPT there one significant fluctuations of domain structure and transition is accompanied by anomalies of the magnetic, magnetoelastic and acoustic properties of MZF, in particular the ultrasonic attenuation coefficient has maximum at varied of temperature. Interaction of elastic waves with unstable domain structure near SOPT is governed by constants of magnetoelastic coupling and depends on the direction and polarization of ultrasonic wave as well as on the orientation of the domain wall. The magnetoelastic term of the thermodynamic potential for crystals MZF belonging to cubic symmetry is determined by two constant of magnetoelastic interaction – $B_{11}$ and $B_{44}$. In the crystals MZF near SOPT there is a change of direction of the axis of easy magnetization from [100] to 111. Thus the most intensive maximum of attenuation of a sound about 30 db/sm is shown for the shear ultrasonic oscillations propagating along an axis [110] with polarization [110], and for the longitudinal ultrasonic waves propagating along basic crystallographic axis this maximum is smaller on 20 db/sm.

In this connection researches of influence of ions of bivalent cobalt on magnetoelastic interaction in crystals MZF was carried out here on the basis of measurements of temperature dependence of attenuation coefficient of shear waves with a direction of propagation 110 both with polarization 110 and also measurements of field dependence of amplitude of ultrasonic oscillations of the same polarization excited in doped samples by the above described method. The magnetoelastic interaction for waves of such polarization in MZF crystals is determined only by of constant $B_{44}$. 
The results and discussion

Figure 1 shows the results of researches of temperature dependence of ultrasonic attenuation coefficient for four samples MZF with the various of cobalt ions content x (in mole %).

![FIGURE 1. The attenuation coefficient of shear ultrasonic waves vs temperature in doped crystals MZF (x₁ = 0, x₂ = 0.05, x₃ = 0.2, x₄ = 0.6).](image)

One can see that the increase of the bivalent cobalt content of in MZF results in shifting of the maximum A value toward higher temperatures in comparison with undoped sample MZF and to decreasing its amplitude down to the total disappearance at the maximal contents of cobalt ions. On figure 2 results of researches of field dependence of amplitude of shear ultrasonic oscillations excited in pure and in doped samples of MZF are given (symbols are the same of samples as in figure 1).

![FIGURE 2. Dependence of B₄₄ value on the external magnetic field in doped crystals MZF](image)

From figure 2 it is possible to see, that with increasing of the cobalt ions content in MZF there is a substantial growth of magnetoelastic interactions and shift of a magnetic field values at which the maximum of B₄₄ value is observed towards larger fields. The similar behaviour A and B₄₄ is caused by the big contributions of Co ions to magnetoelastic coupling strength and to the internal field value which are caused by high anisotropy of ions of bivalent cobalt because of incomplete degeneration of the spin - orbital moment. Slonczewski calculations[6], which has been carried out on the basis of one-ionic model, have shown, that the contribution of ions Co to magnetoelastic interaction is determined by the valence of cobalt ions as well as their position in MZF lattice and also by concentration. The behavior of ultrasonic attenuation coefficient near SOPT vs concentration of Co ions seems possible to explain influence of two factors – 1) change of time of spin-phonon interaction due to the with substantial growth of the anisotropy field in doped samples MZF that results in increasing the temperature of attenuation maximum , and 2) fastening of movement of domain borders by highly anisotropic Co ions that leads to decreasing of amplitude of a maximum with growth of concentration of doping ions.

The experiments carried out allow to develop optimum conditions for increasing the efficiency of application of MZF crystals in devices basing on magnetoelastic interaction.

ACKNOWLEDGMENTS

Authors thank Prof. E.V.Charnaya for useful discussions and Dr. V.I.Ivanova for help in preparation of samples.

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