# Electromagnetic-acoustic conversion—a result of the action of a surface force

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The physical interactions responsible for generation of acoustic waves in a conductor whose surface is exposed to electromagnetic radiation is analyzed. It is shown that the main features of the inertial, induction, and deformation mechanisms of electromagnetic—acoustic conversion can be obtained by assuming that the excitation force is of a purely surface character. Current experimental investigations in this field are reviewed.

## 1. INTRODUCTION

Many diverse electroacoustic and magnetoacoustic effects observed in solids incorporate a range of phenomena which occur at the surface of a metal exposed to electromagnetic radiation and which are encompassed by the general concept of electromagnetic-acoustic conversion (EMAC).[1-3] The crux of EMAC is that when a substance which does not exhibit either piezoelectric or magnetostrictional properties is exposed to electromagnetic radiation, ultrasonic waves with the same frequency (linear response) and with the harmonic frequencies (nonlinear response) are excited in the material. The existence of the boundary of the metal, as a location where the exciting force is concentrated, is of fundamental significance. For a uniform force distribution the problem reduces not to the excitation of acoustic waves in the metal—the subject addressed in the present paper—but rather the displacement of the metal as a whole in the external medium, i.e., the problem of vibrators and membranes. These two problems somehow merge when EMAC is employed for generating ultrasound in semiconductors and dielectrics. In this case a nonconducting solid is coated with a metallic film whose thickness is of the order of the thickness of the skin layer in the metal, and it is this film that plays the role of the source (and detector) of acoustic waves.[79]

The nontriviality of the EMAC as a physical phenomenon is caused, in our opinion, by a number of factors. First, the electromagnetic wave incident on the boundary of the metal excites acoustic waves in an electrically neutral body. In each volume element the electron and ion charges compensate one another, and since the Debye–Hueckel radius of the electron plasma of the metal is of the order of or even several times shorter than the interatomic spacing, this compensation is very precise. Second, we are dealing here with a range of phenomena, since the number of mechanisms giving rise to conversion of electromagnetic and acoustic waves in metals is large. Third, most effects observed in the study of purely acoustic properties (damping and velocity of ultrasound) are sometimes manifested more strongly in EMAC.[4,5]

Electromagnetic-acoustic conversion usually only occurs under the conditions of the skin effect, when the thickness $\delta$ of the skin layer is significantly less than the dimensions $d$ of the sample. This phenomenon is often observed in the electrodynamic properties of the sample with spatial resonance $d=n\lambda/2$, where $n=1, 3,\ldots$, when a half-integer number of wavelengths $\lambda$ of the ultrasound fit within the thickness $d$ of the plate.[6] Figures 1 and 2 display traces of such resonances,[7] corresponding to the excitation of transverse and longitudinal ultrasound ($n=1$) in a tin single crystal. If the order $n$ of the resonance is low, then we are obviously dealing with a situation in which $\lambda$ is much greater than $\delta$. Since both $\lambda$ and $\delta$ are much greater than the interatomic spacing $a$ even under the conditions of the anomalous skin effect, the EMAC is customarily considered to be a volume phenomenon, specially distinguishing the surface force—if it exists—localized at atomic distances from the boundary.[8-10] The restrictions $\lambda,\ d\gg\delta$ make it possible to derive compact equations for the amplitude of the excited ultrasound under the assumption that the exciting force is of a surface character. We underscore the fact that the inequalities $\lambda,\ d\gg\delta$ restrict the problem to comparatively low frequencies $f\ll 1$ GHz (see Sec. 7 of Ref. 1).

## 2. GENERAL SOLUTION OF THE PROBLEM

Due to the high reflectance of the metal only a small fraction of the energy of the electromagnetic waves is dissipated in the metal in the skin layer and, naturally, even less energy is converted into acoustic waves. In this situation there is no need to find a self-consistent solution of the problem of excitation of ultrasound. The EMAC efficiency can be calculated in two steps. In the first step, neglecting the coupling between the electromagnetic and ultrasonic oscillations, all electrodynamic and electronic characteristics of the metal whose surface is exposed to electromagnetic radiation can be calculated, and knowing these characteristics the force density $\mathbf{f}(\mathbf{r},t)=\mathbf{f}(\mathbf{r})e^{-i\omega t}$ acting on the crystal lattice can be calculated. At the second step, knowing the force $\mathbf{f}(\mathbf{r},t)$ exciting the ultrasound makes it possible to solve the acoustic problem, i.e., to calculate the displacement field $\mathbf{U}(\mathbf{r},t)$ in an elastic wave. The elastic field

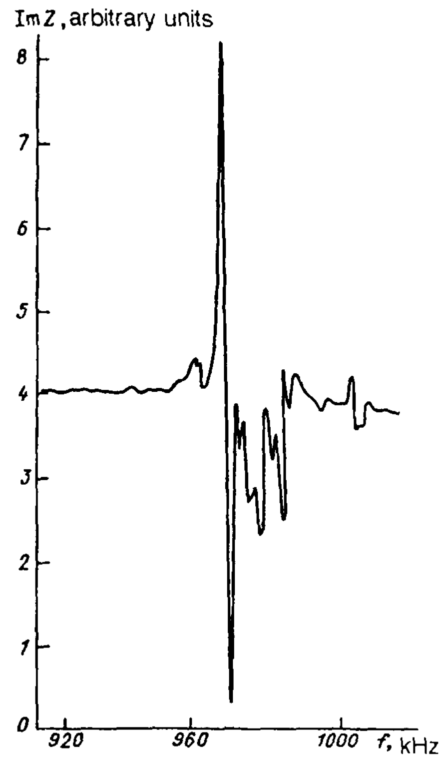



FIG. 1. Transverse-ultrasound resonance in a tin single crystal.[7] $H_0 \| k \| [100]$, $H_0=70$ kOe, $T=4.2$ K, $d=0.1$ cm.

contains both the components concentrated within the skin layer (or, in the more general case, within the mean-free path length of the electrons) and a wave traveling with the sound speed away from the boundary. We assume that the amplitude of this wave $U_\infty$ is the principal characteristic of EMAC in the problem of ultrasound generation in a metal

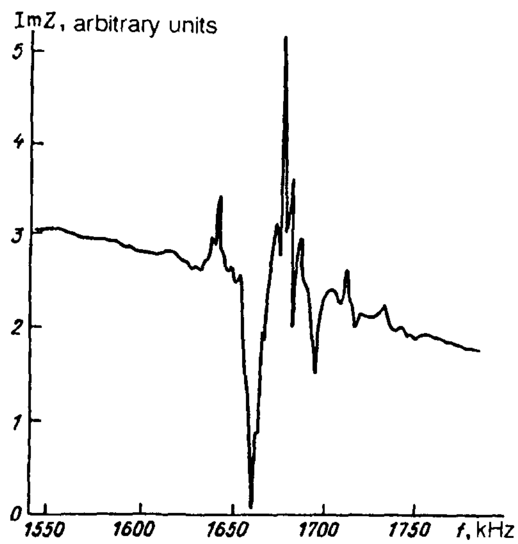



FIG. 2. Longitudinal-ultrasound resonance in a tin single crystal.[7] $H_0 \| k \| [100]$, $H_0=70$ kOe, $T=4.2$ K, $d=0.1$ cm.

occupying the half-space $z>0$. The subscript "$\infty$" in the amplitude actually means that this quantity is measured at a distance much greater than the depth of the skin layer. The damping $\alpha$ of ultrasound can be taken into account by replacing $U_\infty$ by $U_\infty \exp(-\alpha z)$ for a given frequency and polarization.

Three ultrasonic waves with mutually orthogonal polarizations $\varkappa_i (i=1,2,3)$ and phase speeds $S_i$ propagate in each direction in the crystal. Designating by $U_i(r)$ the amplitude of the $i$th wave it is easy to show that this quantity satisfies an inhomogeneous wave equation

$$\Delta U_i(\mathbf{r}) + k_i^2 U_i = f_i(\mathbf{r})/\rho S_i^2,$$
$$f_i(\mathbf{r}) = \mathbf{f}(\mathbf{r}) \varkappa_i, \quad i=1,2,3; \tag{1}$$

here $k_i=\omega/S_i$ is the wave vector of ultrasound, $\rho$ is the density of the metal, and $f_i(\mathbf{r})$ is the projection of the exciting force on the polarization vector of the excited ultrasound. The problem is studied in the harmonic approximation, and the factor $e^{-i\omega t}$ is dropped.

We confine our attention to EMAC with an electromagnetic wave under normal incidence on the boundary of the metallic half-space. Then all quantities depend only on the coordinate $z$, the $z$ axis being normal to the surface of the sample, and the equations (1) simplify:

$$\frac{d^2 U_i(z)}{dz^2} + k_i^2 U_i(z) = \frac{f_i(z)}{\rho S_i^2}. \tag{2}$$

In what follows, we drop the index $i$.

The Newtonian boundary condition

$$U + D \left.\frac{dU}{dz}\right|_{z=0} = 0 \tag{3}$$

makes it possible to investigate the generation of ultrasound when the boundary of the metal is mechanically clamped ($D=0$) and free ($D=\infty$) as well as under intermediate conditions. The first possibility is realized, in particular, in the problem of EMAC in a conducting liquid placed in a dielectric container[11] or with excitation of ultrasound in metallic films on a dielectric substrate.[12-16] The second possibility corresponds to the classical experiment, when the electromagnetic field is produced by an inductance coil placed on the free boundary of a metal.[4,5,17-25] In this case the sample is clamped along the lateral surfaces (Fig. 3); this does not prevent realization of the advantages of the electromagnetic (contactless) method of generation and detection of acoustic waves.

With the help of the Green's function $G_D(z,z')$ satisfying the boundary condition (3), we obtain

$$U(z) = \frac{1}{\rho S^2} \int_0^\infty G_D(z,z') f(z') dz', \tag{4}$$

where

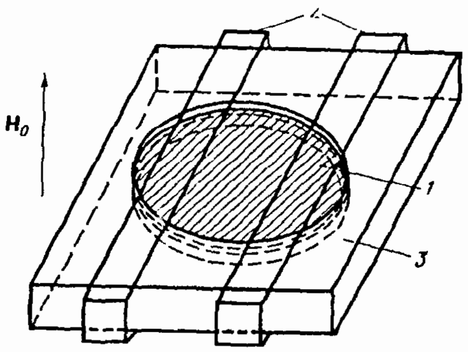



FIG. 3. Arrangement of the experiment on electromagnetic excitation and detection of ultrasound in metals. Sample *1*, encompassed by induction coils *2*, is placed in a dielectric holder *3*. Transverse waves are excited in the sample in a magnetic field $H_0$ perpendicular to the surface.

$$G_D(z,z') = -\frac{e^{ikz'}}{2ik}\left(\frac{1-ikD}{1+ikD}e^{ikz}-e^{-ikz}\right), \quad z<z',$$

$$= -\frac{e^{ikz}}{2ik}\left(\frac{1-ikD}{1+ikD}e^{ikz'}-e^{-ikz'}\right), \quad z>z'. \tag{5}$$

This enables writing the general solution of the acoustic part of the problem in the form

$$U(z) = -\frac{1}{2i\rho S^2 k}\int_0^z\left[\frac{1-ikD}{1+ikD}e^{ikz'}-e^{-ikz'}\right]f(z')\mathrm{d}z' e^{ikz}$$
$$-\frac{1}{2i\rho S^2 k}\int_z^\infty\left[\frac{1-ikD}{1+ikD}e^{ikz}-e^{-ikz}\right]f(z')e^{ikz'}\mathrm{d}z'. \tag{6}$$

The amplitude of the elastic wave at large distances from the boundary is determined by the first terms:

$$U(z)\,|_{z\to\infty} = U_\infty e^{ikz}, \tag{7}$$

where

$$U_\infty = \frac{1}{\rho S^2(1+ikD)}\int_0^\infty\left(D\cos kz - \frac{1}{k}\sin kz\right)f(z)\mathrm{d}z. \tag{8}$$

The surface character of the force ($\delta k \ll 1$) makes it possible to simplify Eq. (8) significantly:

$$U_\infty = \frac{1}{\rho S^2(1+ikD)}\int_0^\infty (D-z)f(z)\mathrm{d}z. \tag{9}$$

Further progress is possible only if the force density $f(z)$ acting on the crystal is specified more precisely.

## 3. NATURE OF THE FORCES

As mentioned above, EMAC is realized by several mechanisms which transform energy from one form (electromagnetic) into another (elastic displacements). Confining attention to EMAC in a normal (nonsuperconducting and nonmagnetic) metal, ignoring nonlinear effects, and separating the thermoelastic mechanism of EMAC into a special paper on the subject,[26] the force density acting on the crystal lattice can be represented in the form of three terms:

$$\mathbf{f} = \mathbf{f}^{\mathrm{ST}} + \mathbf{f}^{\mathrm{L}} + \mathbf{f}^{\mathrm{Def}}, \tag{10}$$

each of which reflects a definite specific feature of the dynamics of the electron gas. Thus, the first term—the Stewart–Tolman force

$$\mathbf{f}^{\mathrm{ST}} = \frac{m}{e}\frac{\partial \mathbf{j}}{\partial t} = -\frac{i\omega m}{e}\mathbf{j} \tag{11}$$

(where $m$ and $e$ are the electron mass and charge, $\omega$ is the frequency, and $\mathbf{j}$ is the current density)—is determined by the noninertial motion of the crystal lattice and, ultimately, it is expressed in terms of the temporal dispersion of the conductivity. A characteristic feature of this mechanism of EMAC is that the true electron mass $m$ appears in Eq. (11). In addition, the inertial force is proportional to the frequency of the incident electromagnetic radiation, and this also identifies the Stewart–Tolman conversion mechanism.

The second component of the exciting force—the Lorentz force

$$\mathbf{f}^{\mathrm{L}} = [\mathbf{j}\mathbf{H}_0]/c \tag{12}$$

is manifested only in the presence of a constant magnetic field $H_0$. This identifies the Lorentzian conversion mechanism. We now call attention to the fact that $\mathbf{f}^{\mathrm{ST}}$ and $\mathbf{f}^{\mathrm{L}}$ are orthogonal to one another. This could be significant for determining the role of the inertial force in EMAC.

Finally, $\mathbf{f}^{\mathrm{Def}}$—the deformation force density—is a vector with the components

$$f_i^{\mathrm{Def}} = \partial\langle\lambda_{ik}\chi\rangle/\partial x_k; \tag{13}$$

here the brackets $\langle ... \rangle$ means integration over the Fermi surface

$$\langle ... \rangle = \frac{2}{(2\pi h)^3}\int_{\varepsilon=\varepsilon_{\mathrm{F}}} ... \frac{\mathrm{d}S}{v},$$

where $\mathrm{d}S$ is an element of the Fermi surface $\varepsilon_{\mathrm{F}}$; $\mathbf{v}$ is the electron velocity; $\lambda_{ik}$ is the deformation potential tensor for free electrons (in the Drude–Lorentz–Sommerfeld model), equal to $mv_iv_k$; $(\partial F/\partial\varepsilon)\chi(\mathbf{p},z)$ is a nonequilibrium correction to the Fermi function $F$, and must be found by solving the corresponding transport equation. The deformation force is produced by direct transfer of the quasimomentum acquired by electrons from the electric field of the radiation to the lattice. This "transfer" occurs, on the average, over a distance equal to the mean-free path $l$ from the location of "acquisition." The electron mean-free path depends on the temperature, and thus the magnitude of the deformation force also depends on the temperature.

The association of the three principal mechanisms of EMAC (inertial, induction, and deformational) to three external parameters of the problem (frequency, magnetic field, and temperature) should not be interpreted literally. It is obvious that the efficiencies of the induction and deformational mechanisms will depend on the frequency, and the efficiency of the deformational interaction will depend on the magnetic field. This division of the sources of the force is more important as a base for further analysis of EMAC at the boundary of a metal.

The arbitrariness of the representation of the force acting on the lattice in the form of three terms is clearly seen in the free-electron model, when under very general assumptions about the collision integral in the transport equation the total force density can be written as

$$\mathbf{f}=ne(\mathbf{E}-\mathbf{j}/\sigma), \tag{14}$$

where $n$ is the electron density, $\mathbf{E}$ is the electric-field density, and $\sigma$ is the specific static conductivity. This simple formula clearly demonstrates the origin of EMAC. It is evident that for $\mathbf{j}=\sigma\mathbf{E}$ the density of the exciting force vanishes, and this in turn is possible only when the conductivity does not exhibit either temporal or spatial dispersion and also when the conductivity does not depend on the constant magnetic field.

Naturally, each term in Eq. (10) can be "recovered" from Eq. (14). It is not difficult to formulate the conditions under which Eq. (14) becomes one of the three terms in Eq. (10). Thus, in order that Eq. (14) to be identical to Eq. (11) only the temporal dispersion of the conductivity need be taken into consideration. In order for Eq. (14) to be identical to Eq. (12) the intensity $\mathbf{E}$ of the electric field must be calculated from the fact that (according to the Drude–Lorentz–Sommerfeld model) for $H_0\neq 0$ the resistance tensor $\hat{\rho}$ in the coordinate system tied to $H_0$ has the form

$$\hat{\rho}=\begin{pmatrix} \rho & H_0/nec & 0 \\ -H_0/nec & \rho & 0 \\ 0 & 0 & \rho \end{pmatrix}, \quad \rho=1/\sigma. \tag{15}$$

Finally, in order for Eq. (14) to be identical to Eq. (13) the electric-field $\mathbf{E}$ and current $\mathbf{j}$ densities must be calculated according to the theory of the anomalous skin effect with $H_0=0$ and with no temporal dispersion of the conductivity ($\omega\tau\rightarrow 0$).

We have discussed the properties of the Drude–Lorentz–Sommerfeld model in such great detail because it can be conjectured that this model is applicable to metals such as potassium, whose Fermi surface is a sphere. If this is so, then the main characteristic of EMAC $U_\infty$ should be expressed in electrodynamic terms and should not contain any parameters describing the interaction of the electrons with the sound. The electron-lattice interaction, which is required for EMAC, is not "discarded" here: The finite mean-free path length and hence the finite conductivity and resistance are a result of the interaction of the electrons with the lattice, to which electrons transfer the momentum that they acquire from the electric field. Identity of Eq. (14) and Eqs. (10)–(13) in the absence of coupling of the electrons with an ideal crystal lattice obtains because in the Drude—Lorentz—Sommerfeld model the deformation-potential tensor $\lambda_{ik}$, as has already been mentioned, degenerates into a bivector $mv_iv_k$ describing momentum transfer to the crystal.

## 4. EMAC. INERTIAL AND INDUCTION INTERACTIONS

A common feature of the Stewart–Tolman and Lorentz interactions is that both are expressed in terms of the alternating current density induced by the electromagnetic radiation in the skin layer. If ultrasound is generated owing to these interactions, then according to Eqs. (1), (8), (11), and (12)

$$U_\infty=\frac{1}{\rho S^2(1+ikD)}\int_0^\infty (D-z)\times\left(\frac{1}{c}[H_0\varkappa]-\frac{i\omega m}{e}\varkappa\right)\mathbf{j}(z)\mathrm{d}z. \tag{16}$$

We recall that $\varkappa$ is the unit vector of polarization excited by the sound wave.

This formula can be transformed with the help of Maxwell's equations without making any assumptions about the character of the conductivity. For this, since $j_z\equiv 0$, we employ the values of the two integrals

$$\int_0^\infty \mathbf{j}(z)\mathrm{d}z=\frac{c}{4\pi}[\hat{z}\mathbf{H}], \tag{17}$$

$$\int_0^\infty z\mathbf{j}(z)\mathrm{d}z=-\frac{c^2}{4\pi i\omega}\mathbf{E}=-\frac{c^2}{4\pi i\omega}\hat{Z}[\mathbf{H}\hat{\mathbf{z}}], \tag{18}$$

where $\mathbf{H}$ and $\mathbf{E}$ are, respectively, the intensity of the magnetic and electric fields of the electromagnetic wave at the boundary of the metal ($z=0$); $\hat{Z}$ is the surface impedance tensor of the conducting half-space; $\hat{z}$ is a unit vector along the $z$ axis; and, $\hat{Z}[\mathbf{z},\mathbf{H}]$ is a vector with the components $Z_{\alpha\beta}[\mathbf{z},\mathbf{H}]_\beta$, where $\alpha$, $\beta\equiv x$, $y$. (We present the following chain of equalities which follow from Maxwell's equations curl $\mathbf{H}=4\pi\mathbf{j}/c$ and curl $\mathbf{E}=i\omega\mathbf{H}/c$:

$$\int_0^\infty zj_x(z)\mathrm{d}z=\frac{c}{4\pi}\int_0^\infty z\,\mathrm{curl}_x H\mathrm{d}z=-\frac{c}{4\pi}\int_0^\infty z\frac{\mathrm{d}H_y}{\mathrm{d}z}\mathrm{d}z$$
$$=\frac{c}{4\pi}\int_0^\infty H_y(z)\mathrm{d}z=\frac{c^2}{4\pi i\omega}\int_0^\infty\frac{\mathrm{d}E_x(z)}{\mathrm{d}z}\mathrm{d}z$$
$$=-\frac{c^2}{4\pi i\omega}E_x(0).$$

Thus

$$U_\infty=\frac{mc}{4\pi e\rho S^2}\frac{1}{1+ikD}([\boldsymbol{\omega}_c\varkappa]-i\omega\varkappa)\times\left(D[\hat{\mathbf{z}}\mathbf{H}]+\frac{ic}{\omega}\hat{Z}[\hat{\mathbf{z}}\mathbf{H}]\right), \tag{19}$$

where $\boldsymbol{\omega}_c=e\mathbf{H}/mc$ is the free-electron cyclotron frequency. The latter formula is very important, especially for $D=\infty$. It turned out that the main characteristic of EMAC—the amplitude $U_\infty$ of the excited sound wave—with a mechanically free boundary ($D=\infty$) does not depend on the electrodynamic characteristics of the conductor, but rather it is determined by the external magnetic field $H_0$ and the magnetic field $H$ of the wave on the surface of the sample as well as by the acoustic properties ($\lambda,\rho,S$) of the sample.

The applicability of the condition of a mechanically free boundary means that $kD \gg 1$ (or $D \gg \lambda$). We note that in the case when sound is excited by a surface force (our case) the inequality $D \gg \lambda$ makes it possible to neglect in Eq. (19) the term containing the impedance, since

$$\delta = |c\mathbf{Z}/i\omega| \tag{20}$$

irrespective of the conduction mechanism, and we assume that $\delta \ll \lambda$.

If $\omega_c \gg \omega$, then the amplitude of the sound excited by the Stewart–Tolman force ($U_\infty^{\mathrm{ST}}$) is small compared to $U_\infty^{\mathrm{L}}$—the amplitude of the sound excited by the Lorentz force. The latter amplitude can be written in a form convenient for calculations:

$$|U_\infty^{\mathrm{L}}| = \frac{\lambda}{4\pi\rho S^2} [\mathbf{H}_0 \varkappa][\mathbf{H}\hat{\mathbf{z}}], \qquad \lambda = \frac{S}{\omega}, \quad D = \infty. \tag{21}$$

It is obvious that under the same conditions

$$|U_\infty^{\mathrm{ST}}| \sim (\omega/\omega_c)|U_\infty^{\mathrm{L}}|. \tag{22}$$

It is easy to show, with the help of the formula (19) and using the expression for the surface impedance, that for EMAC a free metal surface of a metal is preferable to a clamped surface. Indeed, according to Eq. (20)

$$\frac{U_\infty|_{D=0}}{U_\infty|_{D=\infty}} \sim \frac{\delta}{\lambda} \ll 1. \tag{23}$$

## 5. EMAC. DEFORMATION INTERACTION

We now consider the deformation force (13). The components of the deformation-potential tensor are even functions of the quasimomentum $\mathbf{p}$ ($\lambda_{ik}(-\mathbf{p}) = \lambda_{ik}(+\mathbf{p})$). Neglecting the anomalous nature of the skin effect, the nonequilibrium correction to the distribution function of the conduction electrons is an odd function of the quasimomentum. This means that under the conditions of the strictly normal skin effect $l \ll \delta$ the deformation force vanishes (see Eq. (13)). In order to calculate the amplitude of ultrasound excited by the deformation force it is necessary to take into account the finiteness of the mean-free path $l$, even if it is small compared to the thickness $\delta$ of the skin layer. Thus there appears in the problem, together with macroscopic parameters having the dimension of length—the thickness $\delta$ of the skin layer and the wavelength $\lambda$ of the elastic wave, the microscopic parameter $l$, which, in principle, can form any ratio with $\lambda$ and $\delta$. This makes it necessary to consider separately the case of EMAC owing to the deformation force.

As we have already mentioned, the total deformation force acting on the metal is zero, i.e.,

$$\int_0^\infty \mathbf{f}^{\mathrm{Def}}(z)\,dz = 0. \tag{24}$$

It is obvious from the expression for the deformation force (13) that the condition (24) can be satisfied automatically, if owing to the boundary conditions the function $\chi(\mathbf{p},z)$ satisfies the equations

$$\langle \lambda_{iz}(\mathbf{p})\chi(\mathbf{p},0)\rangle = 0. \tag{25}$$

If this is not so, then the surface force $\mathbf{f}^{\mathrm{S}}$ is concentrated directly at the boundary of the metal in a layer, having a thickness of the order of the interatomic separation, where electrons undergo scattering differing significantly from volume scattering.[8–10,27,28] The expression for $\mathbf{f}^{\mathrm{S}}$ can apparently be derived by studying the interaction of the electrons with the boundary. This result is unknown to us, however, even though the interaction of electrons with the boundary of the sample has been studied many times in the derivation of the boundary condition for the function $\chi$. The problem is that there is actually no need for a microscopic derivation, since the exact coordinate-dependence of $\mathbf{f}^{\mathrm{S}}$ is not important. For example, it can be assumed that

$$\mathbf{f}^{\mathrm{S}} = \mathrm{A}\exp(-z/a) \tag{26}$$

and the factor can be chosen so that the condition (24) is satisfied:

$$\mathrm{A} = \langle \boldsymbol{\lambda}(\mathbf{p})\chi(\mathbf{p},0)\rangle / a; \tag{27}$$

here $\boldsymbol{\lambda}$ is a vector with the components $\lambda_{iz}$ ($i = x,y,z$). Then the expression (13) for the deformation force density can be replaced by the formula

$$\mathbf{f}^{\mathrm{Def}}(z) = \frac{\partial}{\partial z}\langle \boldsymbol{\lambda}(\mathbf{p})\chi(\mathbf{p},z)\rangle + \frac{1}{a}\langle \boldsymbol{\lambda}(\mathbf{p})\chi(\mathbf{p},0)\rangle \exp\left(-\frac{z}{a}\right), \tag{27'}$$

and in order to obtain the macroscopic limit the parameter $a$ must be set equal to zero in the expressions obtained.

The unique coordinate dependence of the deformation force makes it necessary to rederive an equation analogous to Eq. (7). Let

$$\langle \varkappa \boldsymbol{\lambda}(\mathbf{p})\chi(\mathbf{p},z)\rangle = \varphi(z). \tag{28}$$

Then

$$f^{\mathrm{Def}} = \frac{\mathrm{d}}{\mathrm{d}z}[\varphi - \varphi(0)\exp(-z/a)].$$

Substituting this expression into Eq. (8) and passing to the limit $a \to 0$ we obtain

$$U_\infty = \frac{1}{\rho S^2(1+ikD)} \int_0^\infty (kD\sin kz + \cos kz)\varphi(z)\,\mathrm{d}z. \tag{29}$$

We shall employ Eqs. (28) and (29) below.

The deformation mechanism of sound excitation has been studied many times.[2,3,29–38] It should be kept in mind, however, that a systematic microscopic theory of this phenomenon is still not available for metals with an arbitrary Fermi surface. First of all, little is known about the deformation-potential tensor. In particular, the components of this tensor have not been measured, even though it is these components (averaged over the Fermi surface) that determine the electronic part of the sound absorption by the metal, which in the limit $kl \gg 1$ does not depend on the dissipative characteristics of the conduction electrons, and for this reason can be employed in order to determine experimentally the components $\lambda_{ik}$ for a metal whose Fermi surface is known. Second, the problem of the skin effect in a metal with a complicated Fermi surface has not been solved outside the $\tau$ approximation, so that $U_\infty$ cannot be calculated for different values of $kl$.

In the present paper we follow the traditional approach and confine our attention to solving the simplest possible problem. Thus we assume that the function $\chi$ satisfies the kinetic equation in the $\tau$ approximation ($\omega\tau \ll 1$, $H_0=0$), i.e.,

$$v_z \frac{\partial \chi}{\partial z} + \frac{\chi}{\tau} = e\mathbf{v}\mathbf{E}. \qquad (30)$$

The interaction with the boundary is described by Fuch's condition

$$\chi(\mathbf{p},0)\,|_{v>0} = Q\chi(\mathbf{p},0)\,|_{v<0}, \qquad (31)$$

where $Q$ is the specularity parameter ($0 \leqslant Q \leqslant 1$), and inside the metal

$$\chi(\mathbf{p},z)\,|_{v<0}. \qquad (32)$$

The most important restriction is associated with the geometry of the problem and the structure of the deformation potential. Let the $z$ axis (more accurately, $p_z$) be the symmetry axis of the Fermi surface of the metal. We assume that the deformation potential has the same structure as the bivector $\tilde{m} v_i v_k$, where the scalar factor $\tilde{m}$ with the dimension of mass (it is often called the effective mass of the electron-phonon interaction) is a phenomenological constant of the problem.

The function $\chi = \chi(\mathbf{p},z)$ does not depend directly on the quasimomentum $\mathbf{p}$, but rather it depends on the components of the velocity $\mathbf{v}$. We assume that the geometry of the problem is such that integrals of the type $\langle v_x v_z \rangle$ and so on computed over the Fermi surface vanish (see Ref. 39 for a discussion of this point). This means that, in particular, in the case of an electromagnetic wave incident normally on the surface of a metal occupying the half-space $z>0$, not only $j_z \equiv 0$ but also $E_z \equiv 0$.

According to Eqs. (30)–(32), having oriented the $x$ axis in the direction of the polarization in the electric field of the electromagnetic wave, we have

$$\chi(\mathbf{p},z) = e\int_z^\infty \frac{v_x E_x(z)}{|v_z|} \exp\left(-\frac{z'-z}{\tau|v_z|}\right) dz', \quad v_z<0,$$
$$= e\int_0^z \frac{v_x E_x(z')}{|v_z|} \exp\left(-\frac{z-z'}{\tau|v_z|}\right) dz'$$
$$+ eQ\int_0^\infty \frac{v_x E_x(z')}{|v_z|} \exp\left(-\frac{z'+z}{\tau|v_z|}\right) dz', \quad v_z>0. \qquad (33)$$

Using Eqs. (14) and (28) it is easy to show that under our assumptions about the geometry of the problem the deformation force excites only a wave for which the projection of the polarization vector $\varkappa$ of ultrasound on the $x$ axis is different from zero. This is obviously the case for a transverse (with respect to the $z$ axis) sound wave in an elastically isotropic body. Then $\varkappa\lambda$ is $\lambda_x = \tilde{m} v_x v_z$. This is the case that we shall study. In calculating $U_\infty$ from Eq. (29) (see also Eqs. (28) and (32)) we shall perform the integration over the Fermi surface separately over the parts where $v_z>0$ and $v_z<0$, i.e., we employ the fact that $\langle ... \rangle = \langle ... \rangle_+ + \langle ... \rangle_-$.

To each point $P$ on the Fermi surface with $v_z<0$ it is possible to associate in a one-to-one manner the point $P'$ on the Fermi surface where $v_z>0$, and $v_x(P') = +v_x(P)$ and $v_z(P') = -v_z(P)$. This correspondence, which is a consequence of our assumption about the geometry of the problem, enables us to switch from integration over the Fermi surface to integration over the part (half) of the Fermi surface where $v_z>0$. Taking this into consideration, it is easy to derive

$$U_\infty^{\mathrm{Def}} = \frac{e}{\rho S^2(1+ikD)} \Bigg\{ k^2 D \Bigg[ \left\langle \frac{\tilde{m} v_x^2 (\tau v_z)^2}{1+(k\tau v_z)^2} \right\rangle_+$$
$$\times \int_0^\infty E_x(z) \cos kz\, dz + (Q-1) \left\langle \frac{\tilde{m} v_x^2 (\tau v_z)^2}{1+(k\tau v_z)^2} \right.$$
$$\left. \times \int_0^\infty E_x(z) \exp\left(-\frac{z}{\tau v_z}\right) dz \right\rangle_+ \Bigg] + (Q-1)$$
$$\times \left\langle \frac{\tilde{m} v_x^2 (\tau v_z)^2}{1+(k\tau v_z)^2} \right\rangle_+ E_x(0) + \frac{i\omega}{c} \Bigg[ (Q+1)$$
$$\times \left\langle \frac{m v_x^2 (\tau v_z)^2}{1+(k\tau v_z)^2} \int_0^\infty H_y(z) \exp\left(-\frac{z}{\tau v_z}\right) dz \right\rangle_+$$
$$-2 \left\langle \frac{m v_x^2 (\tau v_z)^2}{1+(k\tau v_z)^2} \right\rangle_+ \int_0^\infty H_y(z) \cos kz\, dz \Bigg] \Bigg\}. \qquad (34)$$

The equations derived above are convenient for calculations and for obtaining results. The limits of a mechanically free boundary ($D \to \infty$) and a clamped boundary ($D=0$) as well as specular reflection ($Q=1$) and diffuse reflection ($Q=0$) are obvious. It is only necessary to take into account the fact that under the conditions of the anomalous skin effect $E_x(z)$ and $H_y(z)$ depend on the character of the reflection of the electron and, strictly speaking, for $Q=0$ and $Q=1$ different functions $E_x(z)$ and $H_y(z)$ must be substituted into Eq. (34). We neglect the difference in the electric and magnetic fields as a function of the values of the Fuchs' parameter $Q$, since the penetration depth of the fields is apparently insensitive to the value of this parameter.

Since the number of parameters is large, the cases of a mechanically free boundary ($D=\infty$) and a mechanically clamped boundary ($D=0$) must be studied separately. We shall focus on the case $D=\infty$ (as a rule, sound is excited in a conductor with a free boundary[6,40–42]):

$$U_\infty^{\text{Def}}|_{D=\infty} = -\frac{iek}{\rho S^2}\left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\int_0^\infty \left[\cos kz + (Q-1)\right.\right.$$

$$\left.\left.\times \exp\left(-\frac{z}{\tau v_z}\right)\right]E_x(z)\mathrm{d}z\right\rangle_+ , \tag{35}$$

$$U_\infty^{\text{Def}}|_{D=0} = \frac{e}{\rho S^2}\left\{(Q-1)\left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\right\rangle_+ E_x(0)\right.$$

$$+\frac{i\omega}{c}\left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\int_0^\infty \left[(Q+1)\right.\right.$$

$$\left.\left.\left.\times \exp\left(-\frac{z}{\tau v_z}\right) - 2\cos kz\right]H_y(z)\mathrm{d}z\right\rangle_+\right\}. \tag{36}$$

At the beginning of this paper and in the title we declared that EMAC for $k\delta<1$ is a surface phenomenon. The equations (35) and (36) show that when sound is excited by the deformation force (13), strictly speaking, the force can be regarded as a surface force only when $\lambda$ is not only greater than $\delta$ but also greater than the mean free path $l$, i.e., $kl \ll 1$. The opposite limit ($kl \gg 1$), which is often encountered under the conditions of the skin effect ($l \gg \delta$), is also very interesting. For this reason we assume that any values of $kl$ are possible.

The problem contains three parameters having the dimension of length: the wavelength of the sound, the penetration depth of the electromagnetic wave, and the mean free path length. In addition, one should remember that the penetration depth $\delta$ is different in the anomalous and normal skin effects:

$$\delta \approx \delta_n = c/(2\pi\sigma\omega)^{1/2}, \quad l \ll \delta_n,$$

$$\approx (l\delta_n^2)^{1/3}, \quad l \gg \delta \gg \delta_n. \tag{37}$$

By assumption, $\lambda \gg \delta$. The mean free path length $l$ can have any value:

1) $\lambda \gg \delta_n \gg l, \quad \omega\tau \ll S^2l^2/v_F^2\delta_0^2, \quad \delta_0^2/l^2,$

2) $\lambda \gg l \gg (l\delta_n^2)^{1/3}, \quad S/v_F \gg \omega\tau \gg \delta_0^2/l^2,$

3) $l \gg \lambda \gg (l\delta_n^2)^{1/3}, \quad \omega\tau \gg S/v_F, \quad \delta_0^2/l^2.$

Each line specifies the conditions which the frequency of the electromagnetic wave must satisfy ($\tau$ is the relaxation time of the electron gas; $\delta_0 = c/\omega_0$; $\omega_0^2 = 4\pi ne^2/m^*$ is the squared plasma frequency; $v_F$ is the Fermi velocity; $n$ is the conduction-electron density; and, $m^*$ is the effective mass of the conduction electrons). For making estimates we started from the Drude–Lorentz–Sommerfeld theory. The condition of quasistatics $\omega\tau \ll 1$ holds for all of the above inequalities.

The condition $\delta \ll \lambda$ enables replacing in Eq. (35) $\cos kz$ by 1. Then

$$U_\infty^{\text{Def}} = -\frac{iek}{\rho S^2}\left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\int_0^\infty \left[1 + (Q-1)\right.\right.$$

$$\left.\left.\times \exp\left(-\frac{z}{\tau v_z}\right)\right] \times E_x(z)\mathrm{d}z\right\rangle_+ . \tag{38}$$

The integrals containing $\exp(-z/\tau v_z)$ depend strongly on the character of the skin effect. Since they are averaged over the Fermi surface and it is obvious that averaging does not distinguish $v_z=0$, it can be assumed that $v_z\tau \sim l$. For this reason for the normal skin effect

$$\int_0^\infty \exp\left(-\frac{z}{\tau v_z}\right)E_x(z)\mathrm{d}z \approx \tau v_z E_x(0), \tag{39}$$

and for the anomalous skin effect

$$\int_0^\infty \exp\left(-\frac{z}{\tau v_z}\right)E_x(z)\mathrm{d}z \approx \int_0^\infty E_x(z)\mathrm{d}z \approx E_x(0)\delta. \tag{40}$$

From Eqs. (38)–(40) we obtain

$$U_\infty^{\text{Def}}|_{D=\infty} \approx -\frac{iek\delta}{\rho S^2}\left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\right\rangle_+ E_x(0)$$

$$\times 1, \quad l \ll \delta,$$

$$\times Q, \quad l \gg \delta. \tag{41}$$

The case of completely diffuse reflection of electrons by the boundary of the metal ($Q=0$) requires, of course, a more detailed analysis. It may seem unsystematic to employ the formulas of the normal skin effect for calculating $U_\infty^{\text{Def}}$. The derivation of Eq. (34) shows that if the normal-skin-effect expressions are employed for $E_x(z)$ and $H_y(z)$, then the expressions obtained are the result of the method of successive approximations in the ratio $l/\delta$, the first nonvanishing term in the expansion being proportional to $l^2$.

The equation (41) describes the linear acoustic response of the metal. It is evident that the magnitude of the effect is determined by the average value of the $xz$ component of the deformation-potential tensor

$$\left\langle \tau^2 \frac{\lambda_{xz}v_xv_z}{1+(k\tau v_z)^2}\right\rangle_+ = \left\langle \frac{\tilde{m}v_x^2(\tau v_z)^2}{1+(k\tau v_z)^2}\right\rangle_+ .$$

This and similar expressions open up the possibility of determining the role of the local geometry of the Fermi surface in EMAC, but this question has still not been examined in detail. The reason is apparently that the appropriate experimental data are not available. If the Fermi surface is assumed to be spherical (for purposes of making estimates), then

$$\langle \ldots \rangle_+ \approx n\frac{\tilde{m}}{m^*}l^2\Phi(kl),$$

$$\Phi(kl) \approx 1, \quad kl \ll 1, \quad \approx (kl)^{-2}, \quad kl \gg 1. \tag{42}$$

Hence and from Eq. (41), assuming $Q \sim 1$, we obtain an equation that demonstrates the crux of the phenomenon:

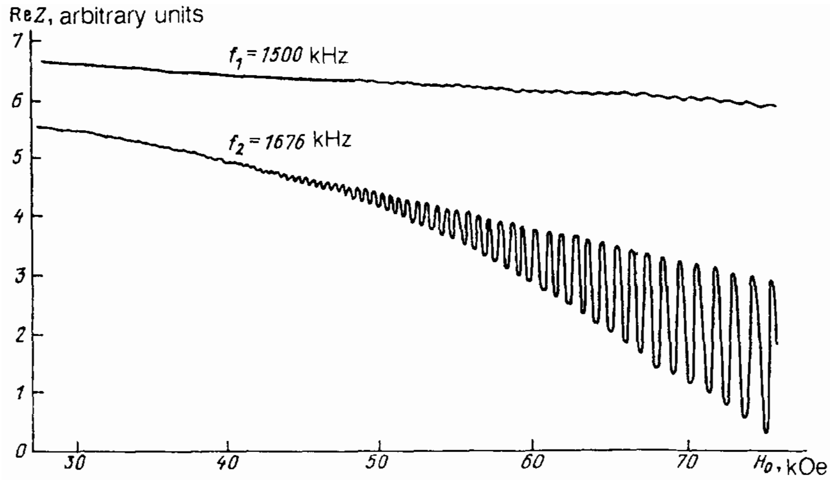



FIG. 4. Quantum oscillations of the real part of the surface impedance of a tin single crystal at a nonresonant frequency ($f_1 = 1500$ kHz) and at the frequency of a standing sound wave in a plate ($f_2 = 1676$ kHz).[7]

$$|U_\infty^{\mathrm{Def}}|_{D=\infty} \simeq l \frac{eE_x(0)l}{\rho S^2} n \frac{\tilde{m}}{m^*} \frac{\delta}{\lambda} \Phi(kl). \qquad (43)$$

For numerical calculations it is convenient to replace $E_x(0)$ by $ZH_y(0)$ and note that $\delta \sim cZ/\omega$.

Comparing Eqs. (36) and (41) shows that Eq. (36) does not contain the factor $\delta/\lambda$, which apparently makes the case of a clamped boundary ($D=0$) more favorable for observing the deformation mechanism of EMAC (at least, for $Q \sim 1$). It seems to us, however, that this assertion requires a more careful analysis that takes into account the structure of the electric field for different values of the parameter $Q$.

## 6. CURRENT EXPERIMENTAL WORK

Ultrasound generation by the Stewart–Tolman force has apparently thus far not been observed in normal metals. Such an experiment would undoubtedly be of fundamental value, since in the inertial mechanism of EMAC electrons can have the "true" mass $m$ and not an effective mass $m^*$ of one or another nature. The question of the manifestation of this conversion mechanism in experiments on surface impedance of superconductors at ultrahigh frequencies[43–47] falls outside the scope of the present paper and essentially requires a separate analysis. As far as ultrasound generation by the Lorentz force is concerned, this question, conversely, has been studied in detail (see, for example, Refs. 1 and 48–50 and the references cited there), and at the present time the induction mechanism of conversion is used mainly as a basis for the contactless method of investigation of the elastic properties of solids and in practical applications (see, for example, Refs. 51 and 52 and the literature cited there). As an example of the application of EMAC in a physical experiment, Fig. 4 displays traces of the quantum oscillations of the surface impedance of a tin single crystal.[7] It is evident that at the frequency of a standing elastic wave over the thickness of the plate the amplitude obtained served oscillations is an order of magnitude higher than the generation amplitude obtained on nonresonance frequencies. The resonance characteristics of the surface impedance of the type shown in Figs. 1 and 2 can also be used for studying quantum oscillations of the velocity and damping of ultrasound.[23,41] For this, the sample, together with the coils encompassing it, must be connected into the positive feedback circuit of a self-excited oscillator.[42] The generation frequency and amplitude of such a device are fixed by the frequency and $Q$ of acoustic resonance in the plate, and these quantities are in turn determined by the speed $S$ and damping $\alpha$ of ultrasound in the metal. Traces of the quantum oscillations of $S$ and $\alpha$ in a tin single crystal are shown in Fig. 5. Contactless methods—EMAC being such a method—yield the most reliable information about subtle physical phenomena, such as quantum oscillations of the elastic properties of metals. In the traditional approach to acoustic measurements it is the presence of sample—transducer contact, unavoidably leading to deformation of the surface, that affects the repeatability of the experimental results.

Electromagnetic acoustic conversion due to deformation interaction is of greatest interest as a phenomenon directly reflecting the dynamics of the electron gas. Analysis of this mechanism of conversion shows that ultrasound generation in the absence of a constant magnetic field occurs at any frequencies (we confine our attention here to the quasistatic case $\omega\tau \ll 1$) and temperatures, and the magnitude of the effect is determined, on the one hand, by the electron-phonon coupling (the tensor $\lambda_{ik}$ is essentially the phonon charge of an electron) and on the other by the relations between parameters having the dimension of length: the thickness $\delta$ of the skin layer, the wavelength $\lambda$ of the elastic wave, and the electron mean free path $l$. The role of nonlocal interactions increases rapidly with increasing $l$ and reaches experimentally measurable values in the regime of the anomalous skin effect. This determines the arrangement of the experiment for studying the deformation mechanism of EMAC: The measurements must be performed on high-quality samples at low temperatures and as high as possible frequencies. Such experiments have been performed on semimetals—bismuth[53–56] and antimony[57]—as well as on many metals—tin,[5,58,59] gallium,[50] aluminum,[61,62] tungsten,[63–65] potassium,[66–72] etc. It is not our purpose to give a complete description of the

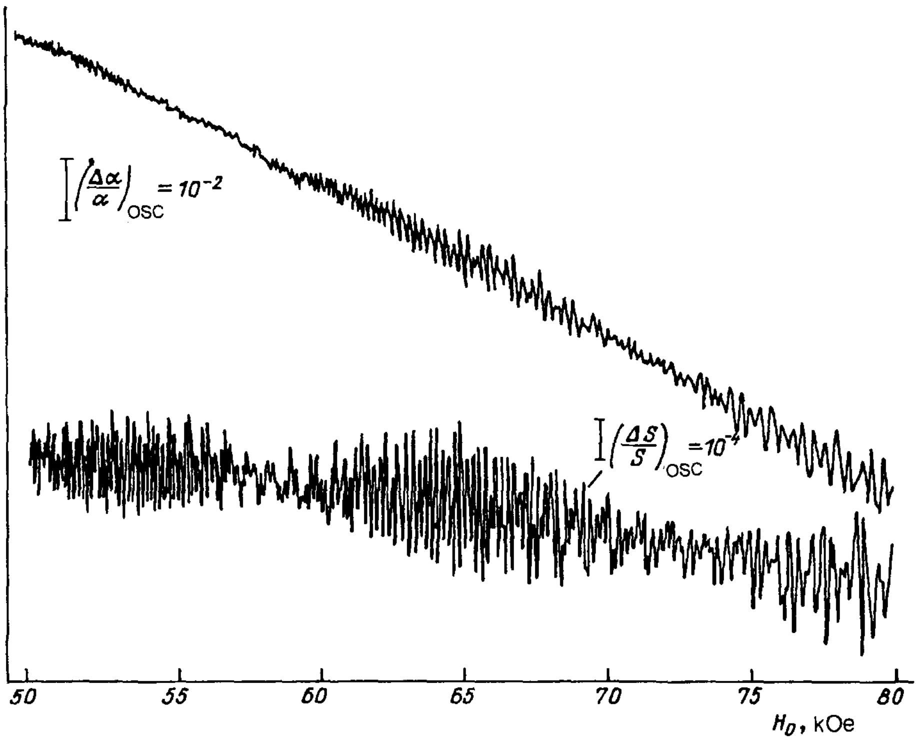



FIG. 5. Quantum oscillations of the speed $S$ and damping $\alpha$ of ultrasound in a tin single crystal.[42]

current experimental situation. We give here only some examples of experimental investigations which exhibit the basic features of the deformation mechanism of EMAC.

*Bismuth.* The peculiarities of the deformation mechanism of electromagnetic excitation of ultrasound in semimetals are determined by the specific nature of the electron spectrum of these materials. The Fermi surface of semimetals consists of electron and hole valleys, separated by a distance in **p** space that is significantly greater than their dimensions. At low temperatures the equilibrium distribution of the carriers in each valley is established over a time much shorter than the time required to establish equilibrium between the valleys. In the case when an electromagnetic wave is incident on the surface of the semimetal the equilibrium distribution of the carriers in the semimetal—both within each valley and between valleys—is destroyed. This means that if the system remains electrically neutral on the whole, the electron density in separate valleys can change. This nonequilibrium, which is specific to semimetals and is associated with the appearance of gradients of the carrier density, leads to the appearance of a deformation force. The specific nature of EMAC in semimetals is most pronounced in the giant quantum oscillations of the damping of ultrasound and the conversion efficiency in bismuth.[4] Traces of these quantities are displayed in Fig. 6. The conversion efficiency $\eta$, defined as the ratio of the energy flux in the excited elastic wave to the energy flux in the incident electromagnetic wave, is proportional to the squared amplitude of the ultrasound $U_\infty$. It was found that, first, additional features are observed in the field dependence of the conversion efficiency—period doubling of the oscillations—and, second, the relative magnitude of the effect in oscillations of the efficiency was an order of magnitude greater than in the oscillations of the damping $\alpha$. These phenomena are explained in Ref. 73.

*Tin.* The temperature dependences (Fig. 7) of the generation amplitude $U_\infty$ and damping $\alpha$ of transverse ultrasound in a tin single crystal indicate that EMAC is closely associated with the acoustic properties of metals.[5] At low temperatures (for the present experiment $T<8$ K) $U_\infty$ and $\alpha$ are virtually independent of the temperature, since the mean free path in this region is determined by the scattering of electrons by impurities. As the temperature increases thermal phonons are included in the scattering processes, and at $T\geqslant 8$ K the plots of $\alpha(T)$ and $U_\infty(T)$ are described by power-law functions of the type $T^{-n}$. The exponent $n=3.3\pm0.1$ for $\alpha$ and $n=6.5\pm0.1$ for $U_\infty$. Thus in the temperature interval studied $U_\infty\sim\alpha^2$, and if it is assumed[74] that the damping is proportional to the mean free path length, then $U_\infty\sim l^2$. The parameters of the ex-

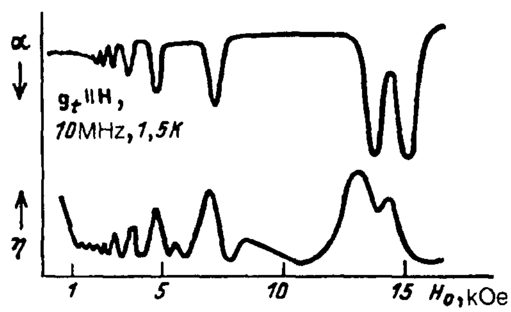



FIG. 6. Giant quantum oscillations of the damping $\alpha$ of ultrasound and the EMAC efficiency $\eta$ in a bismuth single crystal. $\mathbf{k}\|\mathbf{H}_0$, $f=10$ MHz, $T=1.5$ K.[4]

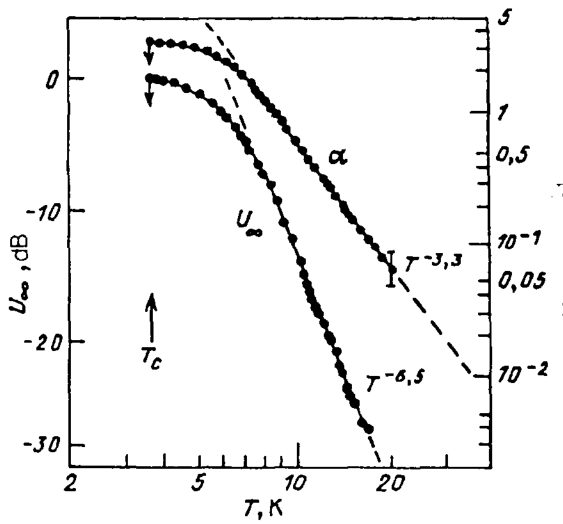



FIG. 7. Temperature dependences of the damping $\alpha$ of transverse ultrasound and generation amplitude $U_\infty$, normalized to the superconducting transition temperature in tin ($T_c=3.73$ K). The dashed lines at temperatures $T>8$ K represent power-law fits to the data points.[5]

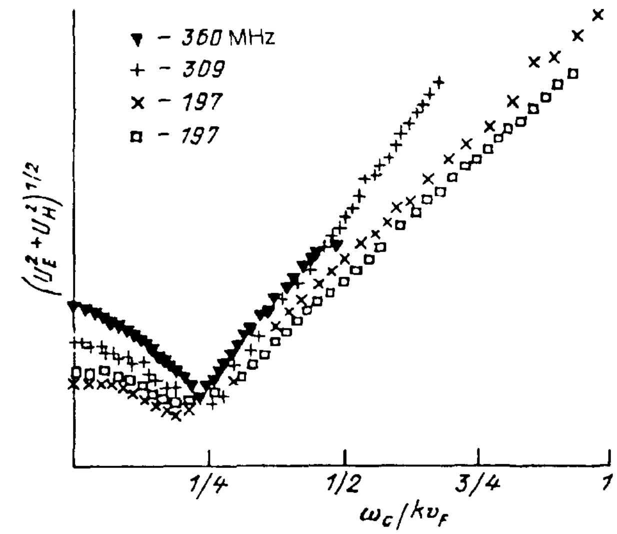


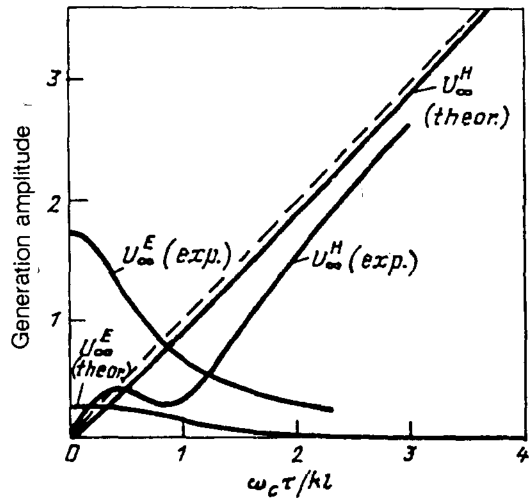

FIG. 8. Modulus of the generation amplitude of transverse ultrasound in aluminum versus the reduced magnetic field at $T=4.2$ K.[62]

perimental sample were characterized by the following values: in the normal state, near the transition into the superconducting state ($T_c=3.72$ K) $l=1.7\cdot10^{-3}$ cm and the thickness of the skin layer $\delta=4\cdot10^{-5}$ cm at the frequency $f=16$ MHz. The parameter $kl$ increases from 0.04 at $T=15$ K up to 1 at $T=T_c$, while $k\delta$ decreases from 0.04 at $T=15$ K to 0.02 at $T=T_c$. Thus in the experiment performed $kl\leqslant1$ and $k\delta\ll1$. The case realized in the experiment corresponds qualitatively to Eq. (43), according to which $U_\infty\sim l^2$, obtained from the top line of Eq. (42).

*Aluminum.* Excitation of transverse ultrasound in aluminum at comparatively high frequencies ($f=90$–$400$ MHz) has been investigated in detail at liquid-helium temperatures.[62] The aim of this work was to determine the absolute values of the efficiency of the deformation mechanism of EMAC, as well as to compare these data to the results obtained in low-frequency measurements.[61] Due to the strong electronic damping of ultrasound in metals at high frequencies the measurements were performed on thin ($d=3.4$ and 2.8 $\mu$m) polycrystalline films, deposited by sputtering Al in vacuum on plane-parallel surfaces of a sapphire single crystal. The main result of the measurements is displayed in Fig. 8. The absolute value of the efficiency of the deformation mechanism of EMAC was determined at the frequency $f=90$ MHz by calibration with respect to measurements in a strong magnetic field ($\omega_c/kv_F\approx2$), when EMAC is due entirely to the induction interaction.

In order to compare the results of different experiments with one another and to the theory the carrier mean free path length or, ultimately, the nonlocality parameter $kl$ must be estimated. The carrier mean free path length in aluminum films was calculated on the basis of the Fuchs–Sondheimer size-effect theory.[75] At liquid-helium temperatures $l=15\pm3$ $\mu$m in both experimental films; this leads to nonlocality parameter $kl=6$ even at $f=200$ MHz. In this situation the amplitude of the excited ultrasound must increase linearly with frequency. It is difficult to determine the frequency dependence from the experimental data presented, but it is important to note that in the absence of a constant field $H_0=0$ the experimental values of the conversion efficiency at all frequencies were an order of magnitude higher than the values estimated from the free-electron model, just as the results of low-frequency measurements.[61] Both the low-frequency[61] and high-frequency[62] experiments were performed at close values of the parameter $kl$. It can be conjectured (compare to Refs. 8 and 9) that the boundaries of the film with the insulating substrate did not significantly affect the EMAC efficiency, so that the EMAC efficiency with $D=0$, $Q=0$ was close to that in the case $D=\infty$, $Q=1$.

*Potassium.* One of the most remarkable pages in the history of EMAC research is the "potassium" problem—the problem of the quantitative disagreement between the experimental and theoretical data on the generation amplitude of transverse ultrasound in potassium. This problem arose soon after the first low-temperature measurements were performed on normal metals, and it continues to attract attention. The crux of the problem is that for a metal with a spherical Fermi surface it is comparatively easy to calculate the generation amplitude[76] and this quantity can be compared to the experimental results. It is also easy to calculate the frequency, temperature, and field dependences in the case when the constant magnetic field is perpendicular to the surface. They can also be compared to experiment. We are accustomed to the fact that in situations when a numerical comparison is made between theory and experiment it most often turns out that experiment "does not reach" the theory and, as a rule, many explanations can be found for why this happens. The opposite situation, when experimental results significantly exceed theoretical estimates, occurs much more rarely. This is what happened in potassium; the discrepancy between the-



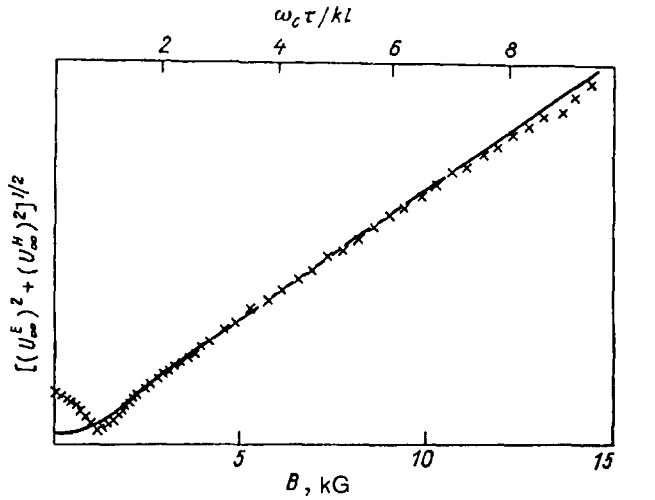




FIG. 9. Field dependences of the generation amplitude of transverse ultrasound in potassium at $T=4.2$ K. $U^E_\infty$(exp) and $U^H_\infty$(exp) are the experimental results. The curves $U^E_\infty$(theor) and $U^H_\infty$(theor) were calculated on the basis of the free-electron model with the nonlocality parameter $kl=1.9$.[68]

FIG. 10. Field dependence of the modulus of the generation amplitude of transverse ultrasound in potassium at $T=4.2$ K and $f=8.97$ MHz. The solid curve was calculated on the basis of the free-electron model with the nonlocality parameter $kl=4.5$.[69]

ory and experiment reached the point where quantitative agreement changes to qualitative agreement.[77]

The main result of the first experimental investigations of EMAC in potassium[66,67] was the very fact of generation of ultrasound in metal in the absence of a constant magnetic field. It was found that in a weak magnetic field the excited transverse elastic waves are polarized along the alternating electric field vector $U^E_\infty$, while in a strong magnetic field they are polarized along the vector of the alternating magnetic field $U^H_\infty$. The results agreed qualitatively with the understanding of the physical processes responsible for EMAC, indicating that in weak magnetic fields and with $H_0=0$ ultrasound is generated by the deformation interaction, while in strong fields ultrasound is generated by the induction mechanism (compare with Eqs. (21) and (43)). The work that actually led to the potassium problem was the experiment of Ref. 68, performed on a single crystal whose [110] axis was perpendicular to the surface. The measurements of the generation amplitude of the fast transverse mode were performed by the echo-pulse method at the frequency $f=9.4$ MHz. The amplitudes of the electric $U^E_\infty$ and magnetic $U^H_\infty$ components of the displacement were measured for two mutually perpendicular positions of the linearly polarized measuring coil. The quality of the sample was characterized by the parameter $kl=1.9$ at $T=4.2$ K. The results of the measurements and the curves calculated on the basis of the free-electron model[76] are displayed in Fig. 9. It is evident that $U^E_\infty$ is maximum with $H_0=0$ and decays rapidly when the magnetic field is switched on. The theoretical curve also has, in principle, the same form, but the experimental value of $U^E_\infty$ with *$H_0=0$ is almost an order of magnitude higher than the* computed value. The experimentally observed field dependence of $U^H_\infty$ is substantially different from the value computed on the basis of the free-electron model. The signal from this polarization at first increases rapidly with increasing magnetic field, passes successively through a maximum and a minimum, and finally, in a strong magnetic field it reaches the linear asymptotic relation characteristic for the induction mechanism of EMAC. We note that it is the generation of ultrasound in strong fields that was employed for calibrating the generation amplitude in weak fields and with $H_0=0$. The next step in the experimental investigation of the potassium problem was Ref. 69, in which the field dependence of the modulus of the amplitude of the excited ultrasound $U_\infty=[(U^E_\infty)^2+(U^H_\infty)^2]^{1/2}$ was measured for the fast transverse mode propagating in the [110] direction. The main result of this work is shown in Fig. 10. This figure also shows the calculation based on the free-electron model, performed using the nonlocality parameter $kl=4.5$. Significant discrepancies between theory and experiment are observed in a weak magnetic field, and with $H_0=0$ in a strong magnetic field the experimental and theoretical dependences are identical. Similar results for the slow transverse mode, propagating in the [110] direction, were obtained in Ref. 70 and very recently in Ref. 72, where it is also shown that the amplitude of nonlocal generation in potassium decays rapidly with increasing temperature. Completing this section, as well as this report, we note that in virtually all works in which theory was compared to experiment it was found that the experimentally determined EMAC efficiency was significantly higher than the theoretical value. Though one can try to explain this discrepancy in aluminum[61–62] and zinc[78] by the structural characteristics of the Fermi surface and the deviation of this surface from the free-electron model, many experiments on potassium—a metal with a spherical Fermi *surface—do not allow for this possibility. This in turn suggests* that the mass of the electron-phonon interaction in metals is significantly higher than not only the carrier effective masses but also the free-electron mass. This is an important result, since electron motion on the Fermi surface and the corresponding effective masses have been well studied for most metals, whereas the electron-phonon interaction still awaits a detailed experimental investigation. The phenomenon of direct electromagnetic-acoustic conversion could be very helpful for such a study.

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