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Thermoelectric mechanism of electromagnetic-acoustic transformation in organic conductors

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Abstract – The thermoelectric mechanism of electromagnetic-acoustic transformation of the energy in an organic conductor with a quasi–two-dimensional electron energy spectrum (Q2D) placed in an external magnetic field has been considered. The amplitude of the acoustic wave excited by the temperature oscillations in a Q2D organic conductor was calculated for both the isothermal and the adiabatic thermal boundary condition. Angular oscillations of the amplitude resulting from the periodic dependence of the electron velocity on the angle between the normal to the layers and the magnetic field has been observed as expected. A comparison with the inductive mechanism of EMAT is made in order to determine the conditions at which the thermoelectric mechanism is dominant over the inductive one in the presence of a magnetic field. The thermoelectric mechanism of EMAT allows new important information on the electronic structure of the organic layered conductors to be obtained.

Introduction. – The search for new materials in the sixties attracted the attention of researches to conductors of organic origin. Most of them have a metal-type electrical conductivity, and their electron energy spectrum can be studied with the help of methods developed for metals. For example, these methods can be rightfully used for studying transport phenomena in organic conductors [1–3], having a layered structure with a sharply pronounced anisotropy in the electrical conductivity: the conductivity across the layers is much smaller than the conductivity in the layer plane. The strong anisotropy of electrical conductivity is apparently associated with a strong anisotropy in the velocity of charge carriers on the Fermi surface, *i.e.*, their energy

$$\varepsilon(\mathbf{p}) = \sum_{n=0}^{\infty} \varepsilon_n(p_x, p_y) \cos\left(\frac{anp_z}{\hbar}\right) \qquad (1)$$

depends weakly on the momentum projection $p_z = \mathbf{pn}$ on the normal n to the layers (a is the distance between adjacent layers, and $\hbar$ is Planck's constant). In tight-binding approximation, the functions $\varepsilon_n(p_x, p_y)$ with $n \geqslant 1$ are much less than the Fermi energy ε_F and fall rapidly with increasing n.

A series of effects exists which are typical for layered conductors of organic origin with metal-type conductivity, but non-existent for ordinary metals. Among these effects, arising from the Q2D nature of the charge carrier energy spectrum, is the high acoustic transparency of an organic conductor in the absence of charge carrier drift along the acoustic wave vector as well as the orientation effect —a strong dependence of kinetic parameters on the orientation of the magnetic field with respect to the layers.

The investigation of transport phenomena in layered conductors is of great interest because it opens new possibilities of determining the electronic structure of these conductors in fine details. The theoretical and experimental investigations of the electrical conductivity and magnetoresistance in Q2D conductors have been subject of an enormous number of studies [4–11]. In ref. [12], a detailed theoretical report is given on the propagation of electromagnetic and acoustic oscillations in organic Q2D conductors. As far as we are aware, there are no experimental results on acoustoelectronic phenomena in Q2D conductors yet. In recent years there has been a rising interest in the research on thermoelectric effects in organic Q2D conductors. A number of publications appeared in which the results of experimental (see ref. [13]) and theoretical studies [14–16] of the thermoelectric coefficients were reported.

It is well known that electromagnetic and acoustic waves in conductors form a coupled system that permits their mutual transformation [17,18]. When an electromagnetic wave of frequency ω is incident on a conducting half-space $z \geqslant 0$ an acoustic wave with the same frequency is excited in the conductor, traveling in the positive z-direction. This phenomenon is known as an electromagnetic-acoustic transformation (EMAT) of the energy [19]. Few mechanisms of EMAT exist in the conducting media. In the absence of an external magnetic field the inertial mechanism is a basic one [20], and in the presence of an external magnetic field there exist deformation, inductive and thermoelectric mechanisms [19,21]. The inertial mechanism is connected with the non-inertiality of the concomitant coordinate system, which moves together with the crystal lattice (Stewart-Tolman effect), the deformation mechanism is associated with the energy renormalization for charge carriers when the crystal is deformed, and the inductive mechanism is due to the Lorentz force acting on the conduction electrons. The thermoelectric mechanism of EMAT occurs when the condition for normal skin effect, $l \ll \delta_s$ (where l is the mean free path of the conduction electrons and δ_s is the skin depth of the electromagnetic field), is fulfilled. In that case when an electromagnetic wave of frequency ω is incident on the surface of the conductor, whose symmetry axis does not coincide with the normal to the surface, nonuniform temperature oscillations of the same frequency appear due to the thermoelectric effect. These oscillations generate thermoelectric stresses which induce acoustic oscillations in the conductor. In ordinary metals the thermoelectric mechanism of EMAT has been well studied [22–24]. For what concerns the Q2D organic conductors, few years ago a detailed publication appeared on EMAT due to the deformation mechanism in these conductors [25]. Assuming the present interest in the research of thermoelectric phenomena in Q2D organic conductors as well as the fact that to date there is no investigation made on the thermoelectric mechanism of EMAT in these conductors, this paper deals with theoretical analyses of the dependence of the acoustic wave amplitude excited by the temperature oscillations in Q2D organic conductors on the magnitude and orientation of an external magnetic field. Two cases of boundary thermal conditions, isothermal and adiabatic will be considered. The amplitude of the acoustic wave associated with the inductive mechanism will also be calculated in order to compare the nature of the attenuation of both the waves in the presence of an external magnetic field.

Formulation of the problem. – In this section we will present the complete system of equations necessary for solving the problem of EMAT due to both the inductive and thermoelectric mechanisms. Suppose that an electromagnetic wave ($E_x = E_z = 0$, $E_y = E$) of frequency ω is incident normally on a surface (along the less conductivity axis (z-axis) *i.e.*, $\mathbf{k} = (0, 0, k)$) of a Q2D organic conductor placed in a magnetic field tilted at an angle $\pi/2 - \theta$ with respect to the conductor's surface, $\mathbf{B} = (B\sin\theta, 0, B\cos\theta)$. The wave is taken to be monochromatic, so the differentiation with respect to the time variable is equivalent to multiplication by $(-i\omega)$. The frequency of the electromagnetic field is constraint with the condition $\omega\tau \ll 1$ (τ is the relaxation time of the conduction electrons), which is, as a rule, always fulfilled when an EMAT is consider.

In a conducting medium, the Lorentz force is a source of longitudinal as well as transversal acoustic oscillations. Apart from that, the temperature oscillations induce only longitudinal acoustic oscillations. In order to make a comparison between both mechanisms, we need to calculate the amplitudes of the longitudinal acoustic waves excited by each of them.

In the absence of the thermoelectric mechanism the basic one is the inductive mechanism of EMAT. In that case the system of equations necessary for describing the propagation of the acoustic wave in a conductor, excited by the incident electromagnetic wave, contains the equation of the theory of elasticity for ionic displacement $\mathbf{U}$:

$$\varrho \frac{\partial^2 U_i}{\partial t^2} = \lambda_{iklm} \frac{\partial U_{lm}}{\partial x_k} + F_i, \tag{2}$$

as well as Maxwell's equations

$$\mathrm{curl}\mathbf{B} = \mu_0 \mathbf{j}; \qquad \mathrm{curl}\mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t}. \tag{3}$$

Here ϱ and λ_{iklm} are the density and elastic tensor of the crystal, μ_0 is the magnetic permeability of the vacuum, U_{lm} is the deformation tensor, and $F_i = (\mathbf{j} \times \mathbf{B})_i$ is the Lorentz force.

Since the electric field of the incident electromagnetic wave is along the y-axis it follows that the only nonzero component of the current density is its y-component, *i.e.*, $\mathbf{j} = (0, j, 0)$. The surface of the conductor is assumed to be free, so that $\frac{\partial U_{zl}}{\partial z}|_{z=0} = 0$. Consequently, the system of equations above can be rewritten in the following form:

$$\omega^2 \varrho U_{zl} = \lambda_{zzzz} k^2 U_{zl} + j_y B \sin\theta, \tag{4}$$

$$\frac{\partial B_x}{\partial z} = \mu_0 j_y; \qquad -\frac{\partial E_y}{\partial z} = i\omega\ B_x, \tag{5}$$

where $\lambda_{zzzz} = s^2 \varrho$, and s is the acoustic-wave velocity.

For the amplitude of the acoustic wave excited by the temperature oscillations in the Q2D organic conductor U_{zl}^{T}, to be calculated, the temperature distribution associated with the oscillations of the electromagnetic field must be determined first.

The complete system of equations describing the propagation of the acoustic wave in the conductor in the presence of the heat flux $\mathbf{Q}$ consists of the thermal conduction equation

$$C \frac{\partial \Theta}{\partial t} + \mathrm{div}\mathbf{Q} = 0, \tag{6}$$

accompanied by the equation for the heat flux, which takes into account the thermoelectric effects,

$$Q_i = k_B T \alpha_{ik} j_k - \kappa_{ik} \frac{\partial \Theta}{\partial x_k}, \tag{7}$$

and the equation of the theory of elasticity including the thermoelectric stress tensor $\sigma^{\mathtt{T}}_{ik} = -\varrho s^2 \beta \Theta(z,t)\delta_{ik}$:

$$\varrho \frac{\partial^2 U_i}{\partial t^2} - \lambda_{iklm} \frac{\partial U_{lm}}{\partial x_k} = -\varrho s^2 \beta \delta_{ik} \frac{\partial \Theta}{\partial x_k}. \tag{8}$$

Here, Θ is linear in the amplitude of the electromagnetic-wave additive term to the equilibrium temperature T, C is the volumetric heat capacity, k_B is the Boltzmann constant, α_{ik} is the thermoelectric power tensor, κ_{ik} is the thermal conductivity tensor, and β_V is the volumetric thermal expansion coefficient.

For the geometry under consideration (the electromagnetic wave is incident normally on the surface of the conductor) all of the quantities depend only on z and the above set of equations take the following form:

$$-i\omega C\Theta + \frac{\partial Q_z}{\partial z} = 0, \tag{9}$$

$$Q_z = k_B T \alpha_{zy} j_y - \kappa_{zz} \frac{\partial \Theta}{\partial z}, \tag{10}$$

$$\frac{\partial^2 U^{\mathtt{T}}_{zl}}{\partial z^2} + q^2 U^{\mathtt{T}}_{zl} = \beta \frac{\partial \Theta}{\partial z}, \tag{11}$$

where $q = \omega/s$.

By substituting eq. (10) in eq. (9) one obtains the following differential equation for the temperature distribution in the conductor:

$$\frac{\partial^2 \Theta}{\partial z^2} + \frac{i\omega C}{\kappa_{zz}}\Theta = \frac{k_B T \alpha_{zy} \frac{\partial j_y(z)}{\partial z}}{\kappa_{zz}}. \tag{12}$$

Equations (11) and (12) must be supplemented with the boundary condition taking into account the thermal conditions at the surface. We shall consider two boundary conditions: isothermal and adiabatic. Consequently, the boundary conditions for the temperature are

$$\Theta(z=0) = 0, \quad \text{and} \quad \frac{\partial \Theta}{\partial z}\mid_{z=0} = 0, \tag{13}$$

and for the amplitude of the acoustic wave in the conductor are

$$\frac{\partial U^{\mathtt{T}}_{zl}}{\partial z}\mid_{z=0} = 0, \quad \text{and} \quad \frac{\partial U^{\mathtt{T}}_{zl}}{\partial z}\mid_{z=0} = -\beta\Theta\mid_{z=0}, \tag{14}$$

respectively.

Calculation of the amplitude $U^{\mathtt{I}}_{zl}$. – The amplitude of the acoustic wave due to the inductive mechanism can be calculated by using eqs. (4) and (5). The current density is obtained from Maxwell's equations, and has the form

$$j_y = -\frac{ik^2}{\omega\mu_0} e^{i(kz-\omega t)}, \tag{15}$$

where

$$k = \frac{1+i}{\delta_s}, \quad \text{and} \quad \delta_s = \sqrt{\frac{2\rho_{yy}}{\omega\mu_0}}. \tag{16}$$

The $\rho_{yy} = 1/\sigma_{yy}$ component of the electric resistivity tensor ρ_{ij} can be derived by means of the σ_{yy} component of the electrical conductivity tensor σ_{ij}. In a magnetic field the components of the conductivity tensor, which relate the current density to the electric field $\mathbf{E}$,

$$j_i = \sigma_{ij} E_j, \tag{17}$$

can be found using the Boltzmann transport equation in the τ approximation for the collision integral [26]. The equation above is valid only in the case of normal skin effect when the relation between current density and electric field is local to a high degree of accuracy.

Without any model assumptions about the electron energy spectrum, the quasi-classical expression for σ_{ij} in the case of periodic motion of a charge with period $T_B = 2\pi/\Omega$ in a magnetic field $\mathbf{B}$, and in the main approximation in the quasi–two-dimensionality parameter $\eta \ll 1$ has the form

$$\sigma_{ij} = \frac{2e^3 B}{(2\pi\hbar)^3} \int \mathrm{d}p_B \int_0^{T_B} \mathrm{d}t v_i(t) \int_{-\infty}^{t} \mathrm{d}t' v_j(t') e^{\frac{t'-t}{\tau}}. \tag{18}$$

Here e is the charge of the conduction electrons, $\Omega = eB/m^*$ is the gyration frequency of an electron in the magnetic field $\mathbf{B}$, m^* is its cyclotron effective mass, $p_B = p_x \sin\theta + p_z \cos\theta = \text{const}$ is the momentum projection in the magnetic-field direction, and t is the time motion of the conduction electrons in the magnetic field under the influence of the Lorentz force $\mathrm{d}\mathbf{p}/\mathrm{d}t = e(\mathbf{v}\times\mathbf{B})$.

In the case of a layered conductor whose electron energy spectrum has the form (only the first two terms in eq. (1))

$$\varepsilon(p) = \frac{p_x^2 + p_y^2}{2m} - \eta \frac{v_F \hbar}{a} \cos\frac{a p_z}{\hbar}, \tag{19}$$

where v_F is the characteristic Fermi velocity of the electrons along the layers, the expression for σ_{yy} in the main approximation in the small parameter of the quasi–two-dimensionality of the electron energy spectrum η takes the form

$$\sigma_{yy} = \rho_{yy}^{-1} = \frac{\sigma_0}{(\Omega\tau)^2 \cos^2\theta + 1}. \tag{20}$$

Here σ_0 coincides in order of magnitude with the electrical conductivity along the layers in the absence of a magnetic field.

Substituting eq. (15) in eq. (4), and using eqs. (16) and (20), one obtains the following expression for the amplitude of the acoustic wave associated with the Lorentz force:

$$U^{\mathtt{I}}_{zl} = \frac{iq^2 B \sin\theta}{\omega\mu_0\rho\left\{\frac{q^2}{i\omega\mu_0\sigma_0}\left(1 + \left(\frac{e\tau B\cos\theta}{m^*}\right)^2\right) - 1\right\}}. \tag{21}$$

Calculation of the amplitude $U^{\mathtt{T}}_{zl}$. – In order to calculate the amplitude of the acoustic wave excited by the temperature oscillations in a Q2D conductor, we must first determine the temperature distribution by means of

$$v_z = \eta v_F \left\{ \sin\xi \left[J_0(\zeta\tan\theta) + 2\sum_{i=1}^{\infty} J_{2i}(\zeta\tan\theta)\cos(2i\Omega\cos\theta t) \right] - \cos\xi \sum_{i=0}^{\infty} J_{2i+1}(\zeta\tan\theta)\sin((2i+1)\Omega\cos\theta t) \right\}. \quad (32)$$

eq. (12). Using the expression for the current density, eq. (15), eq. (12) can be rewritten in the following form:

$$\frac{\partial^2\Theta}{\partial z^2} + k_{\mathrm{T}}^2\Theta = \frac{k_B T \alpha_{zy} k^3}{\omega\mu_0\kappa_{zz}} e^{ikz}, \quad (22)$$

where

$$k_{\mathrm{T}} = \frac{1+i}{\delta_{\mathrm{T}}}, \quad \text{and} \quad \delta_{\mathrm{T}} = \sqrt{\frac{2\kappa_{zz}}{\omega C}}. \quad (23)$$

Here δ_{T} is the depth of penetration of the thermal field in the conductor.

The solution of eq. (22), satisfying the boundary conditions (eq. (13)) is given by

$$\Theta(z) = -\frac{k_B T \alpha_{zy} k}{\omega\mu_0\kappa_{zz}} \frac{1}{\left(1 - \frac{C\rho_{yy}}{\mu_0\kappa_{zz}}\right)} (e^{ikz} - b e^{ik_{\mathrm{T}} z}), \quad (24)$$

where b is 1 for an isothermal boundary condition, and $\delta_{\mathrm{T}}/\delta_s$ for an adiabatic boundary condition.

After the appropriate substitution for the temperature distribution $\Theta(z)$ and the boundary conditions (eq. (14)) in eq. (11), calculations yield the following expression for the amplitude of the acoustic wave excited by the temperature oscillations in a Q2D conductor:

$$U_{zl}^{\mathrm{T}} = \frac{ik_B T\beta\alpha_{zy} k}{\omega\mu_0 q}\Xi_{i,a}, \quad (25)$$

where

$$\Xi_i = \frac{iq\sqrt{2\omega\mu_0\rho_{yy}}}{(1+i)\left(1+\sqrt{\frac{C\rho_{yy}}{\mu_0\kappa_{zz}}}\right)} A, \quad (26)$$

for an isothermal boundary condition, and

$$\Xi_a = \frac{\sqrt{\frac{\mu_0\rho_{yy}\kappa_{zz}}{C}}}{\left(1+\sqrt{\frac{C\rho_{yy}}{\mu_0\kappa_{zz}}}\right)} \left\{ 1 + i\omega\sqrt{\frac{C\rho_{yy}}{\mu_0\kappa_{zz}}} A \right\} \quad (27)$$

for an adiabatic boundary condition. Here the notation is

$$A = \frac{q\left[q - \frac{1}{\delta_s} - \frac{1}{\delta_{\mathrm{T}}}\right] + i\omega\sqrt{\frac{\mu_0}{\rho_{yy}}}\left[\sqrt{\frac{C}{\varkappa_{zz}}} - \frac{q}{\sqrt{2\omega}}\left(1+\sqrt{\frac{C\rho_{yy}}{\mu_0\varkappa_{zz}}}\right)\right]}{(q^2\rho_{yy} - i\omega\mu_0)(q^2\varkappa_{zz} - i\omega C)}. \quad (28)$$

In the frame of the Boltzmann transport theory the thermoelectric power is determined by the electrical conductivity as follows:

$$\alpha_{zy} = \frac{\pi^2 k_B T}{3e}\frac{\mathrm{d}\sigma_{zy}(\varepsilon)}{\mathrm{d}\varepsilon}\Big|_{\varepsilon=\mu}, \quad (29)$$

where μ is the chemical potential of the electron system.

For the model dispersion relation (eq. (19)), the electron velocity along the normal to the layers takes the form

$$v_z = \eta v_F\{\sin\xi\cos(\zeta\tan\theta) - \cos\xi\sin(\zeta\tan\theta)\}, \quad (30)$$

where

$$\xi = \frac{ap_B}{\hbar\cos\theta} \quad \text{and} \quad \zeta = \frac{2a\mu}{\hbar v_F}. \quad (31)$$

We use identities to substitute the Bessel generating functions [27] and to obtain

see eq. (32) *above*

Substituting eq. (32) in eq. (18), we note the terms that survive when $i = n$, and since integrals such as $\int_0^{2\pi} \mathrm{d}\phi\cos(2i\phi)\cos(2n\phi) = \pi\delta_{in}$, where δ_{in}, is the Kronecker delta, one obtains

$$\sigma_{zy} = \frac{1}{2}\eta^2\sigma_0\frac{(\Omega\tau\cos\theta)^2}{1+(\Omega\tau\cos\theta)^2} J_0(\zeta\tan\theta)J_1(\zeta\tan\theta). \quad (33)$$

Performing the derivate over μ yields the final expression for the thermoelectric power α_{zy} of the form

$$\alpha_{zy} = \eta^2\sigma_0\frac{\pi^2 k_B T}{3e\mu}\frac{(\Omega\tau\cos\theta)^2\tan\theta}{1+(\Omega\tau\cos\theta)^2}\left\{ -J_1^2(\zeta\tan\theta) + \frac{1}{2}(J_0(\zeta\tan\theta) - J_2(\zeta\tan\theta))J_0(\zeta\tan\theta)\right\}. \quad (34)$$

An electron system with elastic scattering generally obeys the Wiedemann-Franz law [26]

$$\kappa_{zz} = \frac{\pi^2 k_B^2 T}{3e^2}\sigma_{zz}. \quad (35)$$

The σ_{zz} component of the electrical conductivity can be derived by following the procedure described above when calculating σ_{zy}. After the appropriate substitution, the Wiedemann-Franz formula takes the form

$$\kappa_{zz} = \frac{\pi^2 k_B^2 T}{3e^2}\eta^2\sigma_0\left\{ J_0^2(\zeta\tan\theta) + 2\sum_{i=0}^{\infty}\frac{J_i^2(\zeta\tan\theta)}{1+(i\Omega\tau\cos\theta)^2}\right\}. \quad (36)$$

Discussion. – The amplitudes of the excited acoustic waves by both the inductive and thermoelectric mechanisms are functions of the frequency of the incident electromagnetic wave ω, the magnetic field B, the tangent of the angle between the normal to the layer and the magnetic field $x = \tan\theta$, as well as of the characteristics of the Q2D conductor (resistivity, thermoelectric power and

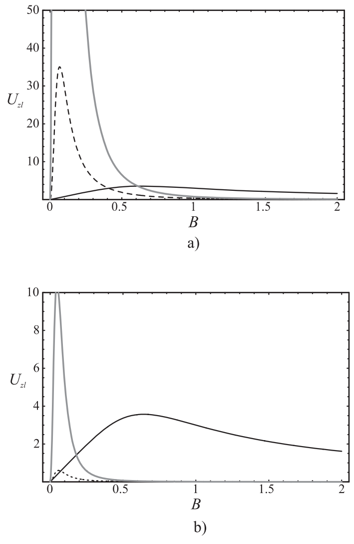

Fig. 1: The dependence of the amplitude of the acoustic wave due to the inductive (solid curve) and thermoelectric mechanism of EMAT for an isothermal (dashed curve) and adiabatic (gray curve) thermal boundary condition on the magnetic field B at $x = 10$, $T = 300\,\mathrm{K}$ and a) $\eta = 0.1$, b) $\eta = 0.01$.

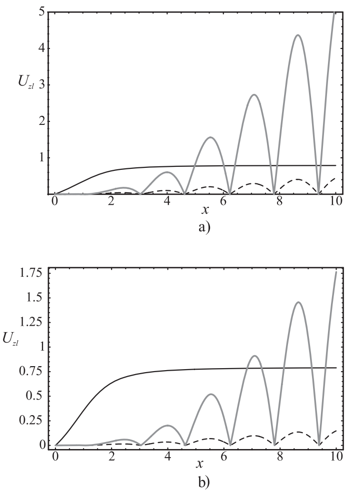

Fig. 2: The dependence of the amplitude of the acoustic wave due to the inductive (solid curve) and thermoelectric mechanism of EMAT for an isothermal (dashed curve) and adiabatic (gray curve) thermal boundary condition on $x = \tan\theta$ at $\eta = 0.01$, $B = 0.1\,\mathrm{T}$, and a) $T = 300\,\mathrm{K}$, b) $T = 100\,\mathrm{K}$.

thermal conductivity). We shall analyze the dependence of the amplitudes $U^{\mathtt{I}}_{zl}$ and $U^{\mathtt{T}}_{zl}$ on the magnitude of the magnetic field and its orientation with respect to the layers. The effectiveness of the excitation of the acoustic wave associated with the inductive mechanism is the highest when $q \ll |k|$ until the condition for normal skin effect $kl \ll 1$ is fulfilled. The acoustic wave is controlled by electrons whose velocities coincide with the Fermy velocity, $v_y = v_F$, *i.e.*, by electrons in the plane of the layer. These are the electrons which are in phase with the acoustic wave as their velocities are normal to the acoustic wave vector $\mathbf{vq} = 0$. At $B = \frac{smx}{e\tau}\sqrt{\frac{\mu_0\sigma_0}{\omega}}$ the interaction of the conduction electrons with the acoustic wave is most effective and the amplitude $U^{\mathtt{I}}_{zl}$ reaches its maximum value. On the other hand, the acoustic wave excited by the thermoelectric effect is controlled by electrons near the extremum v_z of the Fermi surface, reached at $p_z = \pi\hbar/2a$. In Q2D conductors, in which the Fermi surface is a weakly corrugated cylinder open in the direction of the propagation of the acoustic wave, these are the electrons whose velocity component in the direction of propagation of the wave reaches the highest value, $v_z^{max} = \eta v_F$. The interaction of these electrons with the acoustic wave is most effective when $q \ll | k_{\mathtt{T}} |$ until the condition $k_{\mathtt{T}} l \ll 1$ is satisfied and for not very strong corrugation of the Fermi surface as in that case the maximum velocity of the electrons moving along the wave vector $v_z^{max} = \eta v_F$ is larger. The amplitude $U^{\mathtt{T}}_{zl}$ reaches its maximum at $B = \frac{mx}{e\tau}\sqrt{\frac{2\alpha^2\omega\tau\sin^2(x-\pi/4)-\pi x}{\pi x+2\alpha^2\omega\tau\cos^2(x-\pi/4)}}$, where $\alpha = \eta v_F/s$. For larger magnetic fields the acoustic wave is strongly attenuated, and the inductive mechanism dominates over the thermoelectric one as shown in fig. 1.

The angular oscillations of the amplitude of the acoustic wave due to the thermoelectric effect result from the periodic dependence of the electron velocity v_z on the angle θ between the normal to the layers and the magnetic field and take place over the entire range of angles (fig. 2). These oscillations are characteristic for the layered conductors and do not take place in ordinary metals. When the angle θ is noticeably tilted from the normal to the layers but differs from $\pi/2$, all of the sections of the Fermi surface by the plane $p_B = \text{const}$ are closed and at $\eta \ll 1$ they are almost indistinguishable. Because $k_{\mathtt{T}} l \ll 1$ electrons can perform several complete revolutions during the mean free time τ before being scattered, indicating that in the process of the EMAT of the energy in layered conductors a large number of charge carriers is included (due to the weak dependence of $\varepsilon(\mathbf{p})$ on p_z), rather than a small group of electrons as in the case of ordinary metals. This gives rise to the amplitude of acoustic oscillations which makes the thermoelectric mechanism dominant over the inductive one in a wide range of angles at room temperature (fig. 2a). At lower temperatures the free-path length of electrons is increasing and the domination is present only at a larger tilt of the magnetic field (fig. 2b) but is still present because in layered conductors at ultrasonic frequencies the condition $\omega\tau \ll 1$ is well satisfied. Figure 2 shows that the height of the maxima of $U^{\mathtt{T}}_{zl}$ at the adiabatic thermal condition at the boundary exceeds by far that of the maxima at the isothermal one, due to the nonzero temperature distribution at $z = 0$ in the first case $\Theta\,|_{z=0} = -\frac{k_B T\alpha_{zy}k}{\omega\mu_0\kappa_{zz}}\frac{1-\sqrt{\frac{\mu_0\kappa_{zz}}{C\rho_{yy}}}}{1-\frac{C\rho_{yy}}{\mu_0\kappa_{zz}}}$.

The amplitude of oscillations of $U^{\mathtt{T}}_{zl}$ is associated with the periodic dependence of the thermoelectric power α_{zy} and the thermal conductivity κ_{zz} which are functions of

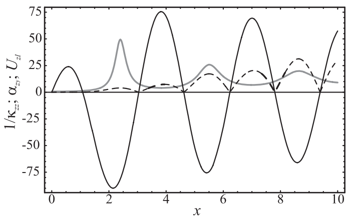

Fig. 3: The positions of the extremes of $1/\kappa_{zz}$ (gray curve), α_{zy} (solid curve) and U_{zl}^{T} (dashed curve) in their angular dependences.

the velocity of the conduction electrons across the layers v_z. This is plotted in fig. 3. The acoustic wave is maximally attenuated when U_{zl}^{T} goes to zero which occurs at angles $\theta = \theta_c = m\pi, m = 0, 1, 2, 3\ldots$, at which α_{zy} is zero. On the other hand, the position of the peaks in the angular dependence of U_{zl}^{T} coincides with the position of the extremes of inverse thermal conductivity $\frac{1}{\kappa_{zz}} = \frac{3e^2}{\pi^2 k_B^2 T}\rho_{zz}$. If $\Omega\tau\cos\theta > 1$, then the first term in eq. (36) is dominant. However, if $\zeta\tan\theta$ equals a zero of the zeroth-order Bessel function, then at that angle κ_{zz} will be a minimum and $1/\kappa_{zz}$ will be a maximum. If $\zeta\tan\theta \gg 1$, then the zeroes occurs at angles $\theta = \theta_n^{max}$ given by $\zeta\tan\theta_n^{max} = \pi(n - \frac{1}{4}), n = 0, 1, 2, 3\ldots$. For $1/\kappa_{zz}$ to be a minimum it should be $\theta = \theta_n^{min}$, where $\zeta\tan\theta_n^{max} = \pi(n + \frac{1}{4})$. Maxima and minima appearing in the angular dependence of $1/\kappa_{zz}$ repeat with a period $\Delta(\tan\theta) = \frac{2\pi\hbar}{aD_p}$, and the period of oscillations of U_{zl}^{T} is $\Delta(\tan\theta) = \frac{\pi\hbar}{aD_p}$, respectively. The diameter of the Fermi surface D_p along the p_y-axis can be determined to a high degree of accuracy from the measured period of the oscillations. The maxima of α_{zy} are shifted with respect to the maxima of $1/\kappa_{zz}$ (and thus of the maxima of U_{zl}^{T}) by the value $2\eta^{1/2}$. The experimental measurement of this value will allow the quasi–two-dimensionality parameter of the electron energy spectrum η to be determined.

Layered organic conductors are very convenient for performing experiments due to their high purity. The observation of the thermoelectric mechanism of EMAT in Q2D conductors at ultrasonic frequencies ($\omega \cong 10^8\,\mathrm{s}^{-1}$) is conditioned by the compliance with certain requirements. In particular, perfect specimens with a free-path length l of electrons of $10^{-3}\,\mathrm{cm}$ (for the conditions $l \ll \delta_s, \delta_{\mathrm{T}}$ to be fulfilled), a relaxation time τ of $10^{-9}\,\mathrm{s}$, (so that $\omega\tau \ll 1$) and magnetic fields up to $0.6\,\mathrm{T}$ must be used. At high frequencies and low enough temperatures δ_s and δ_{T} become much smaller than l. In that case the conditions for the normal-skin effect are violated and the thermoelectric mechanism of EMAT is totally suppressed by both the inductive and deformation mechanisms. Wave processes in layered organic conductors in a magnetic field are quite sensitive to the form of the electron energy spectrum, and their experimental study will provide detailed and reliable information on the relaxation properties of charge carriers.

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