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1998 J. Phys.: Condens. Matter 10 11765

(http://iopscience.iop.org/0953-8984/10/50/015)

J. Phys.: Condens. Matter **10** (1998) 11765–11772. Printed in the UK PII: S0953-8984(98)92321-9

Giant oscillations of the rate of sound attenuation in layered conductors placed in a magnetic field

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Received 9 March 1998, in final form 30 September 1998

Abstract. The propagation of sound waves in layered conductors with quasi-two-dimensional electron energy spectra is studied theoretically. It is shown that in a magnetic field $\boldsymbol{H}$ the rate of sound attenuation G oscillates with a great amplitude as a function of $1/H$ and of the angle between $\boldsymbol{H}$ and the sound wavevector $\boldsymbol{k}$, if the charge carriers are capable of drifting along $\boldsymbol{k}$, and the radius of curvature of their trajectories r is much less than the electron mean free path l but significantly exceeds the wavelength of sound $1/k$. The diameter of the Fermi surface in the direction orthogonal to the vectors $\boldsymbol{k}$ and $\boldsymbol{H}$ can be determined to a high degree of accuracy from the measured period of the oscillations.

Conductors of organic origin are usually thread or layered structures with a sharply pronounced anisotropy of the electrical conductivity. In the layered tetrathiafulvalene salts and halides of tetraselenium tetracene conductors, the conductivity across the layers is substantially less than the conductivity along the layers. It is reasonable to suppose that the electron energy spectrum in such conductors is quasi-two-dimensional in character, i.e. the charge-carrier energy

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

is weakly dependent on the projection of the momentum on the normal to the layers, $p_z = \boldsymbol{p} \cdot \boldsymbol{n}$. The coefficients multiplying the cosines sharply decrease with n, so the maximum value on the Fermi surface $\varepsilon(\boldsymbol{p}) = \varepsilon_F$ of the function $\varepsilon_1(p_x, p_y)$, which is equal to $\max \varepsilon_1(p_x, p_y) = \eta\varepsilon_F$, is much less than ε_F; the functions $\varepsilon_n(p_x, p_y)$ with $n \geqslant 2$ are even smaller. The decrease of the coefficients with the increase of the order of the harmonics is due to the significant distance between the layers. Under such circumstances, the approximation of a strong bond is applicable (the degree of overlapping of the wavefunctions of the electrons belonging to different layers is small). On the other hand, inside the layer a weak-bond approximation is suitable.

By distorting the trajectories of the charges, an external magnetic field significantly affects their motion. If $\boldsymbol{H}$ is not oriented in the plane of the layers, all of the electron orbits in the momentum space are closed and charges gyrate along them with the frequency Ω. In a strong magnetic field $\boldsymbol{H}$ ($\Omega\tau \gg 1$; τ is the charge-carrier relaxation time), all of the kinetic and thermodynamic characteristics of conductors are very sensitive to the form of the electron energy spectrum. A series of effects exist which are typical for layered conductors with metal-type conductivity, but non-existent for ordinary metals [1–6]. Among these effects, arising from the quasi-two-dimensional nature of the charge-carrier energy spectrum, is the high acoustic transparency of a conductor in the absence of charge-carrier drift along the sound wavevector $\boldsymbol{k}$ in the range of the magnetic field for which

$$1 \ll kr \ll kl. \tag{2}$$

In (2), r is the radius of curvature of the electron trajectories, l is the electron mean free path, $1/k$ is the wavelength of sound. Under the conditions of strong anisotropy in the energy–momentum relation for charge carriers, closed electron orbits are almost indistinguishable, allowing the participation of almost all of the conduction electrons on the Fermi surface in the formation of Pippard oscillations [7]. As a result, the amplitude of the periodic variations of the rate of sound attenuation, G, increases sharply as compared to the case for a quasi-isotropic metal. Between the sharp maxima in the dependence of G on $1/H$, the regions of anomalous acoustic transparency are located. It will be shown below that under conditions for which inequalities (2) are satisfied and the charge-carrier drift along $\boldsymbol{k}$ is small (that is, the charge-carrier displacement in the $\boldsymbol{k}$-direction during the period of motion in a magnetic field $T = 2\pi/\Omega$ is less than $1/k$), the rate of sound attenuation oscillates as a function of $1/H$ with a great amplitude.

The rate of sound attenuation can be obtained by means of the solution of the elasticity theory equation for the ionic displacement $\boldsymbol{u}$:

$$-\omega^2 \rho u_i = \lambda_{ijlm} \frac{\partial u_{lm}}{\partial x_j} + F_i. \tag{3}$$

Here ρ and λ_{ijlm} are the density and elastic tensor of the crystal, while $u_{lm} = \partial u_l/\partial x_m$ is the deformation tensor. The wave is taken to be monochromatic with the frequency ω, so the differentiation with respect to the time variable is equivalent to multiplication by $(-\mathrm{i}\omega)$. Equation (3) contains the counterforce $\boldsymbol{F}$ applied to the lattice by the electron system excited by the acoustic wave.

In order to determine the electric field $\boldsymbol{E}$ generated by the sound, equation (3) should be supplemented by the Maxwell equations

$$\Delta \boldsymbol{E} + \mathrm{i}\omega\mu_0 \boldsymbol{j} = 0 \tag{4}$$

and by the electroneutrality condition, which is equivalent to the continuity condition, for the current, i.e.

$$\operatorname{div} \boldsymbol{j} = 0. \tag{5}$$

We have neglected the displacement current. Due to the high density of charge carriers in conductors, a quasi-equilibrium state in the concomitant coordinate system, which moves together with the crystal lattice with the velocity $-\mathrm{i}\omega\boldsymbol{u}$, becomes established in a very short time. That is, during a time of the order of h/ε_F, the distribution function for electrons takes the form $f_0(\varepsilon + \mathrm{i}\omega\boldsymbol{p}\cdot\boldsymbol{u})$. The slow relaxation of electrons in the concomitant coordinate system should be described by means of the non-equilibrium correction $-\psi\, \partial f_0/\partial\varepsilon$ of the Fermi distribution function f. The function ψ satisfies the kinetic equation

$$\begin{gathered} \boldsymbol{v}\cdot\frac{\partial\psi}{\partial\boldsymbol{r}} + \frac{\partial\psi}{\partial t} + \left(\frac{1}{\tau} - \mathrm{i}\omega\right)\psi = g \\ g = -\mathrm{i}\omega\Lambda_{ij}(\boldsymbol{p})u_{ij} + e\widetilde{\boldsymbol{E}}\cdot\boldsymbol{v}. \end{gathered} \tag{6}$$

Here, e and $\boldsymbol{v}$ are the electron charge and velocity. The time t determines the position of a charge on its trajectory in a magnetic field according to the equation of motion

$$\frac{\partial \boldsymbol{p}}{\partial t} = e\boldsymbol{v} \times \boldsymbol{H}. \tag{7}$$

Equation (6) is linearized by the weak perturbation of the electron subsystem, and the collision operator is represented by the approximation of the relaxation time τ. The electric field $\widetilde{\boldsymbol{E}}$ has the form

$$\widetilde{\boldsymbol{E}} = \boldsymbol{E} - \mathrm{i}\omega\mu_0 \boldsymbol{u} \times \boldsymbol{H} + \frac{m\omega^2}{e}\boldsymbol{u}. \tag{8}$$

Here, m is the electron mass. The first two terms in (8) are the relativistic Lorentz transformation of the field $\boldsymbol{E}$, and the third term is connected with non-inertiality of the concomitant coordinate system. The components of the deformation potential tensor $\lambda_{ij}(\boldsymbol{p})$ account for the charge-carrier energy renormalization under strain [8]:

$$\begin{aligned} &\delta\varepsilon = \lambda_{ij}(\boldsymbol{p})u_{ij} \\ &\Lambda_{ik}(\boldsymbol{p}) = \lambda_{ik}(\boldsymbol{p}) - \frac{\langle \lambda_{ik}(\boldsymbol{p})\rangle}{\langle 1\rangle}. \end{aligned} \tag{9}$$

By means of the solution of the kinetic equation, the current density and the force $\boldsymbol{F}$ can be represented as follows:

$$j_i = -\frac{2}{(2\pi\hbar)^3}\int e v_i \psi \frac{\partial f_0}{\partial\varepsilon}\ \mathrm{d}^3 p \equiv \langle e v_i \psi\rangle \tag{10}$$

$$F_i = \mu_0(\boldsymbol{j}\times\boldsymbol{H})_i + \frac{m}{e}\mathrm{i}\omega j_i + \frac{\partial}{\partial x_k}\langle \Lambda_{ik}\psi\rangle. \tag{11}$$

The force $\boldsymbol{F}$ was calculated by Silin for isotropic conductors [9] and generalized by Kontorovich to the case of an arbitrary dispersion law of charge carriers [10].

Let us consider an acoustic wave propagating in the plane of the layers along the x-axis of a conductor, placed in a magnetic field $\boldsymbol{H} = (H\sin\theta, 0, H\cos\theta)$.

Making use of the Fourier method, instead of equations (3)–(5) we obtain the set of linear algebraic equations for the Fourier components $\boldsymbol{u}(k)$ and $\widetilde{\boldsymbol{E}}(k)$ of the ionic displacement and the electric field:

$$\begin{aligned} &\mathrm{i}\omega\mu_0 j_\alpha(k) = k^2[\widetilde{E}_\alpha - \mathrm{i}\omega\mu_0(\boldsymbol{u}(k)\times\boldsymbol{H})_\alpha] \qquad \alpha = y, z \\ &j_x(k) = 0 \\ &-\omega^2\rho u_i(k) = -\lambda_{ixlx}k^2 u_l(k) + (\mathrm{i}m\omega/e)j_i(k) + \mu_0(\boldsymbol{j}(k)\times\boldsymbol{H})_i + \mathrm{i}k\langle\Lambda_{ix}\psi(k)\rangle. \end{aligned} \tag{12}$$

The solution of the kinetic equation in the Fourier representation

$$\psi = \left(\int_{t-T}^{t} \mathrm{d}t'\ g(t')\exp\{\mathrm{i}k[x(t') - x(t)] + \nu(t'-t)\}\right) \Big/ (1 - \exp[-\nu T - \mathrm{i}k\overline{v}_x T]) \equiv \hat{R}g \tag{13}$$

allows us to write the quantities that characterize the system response to the sound wave in the form

$$\begin{aligned} &j_i(k) = \sigma_{ij}(k)\widetilde{E}_j(k) + a_{ij}(k)k\omega u_j(k) \\ &\langle \Lambda_{ix}\psi(k)\rangle = b_{ij}(k)\widetilde{E}_j(k) + c_{ij}(k)k\omega u_j(k) \end{aligned} \tag{14}$$

where

$$\nu = 1/\tau - \mathrm{i}\omega \qquad g(t) = \omega\Lambda_{ji}(t)k_i u_j(k) + e\boldsymbol{v}(t)\cdot\widetilde{\boldsymbol{E}}(k) \qquad x(t) = \int^t v_x(t_1)\ \mathrm{d}t_1.$$

Here, $\bar{v}$ is the velocity of the charge-carrier drift along the x-axis. The Fourier transforms of the electrical conductivity $\sigma_{ij}(k)$ and acoustoelectronic tensors $a_{ij}(k)$, $b_{ij}(k)$, $c_{ij}(k)$ are given by following expressions:

$$\begin{aligned} \sigma_{ij}(k) &= \langle e^2 v_i \hat{R} v_j \rangle \qquad & a_{ij}(k) &= \langle e v_i \hat{R} \Lambda_{jx} \rangle \\ b_{ij}(k) &= \langle e \Lambda_{ix} \hat{R} v_j \rangle \qquad & c_{ij}(k) &= \langle \Lambda_{ix} \hat{R} \Lambda_{jx} \rangle. \end{aligned} \tag{15}$$

The condition for the existence of a non-trivial solution of the set of equations (12) (equating the system determinant to zero) represents the dispersion equation for the problem. The imaginary part of the roots of the dispersion equation determines the attenuation of the acoustic and electromagnetic waves, and the real part describes the renormalization of their velocities related to the interaction between the waves and the conduction electrons. By virtue of the considerable mass difference of ions and electrons, the root k, related to the sound wave, and the root k_e, related to the electromagnetic wave, differ significantly.

Consider the acoustic wave polarized along its wavevector direction. Using (8) and (14), it is easily seen that for $\omega\tau \ll 1$ the root k of the dispersion equation can be represented as follows:

$$k = \frac{\omega}{s} + k_1 \tag{16}$$

where the small correction k_1 takes the form

$$k_1 = \frac{\mathrm{i}k^2}{2\rho s} \frac{1}{(1 - \xi\tilde{\sigma}_{yy})} \left(\xi \left(\tilde{a}_{yx} \tilde{b}_{xy} - \tilde{c}_{xx} \tilde{\sigma}_{yy} \right) + \tilde{c}_{xx} - \mathrm{i} \left(\tilde{a}_{yx} - \tilde{b}_{xy} \right) \frac{H_z \mu_0}{k} + \tilde{\sigma}_{yy} \frac{H_z^2 \mu_0^2}{k^2} \right) \Bigg|_{k=\omega/s} \tag{17}$$

with $s = (\lambda_{xxxx}/\rho)^{1/2}$; also, $\xi = \mathrm{i}\omega\mu_0 c^2/(k^2c^2 - \omega^2)$ and

$$\tilde{\sigma}_{\alpha\beta} = \sigma_{\alpha\beta} - \frac{\sigma_{\alpha x}\sigma_{x\beta}}{\sigma_{xx}} \qquad \tilde{a}_{\alpha j} = a_{\alpha j} - \frac{a_{xj}\sigma_{\alpha x}}{\sigma_{xx}}$$
$$\tilde{b}_{i\beta} = b_{i\beta} - \frac{b_{ix}\sigma_{x\beta}}{\sigma_{xx}} \qquad \tilde{c}_{ij} = c_{ij} - \frac{b_{ix}a_{xj}}{\sigma_{xx}}$$

for $\alpha, \beta = y, z$.

The electron displacement along $\boldsymbol{k}$ during the period of motion $T = 2\pi/\omega$ is

$$\begin{aligned} \overline{v}_x T &= -\tan\vartheta \sum_{n=1}^{\infty} \frac{an}{\hbar} \int_0^T \mathrm{d}t\ \varepsilon_n(t, p_H) \sin\frac{anp_z}{\hbar} \\ &= -\tan\vartheta \sum_{n=1}^{\infty} \frac{an}{\hbar} \int_0^T \mathrm{d}t\ \varepsilon_n(t, p_H) \sin\left\{ \frac{anp_H}{\hbar\cos\vartheta} - \frac{1}{\hbar} anp_x(t, p_H) \tan\vartheta \right\} \end{aligned} \tag{18}$$

where $p_H = p_x \sin\theta + p_z \cos\theta$ is an integral of the motion.

Under the main approximation in a small parameter for the quasi-two-dimensionality η, the velocity of the drift takes the form

$$\overline{v}_x = -\tan\vartheta\ \mathrm{Im} \sum_{n=1}^{\infty} \frac{an}{\hbar} \exp\left\{ \frac{\mathrm{i}anp_H}{\hbar\cos\vartheta} \right\} I_n(\tan\vartheta) \tag{19}$$

where

$$I_n(\tan\vartheta) = \frac{1}{T} \int_0^T \mathrm{d}t\ \varepsilon_n(t) \exp\left\{ -\frac{\mathrm{i}}{\hbar} anp_x(t) \tan\vartheta \right\}. \tag{20}$$

In the case where $kr \gg 1$, the charge carriers that travel in phase with the wave interact most efficiently with the acoustic wave. They give the main contribution to the acoustoelectronic coefficients, which can be easily calculated by means of the stationary-phase method. For example, in the case of a layered conductor whose electron spectrum has the form

$$\varepsilon(p) = \frac{p_x^2 + p_y^2}{2m} + \eta \frac{\hbar}{a} v_0 \cos\left(\frac{ap_z}{\hbar}\right) \qquad v_0 = \sqrt{\frac{2\varepsilon_F}{m}} \tag{21}$$

the asymptotic expression for σ_{yy} at small η is given by

$$\sigma_{yy} = \frac{4Ne}{\pi m v k D}\left\{\frac{1-\sin kD}{(1+\alpha^2)^{1/2}} + \frac{(\pi\gamma)^2}{3}\left(1+\frac{1}{2}\sin kD\right) + \pi\gamma\sin kD\left(1-\frac{1}{(1+\alpha^2)^{1/2}}\right)\right\} \tag{22}$$

where $\alpha = kl\eta\tan\theta J_0(\beta)$, $\beta = (a/\hbar)mv_0\tan\theta$, $D = 2v_0/\Omega$, N is the electron number density, J_0 is the Bessel function.

The rest of the acoustoelectronic coefficients behave in an analogous manner; i.e., they are oscillatory functions of both the inverse magnetic field magnitude and the angle θ.

At $\alpha\gamma \ll 1$, one can easily obtain the following expression for k_1:

$$k_1 = \frac{\mathrm{i}\omega N m v \Omega\tau}{4\pi\rho s^2}\frac{2\pi\gamma\sin^2 kD[1-(1+\alpha^2)^{-1/2}] + (\pi\gamma)^2 + \mathrm{i}\mu(1+\sin kD)}{1-\sin kD + [(\pi\gamma)^2/2](1+\alpha^2)^{1/2} + \pi\gamma\left[(1+\alpha^2)^{1/2}-1\right] + \mathrm{i}\mu} \tag{23}$$

where $\mu = 2\pi^2 v_0\omega^2/s^3\omega_0^2\mu_0\Omega\tau$, ω_0 is the plasma frequency, $\gamma = 1/\Omega\tau$.

Equation (23) is valid for $\Omega\tau \cong (eH\tau\mu_0\cos\theta/m) \gg 1$, i.e. when $\cos\theta$ differs significantly from zero. For the values of $\tan\theta$ at which α becomes equal to zero, there is no drift of the charge carriers along $\boldsymbol{k}$.

At $\alpha \ll 1$,

$$k_1 = \frac{\mathrm{i}\omega N m v \Omega\tau}{4\pi\rho s^2}\frac{\pi\gamma\alpha^2\sin^2 kD + (\pi\gamma)^2 + \mathrm{i}\mu(1+\sin kD)}{1-\sin kD + (\pi\gamma)^2/2 + \pi\gamma\alpha^2/2 + \mathrm{i}\mu} \qquad \alpha \ll 1. \tag{24}$$

When $\gamma^{1/2} \ll \alpha \ll 1$, the oscillating terms significantly exceed the smoothly varying terms not only in the denominator but also in the numerator of expression (24). This yields giant oscillations of the sound attenuation rate $G = \mathrm{Im}\, k$ as a function of both the magnetic field magnitude and the angle θ between $\boldsymbol{H}$ and $\boldsymbol{n}$. In the case where the charge-carrier displacement in the $\boldsymbol{k}$-direction during the relaxation time is much greater than the wavelength of sound, these oscillations take place as well. The quantity k_1 is then

$$k_1 = \frac{\mathrm{i}\omega N m v \Omega\tau}{4\pi\rho s^2}\frac{2\pi\gamma\sin^2 kD + (\pi\gamma)^2 + \mathrm{i}\mu(1+\sin kD)}{1-\sin kD + \pi\gamma\alpha + \mathrm{i}\mu} \qquad 1 \ll \alpha \ll \frac{1}{\gamma}. \tag{25}$$

Thus even a small drift of charge carriers along $\boldsymbol{k}$ affects the sound attenuation G significantly. At $\sin kD = 1$, the function $G(H)$ attains its maximum value:

$$G_{\max} = \frac{2G_0\Omega\tau}{(1+\alpha^2)^{1/2}}. \tag{26}$$

A slight deviation of $\sin kD$ from unity leads to a sharp decrease of G, which reaches its minimum $G_{min} = G_0/\Omega\tau$ at $\sin kD = -1$, if $\alpha^2 \ll \gamma \ll 1$. At $\gamma \leqslant (3\alpha^2/2) \ll 1$ the minimum of $G(H)$ is displaced to the values of H at which $\sin kD$ is close to zero, and for $\sin kD = -1$ the function $G(H)$ has a local maximum $G = G_0\alpha^2$. This maximum increases with increasing α and at $\alpha \geqslant 1$ attains the value G_0—the rate of sound attenuation in the absence of a magnetic field. At the same time, the main maximum decreases with increasing α and draws closer to the local maximum. When $\sin kD = -1$, the rate of sound attenuation oscillates with a great amplitude, exceeding the smoothly varying part of G in times of magnitude $\Omega\tau$.

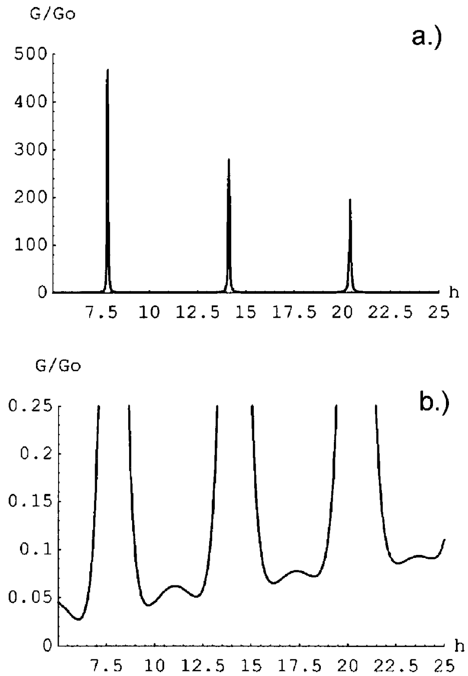

Figure 1. The dependence of the acoustic absorption coefficient G/G_0 on the magnetic field $h = H_0/H$ at $kl = 10^3$, $\eta = 10^{-2}$, $x \equiv \tan(\vartheta) = 1.5 \times 10^{-2}$; (a) $G/G_0 = [0, 500]$, (b) $G/G_0 = [0, 0.25]$.

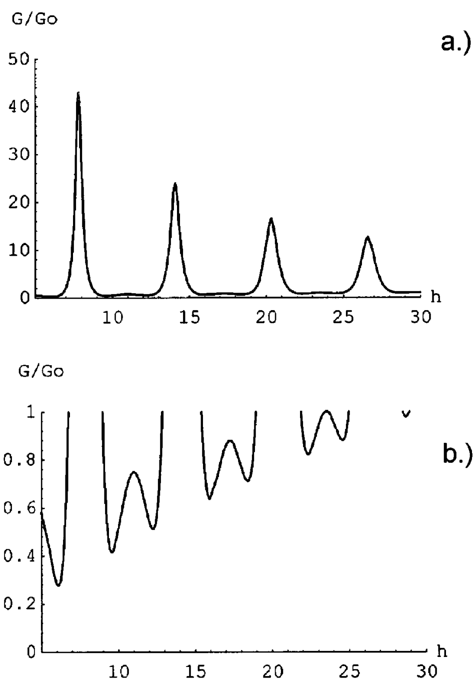

Figure 2. The dependence of the acoustic absorption coefficient G/G_0 on the magnetic field $h = H_0/H$ at $kl = 10^2$, $\eta = 8 \times 10^{-2}$, $x = 8 \times 10^{-2}$; (a) $G/G_0 = [0, 50]$, (b) $G/G_0 = [0, 1]$.

The numerical calculations based on formulae (24) and (25) confirm such analysis. The results for some values of the parameters are shown in figure 1 and figure 2 for the acoustic absorption coefficient G/G_0 versus the magnetic field $h = H_0/H$, with $H_0 = 2\omega v_0 m/es\mu_0$,

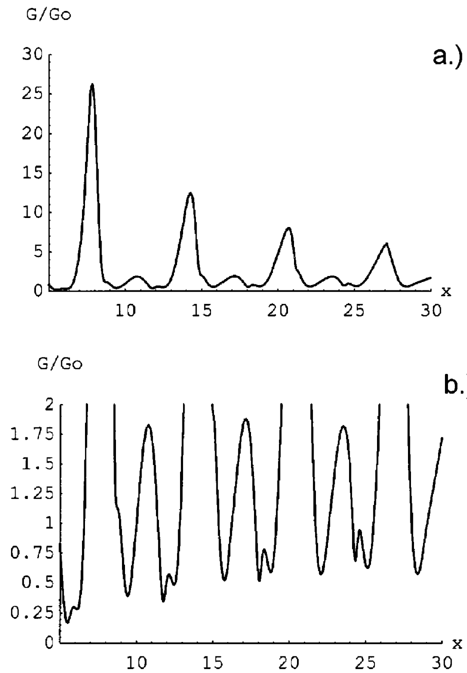

Figure 3. The dependence of the acoustic absorption coefficient G/G_0 on the tangent of the angle x between the magnetic field $\boldsymbol{H}$ and the sound wavevector $\boldsymbol{k}$ at $h = 1$, $kl = 10^2$, $\eta = 10^{-2}$; (a) $G/G_0 = [0, 30]$, (b) $G/G_0 = [0, 2]$.

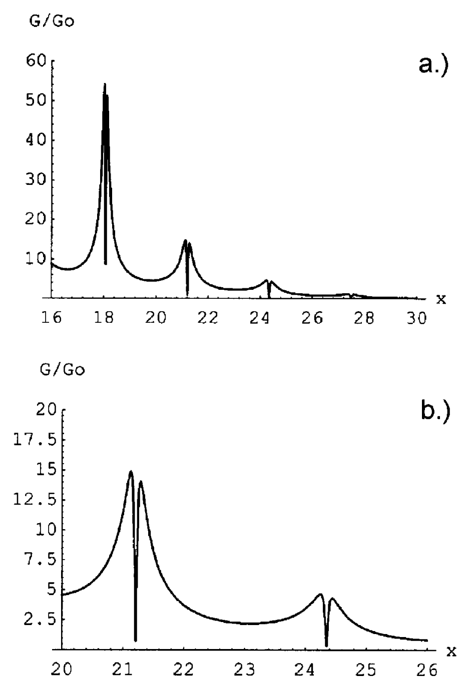

Figure 4. The dependence of the acoustic absorption coefficient G/G_0 on the tangent of the angle x between the magnetic field $\boldsymbol{H}$ and the sound wavevector $\boldsymbol{k}$ at $h = 10^{-1}$, $kl = 10^2$, $\eta = 10^{-1}$; (a) $G/G_0 = [0, 60]$, (b) $G/G_0 = [0, 20]$.

and on figure 3 and figure 4 for the acoustic absorption coefficient G/G_0 versus the tangent of the angle $x = \tan\vartheta$ between the magnetic field $\boldsymbol{H}$ and the sound wavevector $\boldsymbol{k}$.

The oscillating character of the dependence (a), as well as the existence of local maxima

(b) and their locations, are evident.

It is easy to see that the same dependence of G on $1/H$ and $\tan\theta$ remains valid for an arbitrary form of the quasi-two-dimensional electron energy spectrum. If there are only two points of the stationary phase on the electron orbit, the quantity $D = D_p/eH\mu_0\cos\vartheta$ is determined by the diameter D_p of the Fermi surface in the direction orthogonal to the vectors $\boldsymbol{k}$ and $\boldsymbol{H}$.

In much the same way, giant oscillations of the velocity of sound transform in the presence of charge-carrier drift along the wavevector direction.

The observation of the marked influence of the drift on the oscillatory dependence of G on $1/H$ at ultrasonic frequencies ($\omega \cong 10^8$ s^{-1}) is conditioned by compliance with certain requirements. In particular, perfect specimens with a long free-path length of electrons and strong magnetic fields (about 10 T) must be used. For this range of magnetic fields, the Shubnikov–de Haas effect is clearly manifested in the tetrathiafulvalene salts, which proves that the condition $\Omega\tau \gg 1$ is satisfied and, at the same time, the spacing between quantized electron energy levels is much less than both the Fermi energy and also the quantity $\eta\varepsilon_F$. Under these conditions, the quasi-classical description of non-equilibrium processes is valid. For a stronger magnetic field, taking account of the quantization of the electron energy levels becomes essential, but the effects considered above can be observed as well.

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