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Short Communication

Magnetoacoustic resonance in layered conductors

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Abstract. — The resonance absorption of ultrasound wave energy in layered conductors with quasi-two-dimensional electron energy spectrum is predicted. This resonance is connected with sharp Joule loss increase at certain values of an external magnetic field under the Pippard effect conditions.

In a metal placed in a magnetic field \( \mathbf{H} \) ultrasound wave energy is absorbed in a resonance manner, if the radius of the electron trajectory curvature \( r \) is much shorter than the charge carrier mean free path \( l \) but still exceeds significantly the sound wavelength. Apart from the oscillatory dependence of the sound attenuation rate \( \Gamma \) on the inverse magnetic field, predicted by Pippard [1], the function \( \Gamma(\mathbf{H}) \) has sharp resonance maxima, when the charge carrier drift along the wave vector \( \mathbf{k} = (k,0,q) \) during the period of motion \( 2\pi/\Omega = 2\pi m^* c /e H \) is different from zero (\( m^* \) is the charge carrier effective mass, \( c \) the velocity of light, \( e \) the electron charge) [2].

In a metal there is no need to take into account an electromagnetic field generated by sound wave at \( kr \gg 1 \). It is sufficient to consider only the deformation mechanism of sound absorption related to the renormalization of the charge carrier energy \( \delta \varepsilon = \lambda_{ij} u_{ij} \) in deformed crystal lattice. Components of the matrix \( \lambda_{ij} \) connect \( \delta \varepsilon \) with the deformation tensor \( u_{ij} = \partial u_j / \partial x_i \) in the approximation that is linear in the powers of the ionic displacement \( u \). At the Fermi surface (FS) \( \varepsilon(p) = \varepsilon_F \) the components \( \lambda_{ij} \) coincide in order of magnitude with the Fermi energy \( \varepsilon_F \), and they can be calculated easily if the dependence of energy \( \varepsilon \) on quasimomentum \( p \) of electron is known.

In layered conductors with the quasi-two-dimensional energy spectrum of the charge carrier the role of electromagnetic fields is proved to be essential even at \( kr \gg 1 \). Consideration of these fields results in the resonance of \( \Gamma(H) \) at \( k \perp \mathbf{H} \) in the absence of open sections of the FS, when electrons are incapable to drift along \( k \) direction. The reason is that the Pippard oscillations are produced in this case by almost all charge carriers, if \( 1 \ll kr \ll \eta^{-1} \), where \( \eta = v_0 / v_F \) is the ratio of the maximum electron velocity along the normal to the layers, \( v_0 = v_z^{\text{max}} \), to the...
maximum electron velocity $v_F$. As a result, the asymptotic behaviour of the high-frequency conductivity tensor $\sigma_{ij}(k, \omega)$, which connects the electron current density, $j_i$, to the electric field in the system of axes related to the moving crystal lattice, $E_j(k, \omega) = E_j(k, \omega) + [u_i \mathbf{H}] / c$, alters essentially. For instance, the component

$$
\sigma_{yy} = \frac{2 \sigma_0}{\pi k r} \left[ 1 - J_0(\eta kr) \cos \left( kr - \frac{\pi}{2} \right) \right]
$$

becomes sufficiently small for $\eta kr \to 0$ and $kr = \pi/4 + \pi n$, that leads to giant Joule losses. Here, $J_0(z)$, is the Bessel function for the argument $z$, $\sigma_0 = Ne^2 \tau/m^*$, $\tau$ is the mean free path time of electron, $N$ the charge carrier density, and $\omega$ the sound frequency. The Fourier component of the electric field in the laboratory system of axes, $E(k, \omega)$ is to be determined from the Maxwell equations

$$(k^2 - \frac{\omega^2}{c^2})E_\alpha(k) - \frac{4 \pi i \omega}{c^2} j_\alpha(k) = 0, \quad \alpha = y, z
$$

The electric field $E_z$ along $k$-direction can be determined from the electric neutrality condition, which follows from the Poisson equation (see, for example, [3]).

The formula for $\sigma_{yy}$ given above refers to the case when the magnetic field is directed along the normal to the layers and the charge carrier energy spectrum in the plane of the layers is isotropic in the main approximation as to small parameter $\eta$.

The electric current density $j_\alpha(k, \omega)$ can be found readily by using the solution of kinetic equation for the non-equilibrium distribution function of charge carriers $f = f_0 - \Psi \partial f_0 / \partial \varepsilon$ ($f_0(\varepsilon)$ is the Fermi distribution function). In the $\tau$-approximation for the collision integral the solution takes the following form

$$
\Psi = \frac{1}{\Omega} \int_{-\infty}^{\varphi} d\varphi' [\Lambda_{ij}(p_z, \varphi') \dot{u}_{ij} + e \mathbf{E} \mathbf{v}] \exp \left\{ i k [x(\varphi') - x(\varphi)] + \nu(\varphi' - \varphi) \right\}
$$

Here $\varphi$ is the phase variable on the electron orbit in a magnetic field $\varepsilon(\mathbf{p}) = \text{const}$, $p_z = \text{const}$ ; $\nu = 1/\tau - i \omega$, $\Lambda_{ij} = \lambda_{ij} - \overline{\lambda}_{ij}$ where $\overline{\lambda}_{ij}$ is the mean value of the component of the deformation potential tensor $\lambda_{ij}$ on the FS. If one knows $\Psi$ one can find the dissipative function of electron system and, hence, the dependence on $H$ of a part of the rate of sound attenuation

$$
\Gamma = \frac{1}{\rho s |\mathbf{u}|^2} \int_{\varepsilon(p) = \varepsilon_F} dp_z \int_0^{2\pi} d\varphi \frac{2 m^*(p_z) |\psi|^2}{(2\pi \hbar)^3} \frac{1}{\tau}
$$

Here $s$ is the velocity of sound, $\rho$ the ionic density.

In the case under consideration $\Gamma$ has the form

$$
\Gamma = \Gamma_0 \frac{l}{r} \Re \left\{ \frac{1 - J_0^2(\eta kr) + i \beta (1 + J_0(\eta kr) \cos(2kr - \pi/2))}{1 - J_0(\eta kr) \cos(2kr - \pi/2) + i \beta} \right\}
$$

where $\beta = \pi^2 \omega^2 c^2 v_F^2 / 2 \omega_p^2 s^3 \Omega r$, $\omega_p = 2e(\pi N/m^*)^{1/2}$ is the plasma frequency. $\Gamma_0 = N m^* \times v_F \omega / 4 \pi \rho s^2$ of the same order of magnitude as the absorption rate at $H = 0$. Using the expansion for the Bessel function ($J_0(z) \approx 1 - z^2/4$, $z \ll 1$) one can see that the sharp resonance peaks appear at $\eta kr \to 0$, if the parameter $\beta$ is sufficiently small. In a pure two-dimensional case $\Gamma$ is given by

$$
\Gamma = \Gamma_0 \frac{l}{r} \frac{[1 + \cos(2kr(\pi/2)) \beta^2]}{[1 - \cos(2kr - \pi/2)]^2 + \beta^2}
$$
Generalisation to a case of an arbitrary orientation in the \([y, z]\) plane of a magnetic field and arbitrary form of the charge carrier dispersion law presents no difficulty. The simplest geometry, frequently used in experiments, has allowed us to clarify the physical nature of the resonance effect responsible for exclusively by the quasi-two-dimensionality of the electron energy spectrum.

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References

