Magnetoacoustic resonance in metals

A. N. Grishin, V. G. Skobov, L. M. Fisher, and A. S. Chernov
V. I. Lenin All-Union Electrical Engineering Institute

(Submitted 22 March 1982)

Pis'ma Zh. Eksp. Teor. Fiz. **35**, No. 9, 370–372 (5 May 1982)

It is shown that screening of the electron-phonon interaction by vortex fields shifts the lower boundary of the region in which the magnetoacoustic resonance exists by hundreds of kilohertz to several hundred megahertz.

PACS numbers: 76.50. + g

It is customarily assumed that the magnetoacoustic resonance must exist in the frequency range where the wavelength of sound is small compared to the mean free path of electrons (for a mean free path of 1 mm, this range lies above 100 kHz). This Doppler-shifted acoustic cyclotron resonance (DSACR) must appear as a maximum at the edge of the threshold for collisionless absorption.[1] The result in Ref. 1 was obtained assuming that vortex fields created by the sound wave in the metal do not play a significant role. In this paper, we demonstrate that for frequencies that are not too high vortex fields strongly screen the interaction of electrons with sound and the resonance peak, related to DSACR of the main group of carriers, is missing.

Let us examine absorption of transverse sound, propagating in a metal along the magnetic field $\mathbf{H} \parallel z$. For simplicity, we shall assume that the Fermi surface is shaped like a corrugated cylinder and that its axis coincides with the direction of $\mathbf{H}$. We shall assume that the deformation potential has the form $\Lambda_{ik} = -\zeta_0 m v_i v_k$ ($\mathbf{v}$ is the velocity of an electron, m is the cyclotron mass, and ζ_0 is a constant) and we shall calculate the electric current density in the metal and the density of forces acting on the lattice. Thus, the following dispersion equation for the wave in the metal is obtained from the system of equations for the sound and electromagnetic fields (see, for example, Ref. 2)

$$k^2 s^2 - \omega^2 = \frac{k^2 H^2}{4\pi\rho}\left\{\left[1 - \zeta + (4\pi i\omega/k^2c^2)\,\sigma_\pm(0)\zeta\right]^2\left[1 - \frac{4\pi i\omega}{k^2c^2}\,\sigma_\pm(k)\right]^{-1} - \left[(1-\zeta)^2 - \frac{4\pi i\omega}{k^2c^2}\sigma_\pm(0)\zeta^2\right]\right\}, \qquad (1)$$

where ω is the frequency, k is the wave number, s is the velocity of sound in the absence of a magnetic field, c is the velocity of light, ρ is the density of the crystal, $\zeta = \zeta_0(1 \pm i\gamma)$, $\gamma = 1/\Omega\tau$, Ω is the cyclotron frequency, τ is the free flight time, and $\sigma_\pm(k) = \sigma_{xx}(k) \pm i\sigma_{yx}(k)$ is the nonlocal conductivity for circular "plus" and "minus" polarizations. Terms that do not contain a factor ζ are related to the induced interaction of electrons with sound. In deriving (1), we expressed the deformation current and the deformation force in terms of $\sigma_\pm(k) - \sigma_\pm(0)$.

For the Fermi surface being studied, the nonlocal conductivity has the form

 0021-3640/82/09455-04$01.00

$$\sigma_{\pm}(k) = \pm\, i\frac{nec}{H}\left[(1 \pm i\gamma)^2 - \left(\frac{kv}{\Omega}\right)^2\right]^{-1/2}, \tag{2}$$

where n is the concentration of electrons, and v is their maximum velocity along the vector **H**. The conductivity (2) has a singularity at $k = \Omega(1 \pm i\gamma)/v$, corresponding to DSACR. As a result of the screening of the electron–phonon interaction by vortex fields, this singularity is present only in the denominator on the right side of (1). For this reason, the latter vanishes at the singularity if the vortex fields are ignored, the expression $i\omega\sigma_{\pm}(k)(\zeta - 1)^2 H^2/\rho c^2$, which becomes infinite at resonance, would appear on the right side of (1)]. This situation is similar to that occurring in the case of the tilt effect.[3]

Damping of the sound wave is determined by the imaginary part of the expression on the right side of (1). Far from the helicon-phonon and doppleron-phonon resonances, $k = \omega/s$ can be substituted in this expression and the damping coefficient $\Gamma_{\pm}$ is written in the form

$$\Gamma_{\pm}(H) \equiv \mathrm{Im}k = \Gamma_0 \mathrm{Im}\left[\frac{a}{q}(1 - \zeta) \mp \zeta\right]^2 \left[a \pm \frac{q}{\sqrt{(1 \pm i\gamma)^2 - q^2}}\right]^{-1}, \tag{3}$$

where

$$\Gamma_0 = \frac{\omega n m v}{2\rho s^2}, \quad a = \left(\frac{\omega\delta}{s}\right)^3, \quad \delta^3 = \frac{c^2 m v}{4\pi\omega n e^2}, \quad q = \frac{v\omega}{\Omega s} = \frac{H_1}{H}.$$

The parameter a is the cube of the ratio of the anomalous skin layer thickness to the wavelength of sound, while H_1 is the value of the magnetic field corresponding to the collisionless absorption threshold ($q = 1$). The results of the calculation for "plus" polarization, which does not include either a helicon or a doppleron, are presented in Fig. 1. The calculation was carried out for the following values of the parameters: $n = 10^{22}$ cm, $m = 10^{-27}$ g, $v = 10^8$ cm/s, $s = 2 \times 10^5$ cm/s, $\tau = 10^{-9}$ s, $\zeta_0 = 2$ and $\rho = 10$ g · cm^{-3}. In the vicinity of the point H_1 on curve 1, there is a sharp drop due to the disappearance of collionless absorption, but there is no resonance maximum on it. This is a result of the fact that in the case $\omega\delta/s < 1$ vortex fields strongly screen the interaction of electrons with

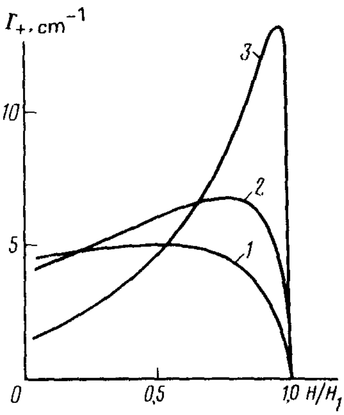

FIG. 1. Magnetic field dependence of the absorption coefficient of sound: curve 1, $a = 0.5$; curve 2, $a = 1$; curve 3, $a = 3$.

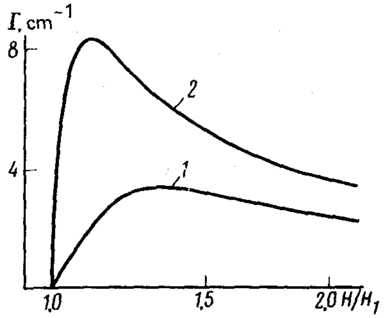

FIG. 2. Curve 1, $a = 3$; curve 2, $a = 5$.

the sound wave (the condition $\omega\delta = s$ corresponds to the frequency 250 MHz). In the opposite case $\omega\delta/s > 1$, the screening is weaker. For this reason, curve 2 has, near the absorption edge, an asymmetric peak, whose height is determined by the quantity a. This situation remains up to very high frequencies $a \sim \gamma^{-1/2}$. For high frequencies, screening disappears and the case examined in Ref. 1 is realized.

For minus polarization, the dispersion equation (1) has a solution corresponding to electromagnetic modes: a helicon and a doppleron. For this reason, for $a = 0.5$ on the curve $\Gamma_-(H)$, there is a strong maximum at the field $H \approx 2H_1$, which is due to a helicon-phonon resonance. For high values of a, this resonance is missing, but a doppleron–phonon resonance (DPR) exists. The corresponding maximum in $\Gamma_-(H)$ is situated to the right of the point H_1 at a distance $\Delta H \approx H_1/2a^2$. For $a = 3$, it is located practically at the threshold of collionless absorption and its height greatly exceeds Γ_0. The magnetoacoustic resonance therefore is not manifested here.

If the Fermi surface is not axially symmetric, then the absorption coefficient Γ has singularities for fields $H_n = H_1/n$, where n is an integer. These singularities in Γ differ from the anomalies near the point H_1, since multiple resonances occur in the region of strong, collionless absorption. For $a \ll 1$, the absorption Γ can be assumed to be inversely proportional to $\text{Re}\sigma(q)$. For this reason, the resonant maxima in $\text{Re}\sigma(q)$ must be manifested as minima in $\Gamma(H)$.

It should also be noted that in metals with anisotropic Fermi surfaces a situation can arise in which the line $\Gamma(H)$ is inverted relative to the point H_1. This must occur when the derivative of the cross-sectional area of the Fermi surface $\partial S/\partial p_z$ has a minimum, i.e., when the longitudinal velocities of all electrons in a given group are greater than some minimum value. For a model Fermi surface of this type, described in Ref. 4, the nonlocal conductivity has the form[1)]

$$\sigma_\pm(K) = \pm i \frac{nec}{H(1 \pm i\gamma)} F\left(\frac{-kv}{\Omega \pm i/\tau}\right), \qquad F(x) = 1 - \frac{x^2}{\sqrt{x^2-1}} \operatorname{arctg} \frac{1}{\sqrt{x^2-1}}, \tag{4}$$

where v is the minimum longitudinal velocity of electrons. The results of a calculation of the damping coefficient for a wave with minus polarization, in which DPR is missing, are illustrated in Fig. 2. Collisionless absorption of the wave in this case exists for fields $H > H_1$.

Thus, the magnetoacoustic resonance,[1] which is caused by the main group of carriers, lies in the hypersonic region. In the transition region, where $1 < a < \gamma^{-1/2}$, experimental observation of this resonance requires careful separation of DPR and DSACR signals, existing in opposite circular polarizations.

We note that the results obtained can explain the shape of the sound absorption curve near DSACR of octahedral holes in tungsten.[5]

We thank L. T. Tsymbal, who pointed out this problem to us.

[1)] In Eq. (8) of Ref. 4, the expression in the radicand should have the opposite sign.

1. É. A. Kaner, V. G. Peschanskiĭ, and I. A. Privorotskiĭ, Zh. Eksp. Teor. Fiz. 40, 214 (1961) [Sov. Phys. JETP 13, 147 (1961)].
2. V. G. Skobov, and É. A. Kaner, Zh. Eksp. Teor. Fiz. 64, 273 (1964) [Sov. Phys. JETP 19, 189 (1964)].
3. É. A. Kaner, L. V. Chebotarev, and A. V. Eremenko, Zh. Eksp. Teor. Fiz. 80, 1058 (1981) [Sov. Phys. JETP 53, 540 (1981)].
4. V. V. Lavrova, S. V. Medvedev, V. G. Skobov, L. M. Fizher, A. S. Chernov, and V. A. Yudin, Zh. Eksp. Teor. Fiz. 66, 700 (1974) [Sov. Phys. JETP 39, 338 (1974)].
5. A. A. Galkin, L. T. Tsymbal, and A. M. Cherkasov, Pis'ma Zh. Eksp. Teor. Fiz. 33, 3 (1981) [JETP Lett. 33, 1 (1981)].

Translated by M. E. Alferieff
Edited by S. J. Amoretty