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ACOUSTIC TRANSPARENCY OF ORGANIC CONDUCTORS IN A MAGNETIC FIELD

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It is predicted that organic conductors have high acoustic transparency in strong magnetic fields. The effect is shown to be due to the two-dimensional nature of the electron energy spectrum.

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Many organic superconductors are layered structures with a pronounced anisotropy of the electrical conductivity in the normal (nonsuperconducting) state: the in-plane conductivity is comparable to the conductivity of rather pure metals and is significantly higher than the conductivity along the normal $\mathbf{n}$ to the layers. It seems reasonable to suppose that the energy of the charge carriers in such a conductor,

$$\epsilon(\mathbf{p}) = \sum_{n=1}^{\infty} \epsilon_n(p_x, p_y)\cos(anp_z/h)\,, \tag{1}$$

depends weakly on the momentum projection $p_z = \mathbf{pn}$ and that the Fermi surface $\epsilon(\mathbf{p}) = \epsilon_F$ is a slightly corrugated cylinder with the "open" direction coinciding with the p_z axis. Here a is the separation between the layers, h is Planck's constant, and the coefficients of the cosines decrease rapidly with increasing n, so that the maximum value of the function $\epsilon_1(p_x, p_y)$ on the Fermi surface is of the order of $\eta\epsilon_F \ll \epsilon_F$ and the functions $\epsilon_n(p_x, p_y)$ with $n \geq 2$ are still smaller.

The experimental observation of Shubnikov–de Haas oscillations in tetrathiafulvalene salts [1,2] and then in halides of tetraseleniumtetracen [3] indicates that it is possible to synthesize organic single crystals in which the mean free path l of the charge carriers is sufficiently large and may be much greater than the diameter $2r$ of their orbits in realizable magnetic fields. Consequently, magnetoacoustic effects [4-6], the observation of which at $r \ll l$ has permitted determination of the Fermi surfaces of almost all metals, should prove useful for investigating the energy spectrum and relaxation properties of charge carriers in layered conductors.

The quasi-two-dimensional character of the electron energy spectrum gives rise to a number of specific effects in a magnetic field $\mathbf{H}$ which do not occur in ordinary metals; in particular, to acoustic transparency when r is less than l while being much greater than the sound wavelength $1/k$.

In metals at

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

the absorption of sound wave energy is determined mainly by the deformation mechanism [7] involving the strain-induced renormalization of the energy of the charge carriers:

$$\delta\epsilon = \lambda_{ij}(\mathbf{p})u_{ij}, \tag{3}$$

where λ_{ij} is the deformation potential tensor and $u_{ij} = \partial u_i/\partial x_j$ is the strain tensor. Allowance for the electromagnetic fields is essential only at higher magnetic fields such that $kr \ll 1$. In quasi-two-dimensional conductors the electromagnetic fields generated by sound play an important role in a significantly wider range of magnetic fields, including fields for which condition (2) holds. This results from the fact that the closed cross sections of the Fermi surface cut by the plane $p_H = \mathbf{pH}$ are almost indistinguishable, and a considerable fraction of the charge carriers contribute to an oscillatory dependence of the acoustoelectronic coefficients which relate the electron fluxes to the strain tensor u_{ij} and the sound-generated electromagnetic fields. As a result the acoustoelectronic coefficients undergo giant oscillations whose amplitude for $kr\eta \ll 1$ is comparable with the slowly varying parts of these coefficients. This is the cause of weak damping of acoustic waves under these conditions. Let us illustrate this for the case in which a sound wave propagates along the layers in a magnetic field $\mathbf{H} = (0, H\sin\vartheta, H\cos\vartheta)$ orthogonal to the wave vector $\mathbf{k}$.

The electromagnetic fields are to be found from Maxwell's equations, which in the Fourier representation have the form

$$E_\alpha(k) = \xi j_\alpha(k); \qquad \alpha = (y, z);$$
$$\mathbf{jk} = 0; \tag{4}$$

where $\xi = 4\pi i\omega/(k^2c^2 - \omega^2), j_\alpha(k)$ and $E_\alpha(k)$ are the Fourier components of the current density and the electric field of a monochromatic wave with the frequency ω.

The relation of the current density

$$\mathbf{j} = 2(2\pi h)^{-3}\int d^3pe\mathbf{v}f(\mathbf{p}, \mathbf{r}, t) \tag{5}$$

to the electric field can be obtained using the solution of the kinetic equation for the charge-carrier distribution function

$$f(\mathbf{p}, \mathbf{r}, t) = f_0(\epsilon(\mathbf{p}) - \mathbf{p}\dot{\mathbf{u}}) - \psi(\mathbf{p}, \mathbf{r})\exp(-i\omega t)\partial f_0/\partial\epsilon. \tag{6}$$

In an approximation to first order in the small perturbation of the electron system, the kinetic equation takes the form:

$$\partial\psi\partial t_H + \mathbf{v}\partial\psi\partial\mathbf{r} - i\omega\psi + W_{\text{col}}(\psi) = e\mathbf{v}\tilde{\mathbf{E}} - i\omega\Lambda_{ij}(\mathbf{p})u_{ij} \tag{7}$$

where $\tilde{\mathbf{E}} = \mathbf{E} - i\omega(\mathbf{v}\times\mathbf{H})/c - \omega^2 m\mathbf{u}/e$. Here $f_0(\epsilon - \mathbf{p}\dot{\mathbf{u}})$ is the equilibrium Fermi distribution function in a reference frame moving with the ion velocity $\dot{u}$, $\Lambda_{ij}(\mathbf{p}) = \lambda_{ij}(\mathbf{p}) - \langle\lambda_{ij}\rangle$, where the angle brackets indicate averaging over the Fermi surface. The collision integral $W_{\text{col}}(\psi)$ in Eq. (7) is a linear integral operator acting upon the function ψ. At low temperatures it can be represented to good accuracy as an operator of multiplication by the collision frequency $1/\tau$, i.e., $W_{\text{col}}(\psi) = \psi/\tau$.

Making use of the solution of equations (4) and (7), one can find a relation between the electric field and ionic displacement $\mathbf{u}$ and finally determine the dissipation function [7]

$$Q = 2(2\pi h)^{-3} \int d^3p\delta(\epsilon - \epsilon_F)\psi W_{\rm col}(\psi). \tag{9}$$

The asymptotic behavior of the acoustic damping rate

$$\Gamma = Q/\rho\omega^2 u^2 s \tag{10}$$

depends essentially on the correlation between the magnetic field and the degree of the corrugation of the Fermi surface. Here ρ is the density of the cystal and s is the velocity of the sound wave. In sufficiently pure specimens, when $kl\eta \gg 1$, there always is a range of magnetic fields in which $1/\eta \ll kr \ll kl$ and the oscillations of the acoustic damping rate are formed only by a small fraction of the conduction electrons, of the order of $(kr\eta)^{-1/2}$. These are charge carriers whose orbit diameters $2r = cD_p/eH_z$ are close to the extreme diameter $D_{\rm ext} = (1 \pm \eta)D$, where D_p is the diameter of the corrugated cylinder in the p_y direction. In this case, as in the case of an ordinary metal, Γ increases with increasing magnetic field in proportion to H:

$$\Gamma = (\omega\tau/r)\{1 + (kr\eta)^{-1/2}\sin(kD - \pi/4)\cos kD\eta\}. \tag{11}$$

However for $kr\eta \ll 1$, almost all charge carriers on the Fermi surface contribute to the Pippard oscillations [4], and the acoustic damping rate depends importantly on kD. For the wave with ionic displacements in the plane of the layers (un $= 0$) the acoustic damping rate

$$\Gamma = \frac{eH\tau\cos\theta\{g_1^2(1+\sin kD) + g_2^2(1-\sin kD) - 2g_1g_2\cos kD\}}{\pi^2h^2avc\{1 + |\xi\tilde{\sigma}_{yy}|^2\}\rho\omega^3u^2} \tag{12}$$

decreases with magnetic field and with the mean free path of the conduction electrons. Exceptions are the resonant values of the magnetic field, at which $kD = 2\pi(n + 1/4)$ (n is an integer) and terms of higher order with respect to the small parameters $1/kD$ and $\gamma = r/l$ should be retained in the asymptotic expression for the component of the high-frequency conductivity tensor

$$\tilde{\sigma}_{yy}(k) = (G/kD)(1 - \sin kD). \tag{13}$$

Here $g_1 = \omega k\Lambda_{jx}u_j, g_2 = \mathrm{k u}H_z avs/c$, and $G = e^2\tau\epsilon_F/a(2\pi h)^2$. At $kD = 2\pi(n + 1/4)$ the damping rate of longitudinal waves increases sharply and the height of the resonance peaks

$$\Gamma_{\rm res} = \frac{\omega\tau/r}{1 + (1/kr^2)^2} \tag{14}$$

is proportional to H if $l \ll kr^2$.

Dissipation of shear waves depends essentially on the form of the off-diagonal components of the tensor $\Lambda_{ij}(\mathrm{p})$. The component Λ_{xx} is apparently of the order of the Fermi energy, while the components Λ_{yx} and Λ_{zx} may be much less than ϵ_F. However, it follows from formula (12) that the damping length of a wave with ionic displacements along the y axis depends on the magnetic field in the same manner as the damping length for longitudinal waves, i.e., the conductor appears to be transparent over a wide range of magnetic fields until $\sin kD$ differs substantially from unity.

Transparency of a layered conductor for a wave with polarization parallel to the normal to the layers occurs only at selected values of a magnetic field for which $\sin kD = -1$. If $\sin kD$ differs substantially from -1, the acoustic damping rate has the form

$$\Gamma = (\omega\tau/r)\eta^2\{1 + \sin kD + kr\eta^2(1 - \sin kD)\}. \tag{15}$$

Thus in an external magnetic field the role of electromagnetic fields generated by sound proves to be essential. Allowance for the fields results in the compensation of a considerable part of the deformational absorption of wave energy and ultimately in acoustic transparency of the layered conductor.

Unfortunately, observation of this effect in HTSC layered structures is faced with difficulties, mainly in connection with the small mean free path of the charge carriers.

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