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MAGNETIC BREAKDOWN GIANT OSCILLATIONS OF ABSORPTION OF SOUND BY METALS

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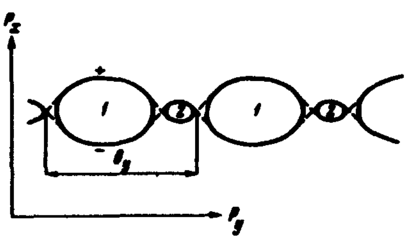
It is well known that the oscillatory dependence of different macroscopic quantities in metals on the magnetic field H is due to its quantizing properties. Usually, owing to the smallness of the quasiclassical parameter $\kappa = e\hbar H/cb_0^2$ (b_0 - characteristic dimension in p-space) the oscillations have a small amplitude $\sqrt{\kappa}$. However, the amplitude of the oscillations for the absorption of a sound wave (with frequency ω and wave vector q) under resonance conditions can become "giant," i.e., $\gtrsim 1$. This circumstance was first pointed out by Gurevich, Skobov, and Firsov (GSF) [1]. The GSF effect is due to oscillations of the density of the number of states on the Fermi surface at the point p_z, where there is resonance (p_z is the projection of the momentum on H = (0, 0, H))[1]. To observe this effect it is necessary to have a very small width of the resonance, namely [1] $\nu_0/qv_0 \lesssim \sqrt{\kappa}$ (ν_0 - collision frequency, v_0 - characteristic velocity).

In the present paper we show, using ultrasound as an example, that the phenomenon of magnetic breakdown [4] of interband tunnel transitions of a conduction electron in a field H leads to the occurrence of new giant (resonant) oscillations of the high-frequency characteristics of the metal, not connected with the oscillations of the density and differing strongly from the GSF effect. We consider as a model a system of periodically repeated symmetrical orbits of two bands, between which (at the locations of closest approach) there occur tunnel transitions with probability $W(\vec{H})$ [4, 5] (see the figure, where 1 and 2 are the numbers of the bands). In addition, it is assumed below that the electron spends most of the time on the trajectories of the first band, i.e., the reinforced inequalities $\alpha = S_1/S_2 \gg 1$ and $\beta = \Omega_1/\Omega_2 \ll 1$ are in force[2] ($S_{1,2}(E, P_z)$ - areas of orbits 1 and 2, $\Omega_{1,2} = eH/m^*_{1,2}c$ - the frequencies of revolutions on the orbits, $m^*_{1,2} = (2\pi)^{-1}(\partial S_{1,2}/\partial E)_{p_z}$, and $\alpha \sim \beta^{-2}$). Under conditions when stationary "magnetic-breakdown" states of the conduction electron are realized, and second orbit, in spite of its smallness, greatly influences the dynamics of the electron. Its "controlling" action (see below), due to coherent effects



[1]A similar phenomenon for electromagnetic waves was considered by I. Lifshitz et al. [2]. Some modification of the GSF effect was considered recently in [3].

[2]For a number of metals exhibiting the property of magnetic breakdown, $\approx < 10^2 - 10^3$ [4].

analogous to those arising in "bleaching" of an optical system, is the cause of the magnetic-breakdown giant oscillations.

The energy spectrum of the considered configuration constitutes a set of "magnetic" bands [4, 5] $E_n(P_x, p_z) = E_n(P_x + 2\pi\sigma/b_y, p_z)$ of broadened Landau levels ($\sigma = e\hbar H/c$, P_x is the x component of the generalized momentum, b_y the period in p space, n the number of the magnetic band). In the zeroth approximation in β, the difference $E_{n+1}(P_x, p_z) - E_n(P_x, p_z) = \hbar\Omega_1(\zeta, p_z)$ (ζ - Fermi energy); i.e., the number of states on the interval $\Delta E >> \hbar\Omega_1$ is determined only by the first (large) band. At the same time, the small band 2 greatly influences the stationary vectors of state $|\mu\rangle$ ($\mu = \{n, P_x, p_z\}$) with accuracy to quantities $\sim\beta$ equal to $c_\mu^{(+)}|\psi_+\rangle + c_\mu^{(-)}|\psi_-\rangle$, where $|\psi_\pm\rangle$ are the usual quasi-classical functions describing motion over the upper (+) and lower (-) sections of the first orbit, and the coefficients $c^\pm$ are connected by the relation

$$|c_\mu^{(+)}|^2 - |c_\mu^{(-)}|^2 = m_1^*(\partial E_n/\partial P_x)2by\ (\langle\psi_\pm/\psi_\pm\rangle = 1),$$

and the average transverse velocity is $\partial E_n/\partial P_x \sim v_0$. It can be stated that the width of the bands E_n (unlike the difference $E_{n+1} - E_n$) depends periodically on the parameter $S_2(\zeta, p_z)/\sigma >> 1$ (the period is equal to 2). Together with it, there oscillate in S_2/σ both $c^\pm$ and the matrix elements $\langle\mu|\hat{\phi}|\mu'\rangle$ of the arbitrary physical quantity $\hat{\phi}$. Under resonance conditions this circumstance leads to magnetic breakdown giant oscillations of the sound absorption coefficient

$$\Gamma(H) = \mathrm{Re}\sum_{\mu\mu'} |\langle\mu|\Lambda(\hat{p}, \hat{r})|\mu'\rangle|^2 (\partial f_0/\partial E_\mu)/\{i \times [(E_\mu - E_{\mu'})\hbar + \omega] - \nu_0\} V_0 I_0 \tag{1}$$

where $f_0(E)$ is the Fermi function, $\Lambda(\hat{p}, \hat{r})$ is the energy of interaction of the electrons with sound, which will not be presented in detailed form. V_0 and I_0 are the volume of the crystal and the density of the sound-wave energy flux, respectively.

We consider first the magnetic breakdown giant oscillations for the case of weak spatial dispersion $\nu_0 << qv_0 << \Omega_1$ and $q_x = 0$. After certain transformations, which are accurate to within $\sim\beta/W$, we can represent formula (2) in the integration with respect to P_x in the form

$$\Gamma(H) = \Gamma^\infty + \frac{1}{2\pi^2\hbar^3 V_0 I_0}\int dE dp_z \times \frac{(\partial f_0/\partial E)\nu_0 m_1^*\ |\langle\Lambda\rangle_+ - \langle\Lambda\rangle_-|^2\, 2\tau\sin(S_2/2\sigma)}{[(\omega - q_z v_z^{(1)})^2 + \nu_0^2]\sqrt{\rho^4 + 4\tau^2\sin^2(S_2/2\sigma)}}, \tag{2}$$

where $v_z^{(1)} = -(\partial S_1/\partial p_z)_E/m_1^*$ is the average velocity along H for the first zone, $\rho^2 = W$, $\tau^2 + \rho^2 = 1$. Γ^∞ is the limiting value of Γ(H) at $H \to \infty$, $\langle\Lambda\rangle_\pm$ are the "classical" mean values of Λ over the section "±." In formula (2) the rapidly oscillating function $\sin(S_2/\sigma)$ "competes" with the resonant denominator, which is small near the resonant point $\overline{p_z}$, where $\omega = q_z v_z^{(1)}$. (Owing to the smallness of ω/qv_0, the derivative $(\partial S_1/\partial p_z)_E$ in p_z is small.) In the most interesting and typical case of closely-located extremal points $S_1(\zeta, p_z)$ and $S_2(\zeta, p_z)$, the function $\sin(S_2/\sigma)$ inside the resonant region can be regarded as constant if $\gamma_q \equiv \nu_0\beta/qv_0\sqrt{\kappa} \ll 1$, $\gamma_T \equiv kT/\hbar\Omega_2 \ll 1$ (T is the temperature). Here

$$\Gamma(H) = \Gamma^{\infty} + (\Gamma_0 - \Gamma^{(\infty)})2r|\sin(S_2/\sigma)| \times [\rho^4 + 4r^2\sin(S_2/2\sigma)]^{-1/2} \tag{3}$$

Γ_0 is $\Gamma(H)$ in the absence of breakdown, and $S_2 = S_2(\zeta, \bar{p}_z)$. According to (3), the oscillations of $\Gamma(H)$ have a period $\Delta(H^{-1}) = 2\pi e\hbar/cS_2(\zeta, \bar{p}_z)$ and an amplitude ~ 1, i.e., they are "gigantic." In the other limiting case, $\gamma_q, \gamma_T \gg 1$, the amplitude of the oscillations is $\sim\exp(-\gamma)$, and an expression for Γ is obtained from (3) by averaging the latter over the argument $\Phi_2 = S_2/\sigma$. The criteria for the occurrence of magnetic breakdown giant oscillations ($\gamma_q, \gamma_T \lesssim 1$) differ from the corresponding criteria for the GSF effect by a small factor β, which improves the possibility of their observation. For $\alpha \lesssim 10^{-2}$, $\nu_0 \sim 10^9$ sec^{-1}, $H \sim 10^4 - 10^5$ Oe. Magnetic-breakdown giant oscillations are possible for $\omega > 10^7$ sec^{-1}. We emphasize that formulas (2) and (3) themselves were obtained for the case when many terms E_n take part in the resonance, and consequently there are no GSF oscillations.

For $qv_0 \gtrsim \Omega_0$ and $q_x \neq 0$, the position of $\bar{p}_z$ depends essentially on

$$P_x(E(P_x + \hbar q_x; \bar{p}_z + \hbar q_z) + E_{n'}(P_x, p_z) = \hbar\omega).$$

Integration with respect to P_x destroys the sharpness of the resonance and therefore the magnetic-breakdown giant oscillations exist only in the vicinity of the points $H_r = cq_x b_y/2\pi e r$, where $\hbar q_x = 2\pi\sigma/b_y$ (r is an integer). The width of these regions

$$\Delta H_r \sim \sqrt{\kappa}\beta^{-1}H_r.$$

The considered effect takes place for all configurations (closed and open), in which all the large trajectories (of the type "±") are almost equal. This is satisfied for all directions of $\vec{H}$ close to the symmetry axis of the crystal. In the opposite case the spectrum becomes "randomized" (see [6]) and an effective broadening of the resonant region takes place. As shown by estimates, for the occurrence of magnetic-breakdown giant oscillations the angle θ between $\vec{H}$ and the symmetry axis should be $\lesssim\sqrt{\kappa}/\beta$ ($\theta \lesssim 10°$).

A detailed analysis of the entire situation using the general method of [6] will be presented in a separate article.

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SPATIAL SELF-MODULATION OF LIGHT AS A RESULT OF ABSORPTION

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We investigate here the process of spatial self-modulation of light upon absorption (SSMA). The self-modulation is the result of the scattering of the interfering coherent waves interacting in the absorbing medium.

The gist of the SSMA phenomenon can be explained in the following manner. Assume that at the initial instant of time t = 0 two plane linearly polarized waves of equal amplitude and frequency are incident from the left on the boundary (x = 0) of an absorbing medium. In the linear approximation the field in it can be represented in the form

$$\frac{1}{2}E_0 e^{-i\omega t} + \text{KC} = \frac{A}{2}e^{-i\omega t + ikx - \alpha x}(e^{iqy} + e^{-iqy}) + \text{KC}, \qquad (1)$$

where $k = (\omega/c)\sqrt{\varepsilon'}$, $\alpha = (1/2)(\omega/c)\varepsilon''/\sqrt{\varepsilon'}$, and $q = k\tan(\theta/2)$. The dielectric constant is $\varepsilon = \varepsilon' + i\varepsilon''$, where $\varepsilon'' \ll \varepsilon'$. As a result of the absorption, a power $Q(r, t)$ is released per unit volume and is proportional to the square of the modulus of the field $|E|^2$. This leads to uneven heating of the medium, which in turn causes a change in the dielectric constant of the medium:

$$\epsilon_1 \approx (\partial\epsilon/\partial T)_p T_1. \qquad (2)$$

Assuming that the entire released power goes over into heat and that the time of observation is small compared with the characteristic time of equalization, due to thermal conductivity, of the temperature of the unevenly heated matter ($t \ll 1/\chi q^2$), we obtain

$$T_1 = \frac{1}{\rho C_p}\int_0^t Q(r, t)dt. \qquad (3)$$

The result is a periodic change, with respect to the coordinate y, of the dielectric constant of the medium. This leads to the appearance of a number of spatial spectra of the scattered field E_1, i.e., to SSMA.

Thus, in essence this effect is close on the one hand to stimulated scattering of light due to absorption (SSA) [1, 2], and on the other hand to the effect of thermal self-focusing and defocusing of light [3]. This is natural. In a macroscopic description, the wave phenomena connected with self-action of light through a medium can be unified under the term self-diffraction of light, taking it, as in the linear theory [4], in a broad sense, i.e., including in self-diffraction the aggregate of all the phenomena occurring during the propagation of the waves in the medium whose inhomogeneities are produced by the wave field itself.

We continue the description of SSMA with an example of two interacting waves. In the given-field approximation ($|E_0| \gg |E_1|$) we obtain for the summary field $E = E_0 + E_1$ from Maxwell's equations and from (2) - (3)