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Combined Resonance of the Interlayer Conductivity in Quasi-Two-Dimensional Conductors

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Analytical expressions have been obtained for the resonance component of the interlayer conductivity of quasi-two-dimensional conductors in an inclined magnetic field in the presence of the Rashba–Dresselhaus spin–orbit coupling. It has been shown that the Dresselhaus interaction makes the main contribution to resonances at combination frequencies in the range of angles between the magnetic field and normal to the conducting layers where angular oscillations of the conductivity occur. The results can be used to experimentally determine the absolute values of the spin–orbit coupling constants.

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INTRODUCTION

The study of magnetic resonance absorption is a verified method for obtaining information on the band structure and other characteristics of conducting systems. In a uniform electromagnetic field and when the spin–orbit coupling is negligible, the orbital and spin dynamics of electrons are independent and resonance absorption is due either to transitions between Landau levels or to spin flip. The spin–orbit interaction couples the orbital and spin motions and allows combination resonance caused by transitions with change in both the number of the Landau level and the spin projection [1, 2].

The spin–orbit coupling in normal metals is usually negligibly weak; for this reason, the combination resonance of the conductivity occurs in semiconductors, semimetals [1, 2], and two-dimensional systems [3]. Layered conducting structures with a quasi-two-dimensional (Q2D) electron energy spectrum constitute an additional class of materials where combination resonance is possible. In particular, low-dimensional organic conductors are strongly anisotropic conductors where high-frequency resonances of different types were experimentally observed (see reviews [4–6] and references therein). The basic structural elements of these materials are organic molecules, e.g., TTF, BEDT–TTF, and BEDO–TTF, which have donor or acceptor properties. Ion radicals of these molecules in Q2D conductors are packed into conducting layers separated by molecular counterion layers. The electrical conductivity along layers in some compounds at room temperature can exceed $10^3 \Omega^{-1}/\text{cm}$ and increases with decreasing tempera-

ture, whereas the conductivity in the transverse direction is about $1 \Omega^{-1}/\text{cm}$ and lower.

Although organic conductors have a complex molecular and crystal structure, their electron band structures are quite simple. The Fermi surface of organic conductors is strongly anisotropic and can consist of quasi-one-dimensional and Q2D sheets. Studies of angular oscillations of the magnetoresistance and quantum magnetic oscillation effects [7] at liquid helium temperatures show that Q2D elements usually have the shape of a weak corrugated cylinder. Tetrathiafulvalene salts β –(BEDT–TTF)₂IBr₂ and β –(BEDT–TTF)₂I₃ exemplify Q2D organic metals whose Fermi surface consists of only one weakly corrugated cylinder. The anisotropy of the electron energy spectrum of a Q2D conductor can be characterized by a small parameter η whose square is the ratio of the conductivities along the normal \mathbf{n} to the layers and in the planes of the layers in the absence of a magnetic field. The energy of the electron in the field of the crystal lattice in the tight binding approximation can be represented in the form of a rapidly convergent series

$$\varepsilon(\mathbf{p}) = \varepsilon_0(p_x, p_y) + \sum_{n=1}^{\infty} \varepsilon_n(p_x, p_y) \cos \frac{np_z}{p_0}. \quad (1)$$

Here, $p_0 = \hbar/a$, a is the distance between the layers, and \hbar is the reduced Planck constant. The functions $\varepsilon_n(p_x, p_y)$ decrease strongly with an increase in their number; the largest function is $\varepsilon_1 = \eta\varepsilon_F$, where ε_F is the Fermi energy. The area of the cross section of the

Fermi surface $S(\varepsilon_F, p_B)$ by the $p_B = \text{const}$ plane depends slightly on the projection of the electron momentum p_B on the direction of the magnetic field \mathbf{B} : such a dependence occurs only in the first order in the anisotropy parameter η . The resonance effects at the absorption of electromagnetic radiation in Q2D systems should be more pronounced than those in quasi-isotropic conductors at comparable mean free paths of charge carriers because almost all electrons on the Fermi surface, rather than a separate group in the extreme cross section of the Fermi surface, are involved in the formation of these effects.

The difference of layered conductors in physical properties from both normal metals and two-dimensional conducting systems is primarily manifested in the transport phenomena in the direction of the normal to the layers, in particular, in the appearance of a series of maxima of the magnetoresistance under the variation of the angle θ between the vectors \mathbf{B} and \mathbf{n} [8, 9]. Angular oscillations of the interlayer magnetoresistance were observed not only in organic metals [7] but also in other low-dimensional inorganic layered conductors, e.g., [10–12]. In this work, the combination resonance of the interlayer conductivity in Q2D conductors in an inclined magnetic field is theoretically studied within the model of Rashba [1, 3] and Dresselhaus [13] spin–orbit coupling. Analytical expressions are obtained for the resonance component of the conductivity depending on the magnitude and direction of the magnetic field. It is shown that the Dresselhaus interaction makes the main contribution to resonances at combination frequencies in the angular range $\tan \theta \gg 1$ where angular oscillations of the conductivity occur.

EQUATION FOR THE CURRENT DENSITY

The current density describing the time and spatial dispersion has the form [14]

$$j_i(\mathbf{r}, t) = \int_{-\infty}^t dt' \int d^3 \mathbf{r}' G_{ik}(\mathbf{r}, \mathbf{r}', t - t') E_k(\mathbf{r}', t'). \quad (2)$$

The kernel of the integral operator (retarded Green's function) in the single-particle approximation can be written as

$$G_{ik}(\mathbf{r}, \mathbf{r}', t - t') = \sum_{\mathbf{v}, \mathbf{v}'} e^{-i(\omega_{\mathbf{v}\mathbf{v}'} - i/\tau_{\mathbf{v}\mathbf{v}'}) (t - t')} \frac{f_{\mathbf{v}'} - f_{\mathbf{v}}}{\varepsilon_{\mathbf{v}} - \varepsilon_{\mathbf{v}'}} \times \langle \mathbf{v} | \hat{j}_k(\mathbf{r}') | \mathbf{v}' \rangle \langle \mathbf{v}' | \hat{j}_i(\mathbf{r}) | \mathbf{v} \rangle. \quad (3)$$

Here, $\mathbf{E}(\mathbf{r}, t)$ is the ac electric field, $f_{\mathbf{v}} = [1 + \exp T^{-1}(\varepsilon_{\mathbf{v}} - \mu)]^{-1}$ is the equilibrium distribution function of quasiparticles with the energy $\varepsilon_{\mathbf{v}}$ in an individual state with the quantum numbers \mathbf{v} and temperature T , μ is the chemical potential, $\omega_{\mathbf{v}\mathbf{v}'} = (\varepsilon_{\mathbf{v}} - \varepsilon_{\mathbf{v}'})/\hbar$,

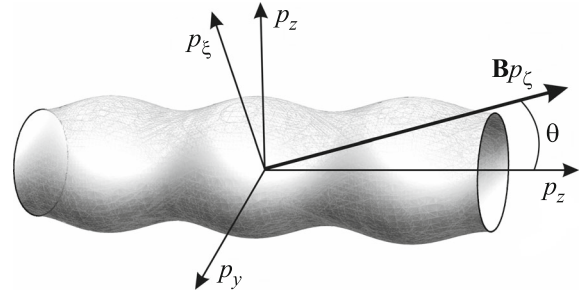


Fig. 1. Fermi surface and coordinate systems.

$\tau_{\mathbf{v}\mathbf{v}'}^{-1} = (\tau_{\mathbf{v}}^{-1} + \tau_{\mathbf{v}'}^{-1})/2$, and $\tau_{\mathbf{v}}$ and $\tau_{\mathbf{v}'}$ are the phenomenological lifetimes of quasiparticles in the states \mathbf{v} and \mathbf{v}' , respectively. For the processes considered below, the width $\hbar/\tau_{\mathbf{v}}$ of the level $\varepsilon_{\mathbf{v}}$ should be much smaller than the distance $\Delta\varepsilon = \varepsilon_{\mathbf{v}} - \varepsilon_{\mathbf{v}_1}$ between adjacent energy levels. In Eq. (3), $|\mathbf{v}\rangle$ are the two-component spinors, which are the eigenfunctions of the single-particle Hamiltonian $\hat{\varepsilon}(\hat{\mathbf{p}})$ and $\langle \mathbf{v} | \hat{j}_i(\mathbf{r}) | \mathbf{v}' \rangle$ are the matrix elements of the current density operator

$$\hat{\mathbf{j}}(\mathbf{r}) = \frac{e}{2} \{ \hat{\mathbf{v}}(\hat{\mathbf{p}}) \delta(\mathbf{r} - \mathbf{r}') + \delta(\mathbf{r} - \mathbf{r}') \hat{\mathbf{v}}(\hat{\mathbf{p}}) \} + c \text{curl} \hat{\boldsymbol{\mu}}_0 \delta(\mathbf{r} - \mathbf{r}'). \quad (4)$$

Here, $\hat{\mathbf{v}} = \partial \hat{\varepsilon} / \partial \hat{\mathbf{p}}$, $\hat{\mathbf{p}} = -i\hbar \nabla / \partial \mathbf{r} - e\mathbf{A}_0(\mathbf{r})/c$ is the kinematic momentum operator, e is the elementary charge, c is the speed of light, $\hat{\boldsymbol{\mu}}_0 = -(g\mu_B/2)\hat{\boldsymbol{\sigma}}$ is the magnetic moment operator of a conduction electron, μ_B is the Bohr magneton, g is the effective g -factor, $\hat{\boldsymbol{\sigma}}$ are the Pauli matrices, and $\mathbf{A}_0(\mathbf{r})$ is the vector potential of a uniform static magnetic field.

Formula (1) is written in the xyz coordinate system, where the z axis is the direction of the minimum conductivity and the y axis can be perpendicular to the magnetic field $\mathbf{B} = (B \sin \theta, 0, B \cos \theta)$. We use another coordinate system $\xi y \zeta$ with the ζ axis along the vector \mathbf{B} and take the gauge of the vector potential $\mathbf{A}_0(\mathbf{r}) = (-By, 0, 0)$ (Fig. 1). The momentum components in both coordinate systems are related to each other by the transformation of rotation by the angle θ between the normal to the layers and the magnetic field.

The Hamiltonian of the electron is taken in the form

$$\hat{\varepsilon}(\hat{\mathbf{p}}) = \hat{\varepsilon}_0(\hat{p}_x, \hat{p}_y) - \hat{\boldsymbol{\mu}}_0 \mathbf{B} + \sum_{n=1}^{\infty} \hat{\varepsilon}_n(\hat{p}_x, \hat{p}_y) \cos \frac{n\hat{p}_z}{p_0} + \hat{V}_{\text{so}}, \quad (5)$$

with the spin–orbit interaction

$$\hat{V}_{\text{so}} = \hat{V}_{\text{R}} + \hat{V}_{\text{D}} = \gamma_{\text{R}} \hat{\boldsymbol{\sigma}}(\hat{\mathbf{p}} \times \mathbf{n}) + \gamma_{\text{D}} \hat{\boldsymbol{\sigma}}(\mathbf{e}_x \hat{p}_x - \mathbf{e}_y \hat{p}_y), \quad (6)$$

which is the sum of the Rashba and Dresselhaus interaction operators. Here, \mathbf{n} is the unit vector of the high symmetry axis of the crystal, which is assumed to coincide with the normal to the layers; γ_R and γ_D are the Rashba and Dresselhaus spin-orbit coupling constants, respectively; \mathbf{e}_x and \mathbf{e}_y are the unit vectors along the x and y axes, respectively; and $\hat{\boldsymbol{\sigma}} = (\sigma_\xi, \sigma_y, \sigma_\zeta)$, where

$$\hat{\sigma}_\xi = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \hat{\sigma}_y = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \hat{\sigma}_\zeta = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}.$$

For the considered conducting systems, the operators of both the energy of motion of the electron along the normal to the layers and the spin-orbit coupling should be treated as perturbations. The complete set of quantum numbers $\mathbf{v} = n, p_\xi, p_B, \sigma$ consists of the number of the Landau level n , the momentum projections p_ξ and p_B , and spin projection $s_\zeta = \sigma_\zeta/2 \equiv \sigma/2$. The spin component of the wavefunction in the zeroth approximation is an eigenfunction of the operator $\hat{\sigma}_\zeta$.

The canonical transformation of the operators $\hat{\varepsilon}$ and $\hat{\mathbf{j}}$

$$\hat{\varepsilon}' = e^{-\hat{S}} \hat{\varepsilon} e^{\hat{S}} = \hat{\varepsilon} + [\hat{\varepsilon}, \hat{S}] + \frac{1}{2!} [[\hat{\varepsilon}, \hat{S}], \hat{S}] + \dots \quad (7)$$

reduces the Hamiltonian to the spin-diagonal form. In the zeroth approximation, the transformed Hamiltonian coincides with the unperturbed operator $\hat{\varepsilon}^{(0)}$ and the matrix elements of the operator \hat{S} in the first order in \hat{V}_{so} are

$$\langle \mathbf{v} | \hat{S} | \mathbf{v}' \rangle = \frac{\langle \mathbf{v} | \hat{V}_{so}^N | \mathbf{v}' \rangle}{\varepsilon_{\mathbf{v}'}^{(0)} - \varepsilon_{\mathbf{v}}^{(0)}}, \quad (8)$$

where \hat{V}_{so}^N is the spin off-diagonal component of the operator \hat{V}_{so} . In this formula, it is assumed that the frequency of the electron paramagnetic resonance (EPR) $\omega_s = g\mu_B B/\hbar$ is not equal to the cyclotron frequency ω_B and its harmonics $l\omega_B$; i.e., combination and cyclotron resonance lines should not coincide; otherwise, the denominator in Eq. (8) can be zero at some n and n' values.

The matrix elements of the operator $\hat{j}_z^{(l)}$ of the component of the orbital part of the current density (4) are diagonal in spin. Therefore, the combination resonance of the interlayer conductivity is determined by the product $\langle \mathbf{v} | \hat{j}_z^{(l)'}(\mathbf{r}) | \mathbf{v}' \rangle \langle \mathbf{v}' | \hat{j}_z^{(l)}(\mathbf{r}) | \mathbf{v} \rangle$ in Eqs. (2) and (3), where $\hat{j}_z^{(l)'}(\mathbf{r}) = [\hat{j}_z^{(l)}(\mathbf{r}), \hat{S}]$.

COMBINATION RESONANCE OF THE INTERLAYER CONDUCTIVITY

Under the assumption that $\mathbf{E}(\mathbf{r}, t)$ and $\mathbf{j}(\mathbf{r}, t) \sim \exp(i\mathbf{k}\mathbf{r} - i\omega t)$, the correction to the interlayer conductivity, which describes the resonances at combination frequencies can be easily obtained from Eqs. (2) and (3) in the form

$$\sigma_{zz}^{(so)} = i \sum_{\mathbf{v}, \mathbf{v}'} \frac{f_{\mathbf{v}'} - f_{\mathbf{v}}}{\varepsilon_{\mathbf{v}'} - \varepsilon_{\mathbf{v}}} \frac{\langle \mathbf{v} | \hat{j}_z^{(l)'}(0) | \mathbf{v}' \rangle \langle \mathbf{v}' | \hat{\mathcal{J}}_z^{(l)'}(\mathbf{k}) | \mathbf{v} \rangle}{\omega_{\mathbf{v}\mathbf{v}'} - \omega - i\tau_{\mathbf{v}\mathbf{v}'}^{-1}}, \quad (9)$$

where $\langle \mathbf{v}' | \hat{\mathcal{J}}_z^{(l)'}(\mathbf{k}) | \mathbf{v} \rangle = \int d^3\mathbf{r} e^{-i\mathbf{k}\mathbf{r}} \langle \mathbf{v}' | \hat{j}_z^{(l)'}(\mathbf{r}) | \mathbf{v} \rangle$. When the electric current flows along the normal to the layers, the depth of the skin layer δ is determined by the conductivity component σ_{zz} , which is smaller than the conductivity in the plane of the layers by a factor of $\eta^{-2} \approx 10^3 - 10^5$, and the spatial dispersion can usually be neglected [15].

To determine the matrix elements $\langle \mathbf{v} | \hat{j}_z^{(l)'}(\mathbf{r}) | \mathbf{v}' \rangle$, it is necessary to use a particular model of the electron energy spectrum. Let the energy of the electron (1) be determined by the zeroth and first Fourier harmonics of the momentum projection on the normal to the layers. Neglecting anisotropy in the plane of the layers and setting $\varepsilon_1(p_x, p_y) = -\varepsilon_\eta$, one can represent Eq. (1) in the form

$$\varepsilon(\mathbf{p}) = \frac{p_x^2 + p_y^2}{2m} - \varepsilon_\eta \cos \frac{p_z}{p_0}, \quad (10)$$

where m is the effective mass, $\varepsilon_\eta = \eta v_F p_0$, and $v_F = \sqrt{2\varepsilon_F/m}$. The case where the energy spectrum in the plane of the layers is specified by an arbitrary positive definite quadratic form of the quasimomentum is reduced to Eq. (10) by means of transformations of rotation and extension of the coordinate axes. The Schrödinger equation for the unperturbed Hamiltonian $\hat{\varepsilon}^{(0)}$ is reduced to the equation of a harmonic oscillator with the frequency $\omega_B = |e|B \cos \theta/(mc)$. If the angle θ between the magnetic field and normal to the layer is close to $\pi/2$, the closed cross sections of the Fermi surface $S(\varepsilon_F, p_B)$ are strongly prolonged, the electron cannot make the complete semiclassical orbit in the momentum space in the mean free path time, and dynamic resonance effects are not manifested. For this reason, the inequality $\eta \tan \theta \ll 1$ is assumed.

The width of the resonance can be represented in the form $\tau_{\mathbf{v}\mathbf{v}'}^{-1} \equiv \tau_s^{-1} = \tau_s^{-1} + \tau_l^{-1}$, where τ_s is the spin-flip time and τ_l^{-1} is the width of the resonance at the transition of the electron from the Landau level

$n' = n + l$ to the level n with the conservation of the spin projection. The simple calculations give

$$\begin{aligned} \sigma_{zz}^{(so)} = & i\eta^2 \frac{\omega_p^2 \omega_B}{\pi^2} w \sum_n \int d\beta \left\{ \frac{f_{n,-1} - f_{n,1}}{\omega_s} |a_{n0}|^2 h_0 \right. \\ & \left. + \sum_{l=1}^{\infty} \left(\frac{f_{n,-1} - f_{n+l,1}}{\Omega_l^{(+)}} h_l^{(+)} + \frac{f_{n,1} - f_{n+l,-1}}{\Omega_l^{(-)}} h_l^{(-)} \right) |a_{nl}|^2 \right\}. \end{aligned} \quad (11)$$

Here, $\omega_p = \sqrt{4\pi n_0 e^2 / m}$ is the plasma frequency, n_0 is the electron density, and $\beta = p_B / (p_0 \cos \theta)$. The subscript of the conserving quantity p_B is omitted in the equilibrium distribution function of electrons f_{n,σ,p_B} . The functions of the frequency

$$h_0 = \frac{\omega + i\tau_s^{-1}}{(\omega + i\tau_s^{-1})^2 - \omega_s^2}, \quad h_l^{(\pm)} = \frac{\omega + i\tau_{ls}^{-1}}{(\omega + i\tau_{ls}^{-1})^2 - \Omega_l^{(\pm)2}} \quad (12)$$

at $\tau_{ls}^{-1} \rightarrow 0$ have sharp maxima at the combination resonance frequencies $\omega_{\text{res}} = |\Omega_l^{(\pm)}|$, where $\Omega_l^{(\pm)} = l\omega_B \pm \omega_s$, whereas the function

$$w = \frac{\gamma_R^2 (\omega_B + \omega_s \cos \theta)^2 + \gamma_D^2 (\omega_s - \omega_B \cos \theta)^2}{a_B^2 (\omega_B^2 - \omega_s^2)^2} \quad (13)$$

of the angle θ determines the contributions to $\sigma_{zz}^{(so)}$ from the Rashba and Dresselhaus interactions, where $a_B = \sqrt{\hbar / (m\omega_B)}$. At large n values, the coefficients

$$a_{nl} = \sqrt{\frac{n!}{(n+l)!}} \frac{\alpha^{l+1}}{2\sqrt{2}^l} e^{-\alpha^2/4} L_n' \left(\frac{\alpha^2}{2} \right) \cos \left(\beta + \frac{\pi l}{2} \right) \quad (14)$$

describe the oscillatory dependence of $\sigma_{zz}^{(so)}$ on the angle θ , where $\alpha = (a_B/r_0) \tan \theta$, $r_0 = cp_0 / (|e|B)$, and $L_n'(x)$ are the associated Laguerre polynomials. The asymptotic representation of Laguerre polynomials [16] allows the expression of a_{nl} in terms of the Bessel functions J_l

$$a_{nl} = \frac{1}{2} \alpha J_l(\sqrt{2n\alpha}) \cos \left(\beta + \frac{\pi l}{2} \right), \quad n \gg 1.$$

In the expression for the energy of the electron

$$\varepsilon_{n,\sigma,p_B} = \hbar\omega_B \left(n + \frac{1}{2} \right) + \frac{\hbar\omega_s \sigma}{2} - \varepsilon_{\perp} \cos \beta,$$

where

$$\varepsilon_{\perp} = \varepsilon_{\eta} e^{-\alpha^2/4} L_n \left(\frac{\alpha^2}{2} \right) \xrightarrow{n \gg 1} \varepsilon_{\eta} J_0(\sqrt{2n\alpha}),$$

terms of the second order in \hat{V}_{so} , i.e., proportional to γ_R^2 and γ_D^2 , are neglected because they give corrections to the energy levels but do not affect the intensity of

the resonance and the angular dependence of the kinetic coefficients.

Each term in the sum over l in Eq. (11) determines the asymptotic behavior of the conductivity $\sigma_{zz}^{(so)}$ near the $\pm l$ th resonance $\omega \approx |\Omega_l^{(\pm)}|$. The first term with $l = 0$ in Eq. (11) corresponds to pure spin transitions and the width of the resonance is determined by the inverse spin flip time τ_s^{-1} .

According to Eq. (14) for a_{nl} , the conductivity $\sigma_{zz}^{(so)}$ in a magnetic field perpendicular to the layers, i.e., at $\theta = 0$, and in the absence of spatial dispersion vanishes within the chosen models of the energy of the electron in the field of the crystal lattice and the spin-orbit coupling. The coefficients a_{nl} in the spatially nonuniform field are also nonzero at $\mathbf{B} \parallel \mathbf{n}$. In particular, when the electric field is $\mathbf{E}(\mathbf{r}) \sim \exp(iky)$ (dispersion in the direction of \mathbf{B} is manifested in the first order in η), at $\theta = 0$ and $\varepsilon_F \gg \hbar\omega_B$, the conductivity $\sigma_{zz}^{(so)}$ is given by Eq. (11) with

$$|a_{nl}|^2 = \frac{1}{4} q^2 J_l^2(\sqrt{2nq}) \sin^2 \beta,$$

where $q = a_B k$, $\sqrt{2nq} = kr_B$, and $r_B = v_F / \omega_B$ is the cyclotron radius. Angular oscillations occur at large $\tan \theta$ values. For the considered geometry of the problem, $kr_B \approx k\delta^{-1} < 1$ and the qualitative picture of the angular dependence of the resonance part of the interlayer conductivity given by Eq. (9) is described by Eq. (11).

If the number of Landau levels below ε_F is large, the conductivity specified by Eq. (11) undergoes de Haas-van Alphen oscillations. The Poisson formula and standard transformations allow the representation of Eq. (11) in the form

$$\sigma_{zz}^{(so)} = \bar{\sigma}_{zz} + \tilde{\sigma}_{zz}, \quad (15)$$

where

$$\begin{aligned} \bar{\sigma}_{zz} = & i\eta^2 \frac{\omega_p^2}{4\pi} \alpha^2 w \\ & \times \left\{ J_0^2(\alpha_1) h_0 + \sum_{l=1}^{\infty} J_l^2(\alpha_1) (h_l^{(+)} + h_l^{(-)}) \right\} \end{aligned} \quad (16)$$

is the smooth component and

$$\begin{aligned} \tilde{\sigma}_{zz} = & i\eta^2 \frac{\omega_p^2 \omega_B}{2\pi} \alpha^2 w \left\{ J_0^2(\alpha_1) \frac{h_0}{\omega_s} \mathcal{Q}_0 \right. \\ & \left. + \sum_{l=1}^{\infty} J_l^2(\alpha_1) \left(\frac{h_l^{(+)}}{\Omega_l^{(+)}} - \frac{h_l^{(-)}}{\Omega_l^{(-)}} \right) \mathcal{Q}_l \right\}. \end{aligned} \quad (17)$$

is the oscillatory component as a function of B^{-1} . Here, $\alpha_l = \alpha\sqrt{2\mu/(\hbar\omega_B)} = (mv_F/p_0)\tan\theta$,

$$Q_l = \sum_{j=1}^{\infty} \frac{(-1)^j}{j} F_l(j\Delta) \psi(j\lambda) \sin \frac{\pi j \omega_s}{\omega_B} \cos \frac{2\pi j \mu}{\omega_B},$$

$$\psi(\lambda) = \frac{\lambda}{\sinh \lambda}, \quad \lambda = 2\pi^2 T / (\hbar\omega_B).$$

The amplitude of oscillating harmonics in the sum Q_l is modulated by the quasiperiodic function

$$F_l(\Delta) = J_0(\Delta) - (-1)^l J_2(\Delta),$$

where the argument $\Delta = 2\pi\varepsilon_\eta J_0(\alpha_l)/(\hbar\omega_B)$ depends not only on B^{-1} but also on $\tan\theta$.

At $\varepsilon_\eta \approx \eta\varepsilon_F \gg \hbar\omega_B$, Eq. (17) describes oscillations of $\sigma_{zz}^{(so)}$ as a function of the inverse magnetic field. In the case $\eta\varepsilon_F \approx \hbar\omega_B$ strictly speaking, it is necessary to take into account the oscillation dependence of the chemical potential μ on B^{-1} , which is determined from the equation of conservation of the number of electrons, and, near the resonances $\omega \approx |\Omega_l^{(\pm)}|$, the dependence of the relaxation time τ_{ls} on B^{-1} . At sufficiently low temperatures $T \leq \hbar\omega_B/(2\pi^2)$, at $\eta\varepsilon_F \gg \hbar\omega_B$ and arbitrary θ values, the oscillatory component of the conductivity $\tilde{\sigma}_{zz} \approx F_l(\Delta)\bar{\sigma}_{zz} \approx \bar{\sigma}_{zz}/\sqrt{\Delta}$ is smaller than the smooth component $\bar{\sigma}_{zz}$ by a factor of $\sqrt{\Delta} \sim \sqrt{\eta\varepsilon_F/(\hbar\omega_B)}$. However, the oscillatory component for $\theta = \theta_i$ at which $\alpha_l \sim \tan\theta_i$ is a root of the Bessel function $J_0(\alpha_l) = 0$ increases to about $\bar{\sigma}_{zz}$. For these directions of the magnetic field, the dependence of the area of the cross section of the Fermi surface $S(\varepsilon_F, p_B)$ on the momentum projection p_B appears in terms quadratic in η .

The resonance conductivity given by Eq. (11) can be represented as the sum $\sigma_{zz}^{(so)} = \sigma_R + \sigma_D$ of contributions from the Rashba and Dresselhaus interactions. According to Eqs. (11) and (13),

$$\frac{\sigma_R}{\sigma_D} = \frac{\gamma_R^2 (\omega_s + \omega_0)^2 \cos^2 \theta}{\gamma_D^2 (\omega_s - \omega_0 \cos^2 \theta)^2}, \quad (18)$$

where $\omega_0 = |e|B/(mc)$. In the angular range $\tan\theta \gg 1$, where angular oscillations of the conductivity are manifested, the asymptotic behavior of $\sigma_{zz}^{(so)}$ is determined by the Dresselhaus spin-orbit coupling, $\sigma_D \gg \sigma_R$. Formulas (11), (13), and (18) can be used to determine the absolute values of the constants γ_R and γ_D . The pure spin transitions with $l = 0$ at the frequency ω_s are of most interest because $\tau_s > \tau_l$ and, in

addition, an increase in θ is accompanied by a decrease in $\omega_B\tau_l$ and, correspondingly, in the intensity of the l th resonance.

CONCLUSIONS

A high-frequency nonuniform electromagnetic field can induce transitions at combination frequencies even in the absence of the spin-orbit coupling [17]. The current density given by Eq. (2) is proportional to the eddy electric field and, correspondingly, includes terms proportional to the magnetic field $\mathbf{B}^- = i(\omega/c)(\mathbf{k} \times \mathbf{E})$.

The spin part $\hat{\mathbf{j}}^{(s)} = c\text{curl}\hat{\boldsymbol{\mu}}_0\delta(\mathbf{r} - \mathbf{r}')$ of operator (4) in Eqs. (2) and (3) determines the high-frequency magnetization excited by the ac field \mathbf{B}^- and the paramagnetic susceptibility $\chi_{ik}(\omega, \mathbf{k})$. In the case $\mathbf{E}(\mathbf{r}), \mathbf{B}^-(\mathbf{r}) \sim \exp(iky)$, resonances of the magnetization at combination frequencies are due to the matrix elements $\langle v | \hat{j}_\xi^{(s)}(\mathbf{r}) | v \rangle \langle v' | \hat{j}_\xi^{(s)}(\mathbf{r}') | v \rangle$.

The structure of resonance absorption of microwave radiation in organic compounds is a superposition of peaks corresponding to different types of resonances. Because of the similarity of the pictures of combination resonance of the conductivity, EPR, and combination resonance of the spin magnetization, there is the problem of identification of experimental data. Unlike $\sigma_{zz}^{(so)}$, the paramagnetic susceptibility does not contain oscillating factors such as $J_l^2((mv_F/p_0)\tan\theta)$. A simple calculation shows that the angular dependence of the smooth part $\bar{\chi}_{\xi\xi}$ is determined by the dependence of resonance functions $h_l^{(\pm)}$ specified in Eqs. (12) on $\omega_B \sim \cos\theta$, and the angular dependence of the quantum part $\tilde{\chi}_{\xi\xi}$ oscillatory as a function of B^{-1} also includes additional factors $J_0(j\Delta)$. Thus, the analysis of the dependence of the absorbed power on the orientation of the magnetic field and the parameters of the Fermi surface makes it possible to certainly distinguish the combination resonance of the interlayer conductivity from EPR and other types of resonances.

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