Low Temperature Physics

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Cite as: Low Temp. Phys. **46**, 1021 (2020); https://doi.org/10.1063/10.0001914
Submitted: 21 August 2020 . Published Online: 21 October 2020

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Cite as: Fiz. Nizk. Temp. **46**, 1206–1215 (October 2020); doi: 10.1063/10.0001914
Submitted: 21 August 2020

View Online Export Citation CrossMark

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ABSTRACT

A theoretical investigation of the combined resonance of interlayer conductivity and spin magnetization, in conductors with quasi-two-dimensional electronic energy spectra. Analytical expressions are obtained for the surface impedance, magnetic susceptibility, and the resonance interlayer conductivity component caused by Rashba–Dresselhaus spin-orbit interaction, with allowance for spatial dispersion.

1. INTRODUCTION

High-frequency resonances are observed in almost all conducting systems that are placed in a strong external magnetic field, if the mean free path of the charge carriers is significant enough for their dynamic properties to manifest. If the spin-orbit interaction and the spatial inhomogeneity of the high-frequency electromagnetic field are neglected, then the electrons' orbital and spin dynamics are independent, and the resonance absorption in nonmagnetic conductors is caused either by transitions between Landau levels, or by the spin-flip. Spin-orbit interaction forms a connection between orbital and spin motion, and enables the resonance that is caused by transitions that occur simultaneously with changes in both the Landau level number, and the spin projection — the combined resonance.[1,2]

In common metals, spin-orbit coupling is usually insignificant. For example, experimental studies concerning the spin-Hall effect in single aluminum crystals[3,4] at helium temperatures yield a potential difference due to spin-orbit interaction that is about equal to 10^{-10}–10^{-9} V. In semiconductors, and two-dimensional electronic systems based on semiconductors, the kinetic and thermodynamic characteristics are very sensitive to features of the charge carrier energy spectrum, and even a small change in the energy bands due to the spin-orbit interaction can lead to noticeable effects.[5–11] For these reasons, combined resonance occurs primarily in semiconductors, semimetals,[1,2] and two-dimensional conducting systems.[5,12] Another type of material in which combined resonance is possible is a layered conducting structure with a quasi-two-dimensional (Q2D) electronic energy spectrum. Layered conductors of organic origin are an example of highly anisotropic conductors, in which different types of high-frequency resonances have been experimentally observed.[13–23] The main structural elements of these substances are the organic molecules or molecular complexes that have donor or acceptor properties. The most well-known examples of such molecules include TTF, BEDT-TTF, BEDO-TTF, etc. In Q2D conductors, the ion-radicals of these molecules are packed into layers that are separated by layers of counterion molecules. These organic molecules are in close proximity to each other, which leads to a significant overlap of the charge carrier wave functions, and as a result, the carriers can move freely from molecule to molecule, forming a conducting plane. In a direction perpendicular to the layers, the organic molecules are separated from each other, and the probability of charge carrier transfer from one conducting plane to another is small. In a number of compounds, the electrical conductivity along the layers at helium temperatures can exceed $10^6\ \Omega^{-1}\times \mathrm{cm}^{-1}$, and decreases with increasing temperature. In the transverse direction it will be 3–5 orders of magnitude lower.

Organic conductors have complex molecular and crystal structures, but their electronic band structures are simple. Their Fermi surface (FS) is strongly anisotropic, and can consist of sheets that are quasi-one-dimensional and Q2D. Studies on the angular oscillations of magnetoresistance, and quantum magnetic oscillations effects[24] at liquid helium temperatures, have shown that the Q2D elements of the FS of well-known organic compounds usually appear as a weakly corrugated cylinder. The anisotropy of the electronic energy spectrum of a Q2D conductor can be characterized by the small parameter η, the square of which is equal to the ratio of the conductivities along the normal $\mathbf{n}$ to the layers and in the plane of the layers, in the absence of a magnetic field. The FS cross section area $S_F(p_B)$ by the $p_B = \mathbf{pB}/B = \text{const}$ plane has a weak dependence on the electron momentum projection p_B in the direction of the magnetic field $\mathbf{B}$, and manifests only in the first-order of the anisotropy parameter η. The resonance phenomena occurring during the absorption of electromagnetic radiation in Q2D conductors should demonstrate more clearly than in ordinary metals with comparable mean free charge carrier paths, since almost all electrons on the FS surface are involved in their manifestation, and not just the select group at the extreme FS cross section.

The difference between the physical properties of layered conductors, ordinary metals, and two-dimensional conducting systems is primarily manifested by the transfer phenomena in the direction normal to the layers, especially in the appearance of a series of magnetoresistance maxima as the angle θ changes between vectors $\mathbf{B}$ and $\mathbf{n}$.[24,25] A prior work,[26] based on the Rashba[1,12]–Dresselhaus[27] spin-orbit interaction model, examines the combined resonance of interlayer conductivity in Q2D conductors with an inclined magnetic field, while neglecting spatial dispersion. It is shown that in the range of θ values in which the angular oscillations $\tan\theta \gg 1$ appear, the main contribution to the resonance at combined frequencies comes from Dresselhaus interaction. This article theoretically examines combined resonance of interlayer conductivity and spin magnetization, and accounts for spatial dispersion. A numerical analysis is performed, which provides a qualitative idea as to the dependence that the kinetic coefficients have on the angle between the magnetic field and the normal to the layers, and their dispersion properties.

2. CURRENT DENSITY EQUATION

According to general current density equations for a system of electrons in an alternating electromagnetic field,[28] the current density in the one-particle approximation, taking into account the time and space dispersion, can be written as

$$j_i(\omega,\mathbf{k}) = i\sum_{\nu,\nu'} \frac{f_0(\varepsilon_\nu)-f_0(\varepsilon_{\nu'})}{\varepsilon_\nu-\varepsilon_{\nu'}} \frac{\langle \nu|\hat{j}_k(0)\nu'\rangle \int d^3\mathbf{r}\, e^{-i\mathbf{kr}}\langle \nu|\hat{j}_i(\mathbf{r})\nu\rangle}{\omega_{\nu\nu'}-\omega-i\tau^{-1}_{\nu\nu'}} E_k(\omega,\mathbf{k}). \tag{1}$$

Here, $\mathbf{E}(\mathbf{r},t) = \mathbf{E}(\omega,\mathbf{k}) \exp(i\mathbf{kr} - i\omega t)$ is the electric field, $f_0(\varepsilon_\nu)$ is the equilibrium distribution function of conduction electrons with energy ε_ν, in an individual state with quantum numbers ν, $\omega_{\nu\nu'} = (\varepsilon_\nu - \varepsilon_{\nu'})/\hbar$, $\tau^{-1}_{\nu\nu'} = (\tau^{-1}_\nu + \tau^{-1}_{\nu'})/2$, and τ_ν and $\tau_{\nu'}$ are the phenomenological quasiparticle lifetimes in the ν and ν' states, $\hbar$ is Planck's constant, $\langle \nu|\hat{j}_i(\mathbf{r})|\nu'\rangle$ are the matrix elements of the current density operator

$$\hat{\mathbf{j}}(\mathbf{r}) = \frac{e}{2}\{\hat{\mathbf{v}}(\hat{\mathbf{p}})\delta(\mathbf{r}-\mathbf{r}') + \delta(\mathbf{r}-\mathbf{r}')\hat{\mathbf{v}}(\hat{\mathbf{p}})\} + c \,\text{rot}\, \hat{\boldsymbol{\mu}}_0\delta(\mathbf{r}-\mathbf{r}'), \tag{2}$$

which is the sum of the orbital $\hat{\mathbf{j}}^{(l)}(\mathbf{r})$ and spin $\hat{\mathbf{j}}^{(s)}(\mathbf{r})$ components, $\hat{\mathbf{v}} = \frac{\partial\hat{\varepsilon}(\mathbf{p})}{\partial\hat{\mathbf{p}}}$, $\hat{\mathbf{p}} = -i\hbar\partial/\partial\mathbf{r} - e\mathbf{A}_0(\mathbf{r})/c$ is the kinematic momentum operator, $A_0(\mathbf{r})$ is the vector potential of the constant uniform magnetic field $\mathbf{B}$, e is the electron charge, c is the speed of light, $\hat{\boldsymbol{\mu}}_0 = \left(\frac{g\mu_B}{2}\right)\hat{\boldsymbol{\sigma}}$ is the electron magnetic moment operator, μ_B is Bohr's magneton, g is the effective g-factor, and $\boldsymbol{\sigma}$ are Pauli matrices. The two-component spinors $|\nu\rangle$ are eigenfunctions of the single-particle Hamiltonian $\hat{\varepsilon}(\hat{\mathbf{p}})$. For the processes considered below, the width $\hbar/\tau_\nu$ of the ε_ν level should be significantly less than the distance $\Delta\varepsilon = \varepsilon_\nu - \varepsilon_{\nu 1}$ between adjacent energy levels.

The electron energy in the field of the crystal lattice of layered conductors is weakly dependent on the momentum projection $p_z = \mathbf{pn}$ on the normal to the layers, and can be written as a rapidly converging series in the tight binding approximation

$$\hat{\varepsilon}(\hat{\mathbf{p}}) = \varepsilon_0(p_x, p_y) + \sum_{n=1}^{\infty}\varepsilon_n(p_x, p_y, \eta)\cos\frac{np_z}{p_0}. \tag{3}$$

Here, the functions $\varepsilon_n(p_x,p_y,\eta)$ decrease significantly as their number increases, the largest being $\varepsilon_1(p_x,p_y,\eta) \simeq \eta\varepsilon_F$, where ε_F is the Fermi energy, $p_0 = \hbar/a$, a is the distance between layers. Formula (3) is written in the xyz coordinate system, in which the z-axis is parallel to the direction of the lowest conductivity, and the y axis can be directed perpendicular to the magnetic field $\mathbf{B} = (B\sin\theta, 0, B\cos\theta)$. In addition, let us use another coordinate system $\xi y\zeta$, in which the ζ axis is parallel to $\mathbf{B}$, and choose the calibration of the vector potential $\mathbf{A}_0(\mathbf{r}) = (-By,0,0)$ (Fig. 1). The momentum components in both coordinate systems are related by the rotation transformation with respect to the angle θ between the normal to the layers and the magnetic field.

In accordance with Formula (3), the electron Hamiltonian is determined by expression

$$\hat{\varepsilon}(\hat{\mathbf{p}}) = \varepsilon_0(\hat{p}_x, \hat{p}_y) - \hat{\boldsymbol{\mu}}_0\mathbf{B} + \sum_{n=1}^{\infty}\varepsilon_n(\hat{p}_x, \hat{p}_y, \eta)\cos\frac{n\hat{p}_z}{p_0} + \hat{V}_{so}. \tag{4}$$

Let us write the spin-orbit interaction operator as the sum

$$\hat{V}_{so} \equiv \hat{V}_R + \hat{V}_D = \gamma_R\hat{\boldsymbol{\sigma}}(\hat{\mathbf{p}}\times\mathbf{n}) + \gamma_D\hat{\boldsymbol{\sigma}}(\mathbf{e}_x\hat{p}_x - \mathbf{e}_y\hat{p}_y) \tag{5}$$

of the Rashba $\hat{V}_R$ and Dresselhaus $\hat{V}_D$ interaction operators. Here, $\mathbf{n}$ is the direction of the crystal's high symmetry axis, which is assumed to coincide with the normal to the layers, γ_R and γ_D are the interaction constants, $\mathbf{e}_x, \mathbf{e}_y$ are the unit vectors along the axes x,y, :

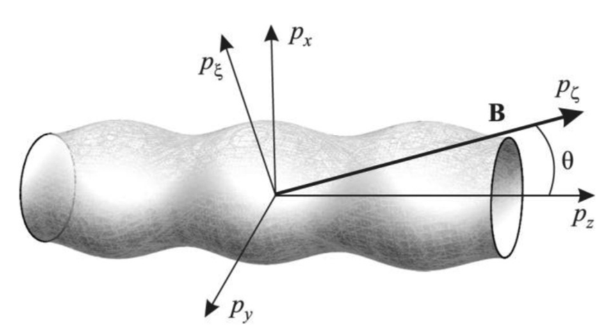

FIG. 1. The Fermi surface and coordinate systems.

$$\hat{\sigma}_\xi = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \hat{\sigma}_y = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \hat{\sigma}_\zeta = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (6)$$

For the considered conducting systems, the operators of the electron motion energy along the normal to the layers, and spin-orbit interaction, should be considered as perturbations. The full set of quantum numbers $\nu = n$, p_ξ, p_ζ,σ consists of the Landau level number n, the projections of the momentum p_ξ, p_ζ and spin $s_\zeta = \sigma_\zeta/2 \equiv \sigma/2$. The spin component $|\sigma\rangle$ of the zeroth-approximation wave function is an eigenfunction of the operator $\hat{\sigma}_\zeta$.

Let us perform the canonical transformation of the operators $\hat{\varepsilon}$ and $\hat{\mathbf{j}}$:

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

which reduce the Hamiltonian (4) to a spin-diagonal form. In the zeroth approximation, the transformed Hamiltonian coincides with the unperturbed operator (4) $\hat{\varepsilon}^{(0)}$, and the matrix elements of the $\hat{S}$ operator in the first order of $\hat{V}_{so}$ are equal to

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

where $\hat{V}_{so}^N$is the spin-off diagonal part of operator $\hat{V}_{so}$. This formula assumes that the spin resonance frequency $\omega_s = g\mu_B B/\hbar$ is not equal to the cyclotron frequency ω_B and its harmonics $l\omega_B$, i.e., the lines of the combined and cyclotron resonances should not coincide; otherwise, at some values of n and n' the denominator in Eq. (8) can vanish.

Using Eqs. (1) and (7), the current density can be written as a series in powers of spin-orbit interaction constants, in which the matrix elements $\langle \nu|\hat{j}_i'(\mathbf{r})|\nu'\rangle$ are calculated based on the eigenfunctions of the unperturbed Hamiltonian.

3. COMBINED RESONANCE OF INTERLAYER CONDUCTIVITY

The spin-diagonal matrix elements of the operator $\hat{j}_z^{(l)}$ of the orbital current density component (2) do not contribute to the current that is caused by transitions at the combined frequencies, and therefore the combined resonance of the interlayer conductivity is determined by the product

$$\langle \nu|\hat{j}_z^{(1)'}(0)|\nu'\rangle\langle \nu'|\int d^3\mathbf{r}\, e^{-i\mathbf{kr}}\hat{j}_z^{(1)'}(\mathbf{r})|\nu\rangle$$

in Eq. (1), where $\hat{j}_z^{(l)'}(\mathbf{r}) = [\hat{j}_z^{(l)}(\mathbf{r}), \hat{S}]$.

In order to find matrix elements $\langle \nu|\hat{j}_z^{(1)'}(\mathbf{r})|\nu'\rangle$, let us use the model conduction electron dispersion law. We restrict ourselves to the zeroth and first Fourier harmonics of the momentum projection onto the normal to the layers in Eq. (3), neglect the anisotropy in the plane of the layers, and set $\varepsilon_1(p_x, p_y, \eta) = -\varepsilon_\eta = -\eta v_F p_0$. As a result, we arrive at

$$\varepsilon(\mathbf{p}) = \frac{p_x^2 + p_y^2}{2m} - \varepsilon_\eta \cos\frac{p_z}{p_0}, \quad (9)$$

here m is the effective mass, $v_F=\sqrt{2\varepsilon_F/m}$. When the inequality

$$\eta \tan\theta \ll 1 \quad (10)$$

is satisfied, the Schrödinger equation for the unperturbed Hamiltonian $\hat{\varepsilon}^{(0)}$ is reduced to the harmonic oscillator equation having the frequency $\omega_B = |e|B\cos\theta/(mc)$.

Let us write the resonance width as $\tau_{\nu\nu'}^{-1} \equiv \tau_{ls}^{-1} = \tau_s^{-1} + \tau_l^{-1}$ where τ_s is the spin-flip time, and τ_l^{-1} is the resonance width at the transition of an electron from the Landau level $n' = n + l$ to the level n with conservation of the spin projection. Assuming that $\mathbf{E}(\mathbf{r}) \sim \exp(iky)$, after simple calculations, we find a correction to the interlayer conductivity, which describes the resonance at combined frequencies

$$\sigma_{zz}^{(so)}(\omega, k) = i\eta^2\frac{\omega_B\omega_p^2}{\pi^2}w(\theta) \times \sum_n \int_{-\pi}^{\pi} d\beta\left\{\frac{f_0(\varepsilon_{n,1}) - f_0(\varepsilon_{n,-1})}{\omega_s}|a_{n,n}(k)|^2 h_0(\omega) + \sum_{l=1}^{\infty}\left(\frac{f_0(\varepsilon_{n+l,1}) - f_0(\varepsilon_{n,-1})}{\omega_B l + \omega_s}h_l^{(+)}(\omega) + \frac{f_0(\varepsilon_{n+l,-1}) - f_0(\varepsilon_{n,1})}{\omega_B l - \omega_s}h_l^{(-)}(\omega)\right)|a_{n,n+l}(k)|^2\right\}, \quad (11)$$

where $\omega_p = \sqrt{4\pi n_0 e^2/m}$ is the plasma frequency, n_0 is the electron density, and $\beta = p_B/(p_0\cos\theta)$. The subscript of the conserved quantity p_B is omitted from the electron energy $\varepsilon_{n,\sigma,p_B} \equiv \varepsilon_{n,\sigma}$. The frequency functions

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

at $\tau_{ls}^{-1} \to 0$ have abrupt maxima at frequencies ω, equal to the combined resonance frequencies $\omega_r = |\Omega_l^{(\pm)}|$, where $\Omega_l^{(\pm)} = l\omega_B \pm \omega_s$, and the function of angle θ

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

determines the contributions from Rashba and Dresselhaus interactions to $\sigma_{zz}^{(so)}$ to $a_B = \sqrt{\hbar/(m\omega_B)}$.

The coefficients

$$a_{n,n'}(k) = \sqrt{\frac{n'+1}{2}}\langle n'+1|e^{ia_Bku}\sin(\beta + \alpha u)|n\rangle - \sqrt{\frac{n}{2}}\langle n'|e^{ia_Bku}\sin(\beta + \alpha u)|n-1\rangle$$

with the help of the formula (see Ref. 29)

$$\langle n|e^{izu}|n+1\rangle\sqrt{\frac{n!}{(n+l)!}}(i\,\mathrm{sgn}\,z)^l\left(\frac{z^2}{2}\right)^{\frac{l}{2}}e^{-\frac{z^2}{4}}L_n^l\left(\frac{z^2}{2}\right)$$

and recurrence relations for the generalized Laguerre polynomials $L_n^l(z)$, can be written as

$$a_{n,n+l}(k)=\frac{i^l}{2\sqrt{2}}\sqrt{\frac{n!}{(n+l)!}}\times\left\{e^{i\beta}z_1^{\frac{l+1}{2}}e^{-\frac{z_1}{2}}L_n^l(z_1)-e^{-i\beta}\mathrm{sgn}^{l+1}(q-\alpha)z_2^{\frac{l+1}{2}}e^{-\frac{z_2}{2}}L_n^l(z_2)\right\}. \quad (14)$$

Here, $|n\rangle$ are the normalized Hermite functions of the dimensionless coordinate u, $z_1=(q+\alpha)^2/2$, $z_2=(q-\alpha)^2/2$, $q=a_Bk$, $\alpha=(a_B/r_0)\tan\theta$, $r_0=p_0/m\omega_B$. At large values of $n\gg 1$, the asymptotic representation of the Laguerre polynomials[30] makes it possible to express Eq. (14) using the Bessel functions J_l:

$$a_{n,n+l}(k)=\frac{i^l}{4}\{e^{i\beta}(q+\alpha)J_l\left[\sqrt{2n}(q+\alpha)\right]-e^{-i\beta}\mathrm{sgn}^{l+1}(q-\alpha)|q-\alpha|J_l\left(\sqrt{2n}|q-\alpha|\right)\}. \quad (15)$$

Equation (15) describes the oscillatory dependence of $\sigma_{zz}^{(so)}$ on the angle θ.

In the expression describing the electron energy spectrum

$$\varepsilon_{n,\sigma}=\hbar\omega_B\left(n+\frac{1}{2}\right)+\frac{\hbar\omega_s\sigma}{2}+\varepsilon_\perp, \quad (16)$$

where

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

we neglect the second-order terms of $\hat{V}_{so}$, i.e., those proportional to γ_R^2 and γ_D^2, since they lead energy level corrections, but do not affect the resonance intensity 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 transition, and the resonance width is determined by the inverse spin-flip time τ_s^{-1}.

According to Eq. (15), the conductivity (11) vanishes in a magnetic field perpendicular to the layers, i.e., at θ=0, if the spatial dispersion is neglected in the chosen models of the electron energy in the field of the crystal lattice (4), and spin-orbit coupling (5). The spatial dispersion of the kinetic coefficients in the presence of a high-frequency electromagnetic field $\omega\tau_l\simeq\omega_B\tau_l\gg 1$ can be characterized by the parameter $\kappa=(\eta\omega_p v_F/\omega_B c)^2$. If the electric field is polarized in the plane of the layers, then η=1. In the case of a normal skin effect, the depth of the skin layer $\delta\simeq k^{-1}\simeq r_B/\sqrt{\kappa}\gg r_B$ is much greater than the cyclotron radius $r_B=v_F/\omega_B$, and in order to implement the conditions of the extremely anomalous skin effect, the inequality $\delta\simeq r_B/\kappa^{1/3}\ll r_B$ must be satisfied. Although for the considered geometry of the problem, when the electric current flows perpendicular to the layers, usually, because of the smallness of the anisotropy parameter η, $\delta>r_B$, the spatial dispersion effect can be significant in structures with high enough conductivity, such as in the organic metal (BEDT-TTF)$_2$IBr$_2$ with a charge carrier density of $n_0\approx 10^{21}$ cm^{-3}.[31]

If the number of Landau levels below ε_F is large, then the conductivity (11) undergoes de Haas–van Alphen oscillations. Using the Poisson formula, we can write Eq. (11) as the sum

$$\sigma_{zz}^{(so)}(\omega,k)=\bar{\sigma}_{zz}(\omega,k)+\tilde{\sigma}_{zz}(\omega,k), \quad (17)$$

of the smooth

$$\bar{\sigma}_{zz}(\omega,k)=i\eta^2\frac{\omega_p^2}{4\pi}w(\theta)\left\{h_0(\omega)U_0+\sum_{l=1}^{\infty}[h_l^{(+)}(\omega)+h_l^{(-)}(\omega)]U_l\right\} \quad (18)$$

and oscillatory with respect to B^{-1} components

$$\bar{\sigma}_{zz}(\omega,k)=i\eta^2\frac{\omega_p^2\omega_B}{2\pi}w(\theta)\times\left\{\frac{h_0(\omega)}{\omega_s}Q_0+\sum_{l=1}^{\infty}\left[\frac{h_l^{(+)}(\omega)}{\Omega_l^{(+)}}-\frac{h_l^{(-)}(\omega)}{\Omega_l^{(-)}}\right]Q_l\right\}. \quad (19)$$

Here

$$Q_l=\sum_{j=1}^{\infty}\frac{(-1)^j}{j}\{U_lJ_0(j\Delta)+V_lJ_2(j\Delta)\}\Psi(j\lambda)\sin\frac{\pi j\omega_s}{\omega_B}\cos\frac{2\pi j\mu}{\omega_B}, \quad (20)$$

$$\Psi(\lambda)=\lambda/\mathrm{sh}\,\lambda,\ \lambda=2\pi^2T/(\hbar\omega_B),$$
$$\Delta=2\pi\varepsilon_\eta J_0(\alpha_1)/(\hbar\omega_B),$$
$$U_1=\frac{1}{2}\{(q+\alpha)^2J_l^2(q_1+\alpha_1)+(q-\alpha)^2J_l^2(q_1-\alpha_1)\},$$
$$V_1=\mathrm{sgn}^{l+1}(q-\alpha)|q^2-\alpha^2|J_l(q_1+\alpha_1)J_l(|q_1-\alpha_1|),$$
$$q_1=\sqrt{2\mu/(\hbar\omega_B)}q=kr_B,$$
$$\alpha_1=\alpha\sqrt{2\mu/(\hbar\omega_B)}=(mv_F/p_0)\tan\theta,$$

T is the temperature, μ is the chemical potential. The amplitudes of the oscillating harmonics in the sum Q_l are modulated by the quasiperiodic functions $J_0(j\Delta)$, $J_2(j\Delta)$, the argument of which depends not only on B^{-1}, but also on $\tan\theta$. If spatial dispersion is neglected, the conductivity $\sigma_{zz}^{(so)}(\omega,0)$ is determined by Eqs. (17)–(19), in which the coefficients U_l and V_l are equal

$$U_l=\alpha^2J_l^2(\alpha_1),\quad V_l=(-1)^{l+1}\alpha^2J_l^2(\alpha_1).$$

Under the condition $\varepsilon_\perp\simeq\eta\varepsilon_F\gg\hbar\omega_B$, Eq. (19) describes the oscillations of interlayer conductivity with changes in the inverse magnetic field. If $\eta\varepsilon_F \simeq \hbar\omega_B$ and $T< \hbar\omega_B/(2\pi^2)$, it is necessary to account for the oscillatory dependence of the chemical potential μ on B^{-1}, and near the resonance $\omega\approx|\Omega_l^{(\pm)}|$, and the relaxation time τ_l as a function of B^{-1}. In the region of sufficiently low temperatures $T< \hbar\omega_B/(2\pi^2)$, at $\eta\varepsilon_F \gg \hbar\omega_B$ and at arbitrary values of θ, the amplitude of the conductivity component that oscillates with the inverse of the magnetic field $\tilde{\sigma}_{zz}$ is $\sqrt{\Delta} \sim \sqrt{\eta\varepsilon_F/(\hbar\omega_B)}$ times smaller than the smooth component $\bar{\sigma}_{zz}^{(so)}$. However, the oscillatory component for the values $\theta= \theta_i$, at which $\alpha_1\sim\tan\,\theta_i$ is a root of the Bessel function $J_0(\alpha_1) = 0$, increases to about $\bar{\sigma}_{zz}^{(so)}$.

The resonance conductivity given by Eqs. (11) and (13) can be represented as the sum $\sigma_{zz}^{(so)} = \sigma_R + \sigma_D$ of the Rashba and Dresselhaus interactions. Having written the cyclotron frequency as $\omega_B= \omega_0 \cos\theta$, where $\omega_0\equiv|e|B/(mc)$ in Eq. (13), it is easy to see that $\sigma_R \simeq \sigma_D$ at $\theta\simeq 1$, but in the range of θ values in which the angular oscillations $\tan\theta\gg 1$ appear, the asymptotic behavior of $\sigma_{zz}^{(so)}$ is determined by the Dresselhaus spin-orbit interaction $\sigma_R \simeq \sigma_D \cos^2\theta \ll \sigma_D$. The angular dependence of the smooth and oscillating components of the interlayer conductivity at characteristic parameter values, is shown in Figs. 2 and 3, respectively.

Let us assume that the conductor occupies the half space $y>0$, and that the condition $\eta\omega_p\gg\omega$ is fulfilled, at which the displacement current can be neglected in Maxwell's equations. An important characteristic of the conducting medium is the surface impedance tensor

$$Z_{ik} = -i\frac{8\omega}{c^2}\int_0^\infty dk[D^{-1}(\omega,k)]_{ik},\;\; i,k=\{x,z\}, \tag{21}$$

which binds the tangential components of the electric field on the surface to the total current, where $D_{ik}=\{k^2\delta_{ik} - (4\pi i\omega c^{-2})(\sigma_{ik} - \frac{\sigma_{ik}\sigma_{iy}}{\sigma_{yy}})\}$. In layered conductors the inequalities $|\sigma_{xz}\sigma_{zx}| \simeq |\sigma_{yz}\sigma_{zy}| \ll |\sigma_{zz}\sigma_{xx}|$, $|\sigma_{zz}\sigma_{yy}|$ are fulfilled,[25] and therefore, the conductivity tensor in the xz,yz planes can be considered diagonal. Let us expand the integrand for the impedance component Z_{zz} into a series in powers of $\sigma_{zz}^{(so)}$, and keep the first two terms:

$$Z_{zz} = Z_{zz}^{(0)} + \Delta Z_{zz}^{(so)} = -i\frac{8\omega}{c^2}\int_0^\infty \frac{dk}{k^2 - 4\pi i\omega c^{-2}\sigma_{zz}(\omega,k)} + \frac{32\pi\omega^2}{c^4}\int_0^\infty \frac{dk\,\sigma_{zz}^{(so)}(\omega,k)}{[k^2 - 4\pi i\omega c^{-2}\sigma_{zz}(\omega,k)]^2}. \tag{22}$$

Here, the first term is the impedance at $\sigma_{zz}^{(so)} = 0$, and the second is the correction due to spin-orbit interaction. The interlayer conductivity in the main approximation $\sigma_{zz}(\omega,k)$ can be easily found with the help of Eq. (1):

$$\sigma_{zz}(\omega,k) = i\frac{\eta^2\omega_p^2}{8\pi^2}\sum_{n,\sigma}\int_{-\pi}^{\pi} d\beta\Bigg\{-\frac{\partial f_0(\varepsilon_{n,\sigma})}{\partial n}\frac{|A_{n,n}(q)|^2}{\omega + i\tau^{-1}} + 2\sum_{l=1}^{\infty}\frac{f_0(\varepsilon_{n,\sigma}) - f_0(\varepsilon_{n+l,\sigma})}{l[(\omega + i\tau_l^{-1})^2 - (\omega_B l)]^2}(\omega + i\tau_l^{-1})|A_{n,n+l}(q)|^2\Bigg\}, \tag{23}$$

where

$$A_{n,n+l}(q) = \langle n|e^{iqu}\sin(\beta + \alpha u)|n+l\rangle\frac{i^{l-1}}{2}\sqrt{\frac{n!}{(n+l)!}} \times\left\{e^{i\beta}z_1^{\frac{l}{2}}e^{-\frac{z_1}{2}}L_n^l(z_1) - e^{-i\beta}\mathrm{sgn}^l(q-\alpha)z_2^{\frac{l}{2}}e^{-\frac{z_2}{2}}L_n^l(z_2)\right\},$$

and the value τ determines the width $\hbar/\tau$ of the n-th energy level of the electron. In the case of $\varepsilon_F \gg \hbar\omega_B$ the summation over n can

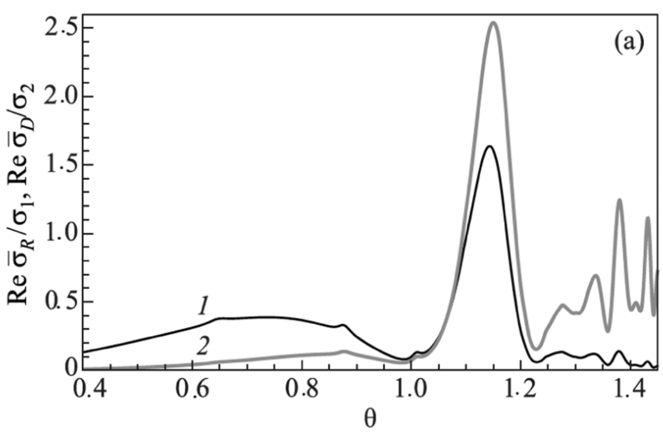

FIG. 2. The dependences of Re$\bar{\sigma}_R/\sigma_1$ (curve 1) and Re$\bar{\sigma}_D/\sigma_2$ (curve 2) on θ, where $\sigma_{1,2} = \eta^2(\omega_p^2/4\pi\omega_0)(\gamma_{R,D}/v_F)^2$, at parameters $mv_F/p_0 = 2.5$, $\omega = 1.8\omega_0$, $\omega_s = 1.4\omega_0$, $\omega\tau_l = 20$, $\omega\tau_s = 200$, $k = 0$ (a) and $kv_F/\omega_0 = 0.2$ (b). The abrupt maximum at $\theta \approx 1.16$ corresponds to resonance at the frequency $\omega = \omega_s + \omega_B$, here and in other figures.

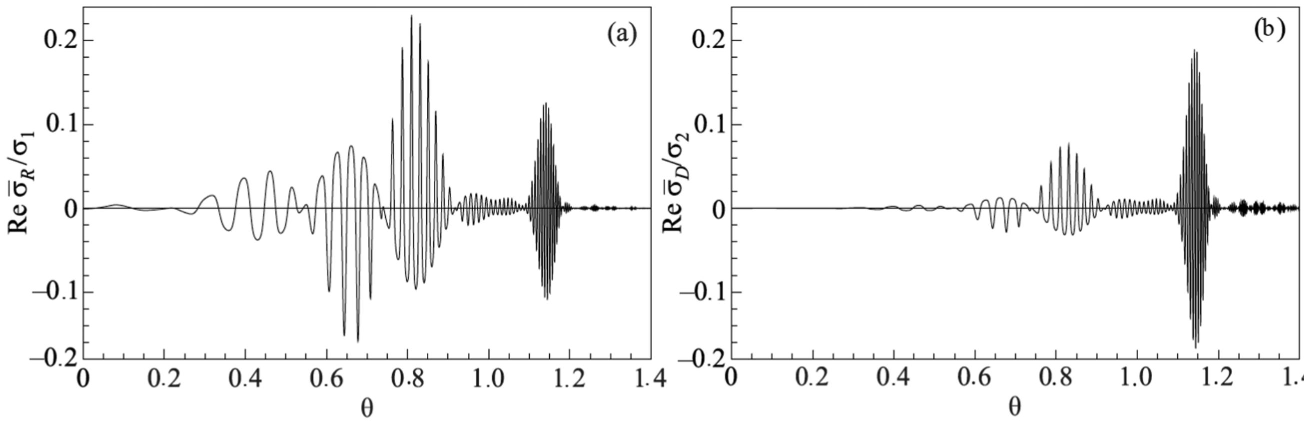

FIG. 3. Qualitative dependences of Re $\tilde{\sigma}_R/\sigma_1$ (a) and Re $\tilde{\sigma}_D/\sigma_2$ (b) on θ, where $\sigma_{1,2} = \eta^2(\omega_p^2/4\pi\omega_0)(\gamma_{R,D}/v_F)^2$, at parameters $mv_F/p_0 = 2.5$, $\omega = 1.8\omega_0$, $\omega_s = 1.4\omega_0$, $2\pi^2 T/(\hbar\omega_0) = 1$, $\omega\tau_l = 10$, $\omega\tau_s = 100$, $\eta\varepsilon_F/(\hbar\omega_0) = 1$, $\eta = 1/30$, and $kv_F/\omega_0 = 0.2$. Slow angular oscillations are caused by the functions $J_0(j\Delta)$, $J_2(j\Delta)$ in Eq. (20).

be replaced by integration, and Eq. (23) can be transformed into

$$\sigma_{zz}(\omega, k) = i\frac{\eta^2\omega_p^2}{8\pi}\left\{\frac{J_0^2(q_1+\alpha_1)+J_0^2(q_1-\alpha_1)}{\omega+i\tau^{-1}} + 2\sum_{l=1}^{\infty}\frac{J_l^2(q_1+\alpha_1)+J_l^2(q_1-\alpha_1)}{(\omega+i\tau_l^{-1})^2-(\omega_B l)^2}(\omega+i\tau_l^{-1})\right\}. \quad (24)$$

By changing the integration variable $k = \lambda(\kappa\omega/\omega_{B)}^{1/2}/r_B$, we write the resonance correction to the impedance as follows:

$$\Delta Z_{zz}^{(so)} = \frac{8\upsilon_F}{c^2}\sqrt{\frac{\omega}{\kappa\omega_B}}\int_0^\infty \frac{d\lambda s_{zz}^{(so)}[\omega, (\kappa\omega/\omega_B)^{1/2}\lambda]}{\{\lambda^2 - is_{zz}[\omega, (\kappa\omega/\omega_B)^{1/2}\lambda]\}}, \quad (25)$$

where s_{zz} is the conductivity attributed to $\eta^2\omega_p^2/(4\pi\omega_B)$, i.e., $s_{zz} = 4\pi\omega_B\sigma_{zz}/(\eta^2\omega_p^2)$, $s_{zz}^{(so)} = 4\pi\omega_B\sigma_{zz}^{(so)}/(\eta^2\omega_p^2)$. If $\kappa \ll 1$, then $\Delta Z_{zz}^{(so)}$ can be expanded into a series in powers of $\sqrt{\kappa}$, and the first term of the series is

$$\Delta Z_{zz}^{(so)} = \frac{2\pi v_F}{c^2}\sqrt{\frac{\omega}{\kappa\omega_B}}\frac{s_{zz}^{(so)}(\omega, 0)}{[-is_{zz}(\omega, 0)]^{3/2}}. \quad (26)$$

In the vicinity of $\omega = \omega_r + \Delta\omega = |l\omega_B \pm \omega_s| + \Delta\omega$, the asymptotic behavior of Eq. (26) looks like

$$\Delta Z_{zz}^{(so)} \approx \frac{\pi v_F}{c^2}\sqrt{\frac{\omega_r}{\kappa\omega_B}}\frac{\tau_{ls}^{-1}+i\Delta\omega}{\Delta\omega^2+\tau_{ls}^{-2}}\alpha_1^2 J_l^2(\alpha_1) \times \frac{\gamma_R^2(\omega_B+\omega_s cos\theta)^2+\gamma_D^2(\omega_s-\omega_B cos\theta)^2}{r_B^2(\omega_B^2-\omega_s^2)^2[-is_{zz}(\omega_r, 0)]^{3/2}}. \quad (27)$$

In the region of $0 < \kappa \leq 1$ the integral in Eq. (25) as a function of κ does not undergo significant changes, therefore, with an increase in κ from $\kappa \ll 1$ to $\kappa \simeq 1$, the correction to the impedance due to spin-orbit interaction decreases according to $1/\sqrt{\kappa}$.

Equations (11), (13), and (26) can be used to experimentally find the absolute values of the constants γ_R and γ_D.[26] Pure spin transitions with $l = 0$ at the frequency ω_s are of particular interest, since $\tau_s > \tau_l$. Additionally, $\omega_B\tau_l$ decreases with increasing θ, as does the intensity of the l-th resonance.

4. COMBINED RESONANCE OF SPIN MAGNETIZATION

Together with spin transitions, non-uniform high-frequency magnetic field $\mathbf{B}^{\sim}(\mathbf{r},t)$ can also excite transitions at combined frequencies, even without taking spin-orbit coupling into account. The current density given by Eq. (1) is proportional to the eddy electric field and, correspondingly, includes terms that are proportional to the magnetic field. The orbital component of operator (2) in Eq. (1) determines the conductivity, and the spin component $\hat{\mathbf{j}}^{(s)}=c$ rot $\hat{\boldsymbol{\mu}}_0\delta(\mathbf{r}-\mathbf{r}')$ determines the high-frequency magnetization and paramagnetic susceptibility.[32]

Setting $\mathbf{k} = (0,k,0)$ from Maxwell's equations

$$\mathbf{j}^{(m)} = c\,\mathrm{rot}\,\mathbf{M}, \quad \mathrm{rot}\,\mathbf{E} = -\frac{1}{c}\frac{\partial\mathbf{B}}{\partial t} \quad (28)$$

we obtain the following expression for the magnetization due to the spin component of the current density operator (2).

$$M_i(\omega, k) = -\frac{\hbar\omega\omega_B}{4\pi}\chi_0\sum_{n,n',\sigma,\sigma'}\int_{-\pi}^{\pi}d\beta \times \frac{f_0(\varepsilon_{n,\sigma})-f_0(\varepsilon_{n',\sigma'})}{\varepsilon_{n,\sigma}-\varepsilon_{n',\sigma'}}\frac{|\langle n'|e^{iqu}|n\rangle|^2\langle\sigma|\hat{\sigma}_k|\sigma'\rangle\langle\sigma'|\hat{\sigma}_i|\sigma\rangle}{(\varepsilon_{n,\sigma}-\varepsilon_{n',\sigma'})/\hbar-\omega-i\tau_{|n-n'|,s}^{-1}} \times B_k^{\sim}(\omega, k). \quad (29)$$

Here, $\mathbf{j}^{(m)}$ is determined by Eq. (1), in which $\mathbf{j} = \mathbf{j}^{(s)}$,

$\chi_0 = \mu_0^2 m p_0/(\pi\hbar^3)$ is the static paramagnetic susceptibility, $\mathbf{B}^{\sim}(\omega,k) = (c/\omega)[\mathbf{k}\times\mathbf{E}(\omega,k)]$, $i,k=\{\xi,\ \zeta\}$ the matrix elements $\langle\sigma'|\hat{\sigma}_k|\sigma\rangle$ are expressed in terms of Kronecker symbols

$$\langle\sigma'|\hat{\sigma}_\xi|\sigma\rangle = \delta_{\sigma,1}\delta_{\sigma',-1} + \delta_{\sigma,-1}\delta_{\sigma',1},\ \langle\sigma'|\hat{\sigma}_\zeta|\sigma\rangle = \delta_{\sigma,\sigma'}.$$

The magnetization component, which is perpendicular to the magnetic field, is caused by transitions with a change in the spin projection, and describes the paramagnetic and combined resonances, is equal to

$$\begin{aligned} M_\xi(\omega,k) &= \chi_{\xi\xi}(\omega,k)B_\xi^{\sim}(\omega,k) \\ &= -\frac{\omega\omega_B}{2\pi}\chi_0\sum_n\int_{-\pi}^{\pi}d\beta\left\{\frac{f_0(\varepsilon_{n,1})-f_0(\varepsilon_{n,-1})}{\omega_s}|b_{n,n}(q)|^2h_0(\omega)\right. \\ &\times\sum_{l=1}^{\infty}\left(\frac{f_0(\varepsilon_{n+l,1})-f_0(\varepsilon_{n,-1})}{\omega_B l+\omega_s}h_l^{(+)}(\omega)\right. \\ &\left.\left.+\frac{f_0(\varepsilon_{n+l,-1})-f_0(\varepsilon_{n,1})}{\omega_B l-\omega_s}h_l^{(-)}(\omega)\right)|b_{n,n+l}(q)|^2\ B_\xi^{\sim}(\omega,k)\right\}, \end{aligned} \tag{30}$$

where

$$b_{n,n+l}(q) = \sqrt{\frac{n!}{(n+l)!}}\left(\frac{q^2}{2}\right)^{\frac{l}{2}}e^{\frac{q^2}{4}}\mathrm{L}_4^l\left(\frac{q^2}{2}\right)\underset{n\gg1}{\longrightarrow}\mathrm{J}_l(\sqrt{2nq}).$$

At $\varepsilon_F \gg \hbar\omega_B$, after standard transformations, the paramagnetic susceptibility can be written as the sum

$$\chi_{\xi\xi}(\omega,k) = \bar{\chi}_{\xi\xi}(\omega,k) + \tilde{\chi}_{\xi\xi}(\omega,k) \tag{31}$$

of the smooth

$$\bar{\chi}_{\xi\xi}(\omega,k) = -\omega\chi_0\left\{\mathrm{J}_0^2(q_1)h_0(\omega) + \sum_{l=1}^{\infty}\mathrm{J}_l^2(q_1)\left[h_l^{(+)}(\omega)+h_l^{(-)}(\omega)\right]\right\} \tag{32}$$

and quantum oscillatory components

$$\begin{aligned} \tilde{\chi}_{\xi\xi}(\omega,k) &= -2\omega\omega_B\chi_0 \\ &\times\left\{\mathrm{J}_0^2(q_1)\frac{h_0(\omega)}{\omega_s} + \sum_{l=1}^{\infty}\mathrm{J}_l^2(q_1)\left[\frac{h_l^{(+)}(\omega)}{\Omega_l^{(+)}} - \frac{h_l^{(-)}(\omega)}{\Omega_l^{(-)}}\right]\right\} \\ &\times\sum_{j=1}^{\infty}\frac{(-1)^j}{j}\mathrm{J}_0(j\Delta)\Psi(j\lambda)\sin\frac{\pi j\omega_s}{\omega_B}\cos\frac{2\pi\mu j}{\hbar\omega_B}. \end{aligned} \tag{33}$$

Equations (30)–(33) are the sum of the magnetic susceptibility's asymptotic behavior near the $\pm l$-th resonance $\omega \approx |\Omega_l^{(\pm)}|$. Upon neglecting the spatial dispersion, i.e., when $q_1=0$, only one term at $l=0$ remains in Eqs. (32) and (33), which describes the electron paramagnetic resonance. The magnetic susceptibility corresponding to the combined resonance is maximized at $q_1 = kr_B \simeq 1$, and at large values $q_1 \gg 1$ it decreases as $1/kr_B$. The angular dependences of the smooth and oscillatory magnetic susceptibility components are shown in Fig. 4. The angular dependence of magnetic susceptibility component $\tilde{\chi}_{\xi\xi}$ that oscillates with changes in B^{-1} is similar to the dependences of the oscillating component of the resonance conductivity considered above and the static magnetic susceptibility in a strong magnetic field.[33] In the region of sufficiently low temperatures, at $\eta\varepsilon_F \gg \hbar\omega_B$ and at arbitrary values of θ, the oscillation amplitude $\tilde{\chi}_{\xi\xi} \simeq J_0(\Delta)\bar{\chi}_{\xi\xi} \simeq \bar{\chi}_{\xi\xi}/\sqrt{\Delta}$ is $\sqrt{\Delta} \simeq \sqrt{\eta\varepsilon_F/(\hbar\omega_B)}$ times less than that of the smooth component $\bar{\chi}_{\xi\xi}$. However, the oscillatory component for $\theta=\theta_i$, at which $\alpha_1\sim\tan\theta_i$ is a root of the Bessel function $\mathrm{J}_0(\alpha_1)=0$, it is equal to $\bar{\chi}_{\xi\xi}$ by order of magnitude. For these directions of the magnetic field, the dependence of the cross section area $S_F(p_B)$ on the momentum projection p_B appears in the second-order terms of η.

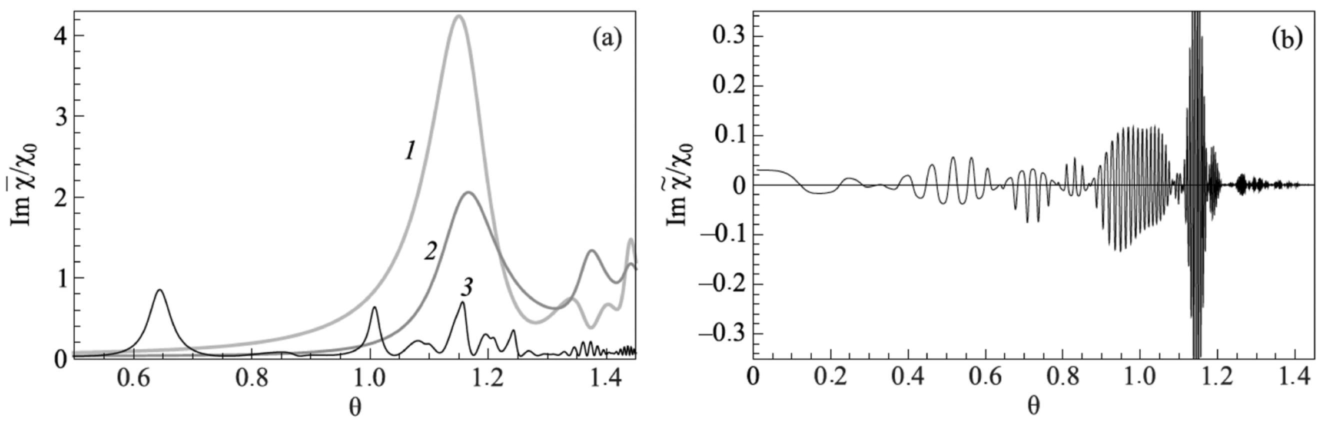

FIG. 4. (a) The dependence of Im $\bar{\chi}/\chi_0$ on θ at $kv_F/\omega_0 = 1$ (1), 0.3 (2), 10 (3). (b) The qualitative dependence of Im $\tilde{\chi}/\chi_0$ on θ $kv_F/\omega_0 = 1$. The slow angular oscillations are caused by the function $\mathrm{J}_0(j\Delta)$ in Eq. (33). The parameter values are: $mv_F/p_0 = 2.5$, $\omega = 1.8\omega_0$, $\omega_s = 1.4\omega_0$, $\omega\tau_l = 10$, $\omega\tau_s = 100$, $2\pi^2T/(\hbar\omega_0) = 1$, $\eta\varepsilon_F/(\hbar\omega_0) = 1$, $\eta = 1/30$.

The correction to the impedance $\Delta Z_{zz}^{(s)}$, under the resonance conditions of spin magnetization, can be found using Eq. (22) for $\Delta Z_{zz}^{(so)}$, into which, instead of $\sigma_{zz}^{(so)}(\omega, k)$, one should substitute the correction to conductivity due to the spin current $\mathbf{j}^{(m)}$:

$$\sigma_{zz}^{(s)}(\omega, k) = -i\frac{k^2c^2}{\omega}\chi_{\xi\xi}(\omega, k)\cos^2\theta. \tag{34}$$

After simple transformations, we get

$$\Delta Z_{zz}^{(s)} = -i\frac{32\pi v_F}{c^2}\sqrt{\frac{\omega}{\kappa\omega_B}}\cos^2\theta \times \int_0^\infty \frac{d\lambda\lambda^2\chi_{\xi\xi}[\omega, (\kappa, \omega/\omega_B)^{1/2}\lambda]}{(\lambda^2 - is_{zz}[\omega, (\kappa\omega/\omega_B)^{1/2}\lambda])^2}. \tag{35}$$

If the incident wave is polarized along the normal to the layers, and normal skin effect conditions are possible, then the principal term of the expansion $\Delta Z_{zz}^{(s)}$ in a series of $\sqrt{\kappa}$ powers

$$\Delta Z_{zz}^{(s)} = -i\frac{8\pi^2 v_F}{c^2}\sqrt{\frac{\omega}{\kappa\omega_B}}\frac{\chi_{\xi\xi}(\omega, 0)}{\sqrt{-is_{zz}(\omega, 0)}}\cos^2\theta \tag{36}$$

describes the electron paramagnetic resonance.

By analogy with Eq. (22), the impedance component Z_{xx} can be written as:

$$Z_{xx} = Z_{xx}^{(0)} + \Delta Z_{xx}^{(s)} = -i\frac{8\omega}{c^2}\int_0^\infty \frac{dk}{k^2 - 4\pi i\omega c^{-2}\hat{\sigma}_{xx}(\omega, k)} - i\frac{32\pi\omega}{c^2}\sin^2\theta\int_0^\infty \frac{dk\, k^2\chi_{\xi\xi}(\omega, k)}{[k^2 - 4\pi i\omega c^{-2}\hat{\sigma}_{xx}(\omega, k)]^2}, \tag{37}$$

where $\hat{\sigma}_{xx} = \sigma_{xx} - \sigma_{xy}\sigma_{yx}/\sigma_{yy}$. In the $\varepsilon_F \gg \hbar\omega_B$ approximation, the components σ_{ij}, $i, j = \{x, y\}$ look like:[34]

$$\begin{aligned}
\sigma_{xx}(\omega, k) &= \frac{i\omega_p^2}{4\pi}\left\{\frac{2J_1(q_1)}{\omega + i\tau^{-1}} + \sum_{l=1}^{\infty}\frac{J_{l-1}^2(q_1) + J_{l+1}^2(q_1) - 2J_{l-1}(q_1)J_{l+1}(q_1)}{(\omega + i\tau_l^{-1})^2 - (\omega_B l)^2}\right\},\\
\sigma_{yy}(\omega, k) &= \frac{i\omega_p^2}{4\pi}\sum_{l=1}^{\infty}\frac{J_{l-1}^2(q_1) + J_{l+1}^2(q_1) - 2J_{l-1}(q_1)J_{l+1}(q_1)}{(\omega + i\tau_l^{-1})^2 - (\omega_B l)^2}(\omega + i\tau_l^{-1}),\\
\sigma_{xy}(\omega, k) &= \frac{\omega_p^2}{4\pi}\sum_{l=1}^{\infty}\frac{J_{l-1}^2(q_1) - J_{l+1}^2(q_1)}{(\omega + i\tau_l^{-1})^2 - (\omega_B l)^2}l\omega_B.
\end{aligned} \tag{38}$$

The impedance correction $\Delta Z_{xx}^{(s)}$ is maximized at $\kappa_1 = (\omega_p v_F/\omega_B c)^2 \simeq 1$, and in the limiting case $\kappa_1 \gg 1$ it decreases as $\kappa_1^{-2/3}$.

In order for combined and cyclotron resonance to manifest, it is necessary for the electron to make several revolutions around its orbit in a magnetic field, over the course of its mean free path. In ordinary metals, achieving this type of mean free path length for charge carriers corresponds to the conditions of the extremely anomalous skin effect, and therefore the magnetic susceptibility (31) is small with respect to the parameter $\delta/r_B \ll 1$. In organic conductors, the fulfillment of the $\omega\tau_l \simeq \omega_B\tau_l \gg 1$ ratio, under the condition that $\delta/r_B \simeq 1$, is much more favorable. This is one more reason why it is possible to experimentally observe the combined resonance of magnetization in layered conductors, in addition to the fact that almost all FS electrons participate in resonance formation, as mentioned in the Introduction.

5. CONCLUSION

The experimentally observed resonance absorption of microwave radiation in organic compounds is a superposition of peaks that correspond to different types of resonances. For example, at helium temperatures, organic metals $(BEDT\text{-}TTF)_2MHg(SCN)_4$ (M = K,Tl) with antiferromagnetic ordering, in a magnetic field perpendicular to the conducting plane, display narrower lines against the background of peaks that correspond to cyclotron resonance. The amplitudes of these lines are 5–10 times smaller, and are presumably caused by the electron paramagnetic and antiferromagnetic resonances.[17] The experimental data in layered conductors can be identified and assigned by analyzing the absorbed power and other conductor characteristics as a function of the direction and magnitude of the magnetic field, as well as the FS parameters. Under normal skin effect conditions, resonance corrections to impedance are proportional to $\sigma_{zz}^{(so)}$ and $\chi_{\xi\xi}$. In contrast to $\sigma_{zz}^{(so)}$, the magnetic susceptibility does not contain any oscillating factors in the form of $J_l^2\left[\left(\frac{mv_F}{p_0}\right)\tan\theta\right]$. The angular dependence of the smooth component $\bar{\chi}_{\xi\xi}$ is determined by the dependence of the resonance functions (12) $h_l^{(\pm)}$ on $\omega_B \sim \cos\theta$, and the angular dependence of the quantum $\tilde{\chi}_{\xi\xi}$. component that oscillates with changes in B^{-1}, as follows from (33), also includes the factors $J_0(j\Delta)$, which describe the slow oscillatory dependence on $\tan\theta$, and the absolute value of the magnetic field. The obtained results can be used to detect the combined resonance of interlayer conductivity and spin magnetization not only in organic conductors, but also in other low-dimensional layered structures of inorganic origin.

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