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Conduction-electron spin resonance in two-dimensional structures

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ABSTRACT

The influence of the conduction-electron spin magnetization density, induced in a two-dimensional electron layer by a microwave electromagnetic field, on the reflection and transmission of the field is considered. Because of the induced magnetization and electric current, both the electric and magnetic components of the field should have jumps on the layer. A way to match the waves on two sides of the layer, valid when the quasi-two-dimensional electron gas is in the one-mode state, is proposed. By following this way, the amplitudes of transmitted and reflected waves as well as the absorption coefficient are evaluated.

1. Introduction

Electron spin resonance has long been used to determine g factors and the longitudinal and transversal relaxation times T_1 and T_2 providing information about electron band structure and allowing us to investigate interactions responsible for spin-flip transitions [1]. This method has acquired an enhanced actuality nowadays because a growing interest in the spin dynamics in two-dimensional (2D) electron systems, which are potentially important for spintronics applications [2,3]. For a long time it was thought that the direct observation of the conduction-electron spin resonance (CESR) in 2D structures is impossible because of small number of current carriers. A breakthrough in this field is recent works [4–10] where the CESR in some 2D semiconductor structures was detected by means of the microwave absorption. The idealness and hence the conductivity of such structures can be high so that the field acting on electron spins can differ appreciably from that of incoming wave because of the field of the electric current excited by the wave. Despite the theory of the spin resonance excitation in bulk conductors is well elaborated (see, e.g., Ref. 11 and references therein), an analogous theory for 2D conductors, to the best of the author knowledge, is still lacking. The purpose of the present note is to fill in this gap.

2. Problem statement and results

A feature of this problem which impedes the immediate application of standard methods, consists in the following. Let the quasi-2D layer aligned along an $x-y$ plane is placed at position $z=0$ between two dielectrics with the permittivities ϵ_1 $(z<0)$ and ϵ_2 $(z>0)$, and z-axis points “upward” to the dielectric 2. Within the frame of classical electrodynamics, properties of a conducting medium enter the Maxwell equations through the material constitutive relations [12], which in the case under study have the form

$$\mathbf{J} = \hat{\sigma}\mathbf{E}, \quad \mathbf{M} = \hat{\chi}\mathbf{H}, \tag{1}$$

where σ and χ are tensors of the electric conductivity and the magnetic susceptibility, respectively. In the following all quantities are assumed to have the time dependence $e^{-i\omega t}$. The great difference between the width of the conducting layer d and the wavelength $\lambda = 2\pi/q_0$ $(q_0 = \omega/c)$ of microwave radiation urges one to treat the layer as strictly two-dimensional sheet so that $\mathbf{J}(\mathbf{r};t) = \delta(z)\mathbf{J}_S(x,y;t)$ and $\mathbf{M}(\mathbf{r};t) = \delta(z)\mathbf{M}_S(x,y;t)$. Then, from the Amper low

$$\nabla \times \mathbf{H}_\omega = -iq_0\epsilon\mathbf{E}_\omega - \frac{4\pi}{c}\mathbf{J}_\omega, \tag{2}$$

it follows the usual expression for the jump of the magnetic field on the sheet

$$\hat{\mathbf{n}} \times (\mathbf{H}_{\omega,2} - \mathbf{H}_{\omega,1}) = \frac{4\pi}{c}\mathbf{J}_{\omega,S} \tag{3}$$

where $\mathbf{n} = \hat{\mathbf{z}}$ and $\mathbf{H}_{\omega,1}$ and $\mathbf{H}_{\omega,2}$ are the values of the magnetic field on the lower and upper sides of the sheet, respectively. Accordingly, from the Faraday low

$$\nabla \times \mathbf{E}_\omega = iq_0(\mathbf{H}_\omega + 4\pi\mathbf{M}_\omega) \tag{4}$$

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it follows

$$\hat{\mathbf{n}} \times (\mathbf{E}_{\omega,2} - \mathbf{E}_{\omega,1}) = 4\pi i q_0 \mathbf{M}_{\omega,S\parallel}, \tag{5}$$

where $\mathbf{M}_{\omega,S\parallel}$ is the parallel component of the 2D magnetization density. Thus both the jumps of the electric and magnetic components of the electromagnetic field on the sheet should be taken into account. If one tries, as usually, to utilize Eqs. (3) and (5) as boundary conditions for matching the fields above and below the sheet, an ambiguity occurs – the jumps make undefined the values of $\mathbf{E}$ and $\mathbf{H}$ which should be used in Eq. (1). Thus, the inequality $d \ll \lambda$ does not allow one to consider the system as a strictly 2D sheet from the very beginning. Therefore, we will first consider d as small but finite quantity, trying to find an additional property of the 2D conductor, which could lift the ambiguity mentioned, and take the limit $d/\lambda \to 0$ on a later stage.

This additional property, in which the following consideration depends on, is the assumption that the electron gas is in the one-mode state, i.e., all electrons occupy only the ground state in the confinement potential forming the 2D gas. Such a situation is usual in semiconductor heterostructures and conducting surfaces and interfaces of oxide insulators. It will be shown below that at the normal incidence of the wave on the one-mode gas the 'averaged' fields $\mathbf{E}_{\omega,av} = \frac{1}{2}(\mathbf{E}_{\omega,1} + \mathbf{E}_{\omega,2})$ and $\mathbf{H}_{\omega,av} = \frac{1}{2}(\mathbf{H}_{\omega,1} + \mathbf{H}_{\omega,2})$, where $\mathbf{E}_{\omega,1,2}$ ($\mathbf{H}_{\omega,1,2}$) are the limit values of the electric (magnetic) field on the lower and upper sides of the layer, respectively, should be substituted into the right-hand sides of Eqs. (3) and (5). Thus, Eqs. (3) and (5) should take the form

$$\hat{\mathbf{n}} \times (\mathbf{H}_{\omega,2} - \mathbf{H}_{\omega,1}) = \frac{4\pi}{c}\hat{\sigma}\left(\frac{\mathbf{E}_{\omega,1} + \mathbf{E}_{\omega,2}}{2}\right), \tag{6}$$

$$\hat{\mathbf{n}} \times (\mathbf{E}_{\omega,2} - \mathbf{E}_{\omega,1}) = 4\pi i q_0 \hat{\chi}\left(\frac{\mathbf{H}_{\omega,1} + \mathbf{H}_{\omega,2}}{2}\right). \tag{7}$$

The standard method supplemented with this matching conditions becomes well defined and straightforwardly gives rise to the following results. The amplitudes of reflection T_{re} and transmission T_{tr} have the form

$$T_{re} = \frac{N_{re}}{D}, \quad T_{tr} = \frac{N_{tr}}{D},$$

$$N_{re} = (n_2 - n_1 + \frac{4\pi\sigma_\omega}{c}) + 2\pi i q_0 \chi_\omega [2n_1 n_2 + \frac{2\pi\sigma_\omega}{c}(n_1 - n_2)],$$

$$N_{tr} = 2n_2 + 2\pi i q_0 \chi_\omega \frac{4\pi\sigma_\omega}{c} n_2$$

$$D = (n_2 + n_1 + \frac{4\pi\sigma_\omega}{c}) - 2\pi i q_0 \chi_\omega [2n_1 n_2 + \frac{2\pi\sigma_\omega}{c}(n_1 + n_2)], \tag{8}$$

while the absorption coefficient, with the accuracy up to terms linear in χ, is

$$A = \frac{N}{Z},$$

$$N = \frac{4\pi\sigma'_\omega}{c} n_1 + 4\pi q_0 n_1 n_2^2 \chi''_\omega \left[1 + \frac{16\pi\sigma'_\omega}{cn_2} + \left(\frac{2\pi\sigma'_\omega}{cn_2}\right)^2 + \left(\frac{2\pi\sigma''_\omega}{cn_2}\right)^2\right],$$

$$Z = \left(n_1 + n_2 + \frac{2\pi\sigma'_\omega}{c}\right)^2 + \left(\frac{2\pi\sigma''_\omega}{c}\right)^2. \tag{9}$$

These equations have been written for the fields with the circular polarization $\mathbf{e}_+ = 1/\sqrt{2}(\mathbf{e}_x + i\mathbf{e}_y)$ when $\mathbf{E} = E_{(-)}\mathbf{e}_+$, $\mathbf{H} = H_{(-)}\mathbf{e}_+$, $\mathbf{M} = M_{(-)}\mathbf{e}_+$, $\mathbf{J} = J_{(-)}\mathbf{e}_+$ and the constitutive relations have the form $M_{(-)} = \chi_\omega^{(+)} H_{(-)av}$ with $\chi_\omega^{(+)} = \chi_{xx}(\omega) + i\chi_{xy}(\omega)$ and $J_{(-)} = \sigma_\omega^{(+)} E_{(-)av}$ with $\sigma_\omega^{(+)} = \sigma_{xx}(\omega) + i\sigma_{xy}(\omega)$. Also the following notations have been used: $n_{1,2} = \sqrt{\epsilon_{1,2}}$ is the refraction index, $\chi''_\omega = \Im\chi_\omega^{(+)}$, $\sigma_\omega = \sigma_\omega^{(+)}$, $\sigma'_\omega = \Re\sigma_\omega^{(+)}$, and $\sigma''_\omega = \Im\sigma_\omega^{(+)}$. Near the frequency ω_{res} of the CESR one gets [13] $\chi_\omega^{(+)} \cong \chi_0(\pi/m) - N(\epsilon_F)\omega/(\omega - \omega_{res} + i/T_2)$, where $\chi_0 = (g\mu_B/2)^2 m/\pi$ is the static susceptibility of 2D degenerate electron gas and $N(\epsilon_F)$ is the density of states for a single spin. Eqs. (8) and (9) show that at $\sigma/c \geq 1$ the effect of the electric current, induced by microwave field, on effective magnetic field acting on electron spins can be appreciable. The derivation of Eqs. (8) and (9) is quite standard and therefore is not given here. The remaining part of the paper presents the proof of the above matching conditions.

3. Matching conditions

So we consider the electron gas which occupies the layer $-d/2 \leq z \leq d/2$. Two facts follow from the assumption about the one-mode state of the gas (see Appendix). The first is that the coordinate dependence of the 3D density of the current and the magnetization has the factorized form

$$\mathbf{J}(\mathbf{r}, t) = \rho(z)\mathbf{J}_S(\mathbf{r}_\parallel, t), \quad \mathbf{M}(\mathbf{r}, t) = \rho(z)\mathbf{M}_S(\mathbf{r}_\parallel, t), \tag{10}$$

where $\mathbf{r} = (x, y, z) = (\mathbf{r}_\parallel, z)$, $\rho(z) = |\psi_0(z)|^2$, $\psi_0(z)$ is the wave-function of the ground state, and $\mathbf{J}_S$ and $\mathbf{M}_S$ are the 2D densities. At the normal incidence of the radiation, $\mathbf{J}_S$ and $\mathbf{M}_S$ loose their coordinate dependence. The second fact is that the constitutive relations (1) take the form

$$J_S^i(\omega) = \sigma_\omega^{ij} \int_z \rho(z) E^j(z, \omega),$$
$$M_S^i(\omega) = \chi_\omega^{ij} \int_z \rho(z) H^j(z, \omega), \tag{11}$$

where $\int_z = \int dz$.

Consider first the question about the value of the electric field which should be used in Ohm's law in the limit $d/\lambda \to 0$. As it is known [and also seen from Eq. (5)], the major reason for a finite difference between the electric field on the upper and lower surfaces of the layer is the magnetization. To make the following explanation more clear the effect of the external magnetic field is omitted for a while. Consider a strictly 2D sheet, uniformly filled with the spin magnetization $\rho(\zeta)\mathbf{m}^{-i\omega t}$, $\mathbf{m} \perp \mathbf{e}_z$ [$\mathbf{m}$ does not depend on $\mathbf{r}_\parallel$ at the normal incidence], which lies inside the layer at $z = \zeta$, $|\zeta| \leq d/2$. By utilizing the fact that the vector-potential created at the point $\mathbf{r}$ by the magnetic dipole $\boldsymbol{\mu}(\mathbf{r}_0)$ placed at the point $\mathbf{r}_0$ is given by $\mathbf{A}(\mathbf{r}) = \boldsymbol{\mu}(\mathbf{r}_0) \times (\mathbf{r} - \mathbf{r}_0)/|\mathbf{r} - \mathbf{r}_0|^3$ [14], one can show that the vector-potential created by the magnetization of the sheet is

$$\mathbf{A}(\mathbf{r}, t) = e^{-i\omega t} 2\pi\rho(\zeta)(\mathbf{m} \times \mathbf{e}_z)\,\mathrm{sign}(z - \zeta), \tag{12}$$

so that the vector-potential created by the total magnetization of the electron layer is given by

$$\mathbf{A}_\omega(z) = 2\pi(\mathbf{m} \times \mathbf{e}_z)\left[\int_{-d/2}^{z} \rho(\zeta)\, d\zeta - \int_{z}^{d/2} \rho(\zeta)\, d\zeta\right]. \tag{13}$$

The corresponding part of the electric field, $\mathbf{E}_\omega = (i\omega/c)\mathbf{A}_\omega$, has the same space dependence. According to Eq. (11), the electric current induced by *this* part of the field is defined by the expression