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Effects of the Electron–Electron Interaction in the Spin Resonance in 2D Systems with Dresselhaus Spin–Orbit Coupling

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Abstract—The effect of the electron–electron interaction on the spin-resonance frequency in two-dimensional electron systems with Dresselhaus spin–orbit coupling is investigated. The oscillatory dependence of many-body corrections on the magnetic field is demonstrated. It is shown that the consideration of many-body interaction leads to a decrease or an increase in the spin-resonance frequency, depending on the sign of the g factor. It is found that the term cubic in quasimomentum in Dresselhaus spin–orbit coupling partially decreases exchange corrections to the spin resonance energy in a two-dimensional system.

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1. INTRODUCTION

Spin–orbit coupling (SOC), being relativistic in nature, leads to the appearance of gyrotropy in two-dimensional (2D) semiconductor systems and spin splitting in the energy spectrum of quasiparticles in the absence of external magnetic field. This splitting for states near the bottom of the conduction band is linear in quasi-momentum and is associated with the lack of an inversion center in the system. Spin splitting in 2D systems based on III–V semiconductors is caused by two main terms in the effective Hamiltonian. This effect is also known as Rashba [1] and Dresselhaus [2] SOC. The former contribution appears in quantum wells (QWs) with different barriers or asymmetric potential profiles and in the presence of an external electric field (structure inversion asymmetry, SIA). The latter contribution is caused by the absence of a spatial inversion center in the unit cell of the materials used to grow the 2D structure (bulk inversion asymmetry, BIA).[1]

The Dresselhaus SOC Hamiltonian in bulk III–V semiconductors for states near the bottom of the conduction band is given by

$$\hat{H}_{SO}(\hat{\mathbf{k}}) = \gamma\{(\hat{k}_y\hat{k}_x\hat{k}_y - \hat{k}_z\hat{k}_x\hat{k}_z)\hat{\sigma}_x + (\hat{k}_z\hat{k}_y\hat{k}_z - \hat{k}_x\hat{k}_y\hat{k}_x)\hat{\sigma}_y + (\hat{k}_x\hat{k}_z\hat{k}_x - \hat{k}_y\hat{k}_z\hat{k}_y)\hat{\sigma}_z\}, \quad (1)$$

where γ is the Dresselhaus coupling constant, $\hat{\sigma}_i$ and $\hat{k}_i$ ($i = x, y, z$) are the Pauli matrices and operators of generalized momentum components. The x, y, z axes are oriented along the cubic crystal axes [100], [010], and [001], respectively. We note that $\hat{H}_{SO}(\hat{\mathbf{k}})$ differs from the Hamiltonian obtained in [2]. The form of Eq. (1) for the Dresselhaus SOC takes into account the regular order for generalized momentum components in the presence of a uniform magnetic field [6].

In 2D systems, due to the size quantization of the momentum component along the growth axis, e.g., the z axis, is replaced in the Hamiltonian by $-i\partial/\partial z$. The SOC Hamiltonian can be obtained by averaging over adiabatically rapid motion along the growth axis. Performing such averaging and taking into account that $\langle\hat{k}_z\rangle = 0$ and $\langle\hat{k}_z^2\rangle \neq 0$, where angle brackets mean quantum-mechanical averaging over the wave function of size quantization, it is easy to obtain the 2D form of the Dresselhaus SOC [7] which, in addition to the term $\hat{H}_D^{(1)}(\hat{k}_x, \hat{k}_y)$ linear in $\hat{k}_x$ and $\hat{k}_y$, contains a component cubic in quasi-momentum, $\hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y)$,

$$\begin{aligned} \hat{H}_{SO}(\hat{\mathbf{k}}) &= \hat{H}_D^{(1)}(\hat{k}_x, \hat{k}_y) + \hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y), \\ \hat{H}_D^{(1)}(\hat{k}_x, \hat{k}_y) &= \beta(\hat{k}_x\hat{\sigma}_x - \hat{k}_y\hat{\sigma}_y), \\ \hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y) &= \gamma(\hat{k}_y\hat{k}_x\hat{k}_y\hat{\sigma}_x - \hat{k}_x\hat{k}_y\hat{k}_x\hat{\sigma}_y), \end{aligned} \quad (2)$$

where $\beta = -\gamma\langle\hat{k}_z^2\rangle$ is the 2D Dresselhaus constant depending on the QW width. The term $\hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y)$ is

[1] In addition to the contributions caused by SIA and BIA, the SOC Hamiltonian in 2D systems also contains a contribution associated with the symmetry and structure of heterojunctions in QWs (interface inversion asymmetry, IIA) [3]. Consideration of the contribution of heterojunctions to the SOC Hamiltonian in the absence of magnetic field leads only to renormalization of the Rashba and Dresselhaus contributions [3–5].

often neglected supposing that $\langle \hat{k}_z^2 \rangle \gg K_F^2$, where $k_F = \sqrt{2\pi n_S}$ is the Fermi wave vector (n_S is the 2D electron density). Recent experiments [8–10] show an important role of the cubic term in the Dresselhaus SOC in 2D systems even at very low electron-gas densities.

The use of electron spin resonance (ESR) for measuring the SOC constant was first proposed in [11, 12]. Due to the extraordinary sensitivity to SOC and high accuracy when determining the absorption-line position, the ESR was efficiently used to determine the Rashba SOC constant in 2D systems [13–17]. The Dresselhaus SOC in QWs placed in a magnetic field leads to the anisotropy of spin splitting of the electronic levels, described as the anisotropy of the effective g-factor in the 2D-system plane [18, 19]. The in-plane anisotropy of the g-factor in the GaAs/AlGaAs QWs, determined by ESR measurements was experimentally observed in [20, 21].

In addition to the fact that the SOC leads to g-factor anisotropy in strong magnetic fields and an increase in the ESR energy in weak magnetic fields,[2] it also causes additional renormalization of the ESR energy in the 2D system, associated with the electron–electron (e–e) interaction (Larmor theorem violation [22]. The effect of the e–e interaction on the ESR energy was studied in detail for 2D systems with Rashba SOC [22–24] and for QWs based on narrow-gap semiconductors such as InAs and InSb [25–30].

The present paper is devoted to studying ESR energy renormalization caused by the e–e interaction in a 2D electronic system with the Dresselhaus SOC cubic, $\hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y)$, and linear, $\hat{H}_D^{(1)}(\hat{k}_x, \hat{k}_y)$, in momentum. In this case, nonparabolicity and disorder effects are disregarded.

2. THEORY

Let us consider a 2D system in the absence of disorder, placed in a magnetic field directed along the z axis perpendicular to the system plane. To calculate the energies and wave functions of single-electron states, instead of momentum operators in the (x, y) plane, it is convenient to introduce the "ladder operators" as follows

$$b^+ = \frac{a_B}{\sqrt{2}}(k_x + ik_y),$$

$$b = \frac{a_B}{\sqrt{2}}(k_x - ik_y),$$

$$bb^+ - b^+b = 1,$$

where $a_B = \sqrt{\hbar c/eB}$ is the magnetic length and B is the magnetic induction. As a result, the Hamiltonian for describing single-particle states can be written as

$$\hat{H}_{(1e)} = \hbar\omega_c\left(b^+b + \frac{1}{2}\right) + \frac{g^*}{2}\mu_B B\sigma_z + \hat{H}_{SO}^{(ax)} + \hat{H}_{SO}^{(wp)},$$
$$\hat{H}_{SO}^{(ax)} = \tilde{\beta}\begin{pmatrix} 0 & ib^+ \\ -ib & 0 \end{pmatrix} + \tilde{\gamma}\begin{pmatrix} 0 & ib^+bb^+ \\ -ibb^+b & 0 \end{pmatrix}, \tag{3}$$

where m^* and g^* are the electron effective mass and g factor, $\mu_B > 0$ is the Bohr magneton, $\omega_c = eB/m^*c$, $\tilde{\beta} = \sqrt{2}\beta/a_B$, and $\tilde{\gamma} = \gamma/\sqrt{2}a_B^3$. The term $\hat{H}_{SO}^{(wp)}$ in Eq. (3) appearing from the anisotropic part of the cubic Dresselhaus SOC term $\hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y)$ has the form

$$\hat{H}_{SO}^{(wp)} = -\tilde{\gamma}\begin{pmatrix} 0 & ib^3 \\ -i(b^+)^3 & 0 \end{pmatrix}. \tag{4}$$

It is easy to see that the dispersion relation of 2D electrons in the absence of term (4) is isotropic and the single-particle problem can be solved analytically. The eigenvalues of the "reduced" Hamiltonian $\hat{H}_{(1e)} - \hat{H}_{SO}^{(wp)}$ have two branches denoted a and b,

$$\tilde{E}_n^{(a)} = \hbar\omega_c n + \sqrt{E_0^2 + \hat{B}_n^2 n}, \quad n = 0, 1, 2, \ldots, \tag{5}$$

and

$$\tilde{E}_n^{(b)} = \hbar\omega_c n + \sqrt{E_0^2 + \hat{B}_n^2 n}, \quad n = 0, 1, 2, \ldots, \tag{6}$$

where

$$E_0 = \frac{1}{2}(\hbar\omega_c + g^*\mu_B B),$$

$$\hat{B}_n = \tilde{\beta} + \tilde{\gamma} n.$$

It is convenient to present the corresponding wavefunctions in the form

$$\tilde{\Psi}_{n,k}^{(a)}(x, y, z) = \sin\varphi_n\begin{pmatrix} |n, k\rangle \\ 0 \end{pmatrix} - i\cos\varphi_n\begin{pmatrix} 0 \\ |n-1, k\rangle \end{pmatrix},$$
$$\tilde{\Psi}_{n,k}^{(b)}(x, y, z) = \sin\varphi_n\begin{pmatrix} 0 \\ |n-1, k\rangle \end{pmatrix} - i\cos\varphi_n\begin{pmatrix} |n, k\rangle \\ 0 \end{pmatrix}, \tag{7}$$

where $|n-1, k\rangle$ corresponds to the normalized wavefunction of the harmonic oscillator [31],

$$\sin\varphi_n = \frac{E_0 + \sqrt{E_0^2 + \hat{B}_n^2}}{\sqrt{\left(E_0 + \sqrt{E_0^2 + \hat{B}_n^2 n}\right)^2 + \hat{B}_n^2 n}},$$
$$\cos\varphi_n = \frac{\hat{B}_n\sqrt{n}}{\sqrt{\left(E_0 + \sqrt{E_0^2 + \hat{B}_n^2 n}\right)^2 + \hat{B}_n^2 n}}. \tag{8}$$

[2] In zero magnetic field, the ESR energy is defined by the spin splitting of the carrier energy spectrum.

The electron spin resonance corresponds to the transition between the (n, a) and ($n + 1$, b) levels. In weak magnetic fields at large n such that $k_F^2 \approx 2n/a_B^2$, the ESR energy takes the form

$$E_{ESR} = \left|\sqrt{(g^*\mu_B B + \hbar\omega_c)^2(\Delta_D^{(1)} + \Delta_D^{(3)}/4)^2} - \hbar\omega_c\right|, \quad (9)$$

where the $\Delta_D^{(1)} = 2\beta k_F$ and $\Delta_D^{(3)} = 2\gamma k_F^2$ are the energy-spectrum splittings in zero magnetic field, corresponding to the linear and cubic parts of the Dresselhaus SOC.

The anisotropic term $\hat{H}_{SO}^{(wp)}$ in the energy and wavefunctions of single-particle states was considered according to perturbation theory with an accuracy up to the second order in wavefunctions (7). The direct diagonalization of $\hat{H}_{(1e)}$ in the basis of eigenfunctions of the "reduced" Hamiltonian $\hat{H}_{(1e)} - \hat{H}_{SO}^{(wp)}$ at m^*, g^*, β, and γ characteristic of 2D structures based on GaAs/AlGaAs shows that the corresponding contribution of $\hat{H}_{SO}^{(wp)}$ does not exceed 1–2% even in weak magnetic fields (at which $k_F^2 \approx 2n/a_B^2$), which validates the use of perturbation theory. We note that the application of perturbation theory in calculating the energies and wavefunctions of single-particle states instead of numerical diagonalization of $\hat{H}_{(1e)}$ makes it possible to save a significant amount of calculation time in solving the many-body problem.

The total Hamiltonian of the 2D system, taking into account the e–e interaction, H_{int}, in the secondary quantization representation is written as

$$\begin{aligned} H &= H_{(0)} + H_{int}, \\ H_{(0)} &= \int d^2\mathbf{r}\Psi^+(\mathbf{r})H_{(1e)}\Psi(\mathbf{r}), \\ H_{int} &= \frac{1}{2}\int d^2\mathbf{r}_1\int d^2\mathbf{r}_2 \\ &\times \Psi^+(\mathbf{r}_1)\Psi^+(\mathbf{r}_2)V(|\mathbf{r}_1 - \mathbf{r}_2|)\Psi(\mathbf{r}_2)\Psi(\mathbf{r}_1), \end{aligned} \quad (10)$$

where $\mathbf{r} = (x, y)$ is the radius vector in the system plane, $V(\mathbf{r}_1 - \mathbf{r}_2)$ is the Coulomb potential, the superscript "+" corresponds to the Hermitian conjugate. In Eq. (10), the field operators $\Psi(\mathbf{r})$ and $\Psi^+(\mathbf{r})$ containing the fermion creation and annihilation operators $a_{n,k,i}$, $a_{n,k,i}^+$, and the wave functions $\varphi_{n,k}^{(a)}(x, y)$ and $\varphi_{n,k}^{(b)}(x, y)$ of single-particle states, calculated taking into account $\hat{H}_{SO}^{(wp)}$,

$$\begin{aligned} \Psi(\mathbf{r}) &= \sum_{n,k,i}\varphi_{n,k}^{(i)}(x, y)a_{n,k,i}, \\ \Psi^+(\mathbf{r}) &= \sum_{n,k,i}\varphi_{n,k}^{(i)+}(x, y)a_{n,k,i}^+, \end{aligned} \quad (11)$$

where $i = a, b$. The use of the Fourier transform for the Coulomb potential,

$$V(|\mathbf{r} - \mathbf{r}'|) = \int\frac{d^2\mathbf{q}}{(2\pi)^2}\tilde{D}(q)e^{i\mathbf{q}(\mathbf{r} - \mathbf{r}')}, \quad (12)$$

makes it possible to reduce calculation of the matrix elements of the e–e interaction H_{int} on the wavefunctions $\varphi_{n,k}^{(a)}(x, y)$ and $\varphi_{n,k}^{(b)}(x, y)$ to calculation of the matrix elements $\langle n_1, k_1|e^{i\mathbf{qr}}|n_2, k_2\rangle$ [24].

The Fourier transform $\tilde{D}(q)$ of the Coulomb potential in the 2D system has the form

$$\tilde{D}(q) = \frac{2\pi e^2}{\varepsilon q}F(q), \quad (13)$$

where ε is the dielectric constant of the system, $F(q)$ is the geometrical form factor taking into account the nonzero thickness of the 2D system along the z axis and the fields of electrostatic images. Then we disregard electron motion along the z axis, setting $F(q) = 1$.

After certain calculations, we can obtain the following expressions for $H_{(0)}$ and H_{int}

$$\begin{aligned} H_{(0)} &= \sum_{n,k,i}E_n^{(i)}a_{n,k,i}^+a_{n,k,i}, \\ H_{int} &= \frac{1}{2}\sum_{n_1\ldots n_4}\sum_{i_1\ldots i_4}\sum_{k_1k_2}\int\frac{d^2q}{(2\pi)^2}\tilde{V}_{n_1,n_2,n_3,n_4}^{(i_1,i_2,i_3,i_4)}(\mathbf{q}) \\ &\times e^{iq_x(k_1 - k_2 + q_y)a_B^2}a_{n_1,k_1,i_1}^+a_{n_2,k_2,i_2}^+a_{n_3,k_2 - q_y,i_3}a_{n_4,k_1 + q_y,i_4}, \end{aligned} \quad (14)$$

where $E_n^{(i)}$ are the energy eigenvalues of the Hamiltonian $\hat{H}_{(1e)}$; in the matrix element $\tilde{V}_{n_1,n_2,n_3,n_4}^{(i_1,i_2,i_3,i_4)}(\mathbf{q}) \propto \tilde{D}(q)$ of the e–e interaction, terms to within the second order of smallness in the matrix element $\hat{H}_{SO}^{(wp)}$ are retained.

To find the ESR energy taking into account the e–e interaction, it is convenient to use the exciton representation [32–34]. During the electron transition between Landau levels in the 2D system, quasielectron–quasihole pairs are generated (quasielectrons over the Fermi level and quasiholes under the Fermi level) which results in the ground-to-excited state transition of the system. To describe the excited state of the 2D system, formed by an electron transferred to an unoccupied or partially occupied Landau level (n, i) and an effective hole appearing at the previous level (n', i'), we define the creation operator of the magnetic exciton with momentum k as

$$A_{n,n',i,i'}^+(\mathbf{k}) = \sum_p e^{ik_x(p + k_y/2)a_B^2}a_{n,p,i}^+a_{n',p + k_y,i'}, \quad (15)$$

which satisfies the following commutation relation

$$[A^{+}_{n_1, n_2, i_1, i_2}(\mathbf{k}_1), A^{+}_{n_3, n_4, i_3, i_4}(\mathbf{k}_2)] = A^{+}_{n_1, n_4, i_1, i_4}(\mathbf{k}_1+\mathbf{k}_2)e^{-\frac{i}{2}a_B^2(\mathbf{k}_1\times\mathbf{k}_2)\big|_z}\delta_{n_2, n_3}\delta_{i_2, i_3} - A^{+}_{n_3, n_2, i_3, i_2}(\mathbf{k}_1+\mathbf{k}_2)e^{\frac{i}{2}a_B^2(\mathbf{k}_1\times\mathbf{k}_2)\big|_z}\delta_{n_1, n_4}\delta_{i_1, i_4}. \quad (16)$$

The energy E_{ex} of this magnetic exciton, measured from the ground state energy $|0\rangle$ of the system satisfies the equation

$$E_{\text{ex}}A^{+}_{n, n', i, i'}(\mathbf{k})|0\rangle = (E^{(i)}_n - E^{(i')}_{n'})A^{+}_{n, n', i, i'}(\mathbf{k})|0\rangle + [H_{\text{int}}, A^{+}_{n, n', i, i'}(\mathbf{k})]|0\rangle. \quad (17)$$

To calculate the commutator on the right-hand side of Eq. (17), it is convenient to rewrite expression (14) in the form

$$\hat{H}_{\text{int}} = \frac{1}{2}\sum_{n_1\ldots n_4}\sum_{i_1\ldots i_4}\int\frac{d^2q}{(2\pi)^2}\tilde{V}^{(i_1, i_2, i_3, i_4)}_{n_1, n_2, n_3, n_4}(\mathbf{q}) \times A^{+}_{n_1, n_4, i_1, i_4}(\mathbf{q})A^{+}_{n_2, n_3, i_2, i_3}(-\mathbf{q}) - \frac{1}{2}\sum_{n_1, n_2, n_3}\sum_{i_1, i_2, i_3}\int\frac{d^2q}{(2\pi)^2}\tilde{V}^{i_1, i_2, i_3, i_4}_{n_1, n_2, n_3, n_4}(\mathbf{q})A^{+}_{n_1, n_2, i_1, i_3}(0). \quad (18)$$

Then, using the commutation relations between the exciton operators and retaining only the terms proportional to the product of the fermion creation and annihilation operators on the right-hand side of expression (17), multiplying this product by the operator of the number of particles, and taking into account that

$$\langle 0|a^{+}_{n_1, p_1, i_1}a_{n_2, p_2, i_2}|0\rangle = \delta_{n_1, n_2}\delta_{p_1, p_2}\delta_{i_1, i_2}\nu^{(i_1)}_{n_1}, \quad (19)$$

where $\nu^{(i)}_n$ is the Landau-level (n, i) filling factor, we find the following expression for $[H_{\text{int}}, A^{+}_{n, n', i, i'}(\mathbf{k})]$ (see also [24–26])

$$[H_{\text{int}}, A^{+}_{n, n', i, i'}(\mathbf{k})]|0\rangle = -\sum_{n_2, i_2}\nu^{(i_2)}_{n_2}(\tilde{E}^{(i, i_2, i, i_2)}_{n, n_2, n, n_2}(0) - \tilde{E}^{(i', i_2, i', i_2)}_{n', n_2, n', n_2}(0))A^{+}_{n, n', i, i'}(\mathbf{k})|0\rangle - (\nu^{(i)}_n - \nu^{(i')}_{n'}) \times \sum_{n_1, n_4, i_1, i_4}\frac{\tilde{V}^{(i_1, i', i, i_4)}_{n_1, n', n, n_4}(\mathbf{k})}{2\pi}A^{+}_{n_1, n_4, i_1, i_4}(\mathbf{k})|0\rangle + (\nu^{(i)}_n - \nu^{(i')}_{n'})\sum_{n_1, n_2, i_1, i_2}\tilde{E}^{(i', i_1, i, i_2)}_{n', n_1, n, n_2}(\mathbf{k})A^{+}_{n_1, n_2, i_1, i_2}(\mathbf{k})|0\rangle. \quad (20)$$

In expression (20), the matrix element $\tilde{E}^{(i_1, i_2, i_3, i_4)}_{n_1, n_2, n_3, n_4}(\mathbf{k})$ is defined as

$$\tilde{E}^{(i_1, i_2, i_3, i_4)}_{n_1, n_2, n_3, n_4}(\mathbf{k}) = \int\frac{d\mathbf{q}}{(2\pi)^2}\tilde{V}^{(i_1, i_2, i_3, i_4)}_{n_1, n_2, n_3, n_4}(\mathbf{q})e^{ia_B^2(\mathbf{q}\times\mathbf{k})_z}. \quad (21)$$

We can see that the second and third terms in expression (20) provide mixing of all possible states of the 2D system, containing magnetic excitons.

Let us dwell on the magnetic-exciton excitation associated with the electron transition between Landau levels (n, a) and $(n+1, b)$, whose energy in the long-wavelength limit corresponds to the ESR energy. It can be shown that in strong magnetic fields such that $B > B_{\text{cr}}$, where B_{cr} is determined from the condition

$$\frac{(\hbar\omega_c)^2}{(\tilde{\beta}+\tilde{\gamma}n)^2} \approx \frac{2n+2+\sqrt{4n(n+2)+A^2}}{4(4-A^2)}, \quad A = 1+\frac{g^*m^*}{2m_0}, \quad (22)$$

mixing of the quasielectron–quasihole pair excitation between Landau levels (n, a) and $(n+1, b)$ with other magnetic excitons can be disregarded (see, e.g., [24]). Thus, the ESR energy taking into account the *e–e* interaction takes the form

$$E_{\text{ESR}} = \left|E^{(a)}_n - E^{(b)}_{n+1} + \Delta^{(e-e)}_{\text{ESR}}\right| \quad \text{at} \quad E^{(a)}_n \geq E^{(b)}_{n+1}, \qquad E_{\text{ESR}} = \left|E^{(b)}_{n+1} - E^{(a)}_n - \Delta^{(e-e)}_{\text{ESR}}\right| \quad \text{at} \quad E^{(a)}_n < E^{(b)}_{n+1}, \quad (23)$$

where the correction to the ESR energy $\Delta^{(e-e)}_{\text{ESR}}$, caused by manyparticle effects, is given by

$$\Delta^{(e-e)}_{\text{ESR}} = (\nu^{(a)}_n - \nu^{(b)}_{n-1})\tilde{E}^{(b, a, a, b)}_{n+1, n, n, n+1}(0) - \sum_{n_2, i_2}\nu^{(i_2)}_{n_2}(\tilde{E}^{(a, i_2, a, i_2)}_{n, n_2, n, n_2}(0) - \tilde{E}^{(b, i_2, b, i_2)}_{n+1, n_2, n+1, n_2}(0)). \quad (24)$$

It is easy to show that, in the absence of SOC, i.e., at $\beta = 0, \gamma = 0$, $\Delta^{(e-e)}_{\text{ESR}} = 0$, the ESR energy, in full compliance with the Larmor theorem, is $g^*\mu_B B$ [35].

3. RESULTS AND DISCUSSION

To illustrate the obtained theoretical results, in this section, we consider a "model" 2D system with the dielectric constant $\varepsilon = 12.5$, the effective electron mass $m^* = 0.067m_0$ (m_0 is the free electron mass), and the g factor varying in the range from –0.4 to 0.4, which are characteristic of 2D structures based on GaAs/AlGaAs [8–10, 20, 21]. The 2D electron density is set to $n_S = 4.0 \times 10^{11}$ cm^{-2}. From the results of various experimental studies of 2D structures based on GaAs/AlGaAs, it is known that the constant γ is in the range between –3 and –35 eV Å^3 [10, 36–40]. In the present study, according to [10], γ is set to –11 eV Å^3.

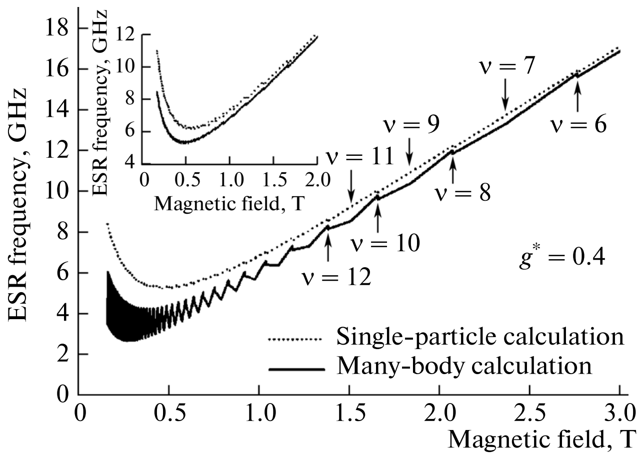

Fig. 1. Spin-resonance frequency as a function of the magnetic field at $\beta = 5 \times 10^3$ eV Å, $\gamma = -11$ eV Å^3, and $g^* = 0.4$ in the single-electron approximation (dotted curve) and with consideration of the e–e interaction (solid curve). Arrows indicate the magnetic fields corresponding to the integer Landau-level filling factor. The inset shows the results of "single-electron" calculations at $\gamma = -11$ (solid curve) and 0 eV Å^3 (dotted curve).

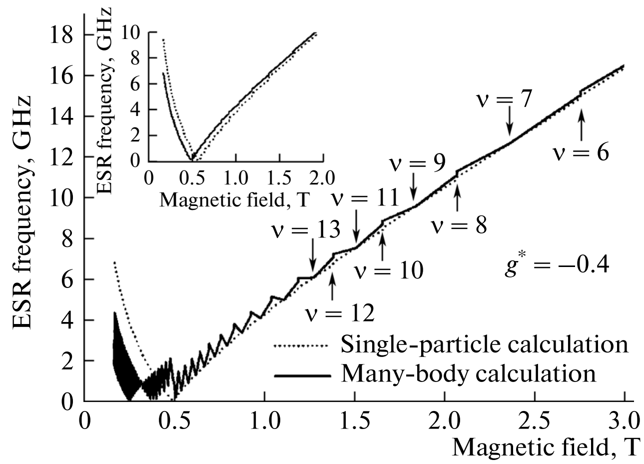

Fig. 2. Spin-resonance frequency as a function of the magnetic field at $\beta = 5 \times 10^3$ eV Å, $\gamma = -11$ eV Å^3, and $g^* = -0.4$ in the single-electron approximation (dotted curve) and with consideration of the e–e interaction (solid curve). Arrows indicate the magnetic fields corresponding to the integer Landau-level filling factor. The inset shows the results of "single-electron" calculations at $\gamma = -11$ (solid curve) and 0 eV Å^3 (dotted curve).

Figures 1 and 2 show the results of calculations of the ESR frequency as a function of the magnetic field at various g factors. The dotted curves correspond to the single-electron approximation, solid curves are the results of calculations taking into account the e–e interaction. The insets in Figs. 1 and 2 show the "single-electron" ESR frequencies with (solid curve) and without (dotted curve) consideration for the cubic term in the Dresselhaus SOC. Arrows indicate the magnetic fields corresponding to integer Landau-level filling factors. We note that the performed calculations disregard excitation mixing between Landau levels (n, a) and (n + 1, b) with other magnetic excitons, i.e., are bounded by the magnetic-field region $B > B_{cr}$, where B_{cr} is determined from condition (22). At the chosen parameters for the model 2D system, the value $B_{cr} \approx 0.16$ T for the Landau level n_F intersecting the Fermi level was reached at $n_F \sim 50$. At a fixed magnetic field, n_F was calculated from the condition

$$2\pi a_B^2 \sum_{n=0}^{n_F} (\nu_n^{(a)} + \nu_{n+1}^{(b)}) = n_S. \quad (25)$$

As seen in Figs. 1 and 2, the behavior of the "single-electron" ESR frequency in the magnetic field is controlled by the 2D electron g-factor sign. At $g^* > 0$, the dependence of the ESR frequency on the magnetic field is U-shaped (Fig. 1); at $g^* < 0$, the dependence is V-shaped (Fig. 2), and the ESR frequency vanishes under the condition

$$\frac{(\hbar\omega_c)^2}{(\tilde{\beta} + \tilde{\gamma} n_F)^2} \approx \frac{2n_F + 1 + \sqrt{4n_F(n_F+1) + A^2}}{4(1 - A^2)}, \quad (26)$$

where A is defined as in expression (22).

In weak magnetic fields, the "single-electron" ESR frequency is defined by the Dresselhaus splitting in zero magnetic field, i.e., by $\Delta_D^{(1)}$ and $\Delta_D^{(3)}$; in strong magnetic fields, the "single-particle" ESR energy tends to the Zeeman energy. The features of the "single-electron" ESR frequency, appearing at even Landau-level filling factors, are related to Fermi-level oscillations in the magnetic field. As the magnetic field increases, the Fermi level jumps from one Landau-level pair split in spin to a lower-lying pair whose spin splitting depends on the Landau-level number n due to the SOC; as a result, a spin-splitting jump at the Fermi level corresponding to the ESR frequency occurs.

Consideration of the e–e interaction leads to significant renormalization of the ESR frequency in the magnetic-field region where the Dresselhaus SOC has a significant effect on the Landau-level splitting. The data of Figs. 1 and 2 show that the contribution to the ESR frequency, caused by the e–e interaction depends on the effective g-factor sign in the 2D system. We can see that the e–e interaction at positive g-factors results in a decrease on the ESR frequency at arbitrary magnetic fields and Landau level filling factors. At $g^* < 0$, the contribution of the e–e interaction to the ESR frequency can be both positive and negative, depending on the magnetic field. In the region of weak magnetic fields corresponding to $B \ll B_0$ (where B_0 is determined from condition (26)), the e–e interaction results in a decrease in the ESR frequency in comparison with "single-electron" values. In strong magnetic fields such that $B \gg B_0$, the e–e interaction leads to an increase in the ESR frequency in comparison with "single-electron" values. In the vicinity of $B \approx B_0$, the

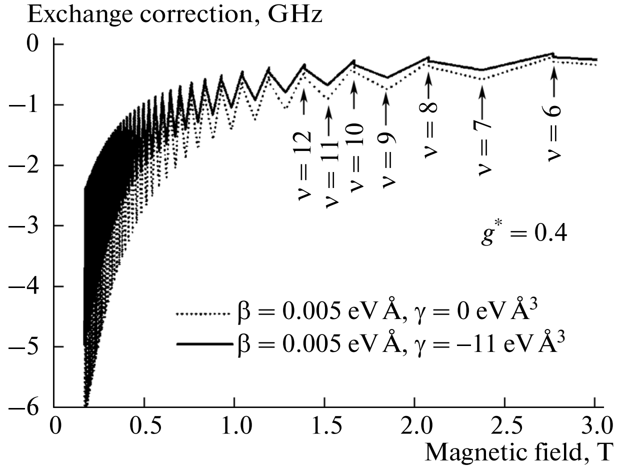

Fig. 3. Many-body corrections to the ESR frequency as functions of the magnetic field at $\beta = 5 \times 10^3$ eV Å and $g^* = 0.4$ with ($\gamma = -11$ eV Å^3, solid curve) and without ($\gamma = 0$ eV Å^3, a dotted curve) consideration of the term cubic in terms of wave vector in the Dresselhaus SOC Hamiltonian.

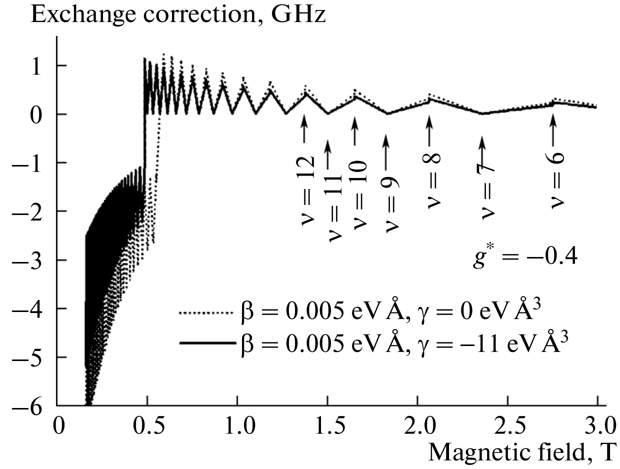

Fig. 4. Corrections associated with the *e*–*e* interaction to the ESR frequency as a function of the magnetic field at $\beta = 5 \times 10^3$ eV Å and $g^* = -0.4$ with ($\gamma = -11$ eV Å^3, solid curve) and without ($\gamma = 0$ eV Å^3, dotted curve) consideration for the term cubic in terms of the wave vector in the Dresselhaus SOC Hamiltonian.

dependence of the ESR frequency on the magnetic field has a complex behavior.

The oscillatory behavior of the "many-body" correction in the magnetic field is associated with oscillations in the difference between the filling factors of the Landau levels involved in the transition corresponding to the spin resonance. At zero temperature in the absence of disorder in the 2D system, the densities of states at Landau levels are described by Dirac δ-functions, and the spin-split Landau levels are not overlapped, which leads to pronounced oscillations of the "many-body" ESR frequency even in weak magnetic fields. In the case of consideration of the finite Landau level width (see, e.g., [41]), ESR frequency oscillations should be spread with decreasing magnetic field due to an increase in the overlap of the densities of states of spin-split Landau levels (n_F, a) and ($n_F + 1$, b). Furthermore, an additional ESR frequency shift can be expected, which is associated with the random potential of impurities in the 2D system with SOC [42].

Figure 3 and 4 show the results of calculations of the corrections to the ESR frequency, associated with the *e*–*e* interaction with (solid curve) and without (dotted curve) consideration of the term cubic in terms of the wave vector in the Dresselhaus SOC Hamiltonian. Arrows indicate the magnetic fields corresponding to integer Landau-level filling factors. We can see that the consideration of $\hat{H}_D^{(3)}(\hat{k}_x, \hat{k}_y)$ in the Dresselhaus Hamiltonian (2) results in a decrease in exchange corrections to the ESR frequency. Since $\gamma < 0$, the cubic term in Eq. (2) is partially compensates the contribution of the linear term $\hat{H}_D^{(1)}(\hat{k}_x, \hat{k}_y)$ and results in a decrease in the total contribution of SOC to the ESR frequency. The decrease in the contribution of the

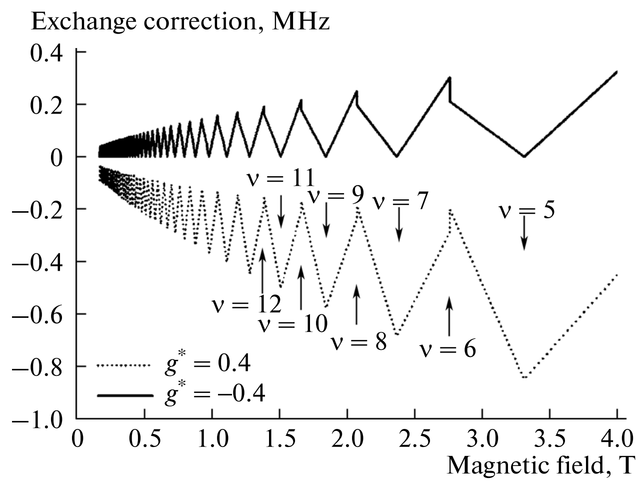

Fig. 5. Manyparticle corrections to the ESR energy at various *g* factors and $\beta = -\gamma k_F^2/4$.

SOC to the ESR frequency according to the Larmor theorem causes a decrease in corrections associated with the *e*–*e* interaction.

It follows from expression (9) that the Dresselhaus spin splitting in zero magnetic field is almost lacking at $\beta = -\gamma k_F^2/4$; in weak magnetic fields, the "single-electron" ESR energy is defined as $g^*\mu_B B$ with an accuracy up to terms $\hat{H}_{SO}^{(wp)}$. As seen in Fig. 5, consideration of the *e*–*e* interaction in this case leads to only insignificant renormalization of the ESR energy whose value does not exceed 1% at magnetic fields of > 0.5 T. In this case, at positive (negative) *g* factors, the ESR frequency decreases (increases) in the entire magnetic-field range.

4. CONCLUSIONS

Violation of the Larmor theorem was theoretically studied, and corrections to the spin-resonance energy in a 2D system with Dresselhaus spin–orbit coupling, caused by the *e*–*e* interaction, were calculated. The oscillatory dependence of the "many-body" ESR frequency on the magnetic field, associated with Fermi-level oscillations in the magnetic field was presented. It was shown that the many-body corrections to the ESR energy depend not only on the interrelation of linear and cubic terms in the Dresselhaus SOC, but also on the electron *g*-factor sign in the 2D system.

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REFERENCES

1. Yu. A. Bychkov and E. I. Rashba, J. Phys. C: Solid State Phys. **17**, 6039 (1984).
2. G. Dresselhaus, Phys. Rev. **100**, 580 (1955).
3. L. E. Golub and E. L. Ivchenko, Phys. Rev. B **69**, 115333 (2004).
4. P. S. Alekseev, JETP Lett. **98**, 84 (2013).
5. Zh. A. Devizorova and V. A. Volkov, JETP Lett. **98**, 101 (2013).
6. P. S. Alekseev, M. V. Yakunin, and I. N. Yassievich, Semiconductors **41**, 1092 (2007).
7. M. I. D'yakonov and V. Yu. Kachorovskii, Sov. Phys. Semicond. **20**, 110 (1986).
8. J. D. Koralek, C. P. Weber, J. Orenstein, B. A. Bernevig, S.-C. Zhang, S. Mack, and D. D. Awschalom, Nature (London) **458**, 610 (2009).
9. M. Kohda, V. Lechner, Y. Kunihashi, T. Dollinger, P. Olbrich, C. Schönhuber, I. Caspers, V. V. Bel'kov, L. E. Golub, D. Weiss, K. Richter, J. Nitta, and S. D. Ganichev, Phys. Rev. B **86**, 081306 (2012).
10. M. P. Walser, U. Siegenthaler, V. Lechner, D. Schuh, S. D. Ganichev, W. Wegscheider, and G. Salis, Phys. Rev. B **86**, 195309 (2012).
11. E. I. Rashba and V. I. Sheka, Sov. Phys. Solid State **3**, 1257 (1961).
12. E. I. Rashba and V. I. Sheka, in *Landau Level Spectroscopy*, Ed. by G. Landwehr and E. I. Rashba (North-Holland, Amsterdam, 1991), p. 131.
13. Z. Wilamowski, N. Sandersfeld, W. Jantsch, D. Többen, and F. Schäffler, Phys. Rev. Lett. **87**, 026401 (2001).
14. Z. Wilamowski, W. Jantsch, H. Malissa, and U. Rössler, Phys. Rev. B **66**, 195315 (2002).
15. G. A. Khodaparast, R. E. Doezema, S. J. Chung, K. J. Goldammer, and M. B. Santos, Phys. Rev. B **70**, 155322 (2004).
16. Z. Wilamowski, H. Malissa, F. Schäffler, and W. Jantsch, Phys. Rev. Lett. **98**, 187203 (2007).
17. Y. Kozuka, S. Teraoka, J. Falson, A. Oiwa, A. Tsukazaki, S. Tarucha, and M. Kawasaki, Phys. Rev. B **87**, 205411 (2013).
18. V. K. Kalevich and V. L. Korenev, JETP Lett. **57**, 571 (1993).
19. P. S. Alekseev, Semiconductors **47**, 1241 (2013).
20. Yu. A. Nefyodov, A. V. Shchepetilnikov, I. V. Kukushkin, W. Dietsche, and S. Schmult, Phys. Rev. B **83**, 041307 (2011).
21. Yu. A. Nefyodov, A. V. Shchepetilnikov, I. V. Kukushkin, W. Dietsche, and S. Schmult, Phys. Rev. B **84**, 233302 (2011).
22. M. Califano, T. Chakraborty, P. Pietiläinen, and C.-M. Hu, Phys. Rev. B **73**, 113315 (2006).
23. A. Shekhter, M. Khodas, and A. M. Finkel'stein, Phys. Rev. B **71**, 165329 (2005).
24. S. S. Krishtopenko, Semicond. Sci. Technol. **29**, 085005 (2014).
25. S. S. Krishtopenko, V. I. Gavrilenko, and M. Goiran, Phys. Rev. B **87**, 155113 (2013).
26. S. S. Krishtopenko, J. Phys.: Condens. Matter **25**, 105601 (2013).
27. S. S. Krishtopenko, J. Phys.: Condens. Matter **25**, 365602 (2013).
28. S. S. Krishtopenko, V. I. Gavrilenko, and M. Goiran, J. Phys.: Condens. Matter **24**, 252201 (2012).
29. S. S. Krishtopenko, V. I. Gavrilenko, and M. Goiran, Solid State Phenom. **190**, 554 (2012).
30. S. S. Krishtopenko, K. P. Kalinin, V. I. Gavrilenko, Yu. G. Sadofyev, and M. Goiran, Semiconductors **46**, 1163 (2012).
31. S. S. Krishtopenko, V. I. Gavrilenko, and M. Goiran, J. Phys.: Condens. Matter **23**, 385601 (2011).
32. S. M. Dikman, V. M. Zhilin, and D. V. Kulakovskii, J. Exp. Theor. Phys. **101**, 892 (2005).
33. S. M. Dickmann and I. V. Kukushkin, Phys. Rev. B **71**, 241310 (2005).
34. S. M. Dickmann and T. Ziman, Phys. Rev. B **85**, 045318 (2012).
35. C. Kallin and B. I. Halperin, Phys. Rev. B **30**, 5655 (1984).
36. W. Knap, C. Skierbiszewski, A. Zduniak, E. Litwin-Staszewska, D. Bertho, F. Kobbi, J. L. Robert, G. E. Pikus, F. G. Pikus, S. V. Iordanskii, V. Mosser, K. Zekentes, and Y. B. Lyanda-Geller, Phys. Rev. B **53**, 3912 (1996).
37. D. Richards, B. Jusserand, G. Allan, C. Priester, and B. Etienne, Solid State Electron. **40**, 127 (1996).
38. W. J. H. Leyland, R. T. Harley, M. Henini, A. J. Shields, I. Farrer, and D. A. Ritchie, Phys. Rev. B **76**, 195305 (2007).
39. S. Faniel, T. Matsuura, S. Mineshige, Y. Sekine, and T. Koga, Phys. Rev. B **83**, 115309 (2011).
40. P. S. Eldridge, J. Hübner, S. Oertel, R. T. Harley, M. Henini, and M. Oestreich, Phys. Rev. B **83**, 041301 (2011).
41. D. Antoniou and A. H. MacDonald, Phys. Rev. B **43**, 11686 (1991).
42. M. Duckheim and D. Loss, Nature Phys. **2**, 195 (2006).

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