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Electrical manipulation of spins in a nanowire with Rashba interaction

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ABSTRACT

We investigate the influence of external electric fields on the spins of a ballistic nanowire in terms of variations of the Rashba parameter and modification of the confinement potential. For a weak Rashba effect, the spins along the confinement direction in a given subband nearly assume full quantization. In the presence of a perpendicular magnetic field, the state of quantization can be manipulated using a transverse electric. This process requires modifications in the spin textures. If an in-plane magnetic field is applied, spins suffer rigid displacement to one edge of the wire and their expectation value becomes independent of the transverse electric field.

1. Introduction

Spintronic devices rely on the manipulation of spins in solid-state systems [1–5]. The spin-precession of propagating electrons at zero magnetic field induced by the Rashba spin–orbit interaction (SOI) represents the basic mechanism of spin dependent field-effect-transistors [6]. This mechanism is plausible since external control on the strength of the Rashba effect is established experimentally [7–9]. The Rashba effect also involves an intersubband coupling that mixes nearest subbands with opposite spins and produces anti-crossings at degenerate points of the energy spectrum of electrons. The reduction of the intersubband coupling is essential to improve the performance of the spin transistor, since it limits the angular distribution of electrons. In principle, this could be achieved by imposing strong transverse confining potential, that is using a quasi-one dimensional system [10,11]. However, when Rashba terms are treated on equal footing, electron spin is not a good quantum number which results in spin textures. Basically, the spin direction shows dependence on the wave vector, k, characterizing the free motion of electrons along the wire axis and the wire transversal coordinate x [12,13].

Also, spins of the nanowire can be manipulated using electric and magnetic fields applied perpendicular or in the plane of the electron gas layer. The magnetic field, $\vec{B}$, strongly alters the subband structure and opens energy gaps, which influences the physical properties of the nanowire including spin [12,14–20]. Electric voltage applied to top gates changes the strength of the SOI as we mentioned above. An electric field applied along the direction of quantum confinement induces a Stark shift in the electron spectra and changes the effective g-factor [21,22]. The effect of the electric field on the spin properties of nanowires is not reported so far.

Here, we study numerically the spin expectation value and its local density and analyze the results analytically using perturbation calculations in the weak Rashba regime. We demonstrate that the spin component along the quantum confinement is mainly quantized in this regime, since local extrema in higher spin branches transform into plateaus. In the presence of transverse electric field, the response of the spin expectation value and its accumulation at the wire edge depend on the direction of the magnetic field. In the case of in-plane magnetic fields, we show that spin density shifts linearly along the confinement direction with an increasing electric field keeping its average constant. For perpendicular magnetic fields, considerable changes in the spin textures occur which displace their average value along the k-axis. This property allows the spin quantization of a given state to be controlled by external electric fields.

2. The model

The nanowire may be realized by confining a two-dimensional electron gas in InAs using external gate voltages. We model such confinement as a parabolic potential, $m\omega x^2/2$, characterized by a frequency ω. The structural inversion asymmetry in InAs leads to the Rashba spin–orbit coupling [15]. The wire axis is taken along the y-direction, which supports free electron motion with a wave function given by $\exp iky$. Application of an electric field, $\vec{F}$, in the lateral direction alters the effective confinement potential. Furthermore, a magnetic field, $\vec{B}$, applied in the z-direction, changes the frequency of the harmonic motion due to orbital effects. However, an in-plane magnetic field only affects the spin degrees of freedom [12].

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The Hamiltonian for perpendicular magnetic fields is given by [15,21]

$$H = \frac{(\vec{p} + e\vec{A})^2}{2m} + \frac{1}{2}m\,\omega^2 x^2 + eFx + \frac{1}{2}g\mu_B \vec{\sigma}\cdot\vec{B} + H_r. \tag{1}$$

Here, $\vec{p} = (p_x, p_y)$, m is the effective mass of the electron, $\vec{\sigma}$ is the Pauli vector operator, μ_B is the Bohr magneton, and g is Lande's g-factor. Using Landau gauge, the vector potential, $\vec{A}$, can be related to the magnetic field $\vec{B} = (0, 0, B)$ through $\vec{A} = xB\hat{e}_y$. The Rashba effect is described by

$$H_r = -\frac{\alpha}{\hbar}\left[(\vec{p} + e\vec{A}) \times \vec{\sigma}\right]_z, \tag{2}$$

where α is the Rashba spin–orbit interaction constant. The strengths of different mechanisms can be compared to each other using the length scales

$$L_0 = \sqrt{\frac{\hbar}{m\omega}}, \quad L_B = \sqrt{\frac{\hbar}{m\omega_c}}, \quad L_{SO} = \frac{\hbar^2}{2m\alpha}, \tag{3}$$

where the cyclotron frequency $\omega_c = eB/m$. We also use the wave vector $k_{EF} = eF/\hbar\omega$ to characterize the action of the external electric field on electrons. After replacing p_y by $\hbar k$, the Hamiltonian in units of $\hbar\omega$ can be expressed as follows:

$$H = H_0 + H', \tag{4}$$

where

$$H_0 = \Omega\left(a^\dagger a + \frac{1}{2}\right) + \frac{1}{2}(L_0 k)^2 - \frac{1}{2}\left(\frac{\Omega x_c}{L_0}\right)^2 + \frac{1}{2}\xi_1\sigma_x + \frac{1}{2}\delta\sigma_z, \tag{5}$$

$$H' = \frac{1}{2}\xi_2(a^\dagger + a)\sigma_x + \frac{i}{2}\xi_3(a - a^\dagger)\sigma_y. \tag{6}$$

We used the shifted ladder operators a and $a^\dagger$, with

$$\Omega = \sqrt{1 + \left(\frac{\omega_c}{\omega}\right)^2} = \sqrt{1 + \left(\frac{L_0}{L_B}\right)^4}, \tag{7}$$

$$x_c = \frac{L_0}{\Omega^2}\left[L_0 k_{EF} + L_0 k\left(\frac{L_0}{L_B}\right)^2\right], \tag{8}$$

$$\xi_1 = \frac{L_0}{L_{SO}}\left(L_0 k - \frac{L_0 x_c}{L_B^2}\right), \tag{9}$$

$$\xi_2 = \frac{1}{\sqrt{2}\Omega}\frac{L_0}{L_{SO}}\left(\frac{L_0}{L_B}\right)^2, \tag{10}$$

$$\xi_3 = \sqrt{\frac{\Omega}{2}}\frac{L_0}{L_{SO}}, \tag{11}$$

and the dimensionless Zeeman splitting

$$\delta = \frac{g}{2}\frac{m}{m_0}\left(\frac{L_0}{L_B}\right)^2 \tag{12}$$

is given in terms of the free electron mass m_0. The first term in Eq. (5) shows that the oscillator frequency assumes a new normalized value Ω, given by Eq. (7), due to the presence of the magnetic field. Coupling of neighboring energy subbands due to Rashba SOI is given by Eq. (6). This complex process can be treated by numerical or perturbation techniques.

In the case of an in-plane magnetic field, making an angle θ with the confinement direction, orbital magnetic effects are absent [12] and the oscillator frequency does not change. The Hamiltonian is given by [12,23,24]

$$H = \frac{p_x^2 + p_y^2}{2m} + \frac{1}{2}m\omega^2 x^2 + eFx + \frac{1}{2}g\mu_B\vec{\sigma}\cdot\vec{B} - \frac{\alpha}{\hbar}(p_x\sigma_y - p_y\sigma_x). \tag{13}$$

The dimensionless form of Eq. (13) reads

$$H = \left(a^\dagger a + \frac{1}{2}\right) + \frac{1}{2}(L_0 k)^2 - \frac{1}{2}(L_0 k_{EF})^2 + \frac{1}{2}\delta(\sigma_x\cos\theta + \sigma_y\sin\theta) + \gamma L_0 k\sigma_x - \frac{i\gamma}{\sqrt{2}}(a^\dagger - a)\sigma_y \tag{14}$$

with $\gamma = L_0/2L_{SO}$.

Because of the free motion of electrons along the y-direction, the eigenfunction may be written in the following form [15]:

$$\Psi_k(x, y) = e^{iky}\begin{pmatrix}\Phi_k^\uparrow(x)\\ \Phi_k^\downarrow(x)\end{pmatrix}. \tag{15}$$

regardless of the direction of the magnetic field. For a given k, the expectation value of the spin is given by

$$\langle\vec{\sigma}(k)\rangle = \int dx\, \Psi_k^\dagger\,\vec{\sigma}\,\Psi_k \tag{16}$$

In either case, the spinor of the wave function can be expanded in terms of the basis, $\psi_{nk\eta}^{(0)}$, of the Hamiltonian obtained after neglecting the Rashba intersubband coupling [22]. For perpendicular magnetic field, we have

$$\psi_{nk\eta}^{(0)}(x, y) = e^{iky}\varphi_n(x)\chi_\eta \tag{17}$$

with

$$\chi_\eta = \frac{1}{\sqrt{1 + c_\eta^2}}(c_\eta|\uparrow\rangle + |\downarrow\rangle). \tag{18}$$

Here, the index n corresponds to the harmonic oscillator wave function $\varphi_n(x)$, and the spin branch is characterized by the index $\eta = \pm 1$. The parameter c_η is given by

$$c_\eta = \frac{1}{\xi_1}\left(\delta + \eta\sqrt{\delta^2 + \xi_1^2}\right). \tag{19}$$

Such decomposition allows the spin to depend on k and the x-coordinate, which gives rise to spin textures [12,17]. We determine the expansion coefficients numerically. In our analysis, we confine ourselves to the weak Rashba regime, where for perpendicular magnetic field the term H' (Eq. (6)) is taken as perturbation with $L_0/L_{SO} < 1$. In this case, $\Psi_k(x)$ can be calculated using the matrix elements $\langle n'\eta'|H'|n\eta\rangle$ and the eigenvalues, $E_{nk\eta}^{(0)}$, corresponding to $\psi_{nk\eta}^{(0)}$. For instance, the first-order wave function correction for the lowest spin branch reads

$$\Psi_k^{(1)}(x) = a_1\varphi_1(x)\chi_{-1} + a_2\varphi_1(x)\chi_{+1}, \tag{20}$$

where

$$a_1 = -\frac{\xi_2 c_{-1}}{\Omega\left(1 + c_{-1}^2\right)}, \tag{21}$$

and

$$a_2 = \frac{-(\xi_2 + \xi_3)c_{-1} + (\xi_3 - \xi_2)c_{+1}}{2\left(\Omega + \sqrt{\delta^2 + \xi_1^2}\right)\sqrt{1 + c_{-1}^2}\sqrt{1 + c_{+1}^2}}. \tag{22}$$