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Fano–Rashba effect in quantum dots

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Abstract
We consider the electronic transport through a Rashba quantum dot coupled to ferromagnetic leads. We show that the interference of localized electron states with resonant electron states leads to the appearance of the Fano–Rashba effect. This effect occurs due to the interference of bound levels of spin-polarized electrons with the continuum of electronic states with an opposite spin polarization. We investigate this Fano–Rashba effect as a function of the applied magnetic field and Rashba spin–orbit coupling.

1. Introduction

Recently, there has been much interest in understanding the manner in which the unique properties of nanostructures may be exploited in spintronic devices, which utilize the spin degree of freedom of the electron as the basis of their operation [1–8]. The main challenge in the field of spintronics is to achieve the injection, modulation and detection of electron spin at the nanometer scale. In 1990, Datta and Das [1] proposed a spin transistor, based on the electron spin precession controlled by an external electric field via spin–orbit coupling. In this proposal, ferromagnetic contacts were used as the spin-polarized source and detector. A natural feature of these devices is the direct connection between their conductance and their quantum-mechanical transmission properties, which may allow their use as an all-electrical means for generating and detecting spin-polarized distributions of carriers.

Enforcing the analogy between quantum dots (QDs) and atomic systems, Fano [9] and Dicke effects [10] were found to be present in several QD configurations. On the other hand, Song *et al* [2] described how a spin filter may be achieved in open QD systems by exploiting Fano resonances that occur in their transmission characteristics. In a QD in which the spin degeneracy of a carrier is lifted, they showed that the Fano effect may be used as an effective means to generate spin polarization of transmitted carriers and that electrical detection of the resulting polarization should be possible.

The Rashba spin–orbit interaction arises from a structure inversion asymmetry resulting from the asymmetry of the in-plane confining potential in semiconductor heterostructures [11, 12]. This effect causes a spin splitting proportional to k. On the other hand, the Fano effect arises from the interference between a localized state and the continuum. In general, the condition for the Fano effect is the presence of two scattering channels at least: the discrete level and the continuum band. The Fano effect in electronic transport through a single-electron transistor allows us to alter the interference between the two paths by changing the voltages on various gates. Kobayashi *et al* [9] reported the first tunable Fano experiment in which a well-defined Fano system is realized in an Aharonov–Bohm ring with a QD embedded in one of its arms. Recently, Shelykh *et al* [13] studied the first Fano-type resonances due to the interaction of electron states with opposite spin orientation. They investigated the electronic transport through Datta and Das spin-modulator devices [1]. They show that the interfaces make the device behave as a Fabry–Perot cavity, so that Breit–Wigner resonances appear in the transmission coefficient. Additionally, interference of quantum-confined electron states with free electron states leads to the appearance of asymmetric Fano lineshapes. Moreover, Sánchez *et al* [14, 15], Wan *et al* [16] and López *et al* [17] predicted the occurrence of Fano lineshapes in a semiconductor quantum wire with local spin–orbit Rashba coupling.

In this work we investigate the electronic transport through a QD with ferromagnetic contacts considering the Rashba spin–orbit interaction. We show that the interference of localized electron states with free electron states leads to the appearance of the Fano–Rashba effect. This effect appears due to the interference of bound levels of spin-polarized electrons with the continuum of electronic states with an opposite spin polarization. We investigate this Fano–Rashba effect as a function of the system parameters.

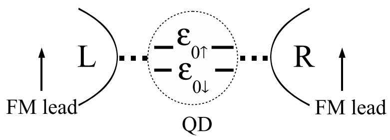

Figure 1. Schematic view of a QD connected to two full polarized ferromagnetic leads.

2. Model

The system under consideration is formed by one QD connected to two full spin-up polarized ferromagnetic leads, as shown schematically in figure 1. The full system is modeled by the Anderson Hamiltonian, namely $H = H_{\mathrm{L}} + H_{\mathrm{D}} + H_{\mathrm{I}}$ with

$$
\begin{aligned}
H_{\mathrm{L}} &= \sum_{i}\sum_{\sigma=\uparrow\downarrow} \varepsilon_\sigma\, c^\dagger_{i\sigma} c_{i\sigma} - v \sum_{\langle i\neq j\rangle}\sum_{\sigma=\uparrow\downarrow} \left(c^\dagger_{i\sigma} c_{j\sigma} + c^\dagger_{j\sigma} c_{i\sigma}\right) \\
H_{\mathrm{D}} &= \sum_{\sigma=\uparrow\downarrow} \varepsilon_{0\sigma} d^\dagger_\sigma d_\sigma + U n_{d\uparrow} n_{d\downarrow} \\
H_{\mathrm{I}} &= -V_0 \sum_{\sigma=\uparrow\downarrow} \left(d^\dagger_\sigma c_{1\sigma} + c^\dagger_{1\sigma} d_\sigma\right) \\
&\quad - V_0 \sum_{\sigma=\uparrow\downarrow} \left(d^\dagger_\sigma c_{-1\sigma} + c^\dagger_{-1\sigma} d_\sigma\right) \\
&\quad - \sum_{\sigma,\sigma'=\uparrow\downarrow} t_{\mathrm{so}}\, [\sigma_x]_{\sigma\sigma'} \left(d^\dagger_\sigma c_{1\sigma'} + c^\dagger_{1\sigma} d_{\sigma'}\right) \\
&\quad - \sum_{\sigma,\sigma'=\uparrow\downarrow} t_{\mathrm{so}}\, [\sigma_x]_{\sigma\sigma'} \left(d^\dagger_\sigma c_{-1\sigma'} + c^\dagger_{-1\sigma} d_{\sigma'}\right),
\end{aligned} \tag{1}
$$

where $c^\dagger_{i\sigma}$ is the creation operator of an electron at site i of the leads in the σ spin state ($\sigma = \uparrow, \downarrow$) and $d^\dagger_\sigma$ is the corresponding operator of an electron with spin σ of the QD. Moreover $n_{d\sigma} = d^\dagger_\sigma d_\sigma$ and V_0 is the coupling between the QD and the leads. U is the Coulomb coupling and it will be neglected hereafter. The potential of the wire is taken to be zero and the hopping in the wire is $-v$. Furthermore we set the site energies in the leads and energy levels of the QD as $\varepsilon_\sigma = 2 + \Delta[\sigma_z]_{\sigma\sigma}$ and $\varepsilon_{0\sigma} = \varepsilon_0 + \mu B[\sigma_z]_{\sigma\sigma}$, respectively, where $\vec{\sigma} = (\sigma_x, \sigma_y, \sigma_z)$ is the Pauli matrix vector, Δ is the ferromagnetic energy, B is the magnetic field in the quantum dot and ε_0 is the energy level in the dot without a magnetic field.

The stationary states of the Hamiltonian H can be written as

$$
|\psi_\sigma\rangle = \sum_{j=-\infty,\, j\neq 0}^{\infty} a_{j\sigma} |j\rangle + b_\sigma |0\rangle, \tag{2}
$$

where $a_{j\sigma}$ and b_σ are the probability amplitudes to find the electron at the site j or at the QD, respectively, with energy $\omega = \varepsilon_\uparrow - 2v\cos k$ or $\omega = \varepsilon_\downarrow - 2v\cosh\kappa$, where $\varepsilon_\uparrow = 2 + \Delta$ and $\varepsilon_\uparrow = 2 - \Delta$. These amplitudes obey the following linear difference equations:

$$
\begin{aligned}
(\omega - \varepsilon_\sigma) a_{j\sigma} &= -v(a_{j+1,\sigma} + a_{j-1,\sigma}), \\
&\quad j \neq -1, 0, 1, \\
(\omega - \varepsilon_\sigma) a_{-1\sigma} &= -v a_{-2\sigma} - V_0 b_\sigma - t_{\mathrm{so}}\, [\sigma_x]_{\sigma\overline{\sigma}}\, b_{\overline{\sigma}}, \\
(\omega - \varepsilon_\sigma) a_{1\sigma} &= -v a_{2\sigma} - V_0 b_\sigma - t_{\mathrm{so}}\, [\sigma_x]_{\sigma\overline{\sigma}}\, b_{\overline{\sigma}}, \\
(\omega - \widetilde{\varepsilon}_{0\sigma}) b_\sigma &= -V_0 (a_{1,\sigma} + a_{-1,\sigma}) \\
&\quad - t_{\mathrm{so}}\, [\sigma_x]_{\sigma\overline{\sigma}}\, (a_{1,\overline{\sigma}} + a_{-1,\overline{\sigma}})
\end{aligned} \tag{3}
$$

where $\widetilde{\varepsilon}_{0\sigma}$ is the renormalized energy level of the QD with spin σ:

$$
\begin{aligned}
\widetilde{\varepsilon}_{0\uparrow} &= \varepsilon_0 + \mu B \\
\widetilde{\varepsilon}_{0\downarrow} &= \varepsilon_0 - \mu B.
\end{aligned} \tag{4}
$$

In order to study the solutions of the above equations, we assume that spin-up electrons are described by a plane wave with unitary incident amplitude, r and t being the reflection and transmission amplitudes. Thus we get

$$
\begin{aligned}
a_{j\uparrow} &= \mathrm{e}^{\mathrm{i}kj} + r\mathrm{e}^{-\mathrm{i}kj}, & j &< 0, \\
a_{j\uparrow} &= t\mathrm{e}^{\mathrm{i}kj}, & j &> 0 \\
a_{j\downarrow} &= A\mathrm{e}^{\kappa j}, & j &< 0 \\
a_{j\downarrow} &= B\mathrm{e}^{-\kappa j}, & j &> 0.
\end{aligned} \tag{5}
$$

Inserting this solution in the equation of motion, we get an inhomogeneous system of linear equations for the unknowns t, r, A and B, leading to the following expression for the transmission amplitude t:

$$
\begin{aligned}
t &= 2\mathrm{i}\alpha_- \sin k[(\omega - \widetilde{\varepsilon}_{0\downarrow}) V_0^2 \\
&\quad - (\omega - \widetilde{\varepsilon}_{0\uparrow}) t_{\mathrm{so}}^2 - 2\alpha_-^2 v \mathrm{e}^{-\kappa}]/[V_0^2(\omega - \widetilde{\varepsilon}_{0\uparrow} + 2\alpha_- \mathrm{e}^{\mathrm{i}k}) \\
&\quad \times (\omega - \widetilde{\varepsilon}_{0\downarrow} - 2\alpha_- \mathrm{e}^{-\kappa}) + t_{\mathrm{so}}^2(\omega - \widetilde{\varepsilon}_{0\downarrow} + 2\alpha_- \mathrm{e}^{\mathrm{i}k}) \\
&\quad \times (-\omega + \widetilde{\varepsilon}_{0\uparrow} + 2\alpha_- \mathrm{e}^{-\kappa})]
\end{aligned} \tag{6}
$$

with $\alpha_\pm \equiv (V_0^2 \pm t_{\mathrm{so}}^2)/v$.

Notice that the above expression reduces to a single resonance when the Rashba spin–orbit coupling is neglected ($t_{\mathrm{so}} = 0$), namely

$$
t = \frac{\mathrm{i}\Gamma_0}{(\omega - \widetilde{\varepsilon}_{0\uparrow} - \Lambda_0) + \mathrm{i}\Gamma_0} \tag{7}
$$

where $\Gamma_0 = 2V_0^2 \sin k$ is the width of the resonance centered at $\widetilde{\varepsilon}_{0\uparrow} + \Lambda_0$ and $\Lambda_0 = (2V_0^2/v^2)\Delta$ is the shift due to the coupling of the QD with the leads. For $V_0 = 0$ the transmission amplitude is also reduced to a single resonance with width $\Gamma_{\mathrm{so}} = 2t_{\mathrm{so}}^2 \sin k$ and centered at $\widetilde{\varepsilon}_{0\downarrow} + \Lambda_{\mathrm{so}}$, where $\Lambda_{\mathrm{so}} = (2t_{\mathrm{so}}^2/v^2)\Delta$ is the shift due to the spin–orbit coupling of the QD with the leads

$$
t = \frac{\mathrm{i}\Gamma_{\mathrm{so}}}{(\omega - \widetilde{\varepsilon}_{0\downarrow} - \Lambda_{\mathrm{so}}) - \mathrm{i}\Gamma_{\mathrm{so}}}. \tag{8}
$$

For this expression we can conclude that the Rashba spin–orbit coupling opens a new channel for tunneling through the QD. Finally, the conductance for spin-up electrons is calculated by means of the Landauer formalism at zero temperature:

$$
\mathcal{G}_\uparrow = \frac{\mathrm{e}^2}{h} T(E_{\mathrm{F}}). \tag{9}
$$

3. Result

Evaluating the transmission probability at $\omega = E_{\mathrm{F}} = 0$ we obtain the spin-dependent conductance

$$
\begin{aligned}
\mathcal{G}_\uparrow &= \frac{\mathrm{e}^2}{h} \\
&\quad \times \frac{4\sin^2 k_{\mathrm{F}}[\alpha_-(\varepsilon_0 + \xi_-) - \mu B \alpha_+]^2}{|(\varepsilon_0 - \xi_-)(\varepsilon_0 - \xi_+) + 2\alpha_+ \mu B(\mathrm{e}^{\mathrm{i}k_{\mathrm{F}}} + \mathrm{e}^{-\kappa_{\mathrm{F}}}) - \mu^2 B^2|^2}
\end{aligned} \tag{10}
$$

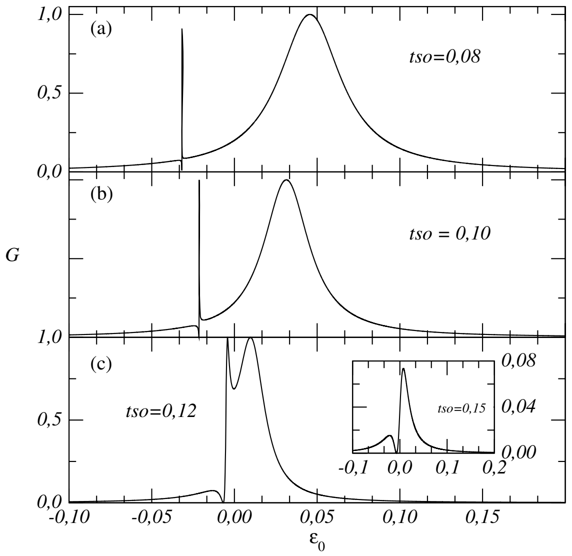

Figure 2. Spin-dependent conductance as a function of the gate voltage ε_0 for different values of the spin–orbit coupling, when $\mu B/2v = 0.001$.

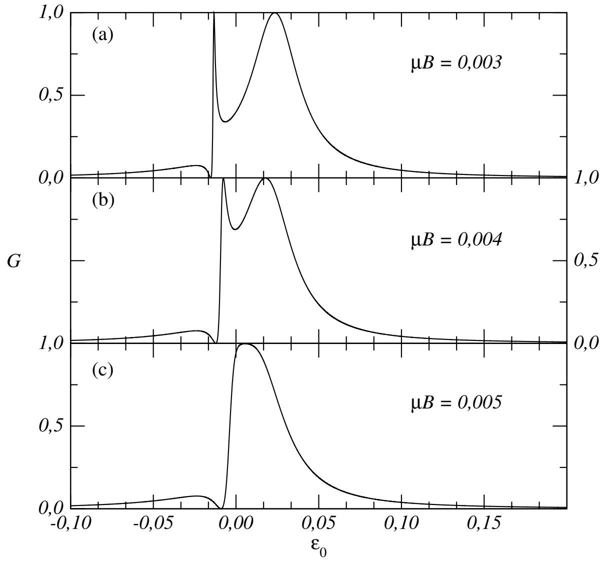

Figure 3. Spin-dependent conductance as a function of the gate voltage for various values of the magnetic field in the dot, when $t_{so}/2v = 0.1$.

where $\xi_+ = 2\alpha_- \cos k_F, \xi_- = 2e^{-\kappa_F}\alpha_-$, $k_F = \cos^{-1}(1 - \Delta/2v)$ and $\kappa_F = \cosh^{-1}(1 + \Delta/2v)$. In what follows we present results for the conductance for $V_0/2v = 0.14$ and $\Delta/2v = 0.1$.

Figure 2 displays the spin-dependent linear conductance versus the gate voltage ε_0 for different values of the spin–orbit coupling at fixed magnetic field. As expected, the linear conductance shows two antiresonances and one Fano antiresonance as a function of the Fermi energy. The antiresonance in the conductance occurs at $\varepsilon_0 = \xi_- + \mu B\alpha_+/\alpha_-$. For small values of magnetic field ($\mu B/v \ll 1$) the conductance of the system can be written approximately as a convolution of a Fano lineshape and a Breit–Wigner lineshape. This is

$$\mathcal{G}_\uparrow \approx \frac{e^2}{h}\frac{(\epsilon_- + q)^2}{\epsilon_-^2 + 1}\frac{1}{\epsilon_+^2 + 1} \tag{11}$$

where $\epsilon_\pm$ are the detuning parameters measuring the energy ε_0 from the resonance centers and normalized by the resonance half-width [$\epsilon_- = (\varepsilon_0 - \xi_-)/\mu B, \epsilon_+ = (\varepsilon_0 - \xi_+)/(2\alpha_- \sin k_F)$] and $q = \alpha_+/\alpha_-$ is the Fano parameter characterizing the lineshape asymmetry.

The spin-dependent conductance versus the gate voltage is displayed in figure 3 for various values of the magnetic field in the QD. We observe that the magnetic field modulates the spin-dependent conductance, allowing for a fine tuning of the system response. Remarkably, as the magnetic field increases the two resonances merge into a single, broad resonance.

Figure 4 displays the behavior of the spin-dependent conductance as a function of the spin–orbit coupling parameter t_{so} at $\varepsilon_0 = 0$ for different values of the magnetic field. We can see that the conductance can also be controlled by selecting a suitable value of t_{so}. Notice that, as t_{so} is varied, the conductance can pass from perfect transmission to perfect reflection. However, t_{so} is dependent on the semiconducting materials used to fabricate the device. Therefore, in contrast

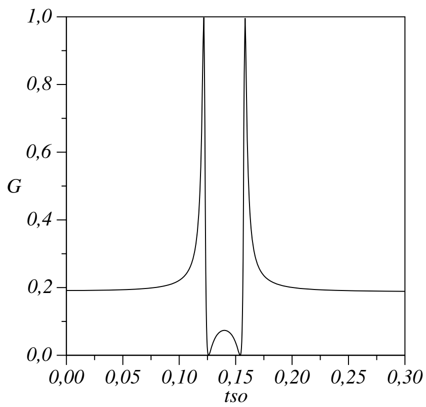

Figure 4. Spin-dependent conductance as a function of the spin–orbit coupling t_{so}, when $\mu B/2v = 0.001$.

to the control gained by changing the gate voltage mentioned in the previous paragraph, one has less control on the spin filtering capabilities of the device by choosing different semiconductors.

Spin-dependent conductance at $\varepsilon_0 = 0$ as a function of the magnetic field in the QD is displayed in figure 5. The conductance as a function of the magnetic field also shows two resonances and one antiresonance. The value of the magnetic field at the antiresonance is $\mu B = 2e^{-\kappa_F}(V_0^2 - t_{s0}^2)^2/(V_0^2 + t_{s0}^2)$.

The Rashba spin–orbit coupling opens a new channel to the conduction that interferes with the direct channel, the key ingredient to produce the destructive interference of the Fano effect. Figure 6 displays a schematic view of the electron tunneling through the QD. The electron with spin up can tunnel directly through the level $\widetilde{\varepsilon}_{0\uparrow}$ without spin-flip processes or can also tunnel indirectly through the level $\widetilde{\varepsilon}_{0\uparrow}$ with two spin-flip processes. The interference between the two tunneling paths gives rise to the Fano–Rashba effect presented in this paper.

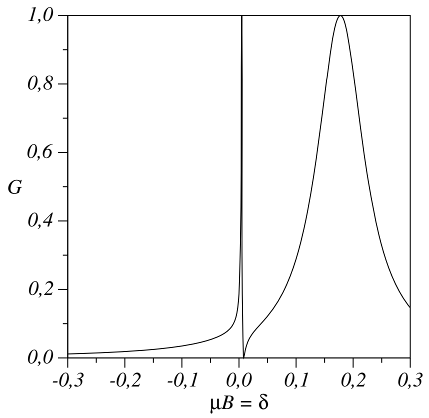

Figure 5. Spin-dependent conductance as a function of the magnetic field in the QD, when $t_{so}/2v = 0.1$.

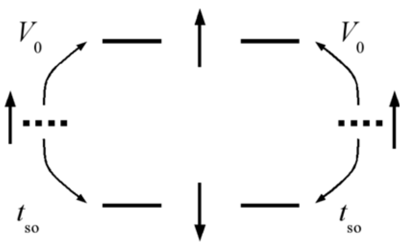

Figure 6. Schematic view of the electron tunneling through the QD.

Finally, we briefly discuss a more realistic situation when the ferromagnetic leads are not fully polarized, as we assumed in this work. In the case of partially polarized leads, a small fraction of electrons with *undesired* spin polarization will be injected in the QD. Since the equations of motion (3) are linear and they lack terms responsible for spin mixing, one could expect a gradual destruction of the Fano–Rashba antiresonance that would be proportional to the loss of the polarization.

4. Summary

Here we investigated the electronic transport through a Rashba QD with ferromagnetic contacts. We have shown that the interference of localized electron states with resonant electron states leads to the appearance of the so-called Fano–Rashba effect. This effect arises from the interference of bound levels of spin-polarized electrons in the QD with the resonant states of opposite spin polarization. We found that the Fano–Rashba effect holds even in the presence of the electron–electron interaction in the QD.

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