Physics Letters A 380 (2016) 1672–1676

Contents lists available at ScienceDirect

Physics Letters A

www.elsevier.com/locate/pla

# Gate-voltage control of equal-spin Andreev reflection in half-metal/semiconductor/superconductor junctions

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ARTICLE INFO

*Article history:*
Received 31 December 2015
Received in revised form 27 February 2016
Accepted 29 February 2016
Available online 2 March 2016
Communicated by L. Ghivelder

*Keywords:*
Andreev reflection
Superconductor
Half-metal

ABSTRACT

With the Blonder–Tinkham–Klapwijk (BTK) approach, we investigate conductance spectrum in Ferromagnet/Semiconductor/Superconductor (FM/Sm/SC) double tunnel junctions where strong Rashba spin–orbit interaction (RSOI) is taken into account in semiconductors. For the half-metal limit, we find that the in-gap conductance becomes finite except at zero voltage when inserting a ferromagnetic insulator (FI) at the Sm/SC interface, which means that the appearance of a long-range triplet states in the half-metal. This is because of the emergence of the unconventional equal-spin Andreev reflection (ESAR). When the FI locates at the FM/Sm interface, however, we find the vanishing in-gap conductance due to the absence of the ESAR. Moreover, the non-zero in-gap conductance shows a nonmonotonic dependence on RSOI which can be controlled by applying an external gate voltage. Our results can be used to generate and manipulate the long-range spin triplet correlation in the nascent field of superconducting spintronics.


## 1. Introduction

Over the last two decades, the proximity effect in hybrid junctions between the superconductor (SC) and ferromagnet (FM) has received much attention [1–14]. In these junctions, the proximity effect would decay rapidly in the FM since the SC with antiparallel spins and FM with parallel spins are antagonistic [15]. When the magnetization of the FM layer is inhomogeneous, however, the long-range proximity effect is observed due to the triplet pair generation in FM through the spin mixing and spin rotation at the FM/SC interface. For example, long-ranged supercurrents which imply triplet pairing have received considerable theoretical [16–25] and experimental [26–30] attention. Microscopically, the superconducting proximity effect is governed by AR, an electron incident on the interface from the normal state material is phase-coherently reflected a hole of opposite spin and momentum to the incident electron, and vice versa [31]. The spin-triplet pairing effects require the formation of ESAR in which the electrons can be Andreev reflected into holes in the same spin band. There are a large number of works on how to create the ESAR at spin-active FM/SC interface including spin flip process, especially, when FM becomes HM that only conduction electrons with equal spin can be paired [32,5,33–37,26,8,38–40].

On the other hand, superconducting spintronics, in analogy with spintronic applications [41,42], have received a boost which aims to use superconductors as active components in spintronic devices [38,43]. One of the topics of the intense research in superconducting spintronics is the effect of spin–orbit (SO) coupling on the charge and spin transports [44–50]. For example, a long-range spin-triplet helix is predicted when the strengths between Rashba and Dresselhaus SO couplings are equal [48]. Jacobsen et al. study the giant triplet proximity effect in $\pi$-biased Josephson junctions with spin–orbit coupling [51]. Furthermore, many efforts have been devoted to the effect of the interfacial SO coupling on the charge and spin transport in the non-magnetic junctions [52–54], especially, including superconductors as active components [55–57]. For instance, Bergeret et al. [50] found that a normal metal with SO coupling can be used as the source of long-range triplet proximity effect in FM/SC hybrid structures. More recently, the study of interfacial SO coupling has been extended to the FM/SC junctions [8] where the magnetoanisotropic Andreev reflection is found in the presence of Rashba and Dresselhaus interfacial SO coupling.

In this article, the tunneling conductance spectrum in FM/Sm/SC double tunnel junctions is studied by means of BTK approach [58] where strong RSOI is taken into account in the semiconductors. For the half-metal limit, we find that the in-gap conductance becomes

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http://dx.doi.org/10.1016/j.physleta.2016.02.053

finite except at zero voltage when inserting a FI at the Sm/SC interface, which means that the appearance of a long-range triplet states in the half-metal. This is because of the emergence of the unconventional ESAR. When the FI locates at the FM/Sm interface, however, we find the vanishing in-gap conductance due to the absence of the ESAR. Moreover, the non-zero in-gap conductance shows a nonmonotonic dependence on RSOI which can be controlled by applying an external gate voltage. Our results are useful to generate and manipulate the singlet–triplet conversion in the nascent field of superconducting spintronics.

This paper is organized as follows. In Sec. 2 we introduce the model and the mode-matching approach. In Sec. 3 we present our numerical results and discussions. The final Sec. 4 is devoted to a brief summary.

## 2. Model and formula

We consider a ballistic FM/Sm/SC double-tunnel junctions where strong RSOI is taken into account in semiconductors. FM and SC are separated from the central Sm by two thin FI layers which can induce the spin-filtering effect. The scattering Hamiltonian of two thin FI is defined by the following expression

$$H_{FI}(x) = (U_1\sigma_0 + U_{M1}\sigma_z)\,\delta(x) + (U_2\sigma_0 + U_{M2}\sigma_z)\,\delta(x-L) \quad (1)$$

Here $\delta(x)$, $U_{1(2)}$ and $U_{M1(2)}$ are $\delta$ function, the barrier amplitude and the exchange potential in the FI [59], respectively.

The effective Hamiltonian (Bogoliubov–de Gennes equation) for quasiparticle states $\Psi(r) = (u_\uparrow(r), u_\downarrow(r), v_\uparrow(r), v_\downarrow(r))^T$ with energy $E$ is given by [60]

$$\begin{pmatrix} H_+ + H_{FI}(x) & i\sigma_y\Delta(x) \\ -i\sigma_y\Delta^*(x) & -H_- - H_{FI}(x) \end{pmatrix}\Psi(\mathbf{r}) = E\Psi(\mathbf{r}) \quad (2)$$

with

$$H_\pm = \left[-\frac{\hbar^2}{2}\nabla\left[\frac{1}{m^*(x)}\right]\nabla - E_F(x)\right]\sigma_0 - h_0(x)\sigma_z + \frac{\lambda_{so}(x)}{\hbar}\left(\pm\sigma_x k_y - \sigma_y k_x\right), \quad (3)$$

with

$$\begin{aligned}
\Delta(x) &= \Delta\Theta(x-L)\\
m^*(x) &= m\Theta(-x) + m'\Theta(x)\Theta(L-x) + m\Theta(x-L),\\
E_F(x) &= E_F\Theta(-x) + (E_F-\delta E_c)\Theta(x)\Theta(L-x)\\
&\quad + E_F\Theta(x-L),\\
\lambda_{so}(x) &= \lambda_{so}\Theta(x)\Theta(L-x),\\
h_0(x) &= h_0\Theta(-x).
\end{aligned} \quad (4)$$

Here, $\Delta$ is the superconducting pair potential, $\lambda_{so}$ is the RSOI constant, $h_0$ is the exchange energy. $m(x)$ and $E_F(x)$ are the position-dependent effective mass and Fermi energy, respectively. The Fermi energy $E_F - \delta E_c$ in the Sm is much smaller than $E_F$ in FM and SC. $\sigma_{x,y,z}$ is the Pauli matrix in the spin space, $\sigma_0$ is the unit matrix, $\Theta(x)$ is the step function and $k = (k_x, k_y)$ is the momentum.

Notice that since the Fermi wave vector $k_F$ in the Sm is much smaller than the Fermi wave vector $k_F$ of the left FM and right SC, the incident angle $\varphi_{FM}$ in the FM has to be virtually zero for transmission according to the translational symmetry along the $y$-axis direction [$k_F^{FM}\sin(\varphi_{FM}) = k_F^{Sm}\sin(\varphi_{Sm}) = k_F^{SC}\sin(\varphi_{SC})$]. Thus, we restrict ourselves to the one-dimensional case in the following.

Consider a beam of spin-up electrons incident from the FM, the general solution of Eq. (2) is of the form aligned

$$\begin{aligned}
\Psi_{\mathrm{I}} &= \left(e^{iq_{e\uparrow}x} + b_\uparrow e^{-iq_{e\uparrow}x}\right)e^{iq_{e\uparrow}x}(1,0,0,0)^T\\
&\quad + b_\downarrow e^{-iq_{e\downarrow}x}(0,1,0,0)^T + a_\uparrow e^{iq_{h\uparrow}x}(0,0,1,0)^T\\
&\quad + a_\downarrow e^{iq_{h\downarrow}x}(0,0,0,1)^T,
\end{aligned} \quad (5)$$

for $x < 0$;

$$\begin{aligned}
\Psi_{\mathrm{II}} &= f_1 e^{ik_{1+}x}(-i,1,0,0)^T + f_2 e^{ik_{1-}x}(i,1,0,0)^T\\
&\quad + f_3 e^{-ik_{1+}x}(i,1,0,0)^T + f_4 e^{-ik_{1-}x}(-i,1,0,0)^T\\
&\quad + f_5 e^{ik_{2+}x}(0,0,-i,1)^T + f_6 e^{ik_{2-}x}(0,0,i,1)^T\\
&\quad + f_7 e^{-ik_{2+}x}(0,0,i,1)^T + f_8 e^{-ik_{2-}x}(0,0,-i,1)^T,
\end{aligned} \quad (6)$$

for $0 < x < L$;

$$\begin{aligned}
\Psi_{\mathrm{III}} &= c_1 e^{ik_{s+}x}(u,0,0,v)^T + c_2 e^{ik_{s+}x}(0,u,-v,0)^T\\
&\quad + c_3 e^{-ik_{s-}x}(v,0,0,u)^T + c_4 e^{-ik_{s-}x}(0,v,-u,0)^T,
\end{aligned} \quad (7)$$

for $x > L$, with

$$\begin{aligned}
q_{e\uparrow(\downarrow)} &= \sqrt{\frac{2m}{\hbar^2}(E_F + E \pm h_0)},\quad q_{h\uparrow(\downarrow)} = \sqrt{\frac{2m}{\hbar^2}(E_F - E \pm h_0)},\\
k_{1\pm} &= \sqrt{\frac{2m'}{\hbar^2}(E_F - \delta E_c + E) + \left(\frac{m'\lambda_{so}}{\hbar^2}\right)^2} \pm \frac{m'\lambda_{so}}{\hbar^2},\\
k_{2\pm} &= \sqrt{\frac{2m'}{\hbar^2}(E_F - \delta E_c - E) + \left(\frac{m'\lambda_{so}}{\hbar^2}\right)^2} \pm \frac{m'\lambda_{so}}{\hbar^2},\\
k_{s\pm} &= \sqrt{\frac{2m}{\hbar^2}\left(E_F \pm \sqrt{E^2-\Delta^2}\right)}.
\end{aligned} \quad (8)$$

Here, the coherence factors in the superconducting region are $u = \sqrt{\frac{1}{2}\left(1 + \frac{\sqrt{E^2-\Delta^2}}{E}\right)}$ and $v = \sqrt{\frac{1}{2}\left(1 - \frac{\sqrt{E^2-\Delta^2}}{E}\right)}$.

The normal reflection $b_\uparrow$, the normal reflection with spin flip $b_\downarrow$, the usual AR $a_\downarrow$, the equal-spin AR $a_\uparrow$ can be determined by imposing the following matching conditions at the left and right interfaces

$$\begin{aligned}
&\Psi_{\mathrm{I}}(0_-) = \Psi_{\mathrm{II}}(0_+),\\
&\Psi_{\mathrm{II}}(L_-) = \Psi_{\mathrm{III}}(L_+),\\
&\hat{v}_x\Psi_{\mathrm{II}}(0_+) - \hat{v}_x\Psi_{\mathrm{I}}(0_-) = \frac{2\hbar k_F}{im}(Z_1\tau_1 + Z_2\tau_2)\Psi_{\mathrm{I}}(0_-),\\
&\hat{v}_x\Psi_{\mathrm{III}}(L_+) - \hat{v}_x\Psi_{\mathrm{II}}(L_-) = \frac{2\hbar k_F}{im}(Z_3\tau_1 + Z_4\tau_2)\Psi_{\mathrm{III}}(L_+),\\
&\tau_1 = \begin{pmatrix} 1 & 0 & 0 & 0\\ 0 & 1 & 0 & 0\\ 0 & 0 & -1 & 0\\ 0 & 0 & 0 & -1 \end{pmatrix},\ \tau_2 = \begin{pmatrix} -1 & 0 & 0 & 0\\ 0 & 1 & 0 & 0\\ 0 & 0 & 1 & 0\\ 0 & 0 & 0 & -1 \end{pmatrix},
\end{aligned} \quad (9)$$

where $Z_{1(3)} = U_{1(2)}/\hbar v_F$ and $Z_{2(4)} = U_{M1(2)}/\hbar v_F$ are dimensionless parameters describing the charge scattering and the magnetic scattering, respectively, with $v_F$ as the Fermi velocity.

The velocity operator $\hat{v}_x$ in the $x$-direction is given by [45]

$$\hat{v}_x = \frac{\partial H}{\hbar\partial k} = \begin{pmatrix} \frac{\hbar}{im(x)}\frac{\partial}{\partial_x} & i\frac{\lambda_{so}(x)}{\hbar} & 0 & 0\\ -i\frac{\lambda_{so}}{\hbar}(x) & \frac{\hbar}{im(x)}\frac{\partial}{\partial_x} & 0 & 0\\ 0 & 0 & -\frac{\hbar}{im(x)}\frac{\partial}{\partial_x} & -i\frac{\lambda_{so}}{\hbar}(x)\\ 0 & 0 & i\frac{\lambda_{so}}{\hbar}(x) & -\frac{\hbar}{im(x)}\frac{\partial}{\partial_x} \end{pmatrix}. \quad (10)$$