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Spin-polarized-current switching mediated by Majorana bound states

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ABSTRACT

We present the study of transport properties of a superconducting wire with strong Rashba spin-orbit coupling for different orientations of an external magnetic field. Using the nonequilibrium Green's functions in the tight-binding approach the crucial impact of the relative alignment of lead magnetization and the Majorana bound state (MBS) spin polarization on the low-bias conductance and shot noise is presented. Depending on this factor the transport regime can effectively vary from symmetric to extremely asymmetric. In the last situation the suppression of MBS-assisted conductance results in current-switch effect allowing the electrical detection of the Majorana fermions.

1. Introduction

The pursuit of experimental observation of a Majorana fermion, originally started in the particle physics, continued then in the solid-state systems [1]. It was shown in the Kitaev model [2] that the emergent quasiparticles appearing in superconducting (SC) system possess the self-Hermitian property of the Majorana fermions. Majorana bound states (MBSs) in low-dimensional structures are spatially separated. The last feature opens a way to utilize the MBSs as the building blocks of fault-tolerant quantum computers [3]. Since the MBSs obey non-Abelian statistics [4] an MBS-based qubit can be manipulated by braiding operations [5].

Among different systems where MBSs were predicted semiconducting wires with strong spin-orbit coupling (SOC) and in proximity to an s-wave SC [6,7] attract considerable attention since the corresponding experimental proofs were provided [8]. Under the influence of an external magnetic field an effective p-wave pairing is realized in the wire and two MBSs appear at its opposite ends [6,7]. In addition to studying the fundamental properties of MBS in the wires, a number of applications was already proposed, e.g. memory cell [9], current switch [10], rectifier and Cooper pair splitter [11,12].

The tunneling spectroscopy measurements reveal the zero-bias conductance peak (ZBP) indicating resonant Andreev transport processes through the zero-energy MBS [13,8]. Furthermore, the shot noise provides supplementary information about the MBSs [14,15]. It is essential that the Majorana nature of the ZBP is not a single interpretation [16]. Therefore, it is necessary to analyze supplementary information concerning quantum transport via the MBSs. One of such properties is an electronic spin polarization in the MBS [17] which we call here the MBS spin polarization. In the present article the influence of MBS spin polarization on nonlocal quantum transport is analyzed. Studying spin-polarized transport we show that the low-bias conductance brings the features allowing both to detect the MBSs and employ that unusual behavior in electronic applications such as a current switch. In Ref. [10] the current switch assisted by the MBSs was already studied. The effect was induced by the combination of quantum phase transition and destructive interference. In contrast, our proposal is based on the features of MBS spin polarization as it is shown in Fig. 1. While the local character of the magnetic field in the proposed scheme can present a certain obstacle for its experimental realization we suggest to apply a magnetic gate to overcome this problem. Note that the scheme containing ferromagnets for the MBS electrical detection was also proposed in Ref. [18]. However, their role differs in principle from the one considered here. In our situation the presence of ferromagnets is not vital to induce the MBSs but is necessary to probe nonlocal transport.

2. The model Hamiltonian

Let us consider a nanowire with the strong Rashba SOC deposited on a grounded s-type SC (see Fig. 1). Hereinafter we name it a 'superconducting wire' taking into account an induced SC pairing in the wire characterized by a parameter Δ. An external magnetic field $\mathbf{B} = (B_x, 0, -B_z) = B(\cos\theta, 0, -\sin\theta)$ is assumed to be oriented at an arbitrary angle in the plane perpendicular to the Rashba field $\mathbf{B}_{SO} \| y$. We suppose in electronic applications the direction and

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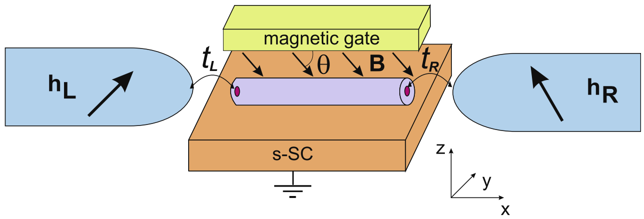

Fig. 1. The superconducting wire between ferromagnetic leads. Red circles indicate the MBSs.

amplitude of magnetic field can be manipulated by a 'magnetic gate', e.g. as it was proposed to manipulate domain-wall motion in [19].

The microscopic Hamiltonian of SC wire is

$$\hat{H}_W = \sum_{j=1}^{N}\left[\sum_{\sigma}\xi_\sigma a_{j\sigma}^{+}a_{j\sigma} + \Delta a_{j\uparrow}a_{j\downarrow} - V_x a_{j\uparrow}^{+}a_{j\downarrow} + H.c.\right] - \sum_{\sigma;j=1}^{N-1}\left[\frac{t}{2}a_{j\sigma}^{+}a_{j+1,\sigma} + \frac{\alpha}{2}\sigma a_{j\sigma}^{+}a_{j+1,\overline{\sigma}} + H.c.\right], \quad (1)$$

where $a_{j\sigma}$ – an electron annihilation operator on jth site of the wire with spin σ; $\xi_\sigma = t + \sigma V_z - \mu$ – an on-site energy of the electron with spin σ taking into account the Zeeman component V_z; μ – a chemical potential of the system; $V_{x(z)} = \mu_B B_{x(z)}$ – x- and z-components of the Zeeman energy; t – a nearest neighbor hopping parameter; α – an intensity of the Rashba SOC.

The SC wire is situated between the ferromagnetic leads characterized by the Stoner-type Hamiltonians:

$$\hat{H}_i = \sum_{k\sigma}\{[\xi_k - eV_i - \sigma M_i\cos\theta_i]c_{k\sigma}^{+}c_{k\sigma} - M_i\sin\theta_i c_{k\sigma}^{+}c_{k\overline{\sigma}}\}, i = L, R, \quad (2)$$

where $c_{k\sigma}^{+}$ – an electron creation operator in ith lead with a wave vector k, spin σ and an energy $\xi_k = \epsilon_k - \mu$; $M_i = \frac{1}{2}g\mu_B h_i$ – an energy of the ith lead magnetization $\mathbf{h}_i$; θ_i – an angle between $\mathbf{h}_i$ and z axis in the xz plane; $\sigma = \pm 1$ or $\uparrow,\downarrow$. The bias voltage $V_{L(R)} = \pm V/2$ is applied to the left (right) lead.

The interaction between the leads and the SC wire is given by a standard tunnel Hamiltonian,

$$\hat{H}_T = t_L\sum_{k\sigma}c_{k\sigma}^{+}a_{1\sigma} + t_R\sum_{p\sigma}c_{p\sigma}^{+}a_{N\sigma} + H.c., \quad (3)$$

where $t_{L(R)}$ – a tunnel parameter between the left (right) lead and the wire. Thus, the total Hamiltonian of the system is $\hat{H} = \hat{H}_L + \hat{H}_R + \hat{H}_T + \hat{H}_W$.

3. Current and noise in terms of the nonequilibrium Green's functions

To calculate spin-dependent transport properties of the SC wire we employ the nonequilibrium Green's functions [20,21] in the spin⊗Nambu space [22]. After some manipulations the general expression describing current in the ith lead can be written as following

$$\langle \hat{I}_i\rangle = e\int_{-\infty}^{+\infty}\frac{d\omega}{\pi}Tr\left[Re\left\{\hat{\sigma}\left(\hat{\Sigma}_i^r\hat{G}_{i,i}^{<} + \hat{\Sigma}_i^{<}\hat{G}_{i,i}^{a}\right)\right\}\right], \quad (4)$$

where $\hat{\sigma} = diag(1,-1,1,-1)$ accounts for the electron and hole transport channels; $\hat{\Sigma}_i^{r,<} = \hat{T}_i^{+}\hat{g}_k^{r,<}\hat{T}_i$ – the Fourier transform of matrix retarded/lesser self-energy function describing the influence of ith lead on the wire; $\hat{g}_k^{r,<}$ – the Fourier transform of matrix retarded/lesser of free-particle Green's function of ith lead. In this work we consider half-metallic leads (e.g. NiMnSb or CrO_2 [23]) where $\theta_i = 0$. Consequently, the time-dependent tunnel coupling matrices are $\hat{T}_i(t) = \frac{t_i}{2}diag\left(e^{-ieV_it}, -e^{ieV_it}, e^{-ieV_it}, -e^{ieV_it}\right)$. $\hat{G}_{i,i}^{<,a} = \hat{P}_i\hat{G}_W^{<,a}\hat{P}_i^{+}$ – the Fourier transforms of the lesser/advanced Green's functions of the wire which are projected on the subspace of site tunnel-coupled with the ith lead. To obtain $\hat{G}_{i,i}^{<,a}$ the projection operators $\hat{P}_L = \left(\hat{I}\hat{O}\right), \hat{P}_R = \left(\hat{O}\hat{I}\right)$ are used, where $\hat{I}$ - 4×4 unit matrix; $\hat{O}$ - $4\times 4N-4$ zero block [24].

The nonequilibrium Green's functions can be obtained from the Dyson and Keldysh equations,

$$\hat{G}_W^r = \left(\omega - \hat{h}_W - \hat{\Sigma}^r\right)^{-1}, \hat{G}_W^a = \left(\hat{G}_W^r\right)^{+}, \quad (5)$$

$$\hat{G}_W^{\lessgtr} = \hat{G}_W^r\hat{\Sigma}^{\lessgtr}\hat{G}_W^a. \quad (6)$$

In expressions (5), (6) $\hat{h}_W$ is the matrix of Hamiltonian (1) in the spin⊗Nambu space. The total self-energy function of the system is $\hat{\Sigma}^n = \hat{P}_L^{+}\hat{\Sigma}_L^n\hat{P}_L + \hat{P}_R^{+}\hat{\Sigma}_R^n\hat{P}_R$. The ith lead components are $\Sigma_i^r = -\frac{i}{2}\Gamma_i, \hat{\Sigma}_i^{<} = (\Sigma_i^a - \Sigma_i^r)\hat{F}_i$, where $\Gamma_i = 2\pi\left(\frac{t_i}{2}\right)^2\rho_i$ – the coupling strength between the wire and the majority subband of ith lead; ρ_i – the DOS of majority subband of ith lead. In the calculations below the leads are treated in the wide-band limit that results in $\Gamma_i = const$. Finally, $\hat{F}_i = diag(n_{i1}, n_{i2}, n_{i1}, n_{i2})$, where $n_{i1,2}(\omega \pm eV_i)$ – the Fermi distribution functions.

Additionally, we analyze the autocorrelations of the current in the leads. In particular, the noise spectral density in the left lead can be written as

$$S_i(\omega) = \int dte^{i\omega t}\langle\delta I_i(t)\delta I_i(0) + \delta I_i(0)\delta I_i(t)\rangle. \quad (7)$$

Substituting (4) into (7) the zero-frequency shot noise is given by [22,25]

$$S_i(\omega) = 2e^2\int_{-\infty}^{+\infty}\frac{d\omega}{2\pi}Tr\Big[\hat{\sigma}\hat{\Sigma}_i^{<}\hat{\sigma}\hat{G}_{i,i}^{>} + \hat{G}_{i,i}^{<}\hat{\sigma}\hat{\Sigma}_i^{>}\hat{\sigma} - \hat{\sigma}\left[\hat{\Sigma}_i\hat{G}_{i,i}\right]^{<}\hat{\sigma}\left[\hat{\Sigma}_i\hat{G}_{i,i}\right]^{>} - \left[\hat{G}_{i,i}\hat{\Sigma}_i\right]^{<}\hat{\sigma}\left[\hat{G}_{i,i}\hat{\Sigma}_i\right]^{>}\hat{\sigma} + \hat{\sigma}\left[\hat{\Sigma}_i\hat{G}_{i,i}\hat{\Sigma}_i\right]^{>}\hat{\sigma}\hat{G}_{i,i}^{<} + \hat{G}_{i,i}^{>}\hat{\sigma}\left[\hat{\Sigma}_i\hat{G}_{i,i}\hat{\Sigma}_i\right]^{<}\hat{\sigma}\Big]. \quad (8)$$

In many works concerning the topologically SC wires the relation between the parameters of system is typically $t \gg \alpha, V_{x,z}, \Delta$ [8,22]. We choose slightly different condition $t \sim \alpha, V_{x,z}, \Delta$ which results in the increase of oscillation period of MBS zero mode energy and considerably simplifies the transport calculations. The transport is analyzed at low temperatures $kT = 10^{-10}$ and low bias $eV \sim E_M$ where E_M is the MBS energy. Therefore and for the sake of simplicity, out of the parametric area where the topologically nontrivial is realized, $\mu^2 + \Delta^2 < V_x^2 + V_z^2 < (2t-\mu)^2 + \Delta^2$ [6,7,26], all physical quantities are set equal to zero. The hopping parameter $t = 1$ is in energy units.

4. Current-switch effect

To understand the features of transport properties we, firstly, analyze the MBS spin-up probability densities at both edges of the SC wire employing the Bogolubov transformation,

$$\beta_l = \sum_{n=1}^{N}\left[u_{ln}a_{n\uparrow} + v_{ln}a_{n\downarrow}^{+} + w_{ln}a_{n\downarrow} + z_{ln}a_{n\uparrow}^{+}\right]. \quad (9)$$

In Fig. 2a, b the MBS spin-up probability densities at the left and right end sites of the wire, $P_{M,L(R)\uparrow}$, are demonstrated as functions of the magnetic-field orientation and amplitude. According to the