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Giant thermopower in superconducting heterostructures with spin-active interfaces

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ABSTRACT

We predict parametrically strong enhancement of the thermoelectric effect in metallic bilayers consisting of two superconductors separated by a spin-active interface. The physical mechanism for such an enhancement is directly related to electron–hole imbalance generated by spin-sensitive quasiparticle scattering at the interface between superconducting layers. We explicitly evaluate the thermoelectric currents flowing in the system and demonstrate that they can reach maximum values comparable to the critical ones for superconductors under consideration.

1. Introduction

It is well known that application of a thermal gradient ∇T to a normal conductor along with electric field $\boldsymbol{E}$ results in the electric current

$$\boldsymbol{j} = \sigma_N \boldsymbol{E} + \alpha_N \nabla T, \quad \alpha_N \sim (\sigma_N/e)(T/\varepsilon_F). \tag{1}$$

Here σ_N defines the Drude conductivity, α_N is the thermoelectric coefficient and ε_F is the Fermi energy. Provided a metal is brought into a superconducting state, Eq. (1) is no longer correct since the electric field cannot penetrate into the bulk of a superconductor. Instead, one finds

$$\boldsymbol{j} = \boldsymbol{j}_s + \alpha_S \nabla T, \tag{2}$$

where $\boldsymbol{j}_s$ is a supercurrent and α_S defines the thermoelectric coefficient in a superconducting state. It turns out that by applying thermal gradient to a uniform superconductor it is not possible to induce and measure any current since thermal current would always be compensated by the supercurrent $\boldsymbol{j}_s = -\alpha_S \nabla T$. The way out is to consider non-uniform superconducting structures in which case no such compensation generally occurs [1,2] and the thermoelectric current can be detected experimentally. Making use of this idea the thermoelectric effect was indeed demonstrated in several experiments with bimetallic superconducting rings [3–5]. Quite surprisingly, the magnitude of the effect was found to be several orders of magnitude bigger than that predicted by theory [6]. The authors of a very recent experimental work [7], also observed a discrepancy between theory and their experimental data.

By now it is well understood that a small theoretical value of the thermoelectric coefficient in ordinary superconductors [6] $\alpha_S \sim \alpha_N$ is directly linked to the assumption that electron–hole symmetry remains preserved in these structures. In this case contributions to the thermoelectric current provided by electron-like and hole-like excitations are of the opposite sign and almost cancel each other. Then, like in a normal metal, one inevitably finds that α_S is controlled by a parametrically small factor $T/\varepsilon_F \ll 1$.

The situation may change if for some reason the electron–hole symmetry gets violated. In this case – as it was demonstrated by a number of authors – a much stronger thermoelectric effect can be expected. The proposed mechanisms for the electron–hole symmetry violation and the related thermoelectric effect enhancement are diverse. In conventional superconductors doped by magnetic impurities, the presence of Andreev bound states formed near such impurities may yield an asymmetry between electron and hole scattering rates which in turn results in a drastic enhancement of the thermoelectric effect [8]. Likewise, the formation of quasi-bound Andreev states near non-magnetic impurities in unconventional superconductors may lead to much larger values of α_S in such systems [9]. Substantial enhancement of thermoelectric currents was also predicted in three terminal hybrid

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ferromagnet–superconductor–ferromagnet (FSF) [10] as well as in FS junctions in the presence of a Zeeman spin-splitting field [11].

In a recent work [12] we argued that the thermoelectric effect can be strongly enhanced also in metallic bilayers consisting of a superconductor and a normal metal (SN) provided these two metals are separated by a thin spin-active interface. By exactly solving the corresponding Bogolyubov–de-Gennes equations we evaluated the wave functions for electron-like and hole-like excitations in such systems demonstrating that spin-sensitive scattering at the SN interface can generate electron–hole imbalance and result in the presence of large thermoelectric currents in such systems. In this paper we will further extend our arguments [12] to superconducting multilayers with spin-active interfaces and demonstrate that thermoelectric properties of such systems may drastically differ from those of bulk superconductors. As a simple example of such systems below we will specifically consider a superconductor with a thin ferromagnetic interlayer. We will show that provided a temperature gradient is applied along this interlayer the system develops a thermoelectric current in which maximum values can be as high as the critical (depairing) current of a superconductor.

The structure of our paper is as follows. In Section 2 we will specify our model and outline our basic quasiclassical formalism of Eilenberger equations to be employed in our analysis of the thermoelectric effect. In Section 3 we will present an efficient method enabling one to derive the solution of these equations for the system under consideration. With the aid of this solution we will then derive a general expression for the thermoelectric current and also briefly discuss our results in Section 4.

2. The model and quasiclassical formalism

In what follows we will consider an extended metallic bilayer consisting of two superconducting slabs S_1 and S_2 as shown in Fig. 1. We will assume that both metals are brought into direct contact with each other via a spin-active interface that is located in the plane $z=0$. Such an interface can be formed, e.g., by an ultrathin layer of a ferromagnet. Our goal is to evaluate an electric current response to a temperature gradient applied to the system along the S_1S_2 interface. This temperature gradient is achieved by setting the temperature T at the left ($x \to -\infty$) and right ($x \to \infty$) edges of the bilayer respectively equal to T_a and T_b, see Fig. 1. For

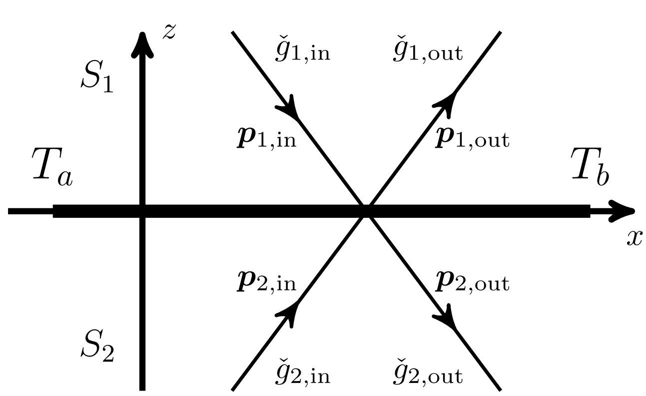

Fig. 1. The system under consideration consisting of two superconducting layers S_1 and S_2 separated by a spin-active interface. The left and right edges of this superconducting bilayer are maintained at temperatures T_a and T_b, respectively. We also schematically indicate the quasiclassical electron Green functions for incoming and outgoing momentum values. These Green functions are matched at the spin-active interface by means of the proper boundary conditions as specified in the text.

the sake of simplicity below we will assume that the temperature depends only on x and does not vary along y- and z-directions.

Within the quasiclassical theory of superconductivity [13], the current density $\boldsymbol{j}(\boldsymbol{r})$ in our system can be evaluated by means of the standard formula

$$\boldsymbol{j}(\boldsymbol{r}) = -\frac{eN_0}{8}\int d\varepsilon \left\langle \boldsymbol{v}_F \mathrm{Sp}[\hat{\tau}_3 \hat{g}^K(\boldsymbol{p}_F, \boldsymbol{r}, \varepsilon)] \right\rangle, \tag{3}$$

where N_0 is the density of state at the Fermi level, $\boldsymbol{p}_F = m\boldsymbol{v}_F$ is the electron Fermi momentum vector, $\hat{\tau}_3$ is the Pauli matrix in the Nambu space, the angular brackets $\langle\cdots\rangle$ denote averaging over the Fermi momentum directions and $\hat{g}^K$ is the Keldysh block of the quasiclassical Green–Eilenberger function matrix:

$$\check{g} = \begin{pmatrix} \hat{g}^R & \hat{g}^K \\ 0 & \hat{g}^A \end{pmatrix}. \tag{4}$$

Here and below the "hat" symbol denotes 4×4 matrices in the Nambu⊗Spin space while the "check" symbol labels 8×8 matrices in the Keldysh⊗Nambu⊗Spin space.

The matrix function $\check{g}$ obeys the transport-like Eilenberger equation [13]:

$$\left[\varepsilon\hat{\tau}_3 - \check{\Delta}(\boldsymbol{r}), \check{g}\right] + i\boldsymbol{v}_F \nabla \check{g}(\boldsymbol{p}_F, \boldsymbol{r}, \varepsilon) = 0 \tag{5}$$

as well as the normalization condition:

$$\check{g}^2 = 1. \tag{6}$$

The order parameter matrix $\check{\Delta}$ has only "retarded" and "advanced" components:

$$\check{\Delta} = \begin{pmatrix} \hat{\Delta} & 0 \\ 0 & \hat{\Delta} \end{pmatrix}, \quad \hat{\Delta} = \begin{pmatrix} 0 & \Delta\sigma_0 \\ -\Delta^*\sigma_0 & 0 \end{pmatrix}, \tag{7}$$

where σ_0 is the unity matrix in the spin space and Δ is the superconducting order parameter. As soon as we are interested in the electronic transport along the interface we set the phase difference between the two superconductors S_1 and S_2 to zero. Under this assumption order parameter can be made to be real everywhere in the system.

As usually, the quasiclassical equations (5) should be supplemented by boundary conditions which describe electron transfer across the *SFS*-interface by matching the Green function matrices $\check{g}$ for incoming and outgoing momentum directions at both sides of this interface, see Fig. 1. In the case of spin-active interfaces the corresponding boundary conditions were derived in [14]. Here we will employ an equivalent approach [15].

The simplest model of the spin-active interface is described by three parameters, i.e. the transmission probabilities for opposite spin directions $D_\uparrow$ and $D_\downarrow$ as well as the so-called spin mixing angle θ which just represents the difference between the scattering phase shifts for spin-up and spin-down electrons. These parameters are assumed to be energy independent which can be justified for sufficiently thin ferromagnetic layers. At the same time, the layer should not be too thin in order to remain in the ferromagnetic state.

Previously we have already made use of this model, e.g., while considering crossed Andreev reflection in three-terminal FSF structures [16] or triplet pairing and dc Josephson effect in SFS junctions [17]. For simplicity we also assume that the above three parameters do not depend on the sign of the quasiparticle momentum along the interface, i.e. $D_\uparrow(\boldsymbol{p}_\parallel) = D_\uparrow(-\boldsymbol{p}_\parallel)$ and so on. Then