Stimulation of a Singlet Superconductivity in SFS Weak Links by Spin–Exchange Scattering of Cooper Pairs

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Josephson junctions with a ferromagnetic metal weak link reveal a very strong decrease of the critical current compared to a normal metal weak link. We demonstrate that in the ballistic regime the presence of a small region with a non-collinear magnetization near the center of a ferromagnetic weak link restores the critical current inherent to the normal metal. The above effect can be stimulated by additional electrical bias of the magnetic gate which induces a local electron depletion of ferromagnetic barrier. The underlying physics of the effect is the interference phenomena due to the magnetic scattering of the Cooper pair, which reverses its total momentum in the ferromagnet and thus compensates the phase gain before and after the spin–reversed scattering. In contrast with the widely discussed triplet long ranged proximity effect we elucidate a new singlet long ranged proximity effect. This phenomenon opens a way to easily control the properties of SFS junctions and inversely to manipulate the magnetic moment via the Josephson current.

Mesoscopic properties of nanometer sized conductors are strongly affected by the injection of correlated electrons (Cooper pairs) from superconducting electrodes. While propagating through normal metal such pairs of electrons are able to preserve their superconducting correlation on mesoscopic lengths providing the superconducting current flow through SNS (superconductor-normal metal-superconductor) weak links1,2. Time reversal of electronic states forming a Cooper pair is an important component of the above correlation. Absence of superconducting pairing interaction in a N part of the superconducting SNS device opens a possibility of easy external manipulation of the spin structure of the propagating Cooper pairs. This offers the means for spintronic manipulation of superconducting weak links. The most efficient “tailoring” of Cooper pairs can be achieved by external in-homogeneity located on a submicron length scale, which is set by superconducting coherence length, determining the scale of spatial correlation of paired electrons. Here we demonstrate that a tip of the magnetic exchange force microscope (MExFM)3,4, which induces localized in space magnetic exchange fields can play the role of such a local probe for a spin state of superconducting Cooper pairs. Existing experimental evidences of the externally induced exchange fields \( h \) in metals of the order of few5,6 or even few tens7 millielectronvolts place the electronic coupling to such field in a range of energies comparable with (or even exceeding) superconducting energy scale \( \Delta \). We will show that the effect of exchange induced gating of Cooper pairs leads to a new phenomenon - stimulation of long- range singlet superconductivity in SFS (superconductor-ferromagnet-superconductor) weak links.

Magnetic exchange interaction in ferromagnetic metals lifts a degeneracy with respect to spin orientation of the electrons, forming a Cooper pair. This leads to different de-Broglie wavelengths of electrons at Fermi surface for spin-up and spin-down orientation and produces a modulation of the Cooper pair wavefunction while propagating along the ferromagnet8. As a result an oscillatory damping of the superconducting ordering is known to appear when a ferromagnetic ordering occurs in a normal metal link connecting two S electrodes9. This phenomenon provides the basis of the \( \pi \)-junction realization10–12. Considering the quantum mechanics of quasiparticle excitations the exchange field leads to phase difference \( \gamma \approx L/\zeta_h \) gained between the electron- and hole- like parts of the total wave function along a path of the length \( L \), where \( \zeta_h = \hbar V_F / 2h \) is a characteristic length determined by the exchange field \( V_E \) is the Fermi velocity).
Fast oscillating contributions by spatial non-homogeneous magnetization. Since they bind electrons with aligned spins (with equal-spin pairs) can be formed in a ferromagnetic metal. It was suggested before that the Cooper pairs do not dephase, thereby leading to long-range proximity due to the destructive trajectory interference. In the second case, the scattered Cooper pair has a reversed spin arrangement and the total phase gain is $\gamma \sim (d_1 + d_2)/\xi_h$, which results in a strong suppression of proximity due to the destructive trajectory interference. To be more precise, we consider the Josephson transport through a normal ballistic nanowire (NW) in contact with a ferromagnetic insulator (FI). The FI turns the NW into an effective ferromagnet with an exchange field $h$. The schematic picture of the SFS device is presented in Fig. 2A. The total length of the constriction $d = d_1 + d_2 + d_3$ is assumed to be large compared to the magnetic coherence length $\xi_h = hV_F/2h: d \gg \xi_h$. For simplicity, we restrict ourselves to the case of a short junction with $d \ll \xi_n$, where $\xi_n = hV_F/T_c$ is the coherence length of normal metal ($T_c$ is the critical temperature of the S layer). The magnetic tip is assumed to bring on localized in space magnetic exchange field inhomogeneity which we model by a stepwise profile:

$$h(x) = \begin{cases} h\xi_0, & \text{in domains } d_1, d_3 \\ h(x_0 \cos \alpha + x_0 \sin \alpha), & \text{in domain } d_2, \end{cases}$$

where $\alpha$ is the angle of the exchange field rotation in the central domain $d_2$ (see Fig. 2B).

**Results**

The current–phase relation of SFS Josephson junction is determined by the quasiclassical relation

$$I = \sum_n I_n = \sum_n a_n \sin \eta \rho \frac{\langle \mathbf{n}, \mathbf{n}_F | \cos \eta \mathbf{n} \rangle}{\langle \mathbf{n}, \mathbf{n}_F \rangle},$$

where $\mathbf{n}$ is the unit vector normal to the junction plane, $\mathbf{n}_F$ is the unit vector along the trajectory, and $a_n$ are the coefficients of the Fourier expansion for the current–phase relation for superconductor–normal metal weak links.
mal metal junction of the same geometry. The angular brackets denote the averaging over different quasiclassical trajectories characterized by a given angle $\theta$ and a certain starting point at the superconductor surface, and for 3D conduction looks as

$$\langle \langle \mathbf{n}, \mathbf{n}_F \rangle \rangle = 2 \int_0^{\pi/2} d\theta \sin\theta \cos\theta \cos(\gamma),$$

where $\cos \theta = (n, n_F)$. At temperatures $T$ close to $T_c$, the current–phase relation (2) is sinusoidal, and the coefficient $a_1$ is determined by the following simple relations$^{13}$:

$$a_1 = \frac{eT_c N f_1}{8\hbar} \left( \frac{\Delta}{T_c} \right)^2.$$  

Here $\Delta$ is the temperature dependent superconducting gap, the factor $N$ is determined by the number of transverse modes in the junction: $N = \frac{1}{s_0} \int ds \int d\mathbf{n}_F (\mathbf{n}_F, \mathbf{n}) \sim S/s_0$, where $S$ is the junction cross-section area, and $s_0^{-1} = (k_F^2/2\pi)^2$, where $k_F$ is the Fermi momentum.

The phase gain $\gamma$ can be conveniently determined from the Eilenberger—type equations$^{22}$ if we use a standard parametrization$^{24}$ of the anomalous quasiclassical Green function $f_1 = f_1^1 + f_1^s$, where $f_1^1$ ($f_1^s$) singlet (triplet) parts of the function, respectively, and $\hat{\theta}$ is a Pauli matrix vector in the spin space. The functions $f_1^1, f_1^s$ satisfy the linearized Eilenberger equations$^{18}$ written for zero Matsubara frequencies

$$-i\hbar V_{ij} \hat{\theta}_j + 2\hbar f_1 = 0, \quad -i\hbar V_{ij} \hat{\theta}_j + 2\hbar f_1 = 0,$$

with the boundary conditions $f_1^1(s = s_0) = 1, f_1^s(s = s_0) = 0$ at the left superconducting electrode (for simplicity we consider the case of the absence of the barriers at the interfaces). The phase gain $\gamma$ along the trajectory in the equivalent SFS junction (Fig. 2B) is determined by the singlet part of the anomalous quasiclassical Green function $f_1^1(s = s_0) = \cos \gamma$ taken at the right superconducting electrode$^{22}$. Solving the equations (5) for the stepwise profile of the exchange field (1) we find (Methods):

$$\cos \gamma = \cos \delta_2 \cos(\delta_1 + \delta_0) - \cos \otimes \sin \delta_2 \sin(\delta_1 + \delta_0) - \sin^2 \delta_1 \sin \delta_3(1 - \cos \delta_0),$$

where $\cos \theta = (n, n_F)$ and $\delta_0 = d_0/\xi_h \cos \theta$ ($i = 1, 2, 3$). Averaging the expression (6) over the trajectory direction $\theta$ and neglecting the terms proportional to $\xi_h/d \ll 1$, which decrease just as for the case of homogeneous ballistic 3D SFS junction, one arrives at the following long–range (LR) contribution:

$$\langle \cos \gamma \rangle_{LR} = -\frac{1}{2} \sin^2 \delta \cos(1 - \cos \delta_2) \cos 2\delta_3,$$

where $\delta_2 = z_0/\xi_h \cos \theta$ and $z_0 = (d_1 - d_2)/2$ is the shift of the central domain with respect to the weak link center. So, the long–range component of the Josephson current at the first harmonic is determined by the relation:

$$I_{LR}^1 = a_1 T_{11}^LR \sin \varphi, \quad T_{11}^LR = \int_0^{\pi/2} d\theta \sin \theta \cos \theta (\cos \gamma)^{LR}.$$

For a very thin central domain $d_2 \ll \xi_h$ in the center of the NW ($z_0 = 0$) one can easily estimate from (7, 8) the critical current of the SFS junction

$$\max \{I_{LR}^1\} = \frac{\hbar}{2 \sin^2 \delta} \left( \frac{d_2}{\xi_h} \right)^2 \ln \frac{\xi_h}{d_2},$$

where $\hbar = (eT_c N/8\hbar) (\Delta/T_c)^2$ is the critical current of the SNS junction for zero exchange field ($\gamma = 0$). We see that the long–ranged critical current reaches the maximum at $\alpha = \pi/2$ and grows with the increase of $d_2$ up to $d_2 \sim \xi_h$. Our numerical calculations show that it is maximum for $d_2 \approx 2.5 \xi_h$ and may reach $\sim 0.7 I_c$. Certainly, the above long–range effect in the first harmonic describes the properties of the SFS constriction if contribution of higher harmonics in the current–phase relation (2) is negligible. We present the analysis of the second harmonic effect in the Supplementary information.

**Discussion**

Figure 3 shows the dependence of the maximal Josephson current $I_{LR}^0 = a_1 T_{11}^0 R$ on the thickness $d_2$ of the 90° domain ($\alpha = \pi/2$) for different values of the shift of the domain $z_0$. Naturally, when the thickness of the central domain $d_2$ goes to zero, the long–range effect disappears. We can see from Fig. 3, that the long–range component of the Josephson current $I_{LR}^0$ coincides approximately with the total supercurrent across the junction (2) until the outer domains are long enough: $d_1, d_3 > \xi_h$. The amplitude of $I_{LR}^0$ depends nonmonotonically on the size of the central domain $d_2$ and has the first maximum at $d_2 \approx 2.5 \xi_h$. Interestingly, that the long–range contribution generates a $\pi$–junction ($I_{LR}^0$ is negative for zero shift of the domain $z_0 = 0$). With the shift of the central domain the junction can be switched from $\pi$– to $0$– state. Figure 4 shows the dependences of the maximal Josephson current $I_c$ on the position of the central domain $z_0$ for different values of the rotation angle $\alpha$. We may see that the critical current is very sensitive to the position of the central domain and the first zero of $I_c$ occurs already at $z_0 \approx 0.5 \xi_h$.

Here we considered a simple step-like model of magnetization distribution in F layer. For the very thin central domain $d_2 \ll \xi_h$, we may easily estimate the long ranged contribution $(\cos \gamma)^{LR} \approx (d_2/\xi_h)^2 (d_2/\hbar)^2$ which is in accordance with the expression (7). For a general profile of the magnetization it may be convenient to use the transfer–matrix method (see the Supplementary information). The smooth (on the scale $\xi_h$) profile of the magnetization decreases the long ranged effect and the proposed mechanism occurs to be most efficient for $d_2 \ll \xi_h$.

Note that the considered phenomenon should generate the oscillating potential profile for the magnetic tip $U(z_0) \sim -I_z(z_0) \cos \phi$ which depends on superconducting phase difference across the junction. This opens an interesting possibility to couple the Josephson current oscillations with mechanical modes of the tip. On the other hand the same effect can produce a change of the orientation of the magnetic moment. Inversely, the precession of the magnetic moment shall modulate the critical current of the junction and provides a direct coupling between the superconducting current and the mag.
The magnetically tunable long-range SFS proximity effect suggested above has a potential to be an important feature of carbon-based superconducting weak links. Graphene sheets and carbon nanotubes are reported to offer a ballistic propagation for electrons on a micrometer length scale. This fact together with appearing reports on a gate tunable magnetism in graphene makes all ingredients of the present theory achievable in experiment. Another possibility is to use the indium antimonide (InSb) nanowires as a superconducting weak link. The indium antimonide nanowires, recently used in the experiments to reveal the signature of Majorana fermions, demonstrated a very high g-factor (g ≈ 50). Anomalously large g-factor reported in such wires offers the possibility to “mimic” a ferromagnetic spin-splitting effect of the order of 10 K by simply applying an external magnetic field of the order of 0.1 Tesla, and then making such nanowire a suitable candidate for a weak link to observe the discussed phenomena. Note that in contrast to the experiments the magnetic field should be applied along the spin-orbit field axis to avoid the interference with the spin-orbit effect.

It should be noted, that new additional functionality of the considered device can be achieved by electric biasing of the magnetic gate. In weakly doped ferromagnetic barriers, such bias (V_e) alters both the charge carrier concentration and the Fermi velocity.

Choosing a polarity of electric gating one can create a depletion region beneath the tip. As a result, both the Fermi velocity V_F and the exchange length ξ_h = hV_F/2h decrease in the spatial region of the domain d_2, and the key parameter responsible for the magnetic exchange scattering δ_2 ≈ d_2/ξ_h grows. For thin domains (δ_2 ≪ 1) the critical current I_c ∝ δ_2^2 increases with the gate voltage V_e and the local depletion of F barrier should result in the stimulation of the superconductivity. This nontrivial interplay between electric and magnetic gating effects can be used to control singlet Josephson current through ferromagnetic nanowires.

To summarize, we studied the interference phenomena originated by the spin–exchange scattering in ferromagnetic ballitic weak link and demonstrated that they provide an efficient way to control the Josephson current and to couple it with a magnetic moment.

**Methods**

A. Transfer–matrix formalism for Eilenberger equations. To consider the Josephson transport through ferromagnetic layer with a non-collinear magnetizations M and exchange field h it is convenient to utilize the transfer–matrix formalism. For this, we need to solve the linearized Eilenberger equations written for zero Matsubara frequencies

\[ -i\hbar \nabla^2 \psi_L f_L + 2h \psi_L = 0, \quad -i\hbar \nabla^2 \psi_R f_L + 2h \psi_R = 0, \]

for the case when the quantization axis is taken arbitrarily in the ferromagnetic layer of a thickness d. We assume that a quasiclassical trajectory s and exchange field h = h (x_0, sin x + x_0, cos x) lie in the plane (x, z), as shown in Fig. 5. The trajectory is characterized by a given angle θ with respect to the z-axis. The triplet part f_0 of the anomalous quasiclassical Green function f = f_0 + f_1 + f_2 consists of two nonzero components and can be written as f_0 = f_0(x_0) + f_0(x_0). Defining the transfer–matrix \( T_s(d, 0) \) that relates the components of the Green function \( \psi_R(x) \) at the left (s = 0) and right (s = x) boundaries of the F layer,

\[ \psi_R(x) = T_s(d, 0) \psi_L(0), \]

we get the following expression:

\[ T_s(d, 0) = \begin{pmatrix} \cos(q_s) & -i \cos \theta \sin(q_s) \\ -i \sin \theta \sin(q_s) & \cos(q_s) \sin^2 \alpha \cos(q_s) - 1 \\ -i \sin(q_s) & \cos \theta \sin(q_s) \end{pmatrix}, \]

where \( q = 1/\xi_h = 2h/V_F \).

In order to elucidate the peculiarities of the Cooper pairs scattering with a spin-flip transition of electrons it is convenient to introduce the new functions \( f_{6} = f_{6} \mp f_{6} \) which describes the pairs with zero spin projection and a reversed spin arrangement. The transfer–matrix \( T_s(d, 0) \) can be drastically simplified if the direction of the exchange field coincides with a spin quantisation axis. In this case, \( x = 0 \) and \( f_{6}(s) = e^{i\alpha} f_{6}(0) \). Calculating the superconducting current at the right electrode \( S_R \) we readily see that it results from the interference with the singlet component coming from the left electrode \( f_{6}(s) = f_{6}(s) + f_{6}(-s)/2 \) (triplet components are irrelevant because the right electrode provides only the singlet component). The oscillating factors \( e^{i\alpha} \) in \( f_{6}(s) \) produce, after the averaging over the trajectories directions (angle θ), a strong damping of the critical current compared to the normal metal (where these factors are absent).

Now we may easily understand the mechanism of the long-ranged proximity effect. Indeed after coming through the first F layer an extra phase factor appears in \( f_{6} \) functions (see Fig. 2): \( f_{6}(s) = e^{i\alpha} f_{6}(0) \). In the absence the middle layer, the \( f_{6} \) at the right electrode would be \( f_{6}(s) + f_{6}(s) = e^{i\alpha} f_{6}(0) \) (and the oscillating factors will strongly damp current. The additional non-collinear middle layer \( d_2 \) will mix up the components \( f_{6} \) and \( f_{6} \) - see the matrix (12) and, for example, \( f_{6}(s) + f_{6}(s) \) in addition to \( e^{i\alpha} f_{6}(0) \) component will have a \( e^{i\alpha} f_{6}(0) \) contribution, i.e. \( f_{6}(s) + f_{6}(s) = a e^{i\alpha} f_{6}(0) + be^{i\alpha} f_{6}(0) \). In fact, namely this mechanism is schematically presented in the Fig. 1(b). Then the resulting \( f_{6} \) function at the right electrode should be \( f_{6}(s) + f_{6}(s) = a e^{i\alpha} f_{6}(0) + be^{i\alpha} f_{6}(0) \) and for \( d_2 \) the oscillating factor at the second term vanishes. This means the emergence of the long-ranged singlet proximity effect discussed in the present report. Note that the
additional noncollinear F layer $d_F$ may strongly increase the critical current, provided that it is placed at the center of the structure.

The transfer–matrix method is very convenient for the calculation of the Josephson transport through the SFS junction containing three ferromagnetic layers with a stepwise profile of exchange field. For this geometry shown in Fig. 2B of the Letter, the anomalous quasiclassical Green function $f(s) = \langle f(x)_{s} f_{0}(x)_{s} f(x) \rangle$ at the right superconducting electrode ($s = s_{0} = d_{c} d_{i}$) can be easily expressed via the boundary conditions $f(0) = (1,0,0)$ at the left superconducting electrode ($s = 0$) as follows:

$$
\begin{align*}
 f(s) &= T_{i}(d_{i}, 0)T_{i}(d_{i}, 0)T_{o}(d_{o}, 0)f(0),
\end{align*}
$$

(13)

where $d = d_{i} + d_{o} + d_{c}$ is the total thickness of the ferromagnetic barrier. As a result, the singlet part $f(s)$ responsible for the Josephson current through the junction, can be written in the form (6).

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