Chirality-controlled spontaneous currents in spin-orbit coupled superconducting rings

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At a superconductor interface with a ferromagnetic insulator (FI), the FI acts to induce a local exchange field within the S layer, which in the presence of spin-orbit interaction promotes a phase-modulated superconducting state. Here we demonstrate that within a thin superconducting loop that is partially proximitized by a FI, spontaneous currents form with a magnetization-orientation-dependent chirality with sizable shifts in Little-Parks oscillations. Furthermore, the critical temperature of the loop is also magnetization-orientation-dependent and conversely, the superconducting transition itself may influence the magnetization direction. More generally, the superconducting region above the FI may serve as a “phase battery” and so offer a new device concept for superconducting spintronics.

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The interaction of interface superconductivity with materials exhibiting strong spin-orbit coupling \[1, 2\] and magnetic exchange fields \[2, 3\] offers enormous potential for the discovery and control of new physical phenomena. For example, a homogeneous magnetic exchange field acting at a superconductor/ferromagnet (S/F) interface induces oscillations in the superconducting order parameter in S/F/S Josephson junctions \[4, 5\] whilst a non-uniform exchange field in such junctions can lead to electron pair conversion from spin-singlet to spin-triplet \[6, 7\] and a dependence of the superconducting critical temperature \(T_c\) on magnetization alignment in S/F/S multilayers \[11\]. In the absence of inversion symmetry, Rashba spin-orbit interaction (SO) in combination with a magnetic exchange field or Zeeman field offers additional physics with an unusual linear over the gradient of the superconducting order parameter \(\Psi\) terms in the Ginzburg-Landau (GL) free energy \((\nabla \Psi) \Psi^*\). Here we call such coupling the “exchange spin-orbit coupling” or “EXSO”. The EXSO effect induces different types of Larkin-Ovchinnikov-Fulde-Ferrell (LOFF)-like helical phases with a non-zero Cooper pair momentum \(\vec{p}\) in the ground state \[12–16\], which play an important role in Majorana physics \[17\] and lead to the formation of \(\varphi_0\)-Josephson junctions with a spontaneous phase difference of \(\varphi_0\) in the ground state \[18, 20\]. Recently, such \(\varphi_0\)-Josephson junctions have been realized experimentally \[21\].

We note that in uniform systems the helical phases do not carry a current; however, such current-carrying states may appear in non-uniform EXSO systems such as close to a magnetic island on a thin film superconductor \[22\] or near an S/F interface within a distance of the London penetration depth of the interface \[23\].

In this letter we demonstrate that the EXSO effect leads to spontaneous currents in a closed superconducting loop in which the superconductor is partially-coupled to a ferromagnetic insulator (FI) as shown in Fig. 1. We further demonstrate that the chirality of the current is controllable through the magnetization alignment of the FI. The center of the ring of radius \(R\) is at \(x = y = 0\). In this geometry, we expect a non-trivial interplay between the Little-Parks effect and helical current carrying states. Even in zero external magnetic field, spontaneous currents are generated in the ground state. The study of Little-Parks oscillations is a powerful tool to probe subtle effects and was recently applied to investigate the superconducting symmetry in \(\text{Sr}_2\text{RuO}_4\) microrings \[24\].

To describe the superconducting loop we apply the general Ginzburg-Landau (GL) approach which is relevant at a temperature \(T\) close to the critical transition \(T_c\) of the loop. In the presence of the EXSO, the density \[
\rho_0 \sim \left( T - T_c \right)^2 \left( \nabla \varphi \right)^2 \]
\[ f(r) \] of the GL free energy \( F = \int f(r) d^2r \) reads
\[
f(r) = a|\psi|^2 + \gamma |D\psi|^2 + \frac{b}{\sqrt{2}}|\psi|^4 \tag{1}
\]
\[ + \left[ \mathbf{n} \times \mathbf{h} \right] \cdot \left[ \psi^* \varepsilon_0(r) D\psi + \text{c.c.} \right].
\]

Here \( a = -\alpha(T_c - T) \), \( \alpha \), \( b \) and \( \gamma > 0 \) are the standard GL coefficients, \( \psi \) is the superconducting order parameter with, \( D = -i\hbar \nabla + (2e/\alpha)A \) is the gauge-invariant momentum operator \((e > 0)\), \( \mathbf{n} \) is the unit vector in the direction along which the inversion symmetry is broken (pointing in the \( z \) direction perpendicular to S/F interface in our case), \( \mathbf{h} \) is the exchange field, and \( \varepsilon_0(r) \) is the EXSO constant. We assume that the exchange field \( \mathbf{h} \) in superconductor is generated by the surface field of the ferromagnetic insulator and as Rashba SO interaction with \( \lambda = \gamma/\hbar \) is the EXSO constant for the bulk superconductor, see for example [13]. Here \( v_{so} \) is the characteristic velocity entering in SO Rashba interaction and \( v_F \) is the Fermi velocity.

For a 1D ring of radius \( R \) it is convenient to write the GL free energy functional (1) using the polar coordinates \((r, \varphi)\), and take into account that the tangential \( \varphi \)-component \( A \) of the vector-potential \( A = (0, A, 0) \) is directly related with the magnetic flux \( \Phi = 2\pi RA \) enclosed by the superconducting loop. Then the GL free energy then reads
\[
F = \sigma R \int_0^{2\pi} d\varphi \left\{ a|\psi|^2 + \frac{\hbar^2}{2R^2} \left| -i \frac{\partial}{\partial \varphi} + \phi \right| \psi \right|^2 + \frac{b}{2}|\psi|^4 \right\}_{\varphi} + \hbar \varepsilon_0 \cos(\varphi - \alpha) \tag{2}
\]
where \( \varepsilon \) is the cross section of the ring, \( \phi = \Phi/\Phi_0 \) is the enclosed magnetic flux in the units of the flux quantum \( \Phi_0 = \pi \hbar c/e \) and
\[
\varepsilon(\varphi) = \varepsilon_0 \left[ \mathbf{n} \times \mathbf{h} \right] = \hbar \varepsilon_0 \cos(\varphi - \alpha)
\]
is a projection on the ring of the Rashba SOC for \(|\varphi| \leq \pi/2 \) and \( \varepsilon(\varphi) = 0 \) otherwise.

To calculate \( T_c \) we are concerned with the linear equation for the superconducting order parameter, which for the functional (2) can be written as
\[
\left( \frac{\partial}{\partial \varphi} + i \phi \right)^2 \psi + 2i\lambda(\varphi) \left( \frac{\partial \psi}{\partial \varphi} + i \phi \psi \right) + i \frac{\partial \lambda}{\partial \varphi} \psi = \tau \psi, \tag{3}
\]
where \( \lambda(\varphi) = \lambda_{so} \cos(\varphi - \alpha) \), \( \lambda_{so} = \varepsilon_0 \hbar R/\gamma \) is the EXSO parameter for \(|\varphi| \leq \pi/2 \) and \( \lambda(\varphi) = 0 \) otherwise. The eigenvalue \( \tau \) determines the shift of the transition temperature
\[
T_c = T_{c0} \left( 1 + \frac{\xi_0^2}{R^2} \tau \right) \tag{4}
\]
with respect to its bulk value \( T_{c0} \). Here \( \xi_0^2 = \gamma^2/\alpha T_{c0} \) is the GL coherence length. Note that averaged exchange field \( h_{av} \sim \hbar a/d \) and to preserve the superconductivity we need \( h_{av} \leq T_c \). However at \( v_{so}/v_F \sim 0.1 \) we may have \( R/\xi_0 \sim 100/10 \) a pretty large spin-orbit parameter \( \lambda_{so} \sim 1/10 \).

The current density corresponding to the functional (2) is determined by the relation
\[
j = -\frac{\partial f}{\partial A} = \frac{2\gamma h}{R} \left[ \psi \frac{\partial \psi^*}{\partial \varphi} - i \psi^* \frac{\partial \psi}{\partial \varphi} + 2 \lambda \psi^* \psi \right] \tag{5}
\]
and we may verify that the Eq. (3) guarantees \( \text{div} j = 0 \).

To explain our main results, we start from a qualitative discussion of the SOC effect on the Little-Parks oscillations. By introducing an effective flux \( \tilde{\phi}(\varphi) = \phi + \lambda(\varphi) \), we write the equation (5) in following form
\[
\left( \frac{\partial}{\partial \varphi} + i \tilde{\phi} \right)^2 \psi = \tilde{\tau} \psi,
\]
where \( \tilde{\tau}(\varphi) = \tau - \lambda^2(\varphi) \). Then the averaged over the ring value
\[
\langle \tilde{\phi} \rangle = \frac{1}{2\pi} \int_0^{2\pi} d\varphi \tilde{\phi}(\varphi) = \phi + (\lambda_{so}/\pi) \cos \alpha,
\]
describes the additional shift
\[
\Delta \phi = - (\lambda_{so}/\pi) \cos \alpha \tag{6}
\]
of the \( T_c \) maximum of the Little-Parks oscillations.

To provide the full solution of (3) we use the Fourier series expansion for the order parameter
\[
\psi = \sum_n \psi_n \exp(-in\varphi) \tag{7}
\]
and EXSO coupling constant
\[
\lambda(\varphi) = \sum_n \lambda_n \exp(-in\varphi), \tag{8}
\]
where
\[
\lambda_n = \lambda_n^r + i \lambda_n^i = \frac{\lambda_{so} \cos \alpha}{2\pi} \int_{-\pi/2}^{\pi/2} d\varphi \cos \varphi \cos(n\varphi) + i \frac{\lambda_{so} \sin \alpha}{2\pi} \int_{-\pi/2}^{\pi/2} d\varphi \sin \varphi \sin(n\varphi) \tag{9}
\]
Then we readily obtain from the Eq. (3) the systems of coupled algebraic equations for the harmonics \( \psi_n \)

\[- (n - \phi)^2 \psi_n + \sum_{m} \lambda_{n-m} (n + m - 2\phi) \psi_m = \tau \psi_n.\] (10)

Let us consider first the case of the flux smaller than the half quantum \( \phi < 1/2 \). In this situation zero harmonics \( \psi_0 \) will dominate ( \( \psi_n \ll \psi_0 \) ) if \( \lambda_{so} \) is small:

\[(\tau + \phi^2 + 2\phi \lambda_0) \psi_0 = \sum_{n \neq 0} \lambda_n^2 (n - 2\phi) \psi_n.\]

On the other hand for \( \psi_n \), we have

\[\tau + (n - \phi)^2 - 2(n - \phi) \lambda_0 \psi_n \approx \lambda_n (n - 2\phi) \psi_0,\]

which leads to the equation for \( \tau \)

\[(\tau + \phi^2 + 2\phi \lambda_0) = \sum_{n \neq 0} \frac{\lambda_n^2 (n - 2\phi)^2}{\tau + (n - \phi)^2 - 2(n - \phi) \lambda_0}.\]

Since \( \tau \approx -\phi^2 - 2\lambda_0 \phi \) and than \( \tau + (n - \phi)^2 - 2\lambda_0 (n - \phi) \approx n [n - 2(\lambda_0 + \phi)] \approx n (n - 2\phi) \) and finally we have for the transition temperature

\[\tau = -\phi^2 - 2\lambda_0 \phi + \sum_{n \neq 0} |\lambda_n|^2 = -(\phi + \lambda_0)^2 + \lambda_{so}^2 / 4,\]

where we took into account that \( \lambda_{-n} = \lambda_n^* \) and

\[\sum_n |\lambda_n|^2 = 2|\lambda_1|^2 + \sum_m |\lambda_{2m}|^2 = \lambda_{so}^2 / 4.\]

In the absence of the external flux \( (\phi = 0) \) we have the following dependence of the critical temperature \( \tau \) on the orientation of the exchange field \( \mathbf{h} \):

\[\tau = \lambda_{so}^2 [1 - (2 \cos \alpha / \pi)^2] / 4.\] (11)

In general case if the flux \( \phi \) is large \( \min |n - \phi| \) occurs for some finite \( n_0 \) and for \( \alpha = 0 \) we have similarly

\[\tau = -(\phi + \lambda_0 - n_0)^2 + \lambda_{so}^2 / 4.\]

The results of the numerical calculations for arbitrary values of \( \lambda_{so} \) are plotted in Fig. 2. We see that the strong EXSO interaction shifts the Little-Parks oscillations with a dependence on the orientation of \( \mathbf{h} \). One can also observe the stimulation of the superconductivity by spin-orbital interaction: in Fig. 2, we compare the dependence of the shifts \( \Delta \phi \) of the \( T_c \) maximum of the Little-Parks oscillations obtained analytically, i.e., using Eq. (8), and numerically. One can see an excellent agreement between the qualitative description (8) and the exact solution of the eigenvalue problem (11). Note that the numerical calculation demonstrates a weak dependence of the maximal \( T_c \) on direction of \( \mathbf{h} \). In Figure 3 we show the dependence of the transition temperature \( \tau \) on the \( \mathbf{h} \)

![Figure 2](image-url)

**FIG. 2:** (Color online) (a) Little-Parks oscillations for different values of EXSO parameter \( \lambda(\phi) = \pm \lambda_{so} \cos(\phi - \alpha) \) for \( |\phi| \leq \pi/2 \): \( \alpha = 0 \) – solid; \( \alpha = \pi \) – dashed. The numbers near the curves denote the corresponding values of the spin-orbit constant \( \lambda_{so} \). (b) Dependence of the shifts of the Little-Parks oscillations \( \Delta \phi (\Delta) \) and the maximal transition temperature \( \tau_{max} \) (—) (see the panel (a) in confusion) on the angle \( \alpha \) for \( \lambda_{so} = 1 \). Solid red line shows the dependence described by the Eq. (3).

![Figure 3](image-url)

**FIG. 3:** (Color online) Dependence of the transition temperature \( \tau \) on the angle \( \alpha \) for \( \lambda_{so} = 0.5 \) (blue solid line) and \( \lambda_{so} = 1.25 \) (red solid line) obtained by the numerical solution of the eigenvalue problem (11) for \( \phi = 0 \). The inserts show the spectra of orbital harmonics for \( \alpha = 0 \) and different values of the spin-orbit constant: left – \( \lambda_{so} = 0.5 \); right – \( \lambda_{so} = 1.25 \).
direction \( \alpha \) in the absence of external flux (\( \phi = 0 \)) for two values of spin-orbital constant \( \lambda_{so} \). The insets in Figure 3 show spectra of orbital harmonics for the selected values of \( \lambda_{so} \) at \( \alpha = 0 \). For small values of the EXSO constant the zero harmonic prevails (left inset in Fig. 3), and \( \tau(\alpha) \) dependence is well described by the expression (11). The spin-orbital interaction results in generation of harmonics with nonzero orbital momenta. As a result the spectrum of orbital harmonics spreads with an increase of the spin-orbital interaction, and a dominant harmonic has a nonzero orbital momenta (right inset in Fig. 3).

We now calculate the current for the case of spin-orbit interaction in the absence of the applied field with \( \phi = 0 \) and the exchange field oriented along the \( x \)-axis (\( \alpha = 0 \)). We calculate the current at \( \varphi = \pi \), where the expression for the current density has the standard form (4) with \( \lambda = 0 \) and

\[
j = -4\epsilon \frac{\gamma |h|}{R} \sum_{n,m} n (-1)^{n-m} \Re \{ \psi_n \psi_n^* \} \approx 4\epsilon \frac{\gamma h \lambda_{so}}{\pi R} \psi_0^2.
\]

Figure 4 shows the dependence the current density \( j \) on the spin-orbital constant \( \lambda_{so} \) for different temperatures \( T < T_c \). One can see that even in the absence of an external magnetic field the superconducting transition occurs in the current-carrying state with nonzero orbital momenta. The sign and magnitude of the spontaneous supercurrent due to spin-orbital interaction is determined by the EXSO constant with a maximum at \( \lambda_{so} \approx 1 \).

We have considered above the case of a thin superconducting ring; however, the effect of the phase accumulation due to EXSO coupling (describing by the constant \( \varepsilon \approx h \varepsilon_0 \)) is very general and can play the role of a superconducting phase battery. Indeed, let us consider the superconducting loop at temperatures below \( T_c \) where modulus \( |\psi| \) of the order parameter is constant. In the case of the infinite wire in the presence of the EXSO coupling the phase modulated solution is realized \( \psi = |\psi| \exp(i\varphi) \) with \( \varphi = q_0 |\psi| \) and \( q_0 = h \varepsilon_0 / \gamma h \) (\( l \)-coordinate along the wire). This solution, however, corresponds to the absence of superconducting current

\[
j = 2i e h \gamma (\psi^* \frac{\partial \psi}{\partial l} - \psi \frac{\partial \psi^*}{\partial l}) - 4 e h \varepsilon_0 |\psi|^2
\]

\[
= 4 e h \gamma |\psi|^2 (\frac{\partial \varphi}{\partial l} - q_0) = 0.
\]

Boundary condition at the interface between modulated (\( \varepsilon \neq 0 \)) and usual superconducting phase (\( \varepsilon = 0 \)) corresponds simply to the continuity of the current. Let us consider as a generic example of the superconducting wire loop consisting of the region of the length \( L_h \) with exchange field and \( L_0 \) without exchange field (see Fig. 5b). The solution in the first region \( |\psi| \exp(iq_1 l) \) and the corresponding current \( j_1 = 4 e h \gamma |\psi|^2 (q_1 - q_0) \). The solution in the second region \( |\psi| \exp(iq_2 l) \) and the corresponding current \( j_2 = 4 e h \gamma |\psi|^2 q_2 \). From the continuity of the current we have \( q_1 - q_0 = q_2 \). The total phase accumulation over the loop will be \( \Delta \varphi = q_1 L_h + q_2 L_0 \) and should be equal to \( 2\pi n \). From this condition we find \( q_2 = (2\pi n - q_0 L_h) / (L_h + L_0) \), what means the spontaneous generation of the current (except the special cases when the phase accumulation in active \( L_h \) region is integer of \( 2\pi \), i.e. \( q_0 L_h = 2\pi n \)). The actual choice of \( n \) is dictated by the minimization of the current \( j_1 = j_2 \) and so the associated magnetic energy, i.e. \( \min \{2\pi n - q_0 L_h \} \). Therefore we may conclude that the part of the wire with EXSO coupling serves as a “phase battery”.

In conclusion, we have demonstrated that the spin-orbit interaction in combination with a magnetic exchange field (EXSO) may induce non-dissipative currents in superconducting loops with a chirality that is controllable through the magnetization alignment. Superconducting elements with EXSO may generate spontaneous phase shifts governed by magnetic moment orientation and by introducing the EXSO wire \( L_h \) in a DC SQUID (see Fig. 5b) we can mimic the role of external flux \( 2\pi n \Phi / \Phi_0 = q_0 L_h \), which opens the way to creating “quiet” qubits [23], similar to that which has been recently fab-
ricated on the basis of the $\pi$–junction. Suitable FI materials for the proposed experiments here include EuS and GdN and potentially oxides such as praseodymium calcium manganese oxide. Finally, we note that the EXSO interaction between magnetic and superconducting subsystems may, in principle, generate an inverse effect in which the superconducting transition provokes reorientation of the magnetic moment.

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