Cooper Pairs Spintronics in Triplet Spin Valves

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We study a spin valve with a triplet superconductor spacer intercalated between two ferromagnets with noncollinear magnetizations. We show that the magnetoresistance of the triplet spin valve depends on the relative orientations of the **d** vector, characterizing the superconducting order parameter, and the magnetization directions of the ferromagnetic layers. For devices characterized by a long superconductor, the effects of a polarized current sustained by Cooper pairs only are observed. In this regime, a supermagnetoresistance effect emerges, and the chiral symmetry of the order parameter of the superconducting spacer is easily recognized. Our findings open new perspectives in designing spintronics devices based on the cooperation of ferromagnetic and triplet correlations.

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Introduction.—The realization of devices with synthetic materials offers interesting technological opportunities. Within this class of devices, the spin valves (SVs) are a notable example. They are made of two ferromagnetic layers separated by a nonmagnetic spacer and show a significant change in the electrical resistance depending on whether the magnetizations of ferromagnetic regions are in a parallel or an antiparallel configuration. This effect, more evident in multilayer structures displaying the giant magnetoresistance effect [1–3], finds application in the information technology industry (e.g., sensors for hard disk drives, magnetic-random-access memory).

The spacer properties strongly affect the SV electric response and, thus, various SV devices have been proposed to study the phase coherent transport through spacers made of nanowires [4], carbon nanotubes [5], ballistic low-dimensional systems [6], and singlet-superconducting regions [7,8]. All these devices, however, contain a scalar spacer, i.e., a middle layer unable to add (nonmagnetic) vectorial quantities relevant in determining the magnetoresistive response of the SV.

Since triplet superconductors (TSCs) can support polarized currents, they represent natural candidates to study SVs having spacers with exotic (i.e., nonscalar) properties. Since the discovery of triplet superconductivity in Sr_2RuO_4 [9], there has been growing interest in the properties of TSC heterostructures [10–19]. Despite this, the study of devices combining TSCs and ferromagnets is still in its infancy and only few unconventional effects have been predicted [13,18–20].

In this Letter, we study the magnetoresistance (MR) properties of a SV whose spacer is a triplet superconductor (see upper panel, Fig. 1). This system is the prototype of a SV having a vectorial spacer. We demonstrate that triplet superconducting spin valves can transmit polarization information by means of dissipationless Cooper pairs (CPs) current in the absence of a quasiparticle (QP) contribution which is instead dominant in SVs with singlet superconducting spacers. The emergence of a magnetoresistance which does not decay with the superconducting spacer length d_s (see Fig. 1) over distances of several times the

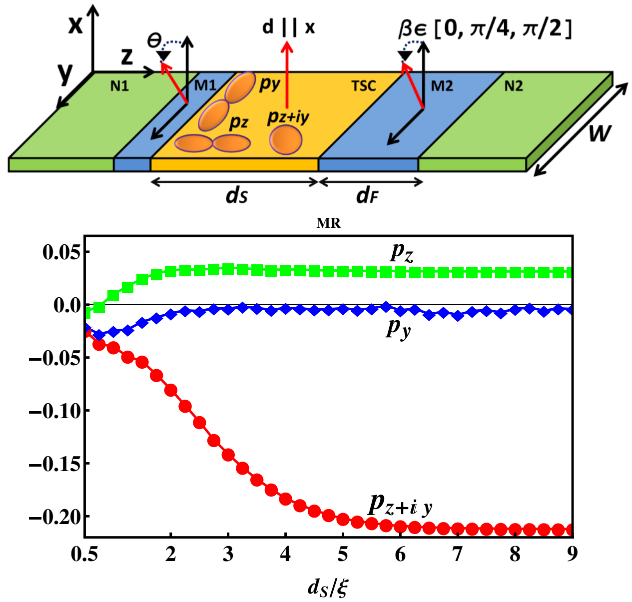

FIG. 1 (color online). (Upper panel) Triplet-spin-valve device. A TSC is intercalated between the nanostructured magnetic regions $M1$ and $M2$ whose magnetic momenta $\vec{M}_1$ and $\vec{M}_2$, belonging to the $x-y$ plane, form the angle θ and $\beta \in [0, \pi/4, \pi/2]$ with respect to the **d** vector (parallel to the x direction) characterizing the superconducting state. The system is biased by means of nonmagnetic leads $N1$ and $N2$ having a transverse dimension W. (Lower panel) Magnetoresistance curves as a function of the superconductor length d_S. The model parameters are fixed as follows: $\varepsilon/\Delta = 0.01$, $\mathbf{h} = 0.65$, $\Gamma = 1.5$, $Z_{\rm BTK} = 1$, $\theta = \pi/2$, $\beta = \pi/4$, $d_F = \xi/10$. Differently from an s-wave spin valve, the MR is a nonvanishing function of d_S.

coherence length ξ is here called the supermagnetoresistance effect (SMRE) [21]. Under the supermagnetoresistance effect, the SV response does not exhibit the conventional Julliere-like [22] behavior, according to which the MR depends on the relative orientation of the polarizations $\vec{M}_1$, $\vec{M}_2$ of the magnetic regions (i.e., MR $\propto \vec{M}_1 \cdot \vec{M}_2$). The violation of the Julliere-like behavior is the distinctive feature of the novel class of SVs having a vectorial spacer.

Theoretical description.—We consider the triplet SV (TSV) of Fig. 1 and combine the Bogoliubov–de Gennes (BdG) approach with the scattering matrix formalism for describing the transport properties along the z direction. The Bogoliubov–de Gennes wave function $\Psi(r)$ describing the system is determined by the eigenvalue problem ($z \neq 0, d_S, d_S + d_F$)

$$\begin{bmatrix} \hat{H}(r) & \hat{\Delta}(r) \\ \hat{\Delta}^\dagger(r) & -\hat{H}^*(r) \end{bmatrix} \Psi(r) = E\Psi(r), \tag{1}$$

where the hat sign indicates 2×2 matrices in spin space, while the single particle Hamiltonian can be written as follows:

$$\hat{H}(r) = \left[-\frac{\hbar^2\nabla^2}{2m} - E_F + V_{\text{int}}(r) \right] \hat{\mathbf{1}} - g\mu_B \hat{\sigma} \cdot \mathbf{M}(r). \tag{2}$$

Here, we introduced an interface potential controlling the barrier transparencies $V_{\text{int}}(r) = U[\delta(z) + \delta(z - d_S) + \delta(z - d_S - d_F)]$ and the magnetic fields describing the ferromagnetic regions $M1$ and $M2$

$$\mathbf{M}(r) = M_1 \delta(z)\mathbf{v}(\theta) + M_2(r)\mathbf{v}(\beta), \tag{3}$$

$$\mathbf{v}(\phi) = \cos(\phi)\mathbf{e}_x + \sin(\phi)\mathbf{e}_y, \tag{4}$$

$\mathbf{e}_{x/y/z}$ being the orthogonal triad of unit vectors. The region $M1$, corresponding to the so-called free layer, is assumed to be very narrow and, thus, is modeled using a Dirac delta potential whose amplitude is proportional to the magnetic momentum M_1. The function $M_2(r)$, describing the magnetization of the region $M2$, is taken spatially homogeneous inside the region and zero elsewhere.

The gap matrix in Eq. (1) is defined as $\hat{\Delta}(r) = i[\hat{\sigma} \cdot \mathbf{d}(r)]\hat{\sigma}_y$, where $\mathbf{d}(r)$ is the vector defining the order parameter of the TSC. Here, we are interested in the case of the equal-spin-pairing unitary state for which $\mathbf{d}(r) = \Delta(r)\mathbf{e}_x$ represents a convenient choice [23]. With this assumption, the triplet Cooper pairs have a z component of the spin $S_z = \pm\hbar$, while the condensate has zero net spin polarization. The magnitude of the gap Δ is assumed to be constant throughout the superconducting region and zero elsewhere. In the following, we consider the three possible orbital pairing states: p_y wave, $\Delta_k = \Delta k_y/k_F$; p_z wave, $\Delta_k = \Delta k_z/k_F$; and the chiral p_{z+iy} wave, $\Delta_k = \Delta[k_z + ik_y]/k_F$ [24]. We assume the translational invariance along the y direction implying the conservation of the wave vector k_y parallel to the interface. Thus, the wave function can be written as $\Psi(r) = e^{ik_y y}\psi(z|E, k_y)$ leading to an effective one-dimensional scattering problem for $\psi(z|E, k_y)$, being the energy E and k_y conserved quantum numbers.

The conductance of the system can be obtained within the scattering field theory [25] where one defines the field

$$\hat{\Psi}_j(z, t|k_y) = \sum_{\beta,\sigma} \int \frac{dEe^{-iEt}}{\sqrt{2\pi\hbar|v_z(E)|}} |\beta\rangle \otimes |\sigma\rangle \times [\hat{a}^\sigma_{j\beta}(E; k_y)e^{ik^{(z)}_\beta z} + \hat{b}^\sigma_{j\beta}(E; k_y)e^{-ik^{(z)}_\beta z}]. \tag{5}$$

The scattering operators $\hat{a}^\sigma_{j\beta}(E; k_y)$ [$\hat{b}^\sigma_{j\beta}(E; k_y)$] destroy an incoming (outgoing) particle of species $\beta \in \{e, h\}$ and spin projection $\sigma \in \{\uparrow, \downarrow\}$ in the lead $j \in \{L, R\}$, while the wave vector $k^{(z)}_\beta = \eta_\beta k_z(E) = \eta_\beta |k(E)| \cos(\alpha_{\text{in}})$ ($\eta_e = -\eta_h = 1$) along the transport direction is written within the Andreev approximation [26], i.e., $v^{j\beta\sigma}_z(E) \approx v_z(E) = \hbar|k(E)|\cos(\alpha_{\text{in}})/m$, where α_{in} is the incidence angle. The scattering matrix $\mathcal{S}$ relates the incoming and the outgoing processes through the equation

$$\hat{b}^\sigma_{j\beta}(E; k_y) = \sum_{\beta'\sigma'j'} \mathcal{S}^{\beta\beta'}_{jj'\sigma\sigma'}(E; k_y)\hat{a}^{\sigma'}_{j'\beta'}(E; k_y), \tag{6}$$

and its elements are obtained by matching the wave functions of the regions $N_{1,2}$, TSC, and $M2$ imposing the boundary conditions (BCs) at the interfaces [27]. For instance, the discontinuity at $z = 0$ described by the local potential $[U - g\mu_B M_1 \hat{\sigma} \cdot \mathbf{v}(\theta)]\delta(z)$ implies the following boundary conditions: (i) $\psi(z = 0^+|E, k_y) = \psi(z = 0^-|E, k_y)$; (ii) $\partial_z\psi(z = 0^+|E, k_y) - \partial_z\psi(z = 0^-|E, k_y) = k_F[1_{4\times4}Z_{\text{BTK}} - \Gamma\mathcal{A}(\theta)]\psi(z = 0^+|E, k_y)$, where we have introduced the matching matrix

$$\mathcal{A}(\theta) = \begin{bmatrix} \hat{\sigma} \cdot \mathbf{v}(\theta) & 0 \\ 0 & \hat{\sigma}^* \cdot \mathbf{v}(\theta) \end{bmatrix}, \tag{7}$$

the Blonder-Tinkham-Klapwijk (BTK) parameter [28] $Z_{\text{BTK}} = 2mU/(\hbar^2 k_F)$ and the spin-active barrier strength $\Gamma = 2m[g\mu_B M_1]/(\hbar^2 k_F)$. The $\mathcal{S}$ matrix depends on the incidence angle α_{in} through the conserved quantity $k_y = |k(E)| \sin(\alpha_{\text{in}})$.

In the limit of low-bias $eV/\Delta \ll 1$, the linear response current I_i flowing through the ith lead is given by the sum of independent contributions of the elementary processes labeled by all possible k_y, i.e., $I_i = \sum_{j,k_y} g_{ij}(k_y)(\mu_j - \mu_s)$, where g_{ij} is the conductance tensor; μ_j is the electrochemical potential of the jth lead while in the superconducting region μ_s is fixed by imposing the electric charge conservation. In the symmetric case ($g_{11} = g_{22}$), the two-terminal conductance is given by $\mathcal{G} = (\bar{g}_{11} - \bar{g}_{12})/2$ with the definition $\bar{g}_{ij} = \sum_{k_y} g_{ij}(k_y)$ [29]; it can be explicitly written in terms of the scattering matrix elements as follows [30]:

$$\mathcal{G} = \frac{e^2 k_F W}{\pi h} \int d\xi d\alpha_{\rm in}[-\partial_\xi f(\xi)]_{eq}\left[\frac{\cos(\alpha_{\rm in})}{2}\right] \times\left[\sum_{\beta\in\{e,h\}} \mathcal{M}_{12}^{\beta\beta}(\xi,\alpha_{\rm in}) + \mathcal{M}_{11}^{\rm he}(\xi,\alpha_{\rm in}) + \mathcal{M}_{11}^{\rm eh}(\xi,\alpha_{\rm in})\right], \quad (8)$$

where $\mathcal{M}_{ij}^{\alpha\beta} = \mathrm{Tr}[\mathcal{S}_{ij}^{\alpha\beta\dagger}\mathcal{S}_{ij}^{\alpha\beta}]$, the trace being performed over the spin. When the incidence angle $\alpha_{\rm in}$ varies in the interval $[-\zeta, \zeta]$, the sum rule $\sum_{j\alpha;k_y}\mathcal{M}_{ij}^{\beta\alpha}(E,k_y) = 2\mathcal{N}_\perp$ is obeyed with the transverse modes number given by $\mathcal{N}_\perp = k_F W \sin(\zeta)/\pi$.

Results.—Here, we define the magnetoresistance of a TSV as $\mathrm{MR} = 1 - \mathcal{G}(\theta,\beta)\mathcal{G}(\theta=\beta,\beta)^{-1}$, being $\mathcal{G}(\theta=\beta,\beta)$ the conductance in parallel configuration of the magnetic momenta of the ferromagnetic regions and θ (β) the angle formed by $M1$ ($M2$) with the $\mathbf{d}$ vector. We consider the zero-temperature limit which is appropriate to study SVs containing TSCs whose typical critical temperatures T_c are quite low. For instance in Sr_2RuO_4, $T_c \approx 0.7$–1.4 K and the zero temperature coherence length $\xi(T=0) \approx 70$ nm [31]. In the computation, we set the maximal incidence angle $\zeta = 35°$ to mimic the angular dispersion of electron waves coming from a remote constriction. The ferromagnetic region $M2$ is made of a weak ferromagnet characterized by normalized Zeeman energy $\mathbf{h} = (g\mu_B M_2)/E_F$ such that $\mathbf{h} \in [0.05, 0.65]$, while the spin-dependent wave vector $k_z^\sigma \approx k_F\sqrt{\cos^2(\alpha_{\rm in}) + \sigma\mathbf{h}}$ is a real quantity for both spin polarizations.

In Fig. 2, the magnetoresistance curves as a function of the angle θ, for short ($d_S = 2\xi$) and long ($d_S = 5\xi$) spacers, are shown. The magnetization direction of the region $M2$ is fixed to $\beta = 0$ [Figs. 2(a) and 2(b)], $\beta = \pi/4$ [Figs. 2(c) and 2(d)], and $\beta = \pi/2$ [Figs. 2(e) and 2(f)]. The magnetoresistance curves pertaining to the symmetries p_z and p_{z+iy} shown in Figs. 2(a), 2(c), and 2(e) are described by the fitting function $\mathrm{MR}(\theta,\beta) = \mathcal{F}(\theta) - \mathcal{F}(\beta)$, where $\mathcal{F}(x) = \sum_{n=1,2}\mathcal{B}_n\cos^n(2x)$ is a symmetry-dependent function. The p_{z+iy} symmetry is well described by the function $\mathcal{F}(x)$ with $\mathcal{B}_1 \approx 0.2$ and $\mathcal{B}_2 = 0$; the case of p_z is instead described by $\mathcal{B}_1 \approx (0.6$–$0.9) \times 10^{-2}$ and $\mathcal{B}_2/\mathcal{B}_1 \approx 5$–$9$. For both cases, $\mathrm{MR}(\theta,\beta)$ is a separable function of θ and β, while for short superconducting regions [see Figs. 2(b), 2(d), and 2(f) for $d_S = 2\xi$] the magnetoresistance curves show a complicated behavior

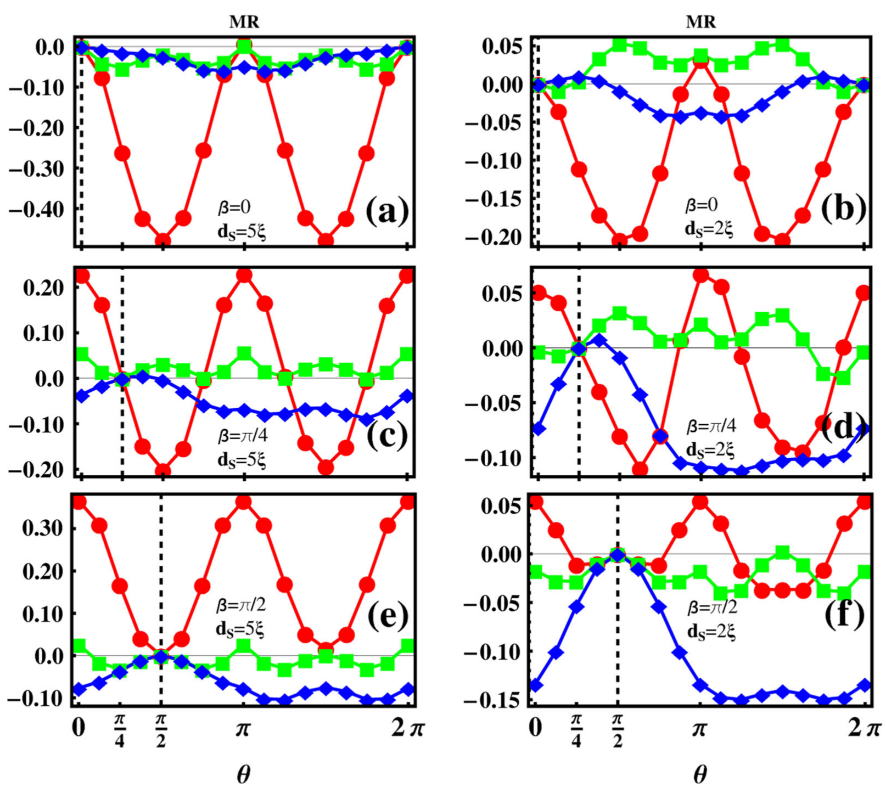

FIG. 2 (color online). Magnetoresistance curves as a function of the angle θ computed for the order parameter symmetry p_z (green filled square), p_y (blue filled diamond), p_{z+iy} (red filled circle). The curves (a), (c), (e) are computed for $d_S = 5\xi$, while the curves (b), (d), (f) are computed by setting $d_S = 2\xi$. The magnetization of the region $M2$ is fixed to form the angle: $\beta = 0$ with the $\mathbf{d}$ vector for the panels (a)–(b); $\beta = \pi/4$ for the panels (c)–(d); $\beta = \pi/2$ for the panels (e)–(f) (see the vertical dashed line). The remaining model parameters are fixed as follows: $\varepsilon/\Delta = 0.01$, $\mathbf{h} = 0.65$, $\Gamma = 1.5$, $d_F = \xi/10$, $Z_{\rm BTK} = 1$. The $\mathrm{MR}(\theta,\beta)$ curves show a period halving from 2π periodicity (for $d_S = 2\xi$) to π (for $d_S = 5\xi$). The period halving is not observed for the p_y symmetry whose periodicity is always 2π.

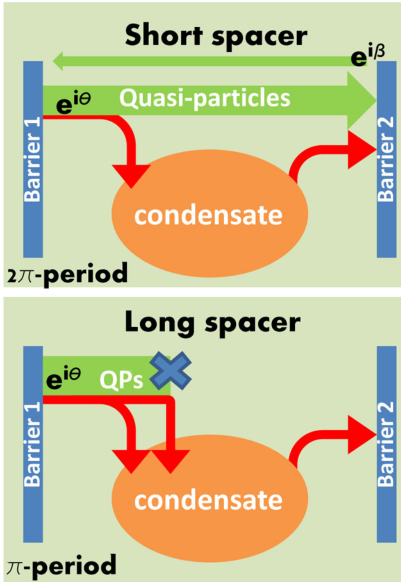

FIG. 3 (color online). A quasiparticle excitation coming from the first spin-active barrier can interfere with the wave reflected from the second barrier only for sufficiently short systems. This because the quasiparticle phase memory is limited to the lifetime $\tau \sim \hbar/\Delta$ of the excitation, while it is lost when the excitation is recombined to form a Cooper pair. Since the quasiparticle phase memory can be retained over a distance $\ell = v_F\tau \sim \xi$, an unconventional π periodicity of the MR curves is observed for spacer length greater than $\sim 2\xi$.

induced by interference effects due to partial conversion of the quasiparticles current into Cooper pairs current.

Moreover, the MR shows a peculiar π periodicity for the chiral and p_z symmetry of the order parameter, while a 2π periodicity is always found in the p_y case. The nonconventional π periodicity of the MR can be explained by the fact that the quasiparticles and the Cooper pairs current contribute differently to the MR: the latter has a π periodicity, while the quasiparticles flux presents a 2π periodicity. The MR curves of the p_y symmetry always present a periodicity of 2π with respect to θ. This is related to the existence of a gapless line ($\alpha_{\text{in}} = 0$) along the transport direction where a quasiparticle's flux always coexists with a Cooper pair current [32].

In the case of p_z and p_{z+iy} symmetry, for short channel ($d_S = 2\xi$) the quasiparticle's current significantly contributes to the conductance determining a dominant 2π periodicity. On the other hand, by increasing the length of the superconducting spacer up to $d_S = 5\xi$ a bulk-like behavior, characterized by a dissipationless Cooper pairs current, is established [32] and a π periodicity is observed. The latter regime is not observable for SVs having a singlet superconducting spacer where the polarized transport is always sustained by a quasiparticle's current. Under this condition, the Julliere-like behavior $\mathrm{MR}(\theta, \beta) \propto \cos(\theta - \beta)$ is obeyed [22].

An additional fingerprint of the triplet correlations is represented by the high magnetoresistance values compared to the s-wave case [8]. The above interpretation can be validated by the analysis of the Andreev reflection (AR) probability [33], i.e., the dominant scattering event at the normal-superconductor interface in which an incoming electron is reflected as a hole. For a single spin active interface (SAI), the AR probability is periodic of π with respect to the magnetization angle θ but the AR coefficient presents a phase factor $\exp(i\theta)$ [33]. This phase can play a role if a second spin active interface is added to the system to form a SV. In fact, for short spacers, the phase difference between the scattering coefficients of the two barriers induces a dominant 2π periodicity. For long spacers, the quasiparticle phase memory is limited to the lifetime $\tau \sim \hbar/\Delta$ of the excitation and, thus, the π period emerges (see Fig. 3). As a consequence, the system behaves like a sequence of two distant spin-active regions ($M1$ and $M2$) interacting only through the condensate. The latter statement explains the separability of the function $\mathrm{MR}(\theta, \beta) = \mathcal{F}(\theta) - \mathcal{F}(\beta)$ and the breakdown of the Julliere relation.

Conclusions.—In conclusion, a SV modified by the inclusion of a triplet superconducting spacer may represent a novel system for spintronics, displaying an unconventional magnetoresistive response. Differently from SVs with normal or s-wave spacers, a TSV shows a magnetoresistive behavior which depends on the relative orientations of the three vectors M_1, M_2, and the $\mathbf{d}$ vector. A nonvanishing symmetry-dependent MR, supermagnetoresistance, has been obtained for long spacers ($d_S \approx 5\xi$) allowing the study of spintronic properties completely determined by the dissipationless spin polarized currents sustained by the Cooper pairs. Experimentally, the regime of pure Cooper pairs spintronics is signaled by a π periodicity of the MR vs θ curves for the chiral and the p_z symmetries; a 2π periodicity is always found in the p_y case.

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[30] In the general nonsymmetric case [8], the conductance is given by $\mathcal{G} = [\bar{g}_{22}\bar{g}_{11} - \bar{g}_{21}\bar{g}_{12}]/(\sum_{ij}\bar{g}_{ij})$, where $\bar{g}_{ik} = (e^2 k_F W/\pi h)\int d\xi d\alpha_{\rm in}[-\partial_\xi f(\xi)]_{eq}[\cos(\alpha_{\rm in})/2] \times [4\delta_{ik} + \mathcal{M}^{he}_{ik}(\xi, \alpha_{\rm in}) + \mathcal{M}^{eh}_{ik}(\xi, \alpha_{\rm in}) - \sum_\beta \mathcal{M}^{\beta\beta}_{ik}(\xi, \alpha_{\rm in})]$.

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[33] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.111.226801 for (i) the analysis of the MR vs **h** curves; (ii) a microscopic description of the transport properties of a normal-TSC interface from the Andreev reflection perspective; (iii) a quantitative explanation of the unconventional π periodicity; (iv) the analysis of the invariance of the transport properties under unitary transformations in spin space.