PHYSICAL REVIEW B **83**, 180504(R) (2011)

Spin Josephson effect with a single superconductor

P. M. R. Brydon,[1,*] Yasuhiro Asano,[2] and Carsten Timm[1,†]
[1]*Institut für Theoretische Physik, Technische Universität Dresden, D-01062 Dresden, Germany*
[2]*Department of Applied Physics, Hokkaido University, Sapporo 060-8628, Japan*
(Received 21 April 2011; published 19 May 2011)

A thin ferromagnetic layer on a bulk equal-spin-pairing triplet superconductor is shown to mediate a Josephson coupling between the spin ↑ and ↓ condensates of the superconductor. By deriving analytic expressions for the bound states at the triplet superconductor-ferromagnet interface, we show that this spin Josephson effect establishes an effective anisotropy axis in the ferromagnetic layer. The associated Josephson spin current is predicted to cause a measurable precession of the magnetization about the vector order parameter of the triplet superconductor.

DOI: 10.1103/PhysRevB.83.180504 PACS number(s): 74.50.+r, 74.20.Rp, 74.70.Xa, 74.70.Pq

Introduction. The complex relationship between superconductivity and magnetism has motivated an enormous effort to understand the properties of heterostructure interfaces between ferromagnets (FMs) and spin-singlet superconductors (SSCs).[1] A remarkable feature of such devices is the existence of proximity-induced spin-triplet superconducting correlations due to the exchange splitting in the FM, which is responsible for the anomalous dynamics of the barrier magnetization in an SSC-FM-SSC Josephson junction.[2] Due to the intimate connection between ferromagnetism and triplet superconductivity, it is natural to consider what results if the SSC was replaced by a triplet superconductor (TSC). This question is of fundamental interest, as the intrinsic spin structure of the Cooper pairs in a TSC allows us to anticipate an unconventional and unique interplay with magnetism, which may be unambiguous signatures of the triplet pairing state in proposed TSCs, such as LiFeAs.[3] For example, bulk spin supercurrents are known to be possible in TSCs,[4–7] and some proposals for their realization require FM elements. The response of the FM component of the device to the spin supercurrent, however, has yet to be investigated. The recent fabrication of superconducting thin films of the suspected TSC Sr_2RuO_4 is an important step toward the creation of TSC-FM heterostructures,[8,9] and so a deeper understanding of the physics of TSC-FM interfaces is timely.

In this Rapid Communication, we show that a thin FM layer on a bulk equal-spin-pairing TSC produces a spin Josephson effect by coupling the spin ↑ and ↓ Cooper pair condensates. The physical mechanism is the spin-dependent phase shift acquired by a Cooper pair undergoing spin-flip reflection at the FM interface, which acts analogously to the phase difference in a Josephson junction. Making only the assumption of spatially constant order parameters, we solve the Bogoliubov-de Gennes (BdG) equations for the bound states at the TSC-FM interface. Using this, we calculate the free energy of the interface, revealing that the spin Josephson effect creates an effective hard or soft axis within the FM layer, depending upon the orbital structure of the TSC gap. Finally, we obtain a general expression for the spin current, thereby showing that it exerts a measurable torque on the FM moment. We propose this effect as a test for triplet pairing.

Model system. We consider a FM layer of width L on a bulk TSC, separated by an atomically thin insulating layer (see Fig. 1). The BdG equation for the quasiparticle states with energy E is

$$\begin{pmatrix} \hat{H}_0(\mathbf{r}) & \hat{\Delta}(\mathbf{r}) \\ \hat{\Delta}^\dagger(\mathbf{r}) & -\hat{H}_0^T(\mathbf{r}) \end{pmatrix} \Psi(\mathbf{r}) = E\Psi(\mathbf{r}), \tag{1}$$

where the caret indicates a 2×2 matrix in spin space. The wave function $\Psi(\mathbf{r})$ is only nonzero for $z<L$. The noninteracting Hamiltonian $\hat{H}_0(\mathbf{r})$ is

$$\hat{H}_0(\mathbf{r}) = \left[\left(-\frac{\hbar^2\nabla^2}{2m} - \mu\right) + U\delta(z)\right]\hat{\mathbf{1}} - \mu_B\hat{\boldsymbol{\sigma}}\cdot\mathbf{H}_{\rm ex}\Theta(z). \tag{2}$$

Here, μ_B is the Bohr magneton, and we make the simplifying but nonessential assumptions that the radius of the Fermi surface k_F and the effective mass m in the normal state of the TSC and the FM are the same. The insulating layer is modeled as a δ function of strength U. The last term in Eq. (2) is the energy due to the exchange field $\mathbf{H}_{\rm ex}$ in the FM. For an incompletely polarized FM, the magnetization $\mathbf{m} = |\mathbf{m}|[\cos(\alpha)\mathbf{e}_x + \sin(\alpha)\cos(\eta)\mathbf{e}_y + \sin(\alpha)\sin(\eta)\mathbf{e}_z]$ is related to $\mathbf{H}_{\rm ex}$ by Luttinger's theorem. We assume a two-dimensional system, so that the majority-spin (parallel to $\mathbf{H}_{\rm ex}$, $s=+$) and minority-spin (antiparallel to $\mathbf{H}_{\rm ex}$, $s=-$) Fermi surfaces have radius $k_{F,s} = \sqrt{(1+s\lambda)}k_F$, where $\lambda = \mu_B|\mathbf{H}_{\rm ex}|/\mu < 1$. This gives $|\mathbf{m}| = \mu_B^2 m|\mathbf{H}_{\rm ex}|/\pi\hbar^2$. For a half-metallic FM, there is a single Fermi surface of radius $\sqrt{2}k_F$, and the exchange splitting is fixed by the details of the system.

The gap matrix is $\hat{\Delta}(\mathbf{r}) = \Theta(-z)i[\hat{\sigma}\cdot\mathbf{d}]\hat{\sigma}^y$, where $\mathbf{d} = \widetilde{\Delta}\hat{\mathbf{d}}$ is the vector order parameter, assumed constant throughout the TSC. $\widetilde{\Delta}$ is an operator, which for Cooper pairs in a relative p-wave orbital state has the real space form $\widetilde{\Delta} = -i\Delta(T)\mathbf{n}\cdot\nabla/k_F$. The gap magnitude $\Delta(T)$ is assumed to have weak-coupling temperature dependence. The unit vector $\mathbf{n}$ defines the orbital state: $\mathbf{n} = \mathbf{e}_z$ for p_z wave; $\mathbf{n} = \mathbf{e}_y$ for p_y wave; and $\mathbf{n} = \mathbf{e}_z + i\mathbf{e}_y$ for $(p_z + ip_y)$ wave. The last choice is of greatest relevance to Sr_2RuO_4 and LiFeAs,[3,8] while the others have been proposed for $(TMTSF)_2X$ (TMTSF stands for tetramethyltetraselenafulvalene, X = PF_6, ClO_4).[10] In the following, it is convenient to express the gap in terms of the Fourier transform of $\widetilde{\Delta}$, which is written $\Delta_{\mathbf{k}} = \Delta(T)\mathbf{n}\cdot\mathbf{k}/k_F$. We fix $\hat{\mathbf{d}} = \mathbf{e}_x$ which defines a TSC where the Cooper pairs have z component of spin $S_z = \pm\hbar$, but the condensed part of the system is unpolarized. Below we show that only the angle

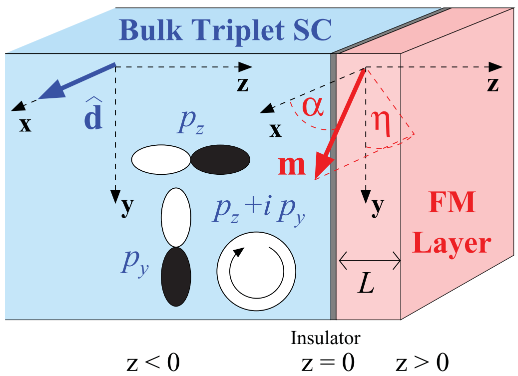

FIG. 1. (Color online) Schematic diagram of the device studied here. The three different choices for the TSC orbital are shown: for the p_z and p_y cases, the white and black lobes indicate opposite signs; the arrow in the $p_z + ip_y$ shows the direction of increasing phase.

between $\hat{\mathbf{d}}$ and $\mathbf{m}$ is relevant for the spin Josephson effect, and so other orientations of $\hat{\mathbf{d}}$ do not result in new physics.

Bound states. We seek solutions of Eq. (1) for states bound to the FM layer. The wave function of such a state has the general form $\Psi(\mathbf{k}_\parallel;\mathbf{r}) = \Psi_{\rm TSC}(\mathbf{k}_\parallel;\mathbf{r})\Theta(-z) + \Psi_{\rm FM}(\mathbf{k}_\parallel;\mathbf{r})\Theta(z)$ and satisfies $\lim_{z\to-\infty}\Psi(\mathbf{k}_\parallel;\mathbf{r}) = 0$ and $\Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=L} = 0$. The momentum component parallel to the interface, $\mathbf{k}_\parallel$, is a good quantum number due to translational invariance. Solving the Andreev equations in the TSC,[11,12] we make the ansatz

$$\Psi_{\rm TSC}(\mathbf{k}_\parallel;\mathbf{r}) = \sum_{\sigma=\uparrow,\downarrow}[a_{1,\sigma}\Psi_\sigma(\mathbf{k}_1;\mathbf{r}) + a_{2,\sigma}\Psi_\sigma(\mathbf{k}_2;\mathbf{r})], \quad (3)$$

where the spinors are given by

$$\Psi_\uparrow(\mathbf{k};\mathbf{r}) = (1,\ 0,\ \gamma(\mathbf{k}),\ 0)^T e^{i\mathbf{k}\cdot\mathbf{r}}e^{\kappa_{\mathbf{k}}z}, \quad (4)$$

$$\Psi_\downarrow(\mathbf{k};\mathbf{r}) = (0,\ 1,\ 0,\ -\gamma(\mathbf{k}))^T e^{i\mathbf{k}\cdot\mathbf{r}}e^{\kappa_{\mathbf{k}}z}, \quad (5)$$

with $\gamma(\mathbf{k}) = -[E + i\,\mathrm{sgn}(k_z)\sqrt{|\Delta_\mathbf{k}|^2 - E^2}]/\Delta_\mathbf{k}$ and $\kappa_\mathbf{k} = (m/\hbar^2|k_z|)\sqrt{|\Delta_\mathbf{k}|^2 - E^2}$. The wave vectors appearing in Eq. (3) are defined by $\mathbf{k}_1 = (\mathbf{k}_\parallel, k_z)$, $\mathbf{k}_2 = (\mathbf{k}_\parallel, -k_z)$. Note that $|\Delta_{\mathbf{k}_1}| = |\Delta_{\mathbf{k}_2}| \equiv |\Delta_{\mathbf{k}_\parallel}|$ for the orbital symmetries considered here.

Depending upon the value of $\mathbf{k}_\parallel$, we have either propagating or evanescent solutions in the FM layer. In the case when there are propagating solutions in both spin channels, we have

$$\Psi_{\rm FM}(\mathbf{k}_\parallel;\mathbf{r}) = \sum_{s=\pm}\{b_{e,s}\sin(k_{e,s}[L-z])e^{i\mathbf{k}_\parallel\cdot\mathbf{r}}\Phi_{e,s} + b_{h,s}\sin(k_{h,s}[L-z])e^{i\mathbf{k}_\parallel\cdot\mathbf{r}}\Phi_{h,s}\}, \quad (6)$$

where the electron and hole spinors are defined by

$$\Phi_{e,s} = (w_s,\ x_s,\ 0,\ 0)^T, \quad \Phi_{h,s} = (0,\ 0,\ w_s^*,\ x_s)^T, \quad (7)$$

with $w_s = s(\cos\alpha - i\sin\alpha\cos\eta)/\sqrt{1 - s\sin\alpha\sin\eta}$, $x_s = \sqrt{1 - s\sin\alpha\sin\eta}$, and the wave vector $k_{e(h),s}$ for electrons (holes) is $k_{e(h),s} = [k_F^2(1+s\lambda) - |\mathbf{k}_\parallel|^2 + (-)2mE/\hbar^2]^{1/2}$. If the radicand is negative, only evanescent solutions are possible; in this case, we replace $k_{e(h),s}$ with $\kappa_{e(h),s}$, where $\kappa_{e(h),s}$ is the inverse decay length.

The coefficients in Eqs. (3) and (6) are chosen so that at the TSC-FM interface, the wave function is continuous, i.e, $\Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=0^-} = \Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=0^+}$, and its derivative obeys $\partial_z\Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=0^+} - \partial_z\Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=0^-} = 2Z\Psi(\mathbf{k}_\parallel;\mathbf{r})|_{z=0^+}$, where $Z = mU/\hbar^2$. The values of E for which the determinant of the resulting system of equations vanishes define the bound-state energies. Explicit expressions for the bound-state energies can be found when the E dependence of the wave vectors is neglected, i.e., $k_{e,s} \approx k_{h,s} \approx k_s$. This approximation is valid for a thin FM layer, such that $(k_{e,s} - k_{h,s})L \approx 2EL/\hbar v_{F,s} \ll 1$.[12,13] For a weakly to moderately polarized FM layer, we have Fermi velocities $v_{F,+} \approx v_{F,-} \sim 10^6$ ms^{-1}, and so for $E \leqslant \max\{|\Delta_\mathbf{k}|\} \sim 0.1$ meV ($T_c \sim 1$ K), we require thin layers less than about 100 unit cells thick. In this limit, we obtain the nondegenerate bound states:

$$E_{\pm,\mathbf{k}_\parallel} = \pm|\Delta_{\mathbf{k}_\parallel}|\sqrt{D_{\mathbf{k}_\parallel}}\,|\cos\alpha|, \quad p_z \text{ wave} \quad (8a)$$

$$E_{\pm,\mathbf{k}_\parallel} = \pm|\Delta_{\mathbf{k}_\parallel}|\sqrt{1 - D_{\mathbf{k}_\parallel}\cos^2\alpha}, \quad p_y \text{ wave} \quad (8b)$$

$$E_{\pm,\mathbf{k}_\parallel} = -|\Delta_{\mathbf{k}_\parallel}|\left[\sqrt{1 - D_{\mathbf{k}_\parallel}\cos^2\alpha}\,\frac{k_y}{k_F} \pm \sqrt{D_{\mathbf{k}_\parallel}}\cos\alpha\,\frac{k_z}{k_F}\right]. \quad (p_z + ip_y) \text{ wave} \quad (8c)$$

Here, we have

$$D_{\mathbf{k}_\parallel} = 4\left[\sum_{s=\pm} s\widetilde{k}_s\cos(k_sL)\sin(k_{-s}L)\right]^2 \times \prod_{s=\pm}\left[1 + 4\widetilde{Z}^2 + \widetilde{k}_s^2 + 4\widetilde{k}_s\widetilde{Z}\sin(2k_sL) + \left(\widetilde{k}_s^2 - 4\widetilde{Z}^2 - 1\right)\cos(2k_sL)\right]^{-1}, \quad (9)$$

where $\widetilde{A} = A/\sqrt{k_F^2 - |\mathbf{k}_\parallel|^2}$ $(A = k_s, Z)$.

The bound-state energies [Eq. (8)] are a central result of our Rapid Communication. They originate due to multiple Andreev reflections within the thin FM layer, which Josephson couple the $S_z = \pm\hbar$ condensates in the TSC. The same physical mechanism is responsible for the formation of Andreev bound states (ABSs) at the tunneling barrier in a Josephson junction.[12,14,15] Remarkably, the states [Eq. (8)] are identical to the (spin degenerate) ABSs in a short Josephson junction of transparency $D_{\mathbf{k}_\parallel}$ between p_z-wave TSCs with phase difference $\Delta\phi = 2\alpha$ [Eq. (8a)], between p_y-wave TSCs with $\Delta\phi = \pi + 2\alpha$ [Eq. (8b)], and between a $(p_z + ip_y)$-wave and $(p_z - ip_y)$-wave TSC with $\Delta\phi = 2\alpha$ [Eq. (8c)]. Since the form of the bound-state energies is fixed by the bulk pairing symmetry, our results should be robust to a self-consistent calculation of the gap.[12]

The spin Josephson coupling can also be understood at a more fundamental level: in the BdG Hamiltonian, the $S_z = \pm\hbar$ condensates in the TSC are independent of one another. Tunneling of a Cooper pair between the two condensates is made possible by the FM layer, where the coupling to the FM moment allows an incident Cooper pair with spin $\sigma\hbar$ to be reflected with spin $-\sigma\hbar$. As a result of this process, the Cooper pair acquires a phase shift $\Delta\theta_{\mathbf{k}_\parallel} + \pi - 2\sigma\alpha$. The last terms are due to the spin flip itself and are primarily

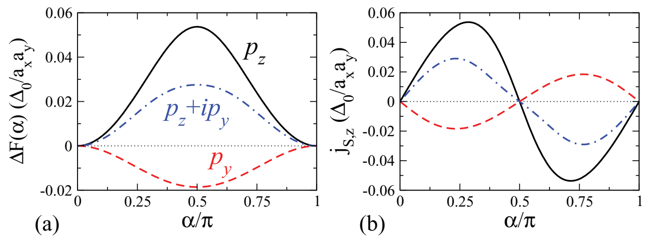

FIG. 2. (Color online) (a) Free energy difference $\Delta F(\alpha)$ and (b) the z component of the spin current $j_{S,z}$ per interface-unit-cell area as a function of α for the three choices of orbital wave function. We take $\lambda = 0.05$, $T = 0.4T_c$, $\eta = 0$, $L = 10a_z$, and $Z = 1$. Δ_0 is the $T = 0$ gap magnitude.

responsible for driving the spin current. $\Delta\theta_{\mathbf{k}_\parallel} = \arg\{\Delta_{\mathbf{k}_2}\} - \arg\{\Delta_{\mathbf{k}_1}\}$ is the phase shift due to the orbital structure of the TSC: for p_z-wave orbitals, we have $\Delta\theta_{\mathbf{k}_\parallel} = \pi$ for all $\mathbf{k}_\parallel$; in the p_y-wave case, we have $\Delta\theta_{\mathbf{k}_\parallel} = 0$; and the superposition of these two orbitals in the $(p_z + ip_y)$-wave TSC gives $\Delta\theta_{\mathbf{k}_\parallel} = \pi - 2\arccos(\sqrt{1 - |\mathbf{k}_\parallel|^2/k_F^2})$. $\Delta\theta_{\mathbf{k}_\parallel}$ accounts for the π phase difference between the bound states in the p_z and p_y cases, and the apparent sign reversal of the p_y component in the $p_z + ip_y$ bound states.

We note that unlike the standard Josephson effect, where the phase difference between the two superconductors drives the supercurrent, the spin Josephson effect is due to phase shifts picked up *during* the tunneling process itself. The phase difference between the spin-up and spin-down condensates is fixed by the orientation of $\hat{\mathbf{d}}$, which is unaffected by the FM layer.

Free energy. The assumption of spatially constant order parameters allows us to write the free energy due to the spin Josephson coupling in the TSC-FM device in terms of the bound states:[15]

$$F = -\frac{1}{2}k_B T \sum_{n=\pm}\sum_{\mathbf{k}_\parallel} \ln[2\cosh(\beta E_{n,\mathbf{k}_\parallel}/2)] + F_0, \quad (10)$$

where F_0 is independent of α and includes the interaction energy in both the TSC and the FM, as well as the contribution from continuum states. We plot the free energy difference $\Delta F(\alpha) = F(\alpha) - F(\alpha = 0)$ per interface-unit-cell area $a_x a_y$ as a function of α in Fig. 2(a). As can be seen, F takes a minimum as a function of the angle α: regarding $\hat{\mathbf{d}}$ as a fixed property of the bulk TSC, the spin Josephson coupling, therefore, establishes a preferred orientation for the magnetization of the FM layer. There is a direct analogy to a short Josephson junction, where the free energy due to the Josephson coupling takes a minimum as a function of $\Delta\phi$.[15]

In the p_z-wave (p_y-wave) case, the formally equivalent Josephson junction with $\Delta\phi = 2\alpha$ ($\Delta\phi = \pi - 2\alpha$) has time-reversal symmetry, and so the Josephson free energy is minimized at $\Delta\phi = 0$. This implies that in the TSC-FM device, the free energy *always* has a minimum at $\alpha = 0$ ($\alpha = \pi/2$). The broken time-reversal symmetry in the $(p_z + ip_y)$-wave case, however, means that the stable value of α is determined by the details of $D_{\mathbf{k}_\parallel}$. Specifically, the p_z and p_y components of the gap favor minima at different values of α: if $D_{\mathbf{k}_\parallel}$ is peaked near $|\mathbf{k}_\parallel| = k_F$ (0), the p_y component (p_z component) dominates and the configuration with $\alpha = \pi/2$ (0) is stable; for more complicated $D_{\mathbf{k}_\parallel}$, the competition between the gap components may stabilize the system at $\alpha \neq 0, \pi/2$. For weak magnetization strengths, the free energy minimum is usually located at $\alpha = 0$.

From Fig. 2(a), we see that to excellent approximation $F \propto \cos 2\alpha$. In writing an effective free energy for the FM layer, we can, therefore, account for the spin Josephson effect by including a term $F_J = f_s(\hat{\mathbf{d}} \cdot \mathbf{m})^2$, i.e., for $f_s < 0$ ($f_s > 0$), $\hat{\mathbf{d}}$ defines an effective easy (hard) axis in the FM layer. Since the sign of f_s is determined by the *orbital* state of the TSC, this reveals an unconventional spin-orbit coupling between the TSC and the FM.

Spin current and magnetization dynamics. The zero-bias charge current I_J in a Josephson junction is given by $I_J = (2e/\hbar)\partial F/\partial\Delta\phi$. We now show that in our device, there is a spontaneous spin current which can be similarly expressed as a derivative of F with respect to α.

Our starting point is the continuity equation for the spin

$$\mathbf{J}_s = \frac{\hbar}{g\mu_B}\frac{d}{dt}\mathbf{M} = \mathbf{M} \times \frac{\partial F}{\partial \mathbf{M}}, \quad (11)$$

where g is the gyromagnetic ratio, and $\mathbf{M}$ is the total moment of the FM layer. The vector notation for the spin current $\mathbf{J}_s$ refers *only* to the polarization; the direction of the spin current is normal to the interface. We omit gradient terms in Eq. (11) as the FM layer is considered to be thin compared to the coherence length of the Cooper pairs. The free energy in our problem has the form $F = F(|\mathbf{M}|, \hat{\mathbf{d}} \cdot \hat{\mathbf{M}})$, allowing us to write

$$\begin{aligned}\mathbf{J}_s &= \mathbf{M} \times \left(\frac{\partial|\mathbf{M}|}{\partial\mathbf{M}}\frac{\partial F}{\partial|\mathbf{M}|} + \frac{\partial\hat{\mathbf{d}}\cdot\hat{\mathbf{M}}}{\partial\mathbf{M}}\frac{\partial F}{\partial\hat{\mathbf{d}}\cdot\hat{\mathbf{M}}}\right)\\ &= \mathbf{M} \times \left(\hat{\mathbf{M}}\frac{\partial F}{\partial|\mathbf{M}|} + \frac{1}{|\mathbf{M}|}[\hat{\mathbf{d}} - (\hat{\mathbf{d}}\cdot\hat{\mathbf{M}})\hat{\mathbf{M}}]\frac{\partial F}{\partial\cos\alpha}\right)\\ &= \hat{\mathbf{p}}\frac{\partial F}{\partial\alpha},\end{aligned} \quad (12)$$

where $\hat{\mathbf{p}} = \hat{\mathbf{d}} \times \mathbf{M}/|\mathbf{M}|\sin\alpha$ is a unit vector which points in the same direction for all $\mathbf{M}$ lying in a fixed plane containing $\hat{\mathbf{d}}$. Inserting Eq. (10) into Eq. (12), we obtain

$$\mathbf{J}_s = -\hat{\mathbf{p}}\frac{1}{4}\sum_{n=\pm}\sum_{\mathbf{k}_\parallel}\frac{\partial E_{n,\mathbf{k}_\parallel}}{\partial\alpha}\tanh(\beta E_{n,\mathbf{k}_\parallel}/2). \quad (13)$$

This closely resembles the Beenakker-van Houten formula for the charge current in a short Josephson junction.[15] It reveals that the spin current in our device is due entirely to resonant tunneling between the two spin condensates through the bound states [Eq. (8)]. Equation (13) gives identical results to the Furusaki-Tsukada technique,[4,7,16] which expresses the spin current in terms of the Andreev reflection coefficients. We show the spin current as a function of α in Fig. 2(b).

If the magnetization is prepared with $0 < \alpha < \pi/2$, Eq. (11) predicts that the spin current will exert a torque on $\mathbf{M}$, causing it to precess about $\hat{\mathbf{d}}$. Writing $\mathbf{M} = AL\mu_B pn\hat{\mathbf{m}}$ and $\mathbf{J}_s = A\mathbf{j}_s$, where A is the area of the TSC-FM interface, and $p = (n_+ - n_-)/n$ is the polarization of the FM, we find the precession frequency to be $\Omega_J = 2g\cos(\alpha)\max\{|\mathbf{j}_s|\}/\hbar npL$. To estimate Ω_J, we assume a weakly polarized FM, $\lambda = 0.05$, with $n = 1$ electron per unit volume $v = a_x a_y a_z$, $\max\{|\mathbf{j}_s|\} =$

$0.025\Delta_0/(a_x a_y)$ [see Fig. 2(b)], $L = 10a_z$, and $T_c = 1$ K. We, hence, find $\Omega_J = 15\cos(\alpha)$ GHz, which is measurable by ferromagnetic resonance (FMR) experiments. As the spin Josephson effect cannot occur for a SSC, the observation of this precession would be very strong evidence of a triplet pairing state, although the precession effects due to multiple FM domains would have to be ruled out.[17] Similarly, it is also necessary to examine the effect of chiral domains of the gap in the candidate material Sr_2RuO_4.[18] Gilbert damping and anisotropy effects in the FM layer must also be included in a complete description of the magnetization dynamics, but do not change the derivation of Eq. (13).

Conclusions. In this Rapid Communication, we have demonstrated that the spin structure of the Cooper pairs in a TSC permits the occurrence of a spin Josephson effect without the need for a second superconductor. We have proposed that a thin FM layer on a bulk TSC can realize this effect. In turn, the spin Josephson coupling establishes an effective easy or hard axis in the FM layer, depending upon the orbital symmetry of the TSC gap. Furthermore, the Josephson spin current causes the magnetization to precess about the $\hat{\mathbf{d}}$ vector with a frequency that is accessible to FMR, realizing a possible experimental signature of the triplet state.

The authors thank D. Manske, J. A. Sauls, and M. Sigrist for useful discussions. Y.A. was supported by KAKENHI on Innovative Areas "Topological Quantum Phenomena" No. 22103002 and KAKENHI No. 22540355 from MEXT of Japan.

---

[*]brydon@theory.phy.tu-dresden.de

[†]carsten.timm@tu-dresden.de

[1]A. I. Buzdin, Rev. Mod. Phys. **77**, 935 (2005); F. S. Begeret, A. F. Volkov, and K. B. Efetov, *ibid.* **77**, 1321 (2005); M. Eschrig, J. Kopu, J. C. Cuevas, and G. Schon, Phys. Rev. Lett. **90**, 137003 (2003).

[2]X. Waintal and P. W. Brouwer, Phys. Rev. B **65**, 054407 (2002); J.-X. Zhu, Z. Nussinov, A. Shnirman, and A. V. Balatsky, Phys. Rev. Lett. **92**, 107001 (2004); V. Braude and Y. M. Blanter, *ibid.* **100**, 207001 (2008); F. Konschelle and A. Buzdin, *ibid.* **102**, 017001 (2009); J. Linder and T. Yokoyama, Phys. Rev. B **83**, 012501 (2011).

[3]P. M. R. Brydon, M. Daghofer, C. Timm, and J. van den Brink, Phys. Rev. B **83**, 060501(R) (2011); S.-H. Baek, H.-J. Grafe, F. Hammerath, M. Fuchs, L. Harnagea, S. Wurmehl, J. van den Brink, and B. Büchner (unpublished).

[4]Y. Asano, Phys. Rev. B **72**, 092508 (2005); **74**, 220501(R) (2006).

[5]J. Linder, M. S. Grønsleth, and A. Sudbø, Phys. Rev. B **75**, 024508 (2007).

[6]P. M. R. Brydon, B. Kastening, D. K. Morr, and D. Manske, Phys. Rev. B **77**, 104504 (2008); P. M. R. Brydon, D. Manske, and M. Sigrist, J. Phys. Soc. Jpn. **77**, 103714 (2008); P. M. R. Brydon and D. Manske, Phys. Rev. Lett. **103**, 147001 (2009).

[7]P. M. R. Brydon, Phys. Rev. B **80**, 224520 (2009).

[8]A. P. Mackenzie and Y. Maeno, Rev. Mod. Phys. **75**, 657 (2003).

[9]Y. Krockenberger, M. Uchida, K. S. Takahashi, M. Nakamura, M. Kawasaki, and Y. Tokura, Appl. Phys. Lett. **97**, 082502 (2010).

[10]A. G. Lebed, K. Machida, and M. Ozaki, Phys. Rev. B **62**, R795 (2000).

[11]C.-R. Hu, Phys. Rev. Lett. **72**, 1526 (1994).

[12]S. Kashiwaya and Y. Tanaka, Rep. Prog. Phys. **63**, 1641 (2000).

[13]Analytic expressions can also be found in the opposite limit $L \gg \hbar v_{F,s}/2E$, although these only hold far from the gap edges [see A. M. Zagoskin, *Quantum Theory of Many-Body Systems* (Springer, New York, 1998)].

[14]Y. S. Barash, A. M. Bobkov, and M. Fogelström, Phys. Rev. B **64**, 214503 (2001); H.-J. Kwon, K. Sengupta, and V. M. Yakovenko, Eur. Phys. J. B **37**, 349 (2004).

[15]C. W. J. Beenakker and H. van Houten, in *Nanostructures and Mesoscopic Systems*, edited by W. P. Kirk and M. A. Reed (Academic, New York, 1992).

[16]A. Furusaki and M. Tsukada, Solid State Commun. **78**, 299 (1991).

[17]R. Grein, M. Eschrig, G. Metalidis, and G. Schon, Phys. Rev. Lett. **102**, 227005 (2009).

[18]F. Kidwingira, J. D. Strand, D. J. Van Harlingen, and Y. Maeno, Science **314**, 1267 (2006); A. Bouhon and M. Sigrist, New J. Phys. **12**, 043031 (2010).