Magneetoanisotropic Andreev Reflection in Ferromagnet-Superconductor Junctions
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Spin-orbit coupling (SOC) is a key interaction in spintronics [1–3], allowing an electrical control of magnetization and, vice versa, a magnetic control of electrical current. In systems lacking space inversion symmetry—be it bulk, hybrid structures, junctions—SOC induces spin-orbit fields [1, 2] as an emergent phenomenon. We are in particular concerned here with interfacial spin-orbit fields which are believed to be behind a wealth of new phenomena, not existent or fragile in the bulk, such as the tunneling anisotropic magnetoresistance (TAMR) [4–7], interfacial spin-orbit torques [8], or skyrmions [9].

Interfacial spin-orbit fields are also important in semiconductor/superconductor [10–13] and F/S junctions [14] for creating Majorana quasiparticle states. It is the latter junctions that we focus on. We investigate the interplay of magnetism and spin-orbit fields. We show that this interplay leads to marked anisotropies in the junction conductance. If the ferromagnet is highly spin polarized—in the half-metal limit—the magnetoanisotropic Andreev reflection depends universally on the spin-orbit fields only. Our results show that Andreev reflection spectroscopy can be used for sensitive probing of interfacial spin-orbit fields in F/S junction.

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is reflected as a hole with the opposite spin, while spin-flip AR implies equal spin of the incoming and reflected particles. These two AR processes, see Figs. 1(b) and (f), introduce, respectively spin-singlet and spin-triplet superconducting correlations at the interface [22, 23].

We consider epitaxial-quality junctions, such as those used in TAMR [6], or point contact geometries [35], in which ballistic transport formalism is applicable. In diffusive tunnel junctions AR could be enhanced by electron-hole coherence [36]. In ferromagnetic junctions such effects would be absent for normal AR due to short coherence length, but spin-flip AR could be enhanced. (Ordinary effects of diffusion could be accounted for by renormalizing the tunneling parameters [32]). We generalize the BTK formalism [37] and solve the Bogoliubov-de Gennes equation [38] for quasiparticle states $\Psi (\mathbf{r})$ with energy $E$,

$$
\left( \hat{H}_F - \Delta \right) \Psi (\mathbf{r}) = E \Psi (\mathbf{r}),
$$

with the single-particle Hamiltonian for electrons $\hat{H}_F = -(\hbar^2/2) \nabla [1/m(\mathbf{r})] \nabla - \mu(\mathbf{r}) - (\Delta_{xc}/2) \Theta(\mathbf{r}) \mathbf{m} \cdot \hat{\sigma} + \hat{H}_B$; for holes $\hat{H}_h = -\hat{\sigma}_y \hat{H}_F \hat{\sigma}_y$. The unit magnetization vector (see Fig. 1) is $\mathbf{m} = [\sin \Theta \cos \Phi, \sin \Theta \sin \Phi, \cos \Theta]$, $\hat{\sigma}$ are Pauli matrices, $\Delta_{xc}$ is the exchange spin splitting in the F region (Stoner model), $m(\mathbf{r})$ is the effective mass, and $\mu(\mathbf{r})$ is the chemical potential. The interfacial scattering is modeled as $\hat{H}_B = (V_0 d + \mathbf{w} \cdot \hat{\mathbf{\sigma}}) \delta(\mathbf{r})$, where $V_0$ and $d$ are the barrier height and width, while $\mathbf{w} = [(\alpha-\beta)k_y, -(\alpha+\beta)k_x, 0]$ is the effective SOC field including Rashba and Dresselhaus terms [1, 2], parametrized by $\alpha$ and $\beta$, respectively, for the crystallographic orientations see Fig. 1. The superconducting pair potential is given by $\Delta = \Delta \Theta(\mathbf{r}) \delta_{\mathbf{r} \cdot \mathbf{u}}$ (the accuracy of such a step-function form of $\Delta$ is discussed in [39]), with the isotropic gap $\Delta$. Similar methodology, for half-metal/S junctions with Rashba coupling inside the superconductor was employed in Ref. [40]. With Rashba-only SOC one should still obtain out-of-plane magnetooisotropy, and this is already implicit in this formalism [40].

Since the in-plane wave vector $\mathbf{k}_\| \in \Sigma$ is conserved, $\Psi(\mathbf{r}) = \Psi(\mathbf{z}) \mathbf{e}^{i\mathbf{k}_\| \cdot \mathbf{r}}$. The solution in the F region for incoming electrons with spin $\sigma$ is

$$
\Psi^{\sigma}_\sigma = \frac{1}{\sqrt{\kappa_\|}} e^{i k_\| z} \chi^e_{\sigma} + r^{e,\sigma}_{\sigma,-\sigma} e^{-i k_\| z} \chi^e_{\sigma,-\sigma} + r^{c,\sigma}_{\sigma,-\sigma} e^{-i k_\| z} \chi^c_{\sigma,-\sigma} + r^{h,\sigma}_{\sigma,-\sigma} e^{-i k_\| z} \chi^h_{\sigma,-\sigma},
$$

with the spinors for the electron-like $\chi^e_{\sigma} = (\chi_{\sigma}, 0)^T$ and hole-like $\chi^h_{\sigma} = (0, \chi_{-\sigma})^T$ quasiparticles, both containing

$$
\chi_{\sigma}^T = (\sigma \sqrt{1 + \sigma^2} \cos \Theta e^{-i \Phi}, \sqrt{1 - \sigma^2} \cos \Theta) / \sqrt{2},
$$

where $\sigma = 1(-1)$ corresponds to the spin parallel (antiparallel) to $\mathbf{m}$. The electron-like (hole-like) quasiparticle wave vectors in the F region are $k_{\sigma}^{(h)} = \sqrt{k_{\|}^2 + 2m_F/h^2 (\Delta E + \sigma \Delta_{xc}/2) - k_{\|}^2}$.

In the S region the scattering states are

$$
\Psi^S_{\sigma} = t^{e,\sigma}_{\sigma,-\sigma} e^{iq \cdot z} \begin{pmatrix} u \\ v \end{pmatrix} + t^{c,\sigma}_{\sigma,-\sigma} e^{iq \cdot z} \begin{pmatrix} 0 \\ u \end{pmatrix} + t^{h,\sigma}_{\sigma,-\sigma} e^{-iq \cdot z} \begin{pmatrix} v \\ 0 \end{pmatrix},
$$

where the quasiparticle wave vectors are given by $q^{(h)}_{\sigma} = \sqrt{q_{\|}^2 + (\Delta E - 2m_S/h^2 \sqrt{E^2 - \Delta^2 - k_{\|}^2})}$. The superconducting coherence factors satisfy $u^2 = 1 - v^2 = (1 + \sqrt{E^2 - \Delta^2}/E)/2$.

Using charge current conservation, the differential conductance at zero temperature, normalized by the Sharvin conductance $[1] G_{sh} = e^2 k_F^2 A/(2\pi h)$ of a perfect contact, is

$$
G = \sum_{\sigma} \int d^2k_{\|} \left[ 1 + R^{h}_{\sigma}(-eV) - R^{c}_{\sigma}(eV) \right],
$$

containing the probability amplitudes in the F region $R^{(h)}_{\sigma}(E, k_{\|}) = \text{Re} \left[ r^{(h)}_{\sigma,-\sigma} r^{(h)}_{\sigma,-\sigma}^2 + k^{(h)}_{\sigma,-\sigma} r^{(h)}_{\sigma,-\sigma}^2 \right]$ which combine the coefficients for the scattering processes with and without spin flip for specular and AR; $V$ is the bias voltage and $A$ is the interfacial area.

To describe our results we introduce dimensionless quantities: $Z = V_0 d/\sqrt{m_F m_S}/(h^2 \sqrt{k_F q_F})$ denotes the barrier strength [31, 37], $\lambda_s = 2\sqrt{m_F m_S}/h^2$ and $\lambda_S = 2\sqrt{m_F m_S}/h^2$ quantify the Rashba and Dresselhaus SOC, and $P = (\Delta_{xc}/2)/\mu_F$ defines the spin polarization in F.

We first examine the influence of SOC on the F/S conductance (see Fig. 2), for a metallic point contact ($Z = 0$) and for a moderate barrier ($Z = 1$). For the former case the conductance tends to decrease with increasing SOC. Even in the half-metallic case ($P = 1$) SOC does not give a finite subgap conductance; spin-flip AR is suppressed. In contrast, for moderate barrier ($Z = 1$), SOC enhances the conductance due to spin-flip AR, even for $P = 1$. Interestingly, at $eV = \Delta$ the conductance is not affected by SOC for any $Z$. Focusing on $G(0)$, Fig. 2 shows that in a metallic contact increasing SOC steadily reduces $G(0)$, while for a moderate barrier $G(0)$ is a nonmonotonic function of SOC, with a (P-dependent) maximum which turns out to be due to spin-flip AR.

The absence of spin-flip AR in metallic contacts can be explained analytically. For $eV \leq \Delta$ quasiparticle transmission is prohibited and subgap conductance $G \sim \sum_{\sigma} \int d^2k_{\|} 2R^{h}_{\sigma}(-eV)$. In the half-metallic case the only contribution to AR comes from spin-flip AR,
FIG. 2. (Color online) Calculated normalized conductance, G(eV/Δ), for different (indicated) spin polarizations P. Rashba SOC is λα = 2 and Dresselhaus SOC is λβ = 0. Magnetization is in-plane (θ = π/2). (a) No interfacial barrier (Z = 0), and (b) modest interfacial barrier (Z = 1) cases are shown. The dashed lines show G without SOC. The insets show the dependence of G(0) on Rashba SOC. In (a), the subgap conductance for P = 1 vanishes for every λα,β; in the inset G(0) = 0 for this case. The illustrations summarize main qualitative impacts of SOC on conductance.

We define the in-plane MAAR as

\[ \text{MAAR}_{[110]}(Φ) = \frac{G(Θ, 0) - G(Θ, Φ)}{G(Θ, Φ)} \bigg|_{Θ=90^\circ}, \]  \hspace{1cm} (6)

and the out-of-plane MAAR as

\[ \text{MAAR}_{[10]}(Θ) = \frac{G(0, Φ) - G(Θ, Φ)}{G(Θ, Φ)} \bigg|_{Φ=-90^\circ}. \]  \hspace{1cm} (7)

The out-of-plane MAAR depends, in general on Φ, but we choose the yz (Φ = −90◦) plane as its reference. The calculated MAAR, in Fig. 3, shows a nonmonotonic dependence on SOC. For metallic contacts (Z = 0) MAAR is determined by the magnetoanisotropy of conventional AR. In the presence of a barrier (exemplified by Z = 1), MAAR gets strongly enhanced due to the additional contribution from spin-flip AR. In-plane MAAR exhibits C2v symmetry due to the interplay of Rashba and Dresselhaus fields, similarly to TAMR [2, 6, 7]. If either of the two fields is absent, in-plane MAAR vanishes. In contrast, out-of-plane MAAR is finite even with the Rashba field alone, which makes it a robust probe of this important interfacial SOC. Interestingly, at eV = Δ MAAR is always absent, as there are no effects of SOC.

While the conductance changes are indicative of interfacial SOC, magnetic anisotropy of the conductance is a true fingerprint. As the main contribution comes from AR, we call this anisotropy effect magnetoanisotropic Andreev reflection. We consider two configurations: in-plane, in which magnetization m changes azimuthally (Φ) in the interfacial plane, and out-of-plane, with polar (Θ) changes of m in a perpendicular plane (see Fig. 1).
on $G$ here; see the discussion to Fig. 2. Additional effects (such as appearance of symmetry lobes) can arise due to the effective mass and Fermi wave vector mismatch [41]. Compared to TAMR, the magnitude of MAAR is giant, varying by orders of magnitude upon changing the spin polarization $P$. (The experimentally measured in-plane TAMR in Fe/GaAs/Au junctions is less than a percent [6].) A detailed model comparison is shown in Fig. 4 for both in- and out-of-plane configurations; TAMR is evaluated by setting $\Delta = 0$. For a typical $P$ of 40%, the ratio MAAR/TAMR is about 10. Moving towards half metals ($P \gtrsim 80\%$), this ratio climbs to more than $10^2$. This giant increase is best illustrated in the half-metallic limit of $P = 1$. For a weak SOC (which is typically the case) an analytical treatment gives [41],

$$MAAR_{[110]}(\Phi = \pi/2) \approx \frac{2\lambda_\alpha \lambda_\beta (1 - \cos 2\Phi)}{\lambda_\alpha^2 + \lambda_\beta^2 + 2\lambda_\alpha \lambda_\beta \cos 2\Phi}.$$ 

(8)

$$MAAR_{[111]}(\Theta = \pi/2) \approx \frac{(\lambda_\alpha + \lambda_\beta)^2 (1 - \cos 2\Theta)}{3(\lambda_\alpha^2 + \lambda_\beta^2 - 2\lambda_\alpha \lambda_\beta + (\lambda_\alpha + \lambda_\beta)^2 \cos 2\Theta)}.$$ 

(9)

Therefore, the in-plane MAAR$_{[\pi/2]}(\Phi = \pi/2) \approx 4\lambda_\alpha \lambda_\beta / (\lambda_\alpha - \lambda_\beta)^2$, and out-of-plane MAAR$_{[\pi/2]}(\Theta = \pi/2) \approx (\lambda_\alpha + \lambda_\beta)^2 / (\lambda_\alpha - \lambda_\beta)^2$, depending universally on the spin-orbit fields only, and diverging as $\lambda_\alpha \approx \lambda_\beta$ (see the in-plane case in Fig. 4). In contrast, TAMR, which is proportional to the product $\lambda_\alpha \lambda_\beta$ [7], has no singular behavior, and is not a universal function of $\lambda_\alpha$ only.

We can trace this giant enhancement of MAAR over TAMR to spin-flip AR. Let us separate phenomenologically the conductance $G = G^{(0)} + G_{so}$ into the sum of SOC independent and dependent parts. In TAMR typically $G^{(0)} \gg G_{so}$, and TAMR $\sim G_{so}/G^{(0)} \ll 1$, even for $P \approx 1$. But in F/S junctions $G^{(0)}$ decreases with increasing $P$, eventually vanishing in the half-metallic limit. For $P \approx 1$ the conductance of the F/S junction is dominated by the spin-flip AR contribution to $G_{so}$. Thus, SOC determines both the conductivity and the magnetotransparency. Furthermore, if $\lambda_\alpha \approx \pm \lambda_\beta$, the spin-flip AR, and so the conductance, can be switched on and off by changing the orientation of $\mathbf{m}$. For $\lambda_\alpha = \lambda_\beta$ and $\Phi = 0$, $\mathbf{m} \parallel \mathbf{w}$ and spin-flip AR yields a finite $G$. However, if $\Phi = \pi/2$, then $\mathbf{m} \parallel \mathbf{w}$ and spin-flip processes are strongly suppressed; $G(eV \leq \Delta)$ at $\Phi = \pi/2$ vanishes. As a result, in-plane MAAR diverges if $\lambda_\alpha = \lambda_\beta$. Similarly for out-of-plane MAAR.

There is one more peculiarity of MAAR in the half-metallic limit. If only Rashba (or only Dresselhaus) SOC is present, MAAR has a fixed universal magnitude of 100%. This case is shown for the out-of-plane configuration in Fig. 4 (in particular the inset for $\lambda_\alpha \lesssim 1$ shows MAAR of 100% for $P \approx 1$). It follows from Eq. (9) that MAAR$_{[\pi/2]}(\Theta = 0, \Theta = \pi/2)$, which gives a universal amplitude of 100% for $\Theta = \pi/2$. In other words, $G(\Theta = 0) = 2G(\Theta = \pi/2)$. The origin of this universal behavior is traced to the spin-flip probability by scattering of spin-polarized electrons off spin-orbit fields. The conductance is determined by spin-flip AR. For out-of-plane magnetization, $\Theta = 0$, two fields, one along $x$ and one along $y$, induce a spin flip. But for an in-plane magnetization, say along $x$, $\Theta = \pi/2$, only the spin-orbit field component along $y$ can flip the spin. This gives the 2:1 ratio in conductances and 100% of MAAR. A more technical and detailed discussion of the differences between MAAR and TAMR can be found in Ref. 41.

Experimental realization of MAAR could follow the measurement geometry of TAMR [6], ideally also the same junction, with the nonmagnetic metal that becomes superconducting at low temperatures. Magnetization of the ferromagnetic layer is typically rotated by an external magnetic field. This field can bring additional anisotropic orbital effects whose presence can be clearly identified from the field magnitude dependence [49]. However, one can avoid these extrinsic effects entirely if one uses dysprosium magnets which can be oriented by the field, but do not need its presence to remain in the rotated posi-
tion [50]. Potential aspects of non-flat tunneling barriers can also be treated [51]. A practical alternative (especially if ballistic junctions are desired) could be a point contact F/S junction geometry [35, 52].

To conclude, we have applied a well established theoretical formalism to systematically explore the magnetic anisotropy of the conductance in F/S junctions due to interfacial SOC. We predict a giant in- and out-of-plane MAAR—when compared with TAMR—exhibiting universal characteristics in the half-metallic regime. The predicted magnetization control of the AR suggests a similar control of the superconducting proximity effect and Majorana states. Our findings reveal an unexplored venue for AR spectroscopy, in the sensitive probing of interfacial SOC and related magnetoanisotropic phenomena.

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[1] I. Žutić, J. Fabian, and S. Das Sarma, Rev. Mod. Phys. 76, 323 (2004).
[2] J. Fabian, A. Matos-Abiague, C. Ertler, P. Stano, and I. Žutić, Acta Phys. Slov. 57, 565 (2007).
[3] S. Maekawa, Concepts in Spin Electronics (Oxford University Press, Oxford, 2006).
[4] L. Brey, C. Tejedor, and J. Fernandez-Rossier, Appl. Phys. Lett. 85, 1996 (2004).
[5] C. Gould, C. Rüster, T. Jungwirth, E. Girgis, G. M. Schott, R. Giraud, K. Brunner, G. Schmidt, and L. W. Molenkamp, Phys. Rev. Lett. 93, 117203 (2004).
[6] J. Moser, A. Matos-Abiague, D. Schuh, W. Wegscheider, J. Fabian, and D. Weiss, Phys. Rev. Lett. 99, 056601 (2007).
[7] A. Matos-Abiague and J. Fabian, Phys. Rev. B 79, 155303 (2009); A. Matos-Abiague, M. Gmitra, and J. Fabian, ibid. 80, 045312 (2009).
[8] K. S. Ryu, L. Thomas, S. H. Yang, and S. Parkin, Nanotech. 8, 527 (2013); A. Chernyshov, M. Overby, X. Liu, J. K. Furdyna, Y. Lyanda-Geller, and L. P. Rokhinson, Nature Phys. 5, 656 (2009).
[9] A. Fert, V. Cros, and J. Sampaio, Nature Nanotech. 8, 152 (2013).
[10] R. M. Lutchyn, J. D. Sau, and S. Das Sarma, Phys. Rev. Lett. 105, 077001 (2010).
[11] Y. Oreg, G. Refael, and F. von Oppen, Phys. Rev. Lett. 105, 177002 (2010).
[12] V. Mourik, K. Zuo, S. M. Frolov, S. R. Plissard, E. P. A. M. Bakkers, and L. P. Kouwenhoven, Science 336, 1003 (2012).
[13] L. P. Rokhinson, X. Liu, and J. K. Furdyna, Nature Phys. 8, 795 (2012).
[14] S. Nadji-Perge, I. K. Drozdov, J. Li, H. Chen, S. Jeon, J. Seo, A. H. MacDonald, B. A. Bernevig, and A. Yazdani, Science 346, 6209 (2014).
[15] Y. A. Bychkov and E. I. Rashba, J. Phys. C 17, 6039 (1984).
[16] G. Dresselhaus, Phys. Rev. 100, 580 (1955).
[17] M. Gmitra, A. Matos-Abiague, C. Draxl, and J. Fabian, Phys. Rev. Lett. 111, 036603 (2013).
[18] G. Deutscher, Rev. Mod. Phys. 77, 109 (2005).
[19] A. I. Buzdin, Rev. Mod. Phys. 77, 935 (2005).
[20] F. S. Bergeret, A. F. Volkov, and K. B. Efetov, Rev. Mod. Phys. 77, 1321 (2005).
[21] A. A. Golubov, M. Y. Kupriyanov, and E. Il’ichev, Rev. Mod. Phys. 76, 411 (2004).
[22] M. Eschrig, Phys. Today 64, 43 (2011).
[23] C. Visani, Z. Sefrioui, J. Tornos, C. Leon, J. Briadico, M. Bibes, A. Barthélémy, J. Santamaría, and J. E. Villégas, Nature Phys. 8, 539 (2012).
[24] S. Kashiwaya and Y. Tanaka, Rep. Prog. Phys. 63, 1641 (2000).
[25] M. J. de Jong and C. W. J. Beenakker, Phys. Rev. Lett. 74, 1657 (1995).
[26] R. J. Soulen, J. M. Byers, M. S. Ososky, B. Nadgorny, T. Ambrose, S. F. Cheng, P. R. Broussard, C. T. Tanaka, J. Nowak, J. S. Moodera, A. Barry, and J. M. D. Coey, Science 282, 85 (1998).
[27] S. K. Upadhyay, A. Palanisami, R. N. Louie, and R. A. Buhrman, Phys. Rev. Lett. 81, 3247 (1998).
[28] J. X. Zhu, B. Friedman, and C. S. Ting, Phys. Rev. B 59, 9558 (1999).
[29] S. Kashiwaya, Y. Tanaka, N. Yoshida, and M. R. Beasley, Phys. Rev. B 60, 3572 (1999).
[30] I. Žutić and S. Das Sarma, Phys. Rev. B 60, R16322 (1999).
[31] I. Žutić and O. T. Valls, Phys. Rev. B 60, 6320 (1999); 61, 1555 (2000).
[32] I. I. Mazin, A. A. Golubov, and B. Nadgorny, J. Appl. Phys. 89, 7576 (2001).
[33] K. Kikuchi, H. Imanura, S. Tulahashi, and S. Maekawa, Phys. Rev. B 65, 020508 (2001).
[34] C. S. Turel, I. J. Guilaran, P. Xiong, and J. Y. T. Wei, Appl. Phys. Lett. 81, 192508 (2011).
[35] P. Stamenov, J. Appl. Phys. 111, 07C519 (2012); 113, 17C718 (2013).
[36] F. W. J. Hekking and Y. V. Nazarov, Phys. Rev. Lett. 71, 1625 (1993).
[37] F. E. Blonder, M. Tinkham, and T. M. Klapwijk, Phys. Rev. B 25, 4515 (1982).
[38] P. G. De Gennes, Superconductivity of metals and alloys (Addison-Wesley, Reading MA, 1989).
[39] C. W. J. Beenakker, Rev. Mod. Phys. 69, 731 (1997).
[40] M. Duckheim and P. W. Brouwer, Phys. Rev. B 83, 054513 (2011).
[41] See Supplemental Material at [ ].
[42] L. P. Gor’kov and E. I. Rashba, Phys. Rev. Lett. 87, 037004 (2001).
[43] Y. Mizuno, T. Yokoyama, and Y. Tanaka, Phys. Rev. B 80, 195307 (2009).
[44] J. Linder and T. Yokoyama, Phys. Rev. Lett. 106, 237201 (2011).
[45] F. Dolcini and L. Dell’Anna, Phys. Rev. B 78, 024518 (2008).
[46] K. Sun and N. Shah, Phys. Rev. B 91, 144508 (2015).
[47] A. V. Balatsky, I. Vekhter, and J.-X. Zhu, Rev. Mod. Phys. 78, 373 (2006).
[48] S. Wu and K. V. Samokhin, Phys. Rev. B 80, 014516
(2009); 81, 214506 (2010).

[49] M. Wimmer, M. Lobenhöfer, J. Moser, A. Matos-Abiague, D. Schuh, W. Wegscheider, J. Fabian, K. Richter, and D. Weiss, Phys. Rev. B 80, 161201 (2009).

[50] C. Betthausen, T. Dollinger, H. Saarikoski, V. Kolkovsky, G. Karczewski, T. Wojtowicz, K. Richter, and D. Weiss, Science 337, 324 (2012).

[51] Z. Y. Chen, A. Biswas, I. Zutic, T. Wu, S. B. Ogale, R. L. Greene, and T. Venkatesan, Phys. Rev. B 63, 212508 (2001).

[52] J. Alicea, Rep. Prog. Phys. 75, 076501 (2012).