Resonant Tunneling Anisotropic Magnetoresistance Induced by Magnetic Proximity

Chenghao Shen, 1 Timothy Leeney, 1 Alex Matos-Abiague, 2 Benedikt Scharf, 1 Jong E. Han, 1 and Igor Žutić 1

1 Department of Physics, University at Buffalo, State University of New York, Buffalo, NY 14260, USA
2 Department of Physics and Astronomy, Wayne State University, Detroit, MI 48201, USA

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We reveal that the interplay between Rashba spin-orbit coupling and proximity-induced magnetization in a two-dimensional electron gas leads to peculiar transport properties and large anisotropy of magnetoresistance. While the related tunneling anisotropic magnetoresistance (TAMR) has been extensively studied before, we predict an effect with a different origin arising from the evolution of a resonant condition with the in-plane rotation of magnetization and having a much larger magnitude. The resonances in the tunneling emerge from a spin parity-time symmetry of the scattering states. However, such a symmetry is generally absent from the system itself and only appears for certain parameter values. Without resonant behavior in the topological surface states of a proximitized three-dimensional topological insulator (TI), TAMR measurements can readily distinguish them from often misinterpreted trivial Rashba-like states inherent to many TIs.

Tunneling magnetoresistance (TMR) has enabled remarkable advances in spintronic applications. While TMR devices require multiple ferromagnetic leads, in the presence of spin-orbit coupling (SOC) even a single ferromagnet (F) yields MR with the change of its magnetization direction. This makes the resulting tunneling anisotropic magnetoresistance (TAMR) a promising effect for scaled-down devices and design simplifications. Since the TAMR originates from the interplay between magnetization, M and SOC, it can also be used for experimentally probing emergent phenomena from interfacial SOC fields in both normal and superconducting heterostructures.

In common TAMR devices, a ferromagnetic lead generates spin polarization, P, which together with the SOC strength determines the transport anisotropy. In vertical tunneling devices the in-plane TAMR is rather small (typically < 1 %) even for large P and exchange energies > eV. In this work we propose a novel geometry in which a spin-polarized lead is not required, but the TAMR is enhanced. We consider tunneling through a gated barrier with proximity-induced magnetism from F on top of a two-dimensional electron gas (2DEG), shown in Fig. 1(a). Magnetic proximity effects offer a versatile method to transform materials with measured proximity-induced exchange energies up to tens of meV.

In the presence of Rashba SOC the transport through the magnetic barrier becomes anisotropic with respect to the magnetization M. Remarkably, the predicted in-plane TAMR is one to two orders of magnitude larger than in most TAMR vertical devices. Surprisingly, this is realized in the proposed device with nonmagnetic leads and the proximity-induced exchange splitting typically two orders of magnitude smaller than the exchange energy in the ferromagnetic lead of other TAMR devices. We find that the enhanced MR sensitivity to changes in M, even for small exchange fields, originates from the emergence of a spin parity-time (PT) symmetry of the scattering states, where T is the time reversal and P, the inversion of both position and spin, generalizing the well-known PT-symmetry. The emergence of the PT symmetry leads to resonances in the transmission, which are highly sensitive to M and eventually result in an enhanced TAMR.

As illustrated in Figs. 1(b) and (c), in the magnetic barrier region the Fermi contour of the 2DEG states is shifted perpendicular to M. The change of the barrier Fermi contour with respect to the Fermi contour of the leads when the in-plane M is varied affects the transmission rates and produces M-dependent changes in the

FIG. 1. (a) Schematic setup. (b) Band structure in the 2DEG and the barrier region (middle) of height V₀ and exchange field Δ. (c) Corresponding Fermi contours, where arrows denote the spin orientations. Dashed lines: the range of a conserved wavevector kₓ in the scattering states. For incident angles exceeding θ₀ backscattering is suppressed. (d) Action of the PT operator on an incident wave with an in-plane spin transforms the incident wave on the left side of the barrier (left panel) into itself, but as a a transmitted wave on the right side of the barrier (right panel). The magnetization orientation M, defined by the in-plane polar angle φ.
device resistance. Furthermore, the barrier Fermi contour undergoes a deformation that increases with the strength of the Rashba SOC. Such a deformation allows for lead-barrier Fermi contour matching enabling multiple states to achieve high-transmission rates. The realization of perfect transmission can be intuitively understood by considering the action of the \( \mathcal{P}_s \mathcal{T} = \mathcal{P} \sigma_z \mathcal{T} \) operator on an incident wave with a given in-plane spin, as schematically shown in Fig. 1(d). \( \mathcal{T} \) reverses both the spin and motion of the incident wave, while \( \mathcal{P}_s = \mathcal{P} \sigma_z \) inverts both the spin (through the action of the Pauli matrix, \( \sigma_z \)) and position (through the action of space inversion \( \mathcal{P} \)) of the wave. As a result, by applying the \( \mathcal{P}_s \mathcal{T} \) operator the incident wave on the left is transformed to itself, but as a transmitted wave on the right. Therefore, scattering states which are eigenfunctions of \( \mathcal{P}_s \mathcal{T} \) experience perfect transmission.

The model Hamiltonian of the system is given by

\[
H = p^2/2m^* + \alpha (\sigma \times p) \cdot \hat{z}/\hbar + [V_0 - \Delta(m \cdot \sigma)]h(x), \tag{1}
\]

where \( m^* \) is the effective mass, \( \alpha \) is the Rashba SOC strength, \( \hat{z} \) is the unit vector along the \( z \)-axis, \( p = (p_x, p_y) \) is the 2D momentum operator, \( \sigma \) is the vector of Pauli matrices, \( V_0 \) describes the potential barrier, \( \Delta \) and \( m \) are the magnitude and direction of the proximity-induced ferromagnetic exchange field. The function \( h(x) = \Theta(d/2 + x)\Theta(d/2 - x) \) describes a square barrier of thickness \( d \). We focus on electrons, not holes [26–29].

Due to Rashba SOC, the wavefunctions can be classified by the helicity index, where \( \lambda = 1 \) (–1) refers to the inner (outer) Fermi contour [recall Fig. 1(c)], and the zero-temperature and zero-bias longitudinal conductance is given by the sum \( G = \sum_{\lambda = \pm 1} G_\lambda \), where

\[
G_\lambda = \frac{e^2}{\hbar} \frac{D}{2\pi} \int -\frac{\Delta}{2} d\theta k_F^\lambda T_\lambda(\theta_\lambda) \cos \theta_\lambda, \tag{2}
\]

is the conductance of the \( \lambda \) channel and its transmission

\[
T_\lambda(\theta_\lambda) = \text{Re}[[t_{\lambda\lambda}]^2 + |t_{\lambda\bar{\lambda}}|^2 (\cos \theta_{\lambda}/\cos \theta_\lambda)], \tag{3}
\]

is the sum of the intra- and inter-channel transmission. \( \theta_\lambda \) is the incident and \( \theta_{\bar{\lambda}} \) the propagation angle of the cross-channel wave with the conservation of the \( k_y \) component. In Eq. (2), \( D \) denotes the width of the sample and \( k_F^\lambda \) refers to the Fermi wavevector of the \( \lambda \) channel.

As show in Fig. 1, the \( \mathcal{P}_s \mathcal{T} \) symmetry leads to perfect transmission. Therefore, transmission resonances occur whenever the scattering states are such that \( \mathcal{P}_s \mathcal{T} \psi(x, y) = \xi \psi(x, y) \), with eigenvalues of the form \( \xi = e^{i\theta} \). However, \( \mathcal{P}_s \mathcal{T} \) does not commute with the Hamiltonian in Eq. (1). Therefore, it is not an intrinsic symmetry of the system. Instead, the \( \mathcal{P}_s \mathcal{T} \) symmetry emerges only for certain specific system parameters and scattering states satisfying,

\[
[H, \mathcal{P}_s \mathcal{T}] \psi_R(x, y) = 0, \tag{4}
\]

where the index \( R \) emphasizes that the relation holds only at resonances.

For illustration we consider an InGaAs/InAlAs 2DEG with \( m^* = 0.05m_0 \), \( m_0 \) is the free-electron mass, and \( \alpha = 0.093 \text{ eVÅ} \). The conductance, normalized to the Sharvin conductance \( G_0 = (e^2/\hbar)(2D/\pi)|\langle 0|\sigma m^*/\hbar^2 \rangle|^2 + 2m^*/\hbar^2|\Delta|^2 \), as a function of the proximity-induced \( \Delta \) is shown in Fig. 2(a) for \( m \parallel \) parallel (dashed line) and perpendicularly (solid line) to the current direction. The conductance exhibits a nonmonotonic behavior with maxima [labeled by (1) and (2)] occurring at different \( \Delta \). The shift between the conductance maxima is largely enhanced when the SOC strength is increased, as illustrated in Fig. 2(b) where \( \alpha = 0.93 \text{ eVÅ} \) has been used [10 times larger than in Fig. 2(a)]. The nonmonotonic behavior of the conductance is a consequence of the collective contributions of multiple resonant states corresponding to different propagation directions of the tunneling carriers. When \( m \) is parallel to the current, the transmission \( (T_\lambda) \) of particles at the Fermi energy exhibits multiple

![FIG. 2. (a) Dependence of conductance on \( \Delta \) with \( E_F = 10 \) meV, \( V_0 = 15 \) meV, \( d = 13 \) nm, and \( \alpha = 0.003 \text{ eVÅ} \). Dashed and solid curves denote the conductance for \( m \parallel x \) and \( m \parallel y \), respectively. Insets (1), (2) show the matching between Fermi contours in the lead (blue and red) and in the barrier (multicolor) at labeled peaks. (b) The same as (a) but with a 10 times greater. (c) and (d) Transmission, \( T_\lambda \), of incident states from dominant \( \lambda = \) –1 band as a function of \( \Delta \) and angle \( \theta \) when \( m \parallel x \) and \( m \parallel y \), with the parameter values from (b).](image-url)
resonances, depending on $\Delta$ and the incident angle, $\theta$, as shown in Figs. 3(c) and (d) for the $\lambda = -1$ incident channels, which are dominant in the total conductance. Their $\lambda = 1$ counterparts are also shown in Ref. [31]. The appearance of a large number of resonant states at the value of $\Delta$ indicated by vertical lines results in the maximum of conductance displayed in Fig. 2(b) and marked with label (3). The parameter space where resonances emerge is shifted, leading to a shift in the $\Delta$ value [see the vertical line in Fig. 2(d)] at which the conductance maximum labelled by (4) in Fig. 2(b) appears.

The origin of the conductance maximum at (3) and (4) in Fig. 2(b) can be understood by examining $T_x(\theta)$ from Figs. 2(c) and (d), where we only showed the dominant $\lambda = -1$ channel contribution which corresponds to the transport via the outer Fermi contour states in the leads. The region of bright colors in Figs. 2(c) and (d) shows strong transmission which approximately satisfies the symmetry condition Eq. (4). Observations of the behaviors in Figs. 2(b)-(d), lead us to conclude that the maximum of the total conductance is determined mainly by the range of $\theta$ for dominant transmission. With $\mathbf{m} \parallel \mathbf{x}$, the Fermi contour inside the barrier shifts vertically and the condition for maximum transmission angle is achieved exactly at $\Delta = V_0$, when the shift of the energy band by the potential $V_0$ inside the barrier is cancelled by the Zeeman shift $-\Delta$ for the $\lambda = -1$ channel. On the other hand, with $\mathbf{m} \parallel \mathbf{y}$, as shown in inset (4) of Figs. 2(b), the best contour matching at large angles is achieved for an enlarged Fermi contour inside the barrier at $\Delta \approx V_0 + \alpha k_F$, as derived in Ref. [31]. This difference for the maximum of the total conductance as a function of $\mathbf{m}$ is the origin of the strong TAMR.

To characterize the strength of the anisotropic response, we introduce the in-plane TAMR coefficient [7,9],

$$\text{TAMR}(\phi) = \frac{R(\phi) - R(0)}{R(0)} = \frac{G(0) - G(\phi)}{G(0)},$$

with $\phi$ defined in Fig. 1(d). Unlike the MR values, which may change considerably from sample to sample, the TAMR coefficient is known to be more robust against specific sample details [32]. Up to the second order in the SOC strength, the extreme values of the TAMR($\phi$) are given by the contrast between the MR measured for $\mathbf{m}$ parallel and perpendicular to the current, i.e., TAMR($\phi = \pi/2$). For brevity, we use TAMR $\equiv$ TAMR($\phi = \pi/2$), unless $\phi$ is explicitly specified.

The dependence of the TAMR on $\Delta$ and $V_0$ is shown in Fig. 3(a). The TAMR exhibits a sharply peaked behavior for $\alpha = 0.093$ eVÅ with extreme values along a line in the vicinity of $V_0 = \Delta$. For a given value of $V_0$, the width of the TAMR peak is estimated as the difference between the values of $\Delta$ corresponding to the maximum of $G(0)$ and $G(\pi/2)$, which, according to our discussion of Fig. 2, in the weak SOC limit is given by $\alpha k_F$. Complementary to our model calculations for an infinite system, we have also performed calculations using Kwant [33] for the case of a finite system with a scattering region discretized into a $100 \times 40$ lattice with a spacing of 2.6 nm [31]. The results in the inset of Fig. 3(a), for the same range of $V_0$ and $\Delta$, reveal that while finite-size effects slightly increase and sharpen the TAMR, the overall qualitative behavior remains similar to the calculations for an infinite system. This confirms the TAMR robustness mentioned above.

Since the helical spin textures in 2DEG systems, depicted in Fig. 1(c), are also inherent to the surfaces states of 3D topological insulators (TIs), it is important to examine if there are any differences between their respective TAMR signatures. In Fig. 3(b) we consider additionally the geometry from Fig. 1(a), but with the 2DEG replaced by a 3D TI. We choose $(\text{Bi}_{1-x}\text{Sb}_x)\text{Te}_3$ with $x = 0.36$, effective mass $m^* = 0.27 m_0$, Rashba SOC $\alpha = 0.36$ eVÅ, Fermi velocity of the surface state $4 \times 10^5$ m/s, the energy difference between the Dirac point and the crossing point of the Rashba bands $\Delta E = 250$ meV, and $E_F = 260$ meV measured from the Dirac point [34,38].

The dependence of the TAMR of the 2DEG on $\mathbf{m}(\phi)$, for different values of $\Delta$ is represented by dashed lines shown in Fig. 3(b). For $\Delta \ll V_0$, the functional form of the conductance is $G(\phi) \approx A + B \cos^2(\phi)$ [17,31], where $A$ and $B$ are functions of system parameters other than $\phi$. It then follows from Eq. (5) that the angular dependence of TAMR is of the form $\text{TAMR} \approx B \sin^2(\phi)$, which is precisely the dependence observed in Fig. 3(b). Interestingly, for the 2DEG the sign of $B$ and even of TAMR can be inverted by changing $\Delta$. While TAMR is positive for $\Delta = 12$ meV and 26 meV, it becomes negative for $\Delta = 17$ meV [see Fig. 3(b)]. In contrast, there is no TAMR sign reversal for the TI surface states, shown by solid lines in Fig. 3(b), which are computed following the same procedure as in Ref. [39]. For both 2DEG and TI the maximum TAMR values in Figs. 3(a) and (b) are quite large compared to...
FIG. 4. TAMR (a) due to topological surface states in the TI/F system as a function of $V_0$ and $\Delta$ and (b) due to both topological and trivial states, where $E_F = 10$ meV, $d = 13$ nm, $m^* = 0.27m_0$ and $\alpha = 0.36$ eVÅ. (c) Dependence of TAMR amplitude on $\Delta$ with $V_0 = 15$ meV for different states.

typical in-plane TAMR $\lesssim 1\%$ in other systems [6, 7, 12, 13]. This predicted magnitude is particularly striking for a 2DEG with commonly found SOC strength and spin-unpolarized leads.

In 3D TIs, depending on the Fermi energy, their transport properties may be dominated by topological surface states with a Dirac-like dispersion, trivial Rashba-like states, or both [40–42]. Despite the different dispersions of the massive carriers in the 2DEG/F and the massless topological states in the TI/F systems, the TAMR in both devices exhibit the same $\sin^2(\phi)$ dependence on $m$. The angular dependence of the TAMR due to topological states in the TI/F structure can be deduced by approximating the barrier region by a Dirac delta function to obtain an approximate analytical TAMR expression [31].

While the TAMR($\phi$) cannot discriminate between the trivial and topological states, we seek if TAMR can still provide their distinguishing signature. Indeed, by comparing TAMR due to (i) only topological states in Fig. 4(a) and (ii) both topological and trivial states in Fig. 4(b), we can see nonmonotonic TAMR trends in either $\Delta$ or $V_0$ arise only from the trivial states. This is better illustrated in Fig. 4(c) with TAMR as a function of $\Delta$. The main $|\text{TAMR}|$ peak at $\Delta = 25$ meV originates from the earlier Fermi contour matching argument. However, the resonant transmission at $\Delta = 9$ meV, arises from a different origin of the standing-wave formation in the barrier due to the constructive interference between the two 2DEG/F interfaces. Such peaks, also resulting from the $P_sT$ symmetry, have resonant conditions analogous to those for simple potential barrier systems [31].

In contrast, with just TI surface states, the Fermi contour shifts the exchange field without changing its diameter due to its linear dispersion relation, and the TAMR lacks such resonance [31]. Therefore, TAMR measurements and their monotonicity in TI/F systems could help to address the controversy [33–39] whether the transport is purely determined by the topological states or if there is also a contribution of trivial states.

To realize magnetic proximity effects for the in-plane transport, magnetic insulators are desirable [34, 50, 51]. This precludes current flow in the more resistive F region [Fig. 1(a)] and minimizes hybridization with the 2DEG or TI to enable a gate-tunable proximity-induced exchange splitting in their respective states. However, as shown in graphene [52–54] for tunable magnetic proximity effects one could also employ ferromagnetic metals, separated by an insulating region from the 2DEG or TI.

While we have focused on a longitudinal transport in a very simple system having no spin-polarized leads, the predicted resonant tunneling behavior emerging from a spin parity-time symmetry of the scattering states is important not just in explaining a surprisingly large TAMR, but also as a sensitive probe to distinguish between trivial and topological states. Our work also motivates several direct generalizations which can be explored in a similar geometry. We expect a rich behavior of the transverse response [39, 55] and unexplored resonant Hall effects as well as detecting different states in magnetic topological insulators [56]. The focus on Rashba spin-orbit coupling can be extended in a growing class of van der Waals materials. For example, transition metal dichalcogenides in addition to their inherent spin-orbit coupling also provide spin-orbit proximity [32, 57–64] and thereby alter spin textures and expected TAMR, while 2D van der Waals ferromagnets support a versatile gate control [65–67].

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