Magnetization dynamics induced by in-plane currents in ultrathin magnetic nanostructures with Rashba spin-orbit coupling

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Recent experiments on ultrathin magnetic layers with broken inversion symmetry reported anomalous current-driven magnetization dynamics. We show that the spin-transfer torque can be significantly modified by Rashba spin-orbit coupling and the modified spin-transfer torque can explain the anomalous magnetization dynamics. This work will be valuable for the development of next generation spintronic devices based on ultrathin magnetic systems.

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Electric control of magnetic states carries a high potential toward device applications, such as magnetic memory and logic. Spin-transfer torque (STT), arising from a spin-polarized current is an efficient way to achieve electric control in magnetic nanostructures. There have been extensive efforts to clarify the properties of STT and enhance the efficiencies of STT both theoretically and experimentally.

Recently the current-driven magnetization dynamics of ultrathin (≤1 nm) magnetic layers has received considerable attention. In particular, experiments on an ultrathin (0.6 nm) magnetic layer (Co) sandwiched between a heavy metal layer (Pt) and an oxide layer (AlOx) revealed a number of anomalous features: (i) When an in-plane current is supplied to the system, a domain wall (DW) in the ultrathin magnetic layer moves against the electron flow direction, which is in clear contrast to STT theories predicting DW motion along the electron flow direction. (ii) DW moves at speeds as high as 400 m/s. This is about four times higher than the highest speed of current-driven DW motion reported previously. The origin of the speed enhancement is not clear. (iii) When an external magnetic field is applied parallel to the in-plane current, the in-plane current induces the magnetization switching of the magnetic layer in a single domain state. Recalling that the trilayer system does not contain a second magnetic layer, this result is again at odds with STT theories, which require a second magnetic layer for current-driven magnetization switching. These anomalies imply that this ultrathin magnetic system is not a mere thin limit of thicker counterparts but is a qualitatively different system governed by different physics. A clear understanding of its core physics will be highly valuable for the development of powerful spintronic devices; higher DW speed implies faster device operation, and switching by an in-plane current opens a possibility to lower the switching energy since the cross-sectional area for an in-plane current can be orders of magnitude smaller than that for a perpendicular current, providing room for the reduction of the switching current threshold.

An experiment reported that conduction electrons in a trilayer system are subject to Rashba spin-orbit coupling (RSOC). The emergence of RSOC (Ref. 19) is reasonable in this ultrathin magnetic layer since its upper and lower layers are made of quite different materials (Fig. 1), breaking the structural symmetry, and the magnetic layer being ultrathin makes the layer more susceptible to symmetry breaking effects. Motivated by this report, we explore the possible relations between RSOC and the anomalies. Previous theoretical studies reported that RSOC generates a contribution to STT that is proportional to −αR × (z × jz), which should be added to conventional STT (Ref. 9) in the absence of RSOC. Here αR is the parameter describing the strength of RSOC, jz is the in-plane current density in the ultrathin magnetic layer, z is the unit vector perpendicular to the layer, and m is the unit vector along the magnetization in the ultrathin magnetic layer. Since this contribution has the same structure as the field torque (−m × H), we call it fieldlike STT (FL-STT) arising from an effective field αRz × jz. FL-STT, however, cannot explain anomalies (i) and (ii), as we demonstrated recently. Here we demonstrate theoretically that when RSOC is combined with electron scattering, it generates still another contribution to STT that is proportional to βα jx (Ref. 10) large, explaining anomaly (iii). To make its magnitude βα large, jx (along the x direction) should be large and m should be forced to acquire a sufficient x component by applying an external magnetic field. By the way, the second term of the effective field is orthogonal to the electron flow direction, which makes the layer more susceptible to symmetry breaking effects.
the PMA direction and is not crucial for the switching. Below we first derive SL-STT and then show that SL-STT explains anomalies (i) and (ii) as well when combined with FL-STT.

The derivation of SL-STT goes as follows. We add the RSOC Hamiltonian \( H_{\text{RSOC}} = (\alpha_0 / \hbar)(\sigma_{\text{op}} \times \mathbf{p}_{\text{op}}) \cdot \mathbf{z} \) to the conventional \( s-d \) Hamiltonian\(^6\) to obtain the total Hamiltonian \( H \) describing conduction electrons in an ultrathin magnetic layer subject to RSOC,

\[
H = \frac{p_{\text{op}}^2}{2m_e} + J_{\text{ex}}\sigma_{\text{op}} \cdot \mathbf{m} + H_{\text{rel}} + H_{\text{RSOC}},
\]

where \( p_{\text{op}} \) is the momentum operator, \( m_e \) is the electron mass, \( J_{\text{ex}} \) is the exchange energy, \( \sigma_{\text{op}} \) is the Pauli spin operator, and \( H_{\text{rel}} \) describes electron scattering responsibility for spin relaxation. Note that \( \mathbf{m} \) appears only as a classical vector field in \( H \). Since STT arises when the local spin direction of conduction electrons deviates from \( \pm \mathbf{m} \), evaluation of the deviation in the presence of RSOC is the central part of this derivation.

Since conduction electron dynamics is much faster than magnetization dynamics, the time dependence of \( \mathbf{m} \) may be ignored in the leading approximation. Corrections to this approximation will be presented below. Then the many-body nonequilibrium state describing current-carrying conduction electrons can be constructed by filling up the eigenstates of the time-independent single-particle Hamiltonian \( H \), following the Landauer-Büttiker description\(^6\) of electron transport. To evaluate the local spin density of this many-body state, it is useful to note that the strong exchange energy \( J_{\text{ex}} \), which is the largest energy scale affecting the conduction electron spin dynamics in conventional metallic ferromagnets (including Co), allows one to classify eigenstates of \( H \) into two groups, majoritylike and minoritylike states. Each group generates majoritylike and minoritylike states. Each group generates

\[
s_{\pm}(r) \equiv (\sigma_{\text{op}} \delta(r_{\text{op}} - r))_{\pm},
\]

where \( r_{\text{op}} \) is the position operator and \( \cdot \) denotes the sum over expectation values over all occupied majoritylike/ minoritylike states. The sum \( s(r) \equiv s_+(r) + s_-(r) \) determines the total local spin density. To evaluate \( s_{\pm}(r) \), it is convenient to derive an equation that it satisfies. From the spin continuity equation determined by \( H \), we obtain the Bloch equation

\[
\nabla \cdot \mathbf{J}_{\pm} = -\frac{s_{\pm}}{\tau_{\text{ex}}} \times \left[ \mathbf{m} + \frac{2\alpha_0 m_e \tau_{\text{ex}}}{\hbar^2} (v_{\pm} \times \mathbf{z}) \right] + \langle \mathbf{\Gamma} \rangle_{\pm},
\]

where \( \tau_{\text{ex}} = \hbar / 2J_{\text{ex}} \), \( \mathbf{\Gamma} = [\sigma_{\text{op}} , H_{\text{rel}}] / i\hbar \), and the spin-current density \( \mathbf{J}_{\pm} = v_{\pm} \times s_{\pm} \). Here \( v_{\pm} \) is the average expectation value of the kinematic velocity operator \( v_{\text{op}} \equiv p_{\text{op}} / m_e + (\alpha_0 / \hbar) \mathbf{z} \times \sigma_{\text{op}} \) over occupied majoritylike/ minoritylike states. In deriving Eq. (2), \( \langle \sigma_{\text{op}} h(v_{\text{op}}), \delta(r_{\text{op}} - r) \rangle + \langle [v_{\text{op}}, \sigma_{\text{op}}], \delta(r_{\text{op}} - r) \rangle \rangle \) is approximated by \( 4[s_{\pm}(r)]_{\pm} \). Equation (2) uniquely fixes \( s_{\pm}(r) \) as a function of \( \mathbf{m} \) and \( v_{\pm} \), where \( v_{\pm} \) carries the current dependence of \( s_{\pm} \). It is convenient to separate \( s_{\pm} \) into longitudinal and transverse components, \( s_{\pm} = (s_{\pm} \cdot \mathbf{m}) \mathbf{m} + s_{\perp} \), since for large \( J_{\text{ex}} \), the longitudinal component is essentially independent of current and \( n_{\perp} \equiv \mp s_{\perp} \cdot \mathbf{m} \) may be identified as the majority/minority number density of conduction electrons. One then makes the relaxation time approximation\(^9\) \( \langle \mathbf{\Gamma} \rangle_{\pm} = -\delta s_{\pm} / \tau_{\text{ex}}, \) where \( \tau_{\text{ex}} \) is the transverse spin relaxation time. In this approximation, the relaxation of the longitudinal spin component is neglected since the transverse relaxation is much faster in conventional metallic ferromagnets and also the longitudinal spin component does not affect STT. From Eq. (2), one then obtains

\[
\delta s_{\pm} = \pm n_{\pm} \tau_{\text{ex}} \left( \frac{\beta + m \times v_{\pm} \times \mathbf{z}}{1 + \beta^2} \right) \left[ D_{\pm}^2 + \frac{2\alpha_0 m_e}{\hbar^2} \mathbf{m} \times (v_{\pm} \times \mathbf{z}) \right],
\]

where \( \beta = \tau_{\text{ex}} / \tau_{\text{ta}} \) and \( D_{\pm} = v_{\pm} \times \mathbf{V}. \) We remark that when the time dependence of \( \mathbf{m} \) is taken into account, the left-hand side of Eq. (2) acquires an additional term \( \delta s_{\pm} / \partial t \) and \( D_{\pm} \) in the above equation is replaced by \( \delta t + v_{\mp} \cdot \mathbf{V} \). Then Eq. (2) becomes the RSOC generalization of the Bloch equation in Ref. 9 without RSOC.

Finally from the relation \( T = \mu_B \tau_{\text{ex}}^{-1} \mathbf{m} \times \mathbf{s} \) (Ref. 9) between the total STT \( T \) and \( \mathbf{s} \), one obtains the Landau-Lifshitz-Gilbert (LLG) equation \( \partial \mathbf{M} / \partial t = -\gamma \mathbf{M} \times \mathbf{H}_{\text{eff}} + (\alpha_0 / M_s) \mathbf{M} \times \partial \mathbf{M} / \partial t + \mathbf{T} \), where \( \mathbf{H}_{\text{eff}} \) is a sum of an external magnetic field and effective magnetic fields due to magnetic anisotropy and magnetic exchange energy, \( \mathbf{M} = M_s \mathbf{m} \) is the magnetization and \( M_s \) is the saturation magnetization in the ultrathin magnetic layer. After grouping together terms of the same structure, one obtains

\[
\frac{\partial \mathbf{M}}{\partial t} = -\gamma \mathbf{M} \times \left( \mathbf{H}_{\text{eff}} + \mathbf{H}_R - \frac{\beta}{M_s} \mathbf{M} \times \mathbf{H}_{\text{R}} \right) + \frac{\alpha_0}{M_s} \mathbf{M} \times \frac{\partial \mathbf{M}}{\partial t} + \frac{\mu_B P}{e M_s (1 + \beta^2)} \mathbf{J}_e \times \mathbf{V} \mathbf{M},
\]

\[
\frac{\partial \mathbf{H}_R}{\partial t} = \frac{\alpha R m_e P}{\hbar e M_s (1 + \beta^2)} (\mathbf{z} \times \mathbf{J}_e),
\]

where \( \mathbf{H}_R \) is the additional effective field due to RSOC, \( \mathbf{J}_e = -e(n_+ v_+ + n_- v_-) \), and \( P \mathbf{J}_e = -e(n_+ v_+ + n_- v_-) \). Note that the first and fourth terms on the right-hand side of Eq. (4) contain the renormalized gyromagnetic ratio \( \gamma \) and the renormalized Gilbert damping \( \alpha \) given by \( \gamma / \gamma_0 = 1 + (n_+ - n_-) / [M_s (1 + \beta^2)] \), and \( \alpha / \gamma_0 = \alpha_0 + \beta (n_+ - n_-) / [M_s (1 + \beta^2)] \). The last and second to the last terms are the nonadiabatic STT (Refs. 9 and 10) and the adiabatic STT.\(^9\) When \( \delta \tau_{\text{ex}} \) is set to zero, Eq. (4) reduces to the LLG equation obtained for thicker magnetic systems.\(^9\)

RSOC effects are contained in the second and third terms. The second term is FL-STT, which has been derived before\(^20,21\) and the third term is SL-STT, whose derivation is one of main results of this Rapid Communication. Note that SL-STT does not contain any space derivative, which is in contrast to the nonadiabatic STT (Ref. 9) being proportional to space derivative of \( \mathbf{m} \). Considering that both STTs
STT determines the nonadiabatic STT, effectively cancelling (enlarging) the effective field and possible for SL-STT to overcancel the nonadiabatic STT in $\pm m_{\alpha}>\beta$ and transverse DW in a PMA system when a current flows to the right should not be uniform for $H_{R}$ can be understood from the direction of the effective field stability equivalence. The effect of FL-STT on the stability (Ref. 22) since it is perpendicular to the nonadiabatic STT arises from electron-scattering-induced spin relaxation and are dynamics (the same DW velocity in ultrathin magnetic systems16 with PMA. Without RSOC, electrons becomes $v_{\perp}p$, this difference is worth clarifying. In the caption of Fig. 3. The inset in Fig. 3(a) shows $v_{DW}$ as a function of $\theta_{SH}$ at $H_{R}=3.0 \times 10^{11} A/m^{2}$ for the stable DW structure in Fig. 2(b). The parameters for the simulation are as follows: $\alpha=0.5$, $\beta=0.25$ (Ref. 26), $M_s=5.0 \times 10^{5} A/m$, the PMA constant $K_{u}=1.0 \times 10^{5} J/m^{3}$, the exchange stiffness constant $A_{ex}=1.0 \times 10^{-11} J/m$, $P=0.7$, $\gamma/2\pi=28.0113$ GHz T$^{-1}$, and $\eta=0.6$ nm. a weaker effect on the stability than FL-STT since, according to Eq. (4), SL-STT is smaller than FL-STT in magnitude by a factor $\beta$, which is usually smaller than 1.9,26 Then combining the above information, we find that there is only one stable DW structure [Fig. 2(a)] when RSOC is sufficiently strong and that it moves fast against the electron flow direction, explaining both anomalies (i) and (ii).

Figures 3(a) and (b) show the micromagnetic simulation results of Eq. (4) for the stable DW structure in Fig. 2(a). Various parameter values used in the simulation are given in the caption of Fig. 3. The inset in Fig. 3(a) shows $v_{DW}$ as a function of a dimensionless parameter $\alpha_{SH} = \alpha / (\gamma M_s/2 \pi)$ at fixed $j_{z}=+3 \times 10^{13} A/m^{2}$, where $\lambda$ is the DW width at $j_{z}=0$. Note that as $\alpha_{SH}$ increases, $v_{DW}$ changes its sign from negative (along the electron flow direction) to positive (against the electron flow direction). The main panel in Fig. 3(a) shows
$v_{DW}$ as a function of $j_x$ at two fixed values of $\alpha_R$, 0 (black squares) and 10 (red circles). For $\alpha_R = 10$, $v_{DW}$ changes from negative to positive at $j_x \approx 3.5 \times 10^{11} \text{ A/m}^2$ and goes above +500 m/s for $j_x > 1.5 \times 10^{12} \text{ A/m}^2$. Thus both anomalies (i) and (ii) can be explained by RSOC if $\alpha_R$ is sufficiently larger than 1. For the value of $\alpha_R = 10^{-10} \text{ eV m}$ (Ref. 18) reported for Pt/Co(0.6 nm)/AlO$_x$, $\alpha_R$ becomes 13 if we assume $\lambda = 3$ nm. The results for the DW structure in Fig. 2(b) are not shown since, when $\alpha_R = 10$, it is unstable for $j_x > 7.4 \times 10^{10} \text{ A/m}^2$.

Lastly, to utilize the RSOC effects for device applications, it is desired to understand when $\alpha_R$ becomes large. In the case of a normal metal in contact with a heavy metal, there are well-known material combinations\textsuperscript{28,29} generating large $\alpha_R$ in the range (0.4–3) × $10^{-10} \text{ eV m}$ for conduction electrons near the interface. In comparison, to the best of our knowledge, experimental data on a ferromagnet in contact with other materials is still quite limited. A recent experiment\textsuperscript{30} on Ta/CoFeB(1.0 nm)/MgO measured $H_R$ in FL-STT and found it to be about a factor of 4 smaller than the corresponding report on Pt/Co(0.6 nm)/AlO$_x$\textsuperscript{18}. A ferromagnetic metal in contact with topological insulators is an interesting combination to explore since topological insulators have strongly spin-orbit coupled states\textsuperscript{31} near their surface. More experiments on various combinations are desired. For certain combinations of nonmagnetic metals [such as Bi/Ag (Ref. 29)], it is known that $\alpha_R$ is drastically enhanced when atoms of neighboring layers get intermixed. A similar enhancement may occur in a ferromagnet in contact with heavy metals and may be responsible for the discrepancy between three experimental groups\textsuperscript{8,27,28} on the values of $H_R$ in Pt/Co(0.6 nm)/AlO$_x$. In particular, the experiment\textsuperscript{27} found $H_R$ to be negligibly small but still observed anomaly (iii). It was argued that the spin Hall effect (SHE) in the Pt layer converts an in-plane charge current to a perpendicular flow of a spin current, which generates the Slonczewski STT. Recalling the structural similarity between the Slonczewski STT and SL-STT, the magnetization switching can be explained by the SHE-induced Slonczewski STT. However, in the absence of FL-STT, the Slonczewski STT alone cannot generate anomalies (i) and (ii), as demonstrated in Figs. 3(c) and 3(d). Note that DW moves along the electron flow direction when $j_x$ is large enough (>5 × $10^{11} \text{ A/m}^2$) and also that large enhancement of the DW speed does not occur.

To conclude, we demonstrated that RSOC generates two contributions (FL-STT and SL-STT) to STT and also that if RSOC is sufficiently strong, they can explain three anomalous features of the magnetization dynamics reported in an ultrathin magnetic system with structural symmetry breaking. This result will be useful for the development of next generation spintronics devices. During the preparation of this manuscript, we were made aware of work\textsuperscript{34} which also derives SL-STT with similar results but does not discuss the implications of SL-STT on the anomalous DW motion.

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MAGNETIZATION DYNAMICS INDUCED BY IN-PLANE . . .

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