Theory of nonequilibrium intrinsic spin torque in a single nanomagnet

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In a single nanomagnet with spin-orbit interactions, the electrical current can generate a nonequilibrium spin density that gives rise to a spin torque on the magnetization. This spin torque does not involve spin transfer mechanism and originates from the band structure itself. We show that this spin torque can be effectively used to switch the direction of the magnetization and the critical switching current density could be as low as 10^4-10^6 A/cm^2 for a number of magnetic systems. Several magnetic systems for possible experimental realization are discussed.

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The manipulation of magnetization in submicronic devices is an important topic related to data storage and magnetic memory technology.[1] Besides field-induced magnetization switching, which is known to be a size and power consuming technique,[2] a current-driven switching based on the spin transfer effect has been realized.[3] This mechanism requires a noncollinear magnetization structure such as spin valves, tunnel junctions, and domain walls. The current density observed in transition-metal-based materials is on the order of 10^7 A/cm^2, which is higher than the current breakdown of semiconductor-based transistors.[2]

In this Brief Report, we propose another current-driven mechanism for the manipulation of the magnetization of a single magnetic two-dimensional electron gas (2DEG) when a Rashba spin-orbit coupling[4] is present. Contrary to the conventional spin-valve structures, this spin torque arises from the band structure of the nanomagnet without the need for noncollinear ferromagnetic layers and does not involve the usual spin transfer mechanism. The transport properties with Rashba interactions in two-dimensional electron gas have been widely studied mainly for paramagnetic systems where there is no spontaneous ferromagnetic magnetization. Inoue *et al.*[5] showed that the presence of an in-plane charge current induces an out-of-equilibrium transverse spin accumulation. This spin accumulation can then exert a torque on neighboring ferromagnetic contacts.[6,7] In contrast, we consider here a single ferromagnetic 2DEG sandwiched between two dissimilar materials, so that the electric potential is highly asymmetric, leading to the presence of the Rashba interaction at the interfaces or within the ferromagnetic layer. The Hamiltonian is thus

$$H = \frac{p^2}{2m} + \frac{\alpha_R}{\hbar}(\mathbf{p}\times\hat{\mathbf{z}})\cdot\boldsymbol{\sigma} - J_{\rm sd}\boldsymbol{\sigma}\cdot\hat{\mathbf{M}}, \tag{1}$$

where we confine the electron motion in two dimensions (x,y), $\hat{\mathbf{M}}=\cos\theta\hat{\mathbf{x}}+\sin\theta\hat{\mathbf{y}}$ is the unit vector for the direction of the local magnetization, $J_{\rm sd}$ is the exchange coupling strength between the conduction electron and the magnetization, and α_R is the Rashba parameter. To immediately see the current-driven effect due to the Rashba term, let us first consider a large $J_{\rm sd}$ such that the conduction-electron spin and $\hat{\mathbf{M}}$ are parallel and the Rashba term is equivalent to the Zeeman energy with the magnetic field $H_{\rm eff}=-(\alpha_R/\hbar)(\langle\mathbf{p}\rangle\times\hat{\mathbf{z}})$. In equilibrium, $H_{\rm eff}=0$ because $\langle\mathbf{p}\rangle$ is zero. In the presence of the current, $\langle\mathbf{p}\rangle$ is proportional to the current $\mathbf{j}_e$ and thus the average magnetic field on the magnetization is proportional to $\alpha_R(\mathbf{j}_e\times\hat{\mathbf{z}})$; it is this current-driven effect that can be used to switch the direction of the magnetization. We now carry out the rigorous calculation below.

The one-electron eigenenergy and wave function are

$$E_{\mathbf{k}}^{\pm} = \frac{\hbar^2k^2}{2m} \pm |\alpha_R(\mathbf{k}\times\hat{\mathbf{z}}) - J_{\rm sd}\hat{\mathbf{M}}| \tag{2}$$

and

$$\Psi_{\mathbf{k}}^{\pm} = \frac{1}{\sqrt{2A}}\begin{pmatrix}\pm e^{i\gamma_k}\\ 1\end{pmatrix}\exp(i\mathbf{k}\cdot\mathbf{r}), \tag{3}$$

where $\tan\gamma_k=(\alpha_R k_x+J_{\rm sd}\sin\theta)/(\alpha_R k_y-J_{\rm sd}\cos\theta)$ and A is the area of the film.

To calculate the nonequilibrium spin density and the spin torque in the presence of the current, we use the Boltzmann equation for the two bands,

$$eE_x\left(-\frac{\partial f_0^\sigma}{\partial k_x}\right) = \sum_{\sigma'}\int d^2k'\,W_{kk'}^{\sigma\sigma'}(f_k^\sigma - f_{k'}^{\sigma'}), \tag{4}$$

where E_x is the electric field in the x direction, f_0^σ is the equilibrium distribution of the two bands ($\sigma=\pm$), and the scattering probability is

$$\begin{aligned} W_{kk'}^{\sigma\sigma'} &= \sum_i |\langle\Psi_k^\sigma|V_{\rm sc}(\mathbf{r}-\mathbf{R}_i)|\Psi_{k'}^{\sigma'}\rangle|^2\delta(E_k^\sigma - E_{k'}^{\sigma'}) \\ &= \frac{2\pi n_i}{\hbar}V^2[1+\sigma\sigma'\cos(\gamma_k-\gamma_{k'})]\delta(E_k^\sigma - E_{k'}^{\sigma'}), \end{aligned} \tag{5}$$

where $V_{\rm sc}(\mathbf{r})=V\delta(\mathbf{r})$ is the impurity scattering potential, $\mathbf{R}_i$ is the impurity position, and n_i is their concentration. The above definition accounts for both intraband ($\sigma=\sigma'$) and interband ($\sigma=-\sigma'$) scatterings, so that Eq. (4) may be solved numerically. However in the limiting case considered below ($E_F \gg J_{\rm sd} \gg \alpha_R k_F$), $\cos(\gamma_k-\gamma_{k'})=1+o(\alpha_R^2)$, so that interband transitions can be neglected to the first order in α_R. This yields to an isotropic scattering probability $W_{kk'}^{\sigma\sigma'}=\pi n_i V^2/\hbar$, which is equivalent to a constant relaxation-time approximation. In this case, the Boltzmann distribution in Eq. (4) has a simple solution,

$$f_k^{\pm} = f_0^{\pm} - eE_x\tau\left(-\frac{\partial f_0^{\pm}}{\partial k_x}\right), \tag{6}$$

where τ is the relaxation time ($e>0$). With the above distribution, the spin density $\delta\mathbf{m}$ and the spin current $\mathbf{j}_s$ can be readily evaluated; i.e.,

$$\delta\mathbf{m} = 2\int d^2k(f_k^+ - f_k^-)(\hat{\mathbf{x}}\cos\gamma_k - \hat{\mathbf{y}}\sin\gamma_k) \tag{7}$$

and

$$\mathbf{j}_s = e\int d^2k(v_x^+ f_k^+ - v_x^- f_k^-)(\hat{\mathbf{x}}\cos\gamma_k - \hat{\mathbf{y}}\sin\gamma_k), \tag{8}$$

where $v_x^{\pm}=(1/\hbar)\partial E_k^{\pm}/\partial k_x$ is the velocity. Once $\delta\mathbf{m}$ is obtained, the spin torque is $\mathbf{T}=-J_{\mathrm{sd}}\delta\mathbf{m}\times\hat{\mathbf{M}}$.

In the limiting case of $J_{\mathrm{sd}}\gg\alpha_R k_F$, where k_F is the Fermi wave vector, one can analytically integrate out the momentum in Eqs. (7) and (8). Up to the first order in α_R/J_{sd},

$$\cos\gamma_k = -\cos\theta + \frac{\alpha_R}{J_{\mathrm{sd}}}(k_y\sin^2\theta + k_x\sin\theta\cos\theta) \tag{9}$$

and

$$\sin\gamma_k = \sin\theta + \frac{\alpha_R}{J_{\mathrm{sd}}}(k_x\cos^2\theta + k_y\sin\theta\cos\theta). \tag{10}$$

By placing the above expressions into Eqs. (7) and (8) and by noticing that the terms linear in k_y average out after integration, we have

$$\delta\mathbf{m} = 2\frac{\alpha_R m}{\hbar e}\frac{j_e}{E_F}(\cos\theta\sin\theta\hat{\mathbf{x}} - \cos^2\theta\hat{\mathbf{y}}) \tag{11}$$

and $\mathbf{j}_s=Pj_e\hat{\mathbf{M}}$, where $P\approx J_{\mathrm{sd}}/E_F$ is the spin polarization of the current. Thus, the spin torque is $\mathbf{T}=-\hat{\mathbf{z}}(2\alpha_R mPj_e/\hbar e)\cos\theta$. If one defines a current-driven magnetic field via $\mathbf{T}=-\gamma\mathbf{M}\times\mathbf{H}_{\mathrm{cd}}$, we have

$$\mathbf{H}_{\mathrm{cd}} = 2\frac{\alpha_R m}{\hbar eM_s}Pj_e(\hat{\mathbf{z}}\times\hat{\mathbf{j}}_e), \tag{12}$$

where M_s is the saturation magnetization and $\hat{\mathbf{j}}_e$ represents the unit vector in the direction of the current.

The current-driven magnetic field or torque discussed above competes with other torques in a ferromagnet. Similar to the current-driven spin torque in spin-valve structures, we can write the dynamic equation of the magnetization by including H_{cd},

$$\frac{d\mathbf{M}}{dt} = -\gamma\mathbf{M}\times(\mathbf{H}_{\mathrm{eff}}+\mathbf{H}_{\mathrm{cd}}) + \frac{\alpha}{M_s}\mathbf{M}\times\frac{d\mathbf{M}}{dt}, \tag{13}$$

where $\mathbf{H}_{\mathrm{eff}}$ is the effective field due to other magnetic interactions. If we consider a single domain ferromagnet with an anisotropy field H_K in $\hat{\mathbf{y}}$ direction, we obtain the critical switching current density by setting $H_{\mathrm{cd}}=H_K$; i.e.,

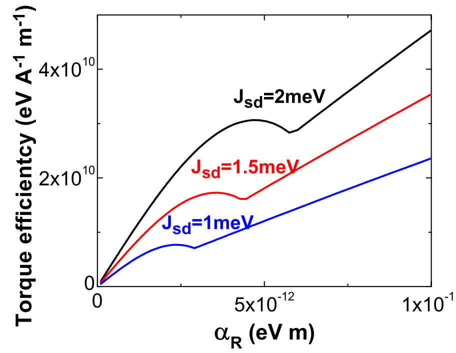

FIG. 1. (Color online) Absolute ratio between the torque and the current density as a function of the Rashba spin-orbit coupling α_R. For these simulations we took $E_F=10$ meV.

$$j_c = \frac{\hbar eH_K M_s}{2\alpha_R mP}. \tag{14}$$

If we take the Rashba constant $\alpha_R=10^{-11}$ eV m for a typical value of a two-dimensional electron gas with a spatial inversion asymmetry, the spin polarization $P=0.5$, the saturation magnetization $M_s=10^4$ (J/T m^3), the effective mass $m^*=0.45$ of heavy holes in GaMnAs,[8] and an anisotropy field of 200 (Oe), we estimate that j_c is about 5×10^6 A/cm^2, which is comparable to the critical current in spin valves and tunneling junctions. Depending on the material parameters (Rashba coupling, effective mass, etc.), we estimate that the current threshold can be as low as 10^4-10^6 A/cm^2.

While the above limiting case $J_{\mathrm{sd}}\gg\alpha_R k_F$ is most interesting and most relevant for observing the spin torque, we can also consider a weak ferromagnet where J_{sd} is smaller than $\alpha_R k_F$. In this case, $\cos(\gamma_k-\gamma_{k'})=\cos(\phi-\phi')+o(J_{\mathrm{sd}})$ [with $\mathbf{k}=k(\cos\phi,\sin\phi)$], and interband scattering is allowed at the zeroth order. Note that, since the Rashba interaction is now dominant, the zeroth order is sufficient to obtain the nonequilibrium torque. In this case, we can use the formalism developed by Schliemann and Loss[9] to solve Eq. (4). Provided that $W_{kk'}^{\sigma\sigma'}=W_{k'k}^{\sigma'\sigma}$, the authors showed that the distribution function can be expressed as a function of longitudinal $\tau_k^{\parallel}$ and transverse $\tau_k^{\perp}$ relaxation times (see Ref. 9 for details),

$$f_k^{\pm} = f_0^{\pm} - eE_x\left[\tau_k^{\parallel}\left(-\frac{\partial f_0^{\pm}}{\partial k_x}\right) + \tau_k^{\perp}\left(-\frac{\partial f_0^{\pm}}{\partial k_y}\right)\right]. \tag{15}$$

Interestingly, we find that $\tau_k^{\parallel}=\tau$ and $\tau_k^{\perp}=0$ at the zeroth order; then we can apply the relaxation-time approximation and we find

$$\delta\mathbf{m} = -\frac{\alpha_R m}{\hbar e}\frac{j_e}{E_F}\hat{\mathbf{y}} \tag{16}$$

and

$$\mathbf{T} = -\hat{\mathbf{z}}\frac{\alpha_R m}{\hbar e}\frac{J_{\mathrm{sd}}}{E_F}j_e\cos\theta. \tag{17}$$

For an arbitrary ratio of $\alpha_R k_F/J_{\mathrm{sd}}$ we show, in Fig. 1, the numerical calculation of the torque efficiency defined as the ratio between the spin torque and the current density $|\mathbf{T}|/j_e$

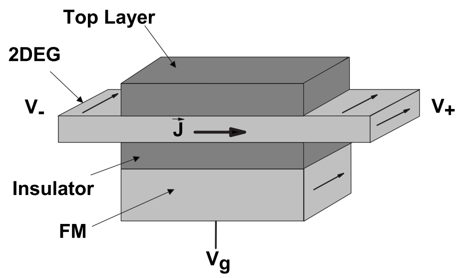

FIG. 2. Schematics of a memory bit based on Rashba-induced magnetization switching. The 2DEG-DMS is asymmetrically embedded between two insulators and the writing current is injected in the plane of the 2DEG. The readout is via the magnetic tunnel junction based on the tunnel magnetoresistance.

as a function of α_R for different values of J_{sd}, maintaining $E_F \gg J_{\mathrm{sd}}, \alpha_R k_F$. This numerical evaluation has been obtained using the Schliemann and Loss formalism.[9] Note that in the case of weak Rashba spin-orbit coupling, the torque arises from the small misalignment of the electron spin from the local magnetization, whereas in the case of strong Rashba coupling, it stems from the anisotropic electron velocity.

It is interesting to compare the present spin torque with the conventional spin transfer torque.[3] The spin transfer torque comes from the absorption of the *transverse* spin current by the magnetization and thus it requires a *noncollinear* magnetization configuration in the direction of the spin current. For the spin-valve structure, the current must be applied perpendicular to the layers and the magnetization of the two magnetic layers must not be collinear. Furthermore, the spin transfer torque involves energy transfer since it competes with the damping torque. The present spin torque is created by the intrinsic spin-orbit coupling in the nonequilibrium condition and thus it does not involve the transfer of the conduction-electron spin to the magnetization. Then, this torque exists for a uniformly magnetized layer with the current in the plane of the layer and acts as an effective field; reducing the magnitude of the anisotropy field to increase the torque efficiency will limit the thermal stability of the magnetic layer. Since the torque does not compete with the damping, it cannot excite current-driven steady magnetization precessions, in contrast with the spin transfer torques.

We now consider the possible realization of our predicted effects in several systems. The essential materials issue is to find a ferromagnet with a sizable Rashba interaction. First note that Rashba spin-orbit coupling arises from a structure inversion asymmetry that restricts the effect to 2DEG. However, a similar torque should exist in three-dimensional layers showing bulk inversion asymmetry leading to a Dresselhaus[10] spin-orbit coupling. Most of the studies on the Rashba interaction have been carried out for nonmagnetic two-dimensional semiconductors. The spatial inversion asymmetry leads to a net potential gradient in the growth direction ($\hat{\mathbf{z}}$), which generates the Rashba Hamiltonian due to spin-orbit coupling $H_R \propto (\mathbf{p} \times \nabla V) \cdot \sigma \propto (\mathbf{p} \times \hat{\mathbf{z}}) \cdot \sigma$. For example, InGaAs/$\mathrm{In}_{0.77}\mathrm{Ga}_{0.23}$As quantum well could produce a Rashba parameter as large as 10^{-11} eV m (Ref. 11) and one would expect that the Rashba parameter should be similar for the Mn-doped dilute magnetic semiconductors (DMSs) in the quantum wells. Recently, 2DEG-DMSs were grown by Bove *et al.*[12] and Teran *et al.*[13] using GaMnAs and CdMnTe quantum well, respectively.

Although Rashba spin-orbit coupling has not been investigated in magnetic 2DEG yet, a number of theoretical and experimental studies have underlined the seminal role of interfacial discontinuity of the potential gradient[14] in the amplitude of the Rashba term. Whereas the confined 2DEG conduction electrons are almost insensitive to any perpendicularly applied electric field,[15] the interfacial discontinuity of the potential gradient is usually very strong and gives rise to a sizable contribution.[14] Consequently, the potential discontinuity responsible for the Rashba spin-orbit interaction should arise in every 2DEG (conventional semiconductors, DMS, and metals) presenting a sufficiently high potential gradient and especially a discontinuity of the band structure at the interface.

Besides conventional semiconductor 2DEG, Rashba-induced spin splitting at metallic surfaces has also been intensively studied both theoretically and experimentally. Ag/Au(111) (Refs. 16 and 17) interface exhibits Rashba interaction of $\alpha_R \approx 4 \times 10^{-12} - 3 \times 10^{-11}$ eV m, while Bi/Ag(111) (Ref. 18) shows Rashba coupling even much stronger than in semiconductors of up to $\alpha_R \approx 3 \times 10^{-10}$ eV m. Magnetic Gd/GdO (Ref. 19) interface also displays a sizable Rashba interaction, depending on the spin projection ($\alpha_R \approx 2.5 \times 10^{-11}$ for minority spins and $\alpha_R \approx 1.5 \times 10^{-11}$ for majority spins). Other systems, e.g., oxide/ferromagnetic interfaces, might possess an important spin-orbit interaction, as shown by tunneling anisotropic magnetoresistance (TAMR) studies in Fe/MgO/Fe junctions[20,21] as well as by the analysis of magnetic anisotropy at AlO_x/Co interface.[22] Recent investigations on electronic reconstruction-induced quasi-2DEG at oxide interfaces[23,24] exhibit a sharp potential gradient[25] that may be asymmetrically designed in order to obtain a large Rashba effective interaction.

These studies are of great interest for our topic since metallic devices seem easier to control than DMS. Although no transition-metals 2DEG have been grown yet, Bihlmayer *et al.*[26] and Henk *et al.*[27] showed that the Rashba coupling extends considerably away from the interface [up to 5 ML for Ag/Au(111)],[26] suggesting that Rashba interaction can be maintained in 2DEG over a few nanometers. Then, even if the electron density is higher in metallic 2DEG than in semiconductor 2DEG, the higher Rashba interaction should lead to a switching current density lower than that determined by the Slowczewski torque.

Finally, we propose a way to probe and use this current-driven magnetization switching. Figure 2 shows the schematics of a memory bit based on Rashba-induced magnetization switching. This device is composed of a conventional magnetic tunnel junction whose free layer is a magnetic 2DEG inserted between two thin insulators. The writing current is injected in the plane of the free layer and can switch the magnetization orientation in either direction depending on the sign of $V_+ - V_-$. The information contained in the free layer magnetization can then be read by applying a bias voltage $V_g - V_-$ and measuring the corresponding resistance across the bottom insulator using the usual tunneling magnetoresistance effect. In this device, the Rashba-induced magnetization switching can be controlled by both the perpendicular applied electric field[14] $\vec{E}$ (through V_g) and the inplane flowing current density (through $V_+ - V_-$). As discussed earlier, the interface between the magnetic 2DEG and the tunnel barrier should have an important interfacial spin-orbit coupling. Systems such as MgO/Fe, AlO_x/Co, or even Au/Fe should provide interesting preliminary results.

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