

Mechanisms of Spin-Polarized Current-Driven Magnetization Switching

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The mechanisms of the magnetization switching of magnetic multilayers driven by a current are studied by including exchange interaction between local moments and spin accumulation of conduction electrons. It is found that this exchange interaction leads to two additional terms in the Landau-Lifshitz-Gilbert equation: an effective field and a spin torque. Both terms are proportional to the transverse spin accumulation and have comparable magnitudes.

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The concept of switching the orientation of a magnetic layer of a multilayered structure by the current perpendicular to the layers was introduced by Slonczewski [1] and Berger [2], and has been followed up by others [3]. The central idea is of a spin transfer from a polarized current to the magnetization of the layer. An alternate mechanism of current induced switching was put forth by Heide *et al.* [4] in which the current across the magnetically inhomogeneous multilayer produces spin accumulation which establishes an energy preference for a parallel or antiparallel alignment of the moments of the magnetic layers; this magnetic “coupling” was posited to produce switching. Recent experiments have reliably demonstrated that the magnetization of a magnetic layered structure is indeed switched back and forth by an applied current [5–7].

In this Letter, we present a spin transfer model in which we solve for the equation of motion of the spin accumulation in order to derive the torque acting on the background magnetization. The salient difference between our model and previous treatments is that we include the effects of spin diffusion which have been found to be important in understanding magnetoresistance of magnetic multilayers for current perpendicular to the plane of the layers (CPP-MR) [8–10]. This allows us to express our results in terms of the same parameters used to interpret CPP-MR. With respect to the model of Heide *et al.* [4] our effect arises from transverse spin accumulation, i.e., perpendicular to the local magnetization, and not from the longitudinal spin accumulation.

Let us consider a magnetic multilayer with the current perpendicular to the plane of the layer (defined as x direction). The linear response of the current to the electrical field can be written as a spinor form,

$$\hat{j}(x) = \hat{C}E(x) - \hat{D} \frac{\partial \hat{n}}{\partial x}, \quad (1)$$

where $E(x)$ is the electric field, \hat{j} , \hat{C} , \hat{D} , and \hat{n} are the 2×2 matrices representing the current, the conductivity, the diffusion constant, and the *accumulation* at a given position. The diffusion constant and the conductivity are related via the Einstein relation $\hat{C} = e^2 \hat{N}(\epsilon_F) \hat{D}$ for a degenerate metal, where $\hat{N}(\epsilon_F)$ is the density of states at the Fermi level. In general, one can express these matrices in terms of the Pauli spin matrix $\boldsymbol{\sigma}$,

$$\hat{C} = C_0 \hat{I} + \boldsymbol{\sigma} \cdot \mathbf{C}, \quad (2)$$

$$\hat{D} = D_0 \hat{I} + \boldsymbol{\sigma} \cdot \mathbf{D}, \quad (3)$$

and

$$\hat{n} = n_0 \hat{I} + \boldsymbol{\sigma} \cdot \mathbf{m}, \quad (4)$$

where $2n_0$ is the charge accumulation and \mathbf{m} is the spin accumulation. By placing Eqs. (2)–(4) into (1), we rewrite the linear response in terms of the electric current j_e and magnetization current \mathbf{j}_m as

$$j_e \equiv \text{Re}(\text{Tr} \hat{j}) = 2C_0 E(x) - 2D_0 \frac{\partial n_0}{\partial x} - 2\mathbf{D} \cdot \frac{\partial \mathbf{m}}{\partial x} \quad (5)$$

and

$$\mathbf{j}_m = \text{Re} \text{Tr}(\boldsymbol{\sigma} \hat{j}) = 2\mathbf{C}E(x) - 2\mathbf{D} \frac{\partial n_0}{\partial x} - 2D_0 \frac{\partial \mathbf{m}}{\partial x}. \quad (6)$$

It is noted that we have chosen the unit $e = \mu_B = 1$ for the notation convenience so that the electrical current and the magnetization current have the same unit. For a transition ferromagnet, one defines the spin polarization parameter β as $\mathbf{C} = \beta C_0 \mathbf{M}_d$, where \mathbf{M}_d is the *unit* vector to represent the direction of the local magnetization. Similarly, we can introduce a spin polarization β' for the diffusion constant $\mathbf{D} = \beta' D_0 \mathbf{M}_d$. These two polarization parameters are not necessarily the same, i.e., when the density of states are

different for spin up and down electrons, $\beta \neq \beta'$. By inserting these relations into Eqs. (5) and (6), and eliminating the electric field and charge density, we obtain

$$\mathbf{j}_m = \beta j_e \mathbf{M}_d - 2D_0 \left[\frac{\partial \mathbf{m}}{\partial x} - \beta \beta' \mathbf{M}_d \left(\mathbf{M}_d \cdot \frac{\partial \mathbf{m}}{\partial x} \right) \right], \quad (7)$$

where we have dropped an uninteresting term proportional to the derivative of the charge accumulation $\partial n_0 / \partial x$. The diffusion term in this equation is strictly valid only when the time for momentum relaxation τ is shorter than the characteristic time for the motion of the transverse component of the spin accumulation which is of the order of h/J . The opposite limit, i.e., the ballistic regime for the transverse component, will be addressed in a forthcoming publication; it leads to another expression for the length scale λ_J introduced in Eq. (10) below.

We now describe the equations of motion for the spin accumulation and local magnetization when we turn on the interaction between the spin accumulation and the local moment via the contact interaction,

$$H_{\text{int}} = -\mathbf{J} \mathbf{m} \cdot \mathbf{M}_d. \quad (8)$$

With this interaction, the equation of motion for the spin accumulation is

$$\frac{d\mathbf{m}}{dt} + (J/\hbar) \mathbf{m} \times \mathbf{M}_d = -\frac{\mathbf{m}}{\tau_{\text{sf}}}, \quad (9)$$

where τ_{sf} is the spin-flip relaxation time of the conduction electron. The second term on the left hand side represents the precessional motion of the accumulation due to the sd interaction when the magnetization directions of the spin accumulation and the local moments are not parallel. Since the conduction electrons carry a spin current given by Eq. (7), we replace $\frac{d\mathbf{m}}{dt}$ by $\frac{\partial \mathbf{m}}{\partial t} + \frac{\partial \mathbf{j}_m}{\partial x}$. By using Eq. (7), we find

$$\begin{aligned} \frac{1}{2D_0} \frac{\partial \mathbf{m}}{\partial t} &= \frac{\partial^2 \mathbf{m}}{\partial x^2} - \beta \beta' \mathbf{M}_d \left(\mathbf{M}_d \cdot \frac{\partial^2 \mathbf{m}}{\partial x^2} \right) \\ &\quad - \frac{\mathbf{m}}{\lambda_{\text{sf}}^2} - \frac{\mathbf{m} \times \mathbf{M}_d}{\lambda_J^2}, \end{aligned} \quad (10)$$

where we have defined $\lambda_{\text{sf}} \equiv \sqrt{2D_0\tau_{\text{sf}}}$ and $\lambda_J \equiv \sqrt{2\hbar D_0/J}$. As pointed out just after Eq. (7), this expression for the length scale λ_J holds only when $\tau < h/J$. In the opposite limit the length scale is $v_F h/J$. For transition-metal ferromagnets τ and h/J have the same order of magnitude.

The equation of motion for the local magnetization is

$$\frac{d\mathbf{M}_d}{dt} = -\gamma_0 \mathbf{M}_d \times (\mathbf{H}_e + \mathbf{J} \mathbf{m}) + \alpha \mathbf{M}_d \times \frac{d\mathbf{M}_d}{dt}, \quad (11)$$

where γ_0 is the gyromagnetic ratio, \mathbf{H}_e is the magnetic field including the contributions from the external field, anisotropy and magnetostatic field, the additional effective field $\mathbf{J} \mathbf{m}$ is due to coupling between the local moment

and the spin accumulation, and the last term is the Gilbert damping term.

To solve for the dynamics of the spin accumulation and the local moment, we need to simultaneously determine them using Eqs. (10) and (11). The time scales are very different for the spin accumulation and the local moments. The characteristic time scales of the former are of the order of τ_{sf} and h/J , see Eq. (9), i.e., of the order of picoseconds (10^{-12} sec). For the local moment, the time scale is $\gamma_0^{-1} (H_e + J m_{\perp})^{-1}$. For a magnetic field of the order of 0.1 T, this time scale is of the order of nanoseconds. Therefore, as long as one is interested in the magnetization process of the local moments, one can always treat the spin accumulation in the limit of long times. The two dynamic equations are then simply decoupled: we first solve Eq. (10) with fixed local moments (independent of time) and set the left hand side of Eq. (10) to zero. Once the spin accumulation is obtained, we substitute it into Eq. (11) to solve the dynamics of the local moments.

Before we proceed to solve for the *stationary* solution of Eq. (10), let us first discuss the general features derived from Eq. (10). We separate the spin accumulation into longitudinal (parallel to the local moment) and transverse (perpendicular to the local moment) modes. Equation (10) can now be written as

$$\frac{\partial^2 \mathbf{m}_{\parallel}}{\partial x^2} - \frac{\mathbf{m}_{\parallel}}{\lambda_{\text{sdl}}^2} = 0, \quad (12)$$

where $\lambda_{\text{sdl}} = \sqrt{1 - \beta \beta'} \lambda_{\text{sf}}$, and

$$\frac{\partial^2 \mathbf{m}_{\perp}}{\partial x^2} - \frac{\mathbf{m}_{\perp}}{\lambda_{\text{sf}}^2} - \frac{\mathbf{m}_{\perp} \times \mathbf{M}_d}{\lambda_J^2} = 0. \quad (13)$$

The longitudinal spin accumulation \mathbf{m}_{\parallel} decays at the length scale of the spin diffusion length λ_{sdl} while the transverse spin accumulation \mathbf{m}_{\perp} decays as λ_J . The spin diffusion length λ_{sdl} has been measured to be about 60 nm in Co [11]. We estimate λ_J by taking the typical diffusion constant of a metal to be 10^{-3} (m^2/s) and $J = 0.1\text{--}0.4$ (eV) so that λ_J is about 1.2 to 2.4 nm. Thus, the transverse spin accumulation has a much shorter length scale compared to the longitudinal one.

Before we apply Eqs. (12) and (13) to a multilayer structure, we take a look at the influence of the spin accumulation on the local moment. As seen from Eq. (11), the longitudinal spin accumulation has no effect on the local moment. We may rewrite Eq. (11) in terms of the transverse spin accumulation only,

$$\frac{d\mathbf{M}_d}{dt} = -\gamma_0 \mathbf{M}_d \times (\mathbf{H}_e + \mathbf{J} \mathbf{m}_{\perp}) + \alpha \mathbf{M}_d \times \frac{d\mathbf{M}_d}{dt}. \quad (14)$$

To discuss the transverse accumulation we introduce a vector \mathbf{A} such that $\mathbf{J} \mathbf{m}_{\perp} = \mathbf{A} \times \mathbf{M}_d$. If one considers a system with two ferromagnetic layers whose magnetization directions are not parallel to each other, the spin accumulation at one layer depends on the orientation of the other. Let us suppose that the above equation is used for

the layer $F1$, i.e., denote $\mathbf{M}_d = \mathbf{M}_d^{(1)}$. The magnetization of the other layer is labeled as $\mathbf{M}_d^{(2)}$. Without loss of generality, we can write the two components of the accumulation in the plane transverse to $\mathbf{M}_d^{(1)}$ as $J\mathbf{m}_\perp = a\mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)} + b(\mathbf{M}_d^{(1)} \times \mathbf{M}_d^{(2)}) \times \mathbf{M}_d^{(1)}$, where a and b are determined by geometric details of the multilayer. Placing this into Eq. (14), we find

$$\frac{d\mathbf{M}_d^{(1)}}{dt} = -\gamma_0\mathbf{M}_d^{(1)} \times (\mathbf{H}_e + b\mathbf{M}_d^{(2)}) - \gamma_0a\mathbf{M}_d^{(1)} \times (\mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)}) + \alpha\mathbf{M}_d^{(1)} \times \frac{d\mathbf{M}_d^{(1)}}{dt}. \quad (15)$$

Thus the transverse spin accumulation produces two effects simultaneously: the term $b\mathbf{M}_d^{(1)} \times b\mathbf{M}_d^{(2)}$ is the torque due to an effective field $b\mathbf{M}_d^{(2)}$ and the other is $a\mathbf{M}_d^{(1)} \times (\mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)})$ which is called the ‘‘spin torque.’’ Both terms lead to corrections to the original Landau-Lifshitz-Gilbert equation. It has been shown that both terms are capable to switch the magnetic moments [12]. Note the effective field introduced here looks as if it arises from the current induced coupling named NEXI; however, it is different as NEXI was attributed to the longitudinal component of the spin accumulation [4]. In contrast we have shown that only the transverse spin accumulation must be taken into account and that the longitudinal accumulation does not produce any effect on the motion of local moments. An even more striking difference is Heide’s finding that ‘‘the presence of a second ferromagnetic layer is not necessary.’’ This is because his longitudinal accumulation exists for a single F layer, while a second F layer with tilted magnetization is required for transverse accumulation and for our mechanism. It is notable that the ‘‘spin torque’’ term, first introduced by Slonczewski, appears on an equal footing with the torque due to the effective field $b\mathbf{M}_d^{(2)}$ as both are related to the *transverse* spin accumulation.

We now explicitly verify that the solution of the transverse accumulation \mathbf{m}_\perp indeed has our proposed general form and we quantitatively determine the magnitude of the effective field (proportional to b term) and the ‘‘spin torque’’ (proportional to a term) entering Eq. (15). To obtain a physically transparent solution of the spin accumulation, we choose an oversimplified case to perform our calculation so that the effective field and spin torque can be analytically derived. We consider a system consisting of a very thick ferromagnetic layer which is assumed to be pinned, a spacer layer which is infinitely thin so that the spin current is conserved across the layer when there is no spin flip scattering in this region, and a thin ferromagnetic layer backed by an ideal paramagnetic layer. In addition we make our calculation simpler by neglecting spin-dependent reflection at the interfaces. In such a system, we look for the solution of the spin accumulation in the thin $F1$ layer whose magnetization direction is at the positive z direction. The magnetization direction of the

pinned layer is $\mathbf{M}_d^{(2)} = \cos\theta\mathbf{e}_z - \sin\theta\mathbf{e}_y$, where θ is the angle between $\mathbf{M}_d^{(2)}$ and $\mathbf{M}_d^{(1)} = \mathbf{e}_z$. From Eqs. (12) and (13), and by assuming the same λ_{sdl} for the thin magnetic layer, $F1$, as for the nonmagnetic layer which backs it, we write the solution for the $F1$ layer as

$$m_z(x) = G_1 \exp(-x/\lambda_{\text{sdl}}), \quad (16)$$

$$m_x(x) = G_2 \exp(-x/l_+) + G_3 \exp(-x/l_-), \quad (17)$$

$$m_y(x) = -iG_2 \exp(-x/l_+) + iG_3 \exp(-x/l_-), \quad (18)$$

where $l_\pm^{-1} = \sqrt{(1/\lambda_{\text{sf}}^2) \pm (i/\lambda_j^2)}$. To determine the constants of integration, we assume the thick magnetic layer $F2$ is half metallic so that the current is fully spin polarized and we demand that the spin current is continuous across $F2/N/F1$ interface [13]; we find

$$\beta j_e - 2D_0(1 - \beta\beta') \left(-\frac{G_1}{\lambda_{\text{sdl}}} \right) = j_e \cos\theta, \quad (19)$$

$$-2D_0 \left(\frac{G_2}{l_+} + \frac{G_3}{l_-} \right) = 0, \quad (20)$$

and

$$-2D_0(-i) \left(-\frac{G_2}{l_+} + \frac{G_3}{l_-} \right) = -j_e \sin\theta. \quad (21)$$

Thus we determine the constants to be

$$G_1 = -\frac{j_e \lambda_{\text{sdl}} (\beta - \cos\theta)}{2D_0(1 - \beta\beta')}, \quad (22)$$

$$G_2 = \frac{j_e l_+ \sin\theta}{4iD_0}, \quad (23)$$

and

$$G_3 = -\frac{j_e l_- \sin\theta}{4iD_0}. \quad (24)$$

Therefore, we find the transverse spin accumulation

$$\mathbf{m}_\perp = -\left(\frac{j_e}{2D_0} \right) [\text{Im}(l_+ e^{-x/l_+}) \mathbf{M}_d^{(2)} + \text{Re}(l_+ e^{-x/l_+}) \times (\mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)})] \times \mathbf{M}_d^{(1)}, \quad (25)$$

where we have used $-\sin\theta\mathbf{e}_x = \mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)}$ and $\sin\theta \times \mathbf{e}_y = (\mathbf{M}_d^{(2)} \times \mathbf{M}_d^{(1)}) \times \mathbf{M}_d^{(1)}$. We immediately see that the form of the spin accumulation given above is precisely the form we used in deriving Eq. (15). To obtain the coefficients a and b entering Eq. (15) we average this spin accumulation over $0 \leq x \leq t_F$ where t_F is the thickness of the $F1$ layer and find

$$a = -\frac{Jj_e}{2D_0 t_F} \text{Im}[l_+^2 (1 - e^{-t_F/l_+})] \quad (26)$$

and

$$b = \frac{Jj_e}{2D_0 t_F} \text{Re}[l_+^2 (1 - e^{-t_F/l_+})]. \quad (27)$$

It is noted that both a and b change sign under time reversal. The former agrees with that found in [1–3], while the latter has not been considered by these authors. To

estimate a and b , we take the limit $\lambda_{sf} \gg \lambda_J$ in which case $l_+ = (1 + i)\lambda_J/\sqrt{2}$. By placing this into Eqs. (26) and (27), we find

$$a = -\frac{hj_e a_0^3}{\sqrt{2} e \mu_B \lambda_J} \left(\frac{1 - \cos \xi e^{-\xi}}{\xi} \right) \quad (28)$$

and

$$b = \frac{hj_e a_0^3}{\sqrt{2} e \mu_B \lambda_J} \left(\frac{\sin \xi e^{-\xi}}{\xi} \right), \quad (29)$$

where $\xi = t_F/(\sqrt{2} \lambda_J)$, a_0 is the lattice constant, and we have reinserted the electric charge and Bohr magneton so that a and b have units of a magnetic field. If we take $\lambda_J = 20 \text{ \AA}$, $a_0 = 2 \text{ \AA}$, $j_e = 10^{11} \text{ A/m}^2$, we find $a = -1056 \text{ (Oe)}$ and $b = 457 \text{ (Oe)}$ for a typical experiment with $t_F = 25 \text{ \AA}$.

In conclusion, we have found that by considering the exchange forces between the conduction electron spin and the background magnetization for the spin current perpendicular to the layers of a magnetic multilayer there exists the effective field *and* torque, *both* of which contribute to current driven reversal of the magnetization. We treat both terms on an equal footing and demonstrate that they have a common origin. Our solution differs in two important aspects from previous work: we find the *longitudinal* spin accumulation does not play a role in the switching, and the spin torque, as well as the effective field, arises from a region in the magnetic layer within $\sim \lambda_J$ of the interface. Therefore, in the diffusive regime, the decay length in our theory is related *neither* to the phase of the wave function [1,2], *nor* to the spin diffusion length as in the effective field concept of switching [4].

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 - [13] To obtain a simple and transparent result we have neglected the spin accumulation in the thick layer, i.e., we have assumed the spin current in the thick layer is unchanged up to the interface. This is strictly valid only for half-metallic ferromagnetic layers. A more realistic calculation should treat the spin accumulation in the thick layer on an equal footing as the thin layer, i.e., in a self-consistent manner; one should also include the interface spin dependent scattering (reflections and diffuse) which is omitted in our simple illustration.