Domain-wall resonance induced by spin-polarized current in metal thin films with stripe structures

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The dynamics of a domain wall induced by spin-polarized current is studied in ferromagnetic metallic thin films with stripe structures. The solution of the generalized Landau-Lifshitz-Gilbert equation including adiabatic and nonadiabatic spin torque is used to describe motion and distortion of the domain wall. We show that ac currents facilitate the manipulation of the domain wall since domain wall resonance decreases the critical current density for depinning. In contrast with dc current, both the two spin torques induced by ac currents can produce wall motion and distortion.

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I. INTRODUCTION

The interaction between electronic transport and domain wall motion has received much attention due to the promise of applications in nonvolatile memory devices. In a ferromagnetic metal, the electron scattering by a domain wall due to *s*-*d* exchange coupling is a source of giant magnetoresistance.[1,2] Conversely, the spin-polarized current can induce the domain wall motion and distortion in the same way as magnetic fields.[3–12] These effects can be used to change the magnetic configuration in confined structures.[13–15] When a spin-polarized current traverses a domain wall, the conductance electrons precess in an exchange field and effect a torque on the local moments due to the exchange coupling between the conductive electrons and the local moments in the domain wall. This spin-transfer torque was proposed by Slonczewski[16] and Berger[17] in 1996. They consider the case of magnetic multilayers and the spin torque is exerted on single-domain ferromagnetic films. Other spin-transfer torque effects have been observed for ballistic[18,19] and diffusive[3,4,20] transport in multidomain ferromagnets.

Stripe-domain materials have the potential for memory device application. Field-driven domain wall dynamics in insulating, such as yttrium iron garnet (YIG),[21] and metallic[22] magnetic thin films with stripe-domain structure have been extensively studied. Spin-polarized current are more easily localized than magnetic fields. Thus, it is interesting to explore the effect of spin-polarized current on these confined structures.

Here, we investigate the action of spin-polarized current on the stripe-domain structure of metallic thin films. We discuss the roles of adiabatic and nonadiabatic torques in driving the wall motion. In the case of dc current, the former distorts the wall's internal structure, while the latter produces a pressure on the wall. However, for ac current, these two torques both have an effect on the motion and distortion. We describe the current-induced domain-wall motion and distortion by a forced oscillating equation with friction. We show that the critical current density for displacing domain wall from a restoring potential is decreased by using ac current due to resonance.

In the following, we display the equations of domain-wall motion in the presence of current in Sec. II. In Sec. III, we calculate critical current density for depinning domain wall by direct and alternating current, respectively. A summary is given in Sec. IV.

II. EQUATION OF DOMAIN-WALL MOTION

We consider a ferromagnetic metallic film in the xy plane, with the easy axis z perpendicular to the film, as shown in Fig. 1. An in-plane current along the y direction goes through the film. This current exerts spin-transfer torques on the local magnetization. In the presence of adiabatic and nonadiabatic spin-transfer torques, the generalized Landau-Lifshitz-Gilbert (LLG) equation describing dynamics of magnetization in film can be written as

$$\frac{\partial \mathbf{M}}{\partial t} = \gamma \mathbf{M} \times \frac{\delta W}{\delta \mathbf{M}} + \frac{\alpha}{M_s} \mathbf{M} \times \frac{\partial \mathbf{M}}{\partial t} + \mathbf{T}_a + \mathbf{T}_n, \qquad (1)$$

where $\mathbf{M}$ is the magnetization of ferromagnet, M_s is the saturation magnetization, γ is the gyromagnetic ratio, α is the Gilbert damping parameter, and W is the energy density. $\mathbf{T}_a$ and $\mathbf{T}_n$ are the adiabatic and nonadiabatic spin torque.[3,4] $\mathbf{T}_a = -(a_J/M_s^2)\mathbf{M} \times [\mathbf{M} \times (\partial \mathbf{M}/\partial y)]$, where $a_J = a j_e$, $a = P\mu_B/[eM_s(1+\xi^2)]$, $\xi = \tau_{ex}/\tau_{sf}$, $\tau_{ex} = \hbar/(SJ_{ex})$, J_{ex} is the *s*-*d* exchange coupling strength, τ_{sf} is the spin-flip relaxation time, P is the spin polarization of the current, μ_B is the Bohr magneton, and j_e is the electric current density. This torque is caused by the nonequilibrium conduction electrons whose polarizations coincide with the local moments. $\mathbf{T}_n =$

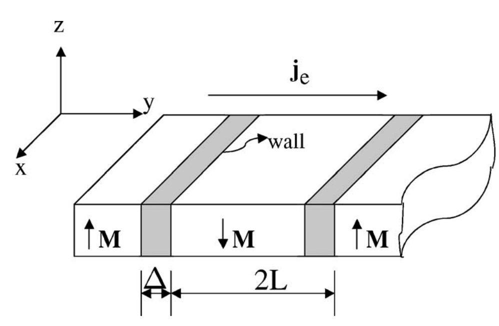

FIG. 1. Geometry of the film with stripe-domain structures.

$-(n_J/M_s)\mathbf{M}\times(\partial\mathbf{M}/\partial y)$, with $n_J=nj_e$ and $n=P\mu_B\xi/[eM_s(1+\xi^2)]$. This torque is the effect of the nonequilibrium conductive electrons which polarizations deviate from the local moments.

Taking into account $\mathbf{M}^2=M_s^2=const$ below the Curie temperature, Eq. (1) can be transformed into two coupled equations for the polar angle θ and azimuth angle ϕ defined relative to the z axis

$$\frac{\partial\theta}{\partial t}+\alpha\sin\theta\frac{\partial\phi}{\partial t}=-\frac{\gamma}{M_s\sin\theta}\frac{\delta W}{\delta\phi}+a_J\frac{\partial\theta}{\partial y}+n_J\sin\theta\frac{\partial\phi}{\partial y},$$

$$\sin\theta\frac{\partial\phi}{\partial t}-\alpha\frac{\partial\theta}{\partial t}=\frac{\gamma}{M_s}\frac{\delta W}{\delta\theta}+a_J\sin\theta\frac{\partial\phi}{\partial y}-n_J\frac{\partial\theta}{\partial y}. \quad (2)$$

In order to solve above LLG equations, it is convenient to write them in an equivalent variational form

$$\begin{aligned}\delta W&=\left(\frac{\delta W}{\delta\theta}\right)\delta\theta+\left(\frac{\delta W}{\delta\phi}\right)\delta\phi\\&=\frac{M_s}{\gamma}\Bigg[\left(-\alpha\frac{\partial\theta}{\partial t}+\sin\theta\frac{\partial\phi}{\partial t}+n_J\frac{\partial\theta}{\partial y}-a_J\sin\theta\frac{\partial\phi}{\partial y}\right)\delta\theta\\&\quad+\Bigg(-\sin\theta\frac{\partial\theta}{\partial t}-\alpha\sin^2\theta\frac{\partial\phi}{\partial t}\\&\quad+a_J\sin\theta\frac{\partial\theta}{\partial y}+n_J\sin^2\theta\frac{\partial\phi}{\partial y}\Bigg)\delta\phi\Bigg]. \quad (3)\end{aligned}$$

We introduce the dynamic wall structure as follows[23]

$$\theta(y,t)=2\arctan\exp\left[\frac{y-q(t)}{\Delta(t)}\right],$$

$$\phi(y,t)=\psi(t), \quad (4)$$

where $\Delta(t)$ represents the wall width, and $q(t)$ is the normal-displacement coordinate of the wall.

Including the local exchange energy, the anisotropy energy, and the demagnetizing energy, we write the total instantaneous energy density in terms of θ and ϕ

$$\begin{aligned}W_{inst}&=A\left[\left(\frac{\partial\theta}{\partial y}\right)^2+\left(\sin\theta\frac{\partial\phi}{\partial y}\right)^2\right]+K\sin^2\theta\\&\quad+2\pi M_s^2\sin^2\theta\sin^2\phi, \quad (5)\end{aligned}$$

where A is the exchange stiffness coefficient, and K is the easy-axis anisotropy constant. We consider the one-dimensional model in which θ and ϕ vary only in the y direction and the thickness of film is larger than the wall width. This condition ignores the surface effects.[24]

For the special case of no current and driving field, taking the static solution for a π domain wall with $\theta(y)=2\arctan(e^{y/\Delta})$ and $\phi(y)=\psi=const$ in Eq. (5), and assuming the wall width Δ as a variable to be determined, the instantaneous wall energy is given[25]

$$\sigma_{inst}=\int W_{inst}dy=\int_0^\pi W_{inst}\left(\frac{\partial\theta}{\partial y}\right)^{-1}d\theta=2A/\Delta+2\Delta\kappa, \quad (6)$$

where $\kappa=K+2\pi M_s^2\sin^2\psi$. Minimizing σ_{inst} with respect to Δ, the instantaneous wall energy is $\sigma_{inst}=4\sqrt{A\kappa}$, and the wall width is $\Delta=\sqrt{A/\kappa}$.

For the case of nonzero current, we assume the wall-magnetization angle ϕ is independent of y. In general, $\Delta(t)$ varies less than $q(t)$, namely, the distortion of wall width is small during the wall motion, $|d\Delta/dt|\ll|dq/dt|$. So from Eq. (4), one can obtain

$$\frac{\partial\theta}{\partial y}=\Delta^{-1}\sin\theta,$$

$$\delta\theta=-\Delta^{-1}\sin\theta dq,$$

$$\frac{\partial\theta}{\partial t}=-\Delta^{-1}\sin\theta\frac{\partial q}{\partial t}. \quad (7)$$

Integrating Eq. (3) over y and substituting Eq. (7), a pair of partial differential equations about wall energy is obtained

$$\frac{\partial\sigma}{\partial\psi}=\frac{2M_s}{\gamma}\left(\frac{dq}{dt}-\alpha\Delta\frac{d\psi}{dt}+a_J\right),$$

$$\frac{\partial\sigma}{\partial q}=\frac{2M_s}{\gamma}\left(-\alpha\Delta^{-1}\frac{dq}{dt}-\frac{d\psi}{dt}-n_J\Delta^{-1}\right). \quad (8)$$

Equation (8) expresses the change of domain-wall energy arising from local magnetization precession, vicious damping, and spin-polarized current. In high-Q ($K>2\pi M_s^2$) stripe-domain structure, the domain wall sits in a potential well produced by the demagnetization field.[21] The restoring energy density is written as $\sigma_p=kq^2/2$ with $k=16\pi M_s^2/L$ and L is half the length of domain. For materials with sufficiently low coercivity, it is reasonable to only consider the restoring force generated by demagnetization field. So, $\sigma=\sigma_{inst}+\sigma_p$. We notice that $\partial\sigma/\partial\psi$ represents the torque on the wall magnetization in the xz plane and the adiabatic spin torque can distort the wall's internal structure, because spin transfer is dominant in the adiabatic approximation.[6] From Eq. (8), we obtain

$$\frac{dq}{dt}=2\pi\gamma M_s\Delta\sin 2\psi+\alpha\Delta\frac{d\psi}{dt}-a_J. \quad (9)$$

$\partial\sigma/\partial q$ is the pressure on the wall. The nonadiabatic torques generate global pressure on the walls due to the reigning momentum transfer for the nonadiabatic electrons.[6] The instantaneous pressure arises from the applied field and magnetostatic restoring potential with a restoring force constant k. Thus,

$$\frac{d\psi}{dt}=-\gamma\frac{k}{2M_s}q-\alpha\Delta^{-1}\frac{dq}{dt}-n_J\Delta^{-1}. \quad (10)$$

We assume that at $t=0$, the current and perpendicular magnetic field is applied and the wall lies in the pinning center ($q=0$). Then, from Eqs. (9) and (10), the initial velocity of wall is $(-\alpha n_J - a_J)/(1+\alpha^2)$. Like the driving field, both the nonadiabatic and adiabatic torque give an initial velocity to the domain wall, but the former is damped. Also, ψ starts precessing at the Larmor frequency $-(n_J-\alpha a_J)/[\Delta(1+\alpha^2)]$. As ψ increases, The velocity decreases.

III. DOMAIN-WALL RESONANCE INDUCED BY CURRENT

Here, we discuss the case of small deviations of ψ from equilibrium, and the linear variation of the wall velocity with ψ. Linearizing Eq. (9) and eliminating ψ using Eq. (10), we obtain

$$m\frac{d^2q}{dt^2}+b\frac{dq}{dt}+kq=-\frac{2M_s}{\gamma}\left(\frac{n_J}{\Delta}+\frac{\alpha dn_J/dt+da_J/dt}{4\pi\gamma\Delta M_s}\right), \tag{11}$$

where

$$m=(1+\alpha^2)\frac{1}{2\pi\gamma^2\Delta},\quad b=\alpha\left(\frac{2M_s}{\gamma\Delta}+\frac{k}{4\pi\gamma M_s}\right), \tag{12}$$

with $\Delta=\sqrt{A/K}$.

In dc current case, adiabatic torque only twists the magnetization of wall, or transfers magnetic momentum to the wall, while nonadiabatic torque produces a global pressure on the wall.[4,6,8] However, in the ac current case, Eq. (11) indicates that both the adiabatic torque and nonadiabatic one contributes to the global motion of wall.

For an applied direct current, Eq. (11) describes damped oscillation of the wall. It is easy to see that only nonadiabatic spin torque affects wall motion in the dc current case. This is due to considering the domain wall as a particle and only nonadiabatic torque produces a pressure on it. The current changes the equilibrium location of the oscillation. In the absent of driving field, the change is evaluated as $\Delta q_0 = j_e\xi P\mu_B L/[8\pi e\gamma\Delta M_s^2(1+\xi^2)]$. If $\Delta q_0 \geqslant L$ (the range of the restoring potential), the wall is depinned. So, the current density for depinning is

$$j_{dc}^{cri}=8\pi(1+\xi^2)e\gamma\Delta M_s^2/(\xi P\mu_B). \tag{13}$$

Taking $J_{ex}\approx 1$ eV, $S=2$, $\tau_{sf}\approx 10^{-12}$ s, $\alpha\approx 10^{-2}$ for typical ferromagnetic metal, and $M_s=1125$ G, $A=1.0\times 10^{-6}$ erg/cm, $K=1.9\times 10^{7}$ erg/cm^3, $P=0.35$ for Co_3Pt, we can calculate the critical current density $j_{dc}^{cri}=6.46\times 10^{11}$ A/cm^2.

For the case of alternating current, taking $j_e=j_{ac}\cos\omega t$ with oscillating frequency ω, Eq. (11) describes forced damped oscillations. Considering $b^2<4mk$, we obtain

$$q=C\exp\left(-\frac{bt}{2m}\right)\cos\left(\frac{\sqrt{4mk-b^2}}{2m}t\right)-\frac{\frac{2M_s}{\gamma}\sqrt{\frac{n^2}{\Delta^2}+\frac{(\alpha n+a)^2}{(4\pi\gamma M_s)^2}\omega^2}\,j_{ac}}{\sqrt{(k-m\omega^2)^2+b^2\omega^2}}\cos(\omega t-\delta-\delta'), \tag{14}$$

where

$$\delta=\arctan\frac{b\omega}{k^2-m\omega^2},\quad \delta'=\arctan\frac{(\alpha n+a)\omega}{4\pi n\gamma M_s}, \tag{15}$$

and C is a real constant which can be determined from the initial conditions. The first term in Eq. (14) rapidly approaches to zero because of the large value of $b/2m$. From Eq. (14), the resonance frequency is $\omega=4\pi\gamma M_s[(4\Delta^2/L^2+\xi^2+\xi^4)^{1/2}-\xi^2]^{1/2}$, and smaller current leads to a relative larger displacement. Taking $q_{max}=L$, we can obtain the critical ac current density that can depin domain wall in the presence of resonance, which is

$$j_{ac}^{cri}=\frac{\xi}{2}\left\{\frac{L}{\Delta}\left[2\sqrt{\pi}\sqrt{(\pi\xi^4-\alpha^2\xi^2)\frac{L^2}{\Delta^2}+4\xi^2\frac{L}{\Delta}+4}+(\alpha^2-2\pi\xi^2)\frac{L}{\Delta}-4\right]\right\}^{1/2} j_{dc}^{cri}. \tag{16}$$

Using the material parameters of Co_3Pt, and taking $L=1\ \mu$m, the resonant frequency is $\omega=12.8$ GHz and $j_{ac}^{cri}=0.18 j_{dc}^{cri}$. Without the currents, to counter the restoring force, the driving field is not less than $8\pi M_s$. For example, the value is 2.8 T for Co_3Pt. Combining the magnetic field and spin current, the critical current density can also be reduced. Applying a magnetic field perpendicular to the film, adds a term $2M_sH_\perp$ to the right of Eq. (11). We obtain $\tilde{j}_{dc}^{cri}=[1-H_\perp/(8\pi M_s)]j_{dc}^{cri}$, and $\tilde{j}_{dc}^{cri}$ denotes the critical current density due to driving field and current.

IV. CONCLUSION

In summary, the dynamics of domain wall under the spin-polarized current in stripe-domain metallic films is investigated by means of the LLG equations. The action of adiabatic and nonadiabatic spin torque on motion and distortion of domain wall are different for dc and ac current. In the dc case, the two torques drive wall motion and distortion, respectively. Whereas, the roles are mixed for the ac current case. Furthermore, in order to depin the domain wall from a potential generated by demagnetization field, an ac current is preferred to the dc one because of decrease of critical current density induced by the resonance. Spin-polarized current control of magnetization configurations has the advantage of localizability. Stripe-domain ferromagnetic metallic films are candidates of information storage materials. Our results are potentially useful for applications in magnetic reading and writing.

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