Current-Induced Control of Spin-Wave Attenuation

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The current-induced modification of the attenuation of a propagating spin wave in a magnetic nanowire is studied theoretically and numerically. The attenuation length of spin wave can increase when the spin waves and electrons move in the same direction. It is directly affected by the nonadiabaticity of the spin-transfer torque and thus can be used to estimate the nonadiabaticity. When the nonadiabatic spin torque is sufficiently large, the attenuation length becomes negative, resulting in the amplification of spin waves.

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The spin wave (SW) is ubiquitous in magnetic systems and has long been a fundamental research topic in magnetism [1]. Much effort has been extended in understanding its wave properties such as the dispersion relation, reflection, and tunneling [2]. Thanks to recent progress in the fabrication of magnetic nanostructures, it is now possible to use SWs to deliver signal information. Several device concepts have been proposed such as the SW logic device [3] and the SW bus interconnect [4]. Despite the soundness of the concepts, however, the SW amplitude measured at a distance from a source is very marginal because of a substantial attenuation of SW amplitude, which is the fundamental limitation of SW for applications.

The amplitude attenuation is caused by the dissipation of the magnetic energy into the environment and thus related to the intrinsic damping. Recently, the current-induced manipulation of local magnetization via spin-transfer torque (STT) [5] has received considerable attention because of its significance in the fundamental understanding of spin transport and potential for applications in nanoscale magnetic devices. STT provides the antidamping effect and enables a new class of current-induced magnetization dynamics such as full magnetization reversal [6], steady-state precession [7], and domain wall motion [8]. It also occurs for SWs and causes the current-induced SW Doppler shift in the adiabatic limit as theoretically predicted [9] and experimentally confirmed [10]. Until now, however, the antidamping effect of STT on the attenuation of propagating SWs has not been investigated yet.

In continuously varying magnetization such as SWs, the spin-transfer torque τ_s is composed of the adiabatic and the nonadiabatic terms, and given by [11]

$$\tau_s = \mathbf{u} \cdot \nabla \hat{m} - \beta[\mathbf{u} \cdot (\hat{m} \times \nabla)]\hat{m} \tag{1}$$

where $\hat{m}$ is the unit vector of magnetization, $\mathbf{u} = u_0\hat{x}$, $u_0(= \mu_B j_e P/eM_S)$ is the magnitude of adiabatic spin torque, P is the spin polarization, j_e is the current density, M_S is the saturation magnetization, and β is the ratio of nonadiabatic spin torque to adiabatic one which is related to spin relaxation or momentum transfer [11]. Although the magnitude of β is predicted to be small [i.e., the order of the damping constant $\alpha(\sim 0.01)$], it is crucially important to understand the spin transport in the mesoscopic system [12–15] and the damping mechanism [16]. However, the magnitude of β is still highly controversial. For instance, the experimental estimations of β are widely distributed, i.e., $\beta = 8\alpha$ [17], $\beta = \alpha$ [18], $\beta = 2\alpha$ [19], $\beta \neq \alpha$ [20], and $\beta > 1$ [21]. Note that all the experiments have estimated β using the current-induced domain wall motion. This wide distribution may be caused by the fact that the domain wall is a localized object and very sensitive to local defects such as edge roughness of nanowire. In contrast, SWs are not localized and thus can be an attractive tool to estimate β.

In this Letter, we theoretically show that by injecting an electric current through a magnetic nanowire where SW is *externally* excited at a local region, it is possible to suppress the SW attenuation and even amplify the SWs when the nonadiabatic torque overcomes the intrinsic damping

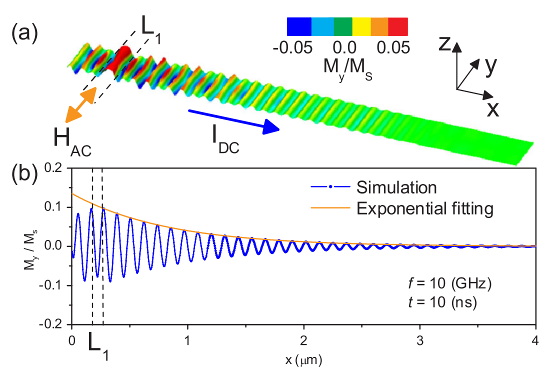

FIG. 1 (color online). (a) Schematic of the propagating spin waves along the nanowire (width = 120 nm and thickness = 6 nm). Perspective-view images for the normalized M_z components and the color contour for the normalized M_y components. (b) The spatial variation of the normalized M_y component is shown for $|H_{ac}| = 100$ Oe and $u_0 = 0$ m/s.

torque. Therefore, β can be experimentally determined by measuring the SW amplitude depending on the current polarity. Our finding is related to the current-induced SW instability [14,22,23]. Previous studies on the SW instability have focused on current-induced domain nucleations from uniformly magnetized state without an external SW excitation. The domain nucleation occurs in a chaotic manner when the current increases above a threshold [24]. In contrast, when the SW is externally excited at a local region, the SW dynamics is not chaotic for a certain distance from the source even at the current above a threshold, and thus experimentally detectable. We verify our theoretical prediction by using micromagnetic study. Finally, we show how high accuracy in measuring the SW amplitude is needed to experimentally estimate β.

The model system is shown in Fig. 1(a). An alternating field with the frequency of $\omega_f/2\pi$ and the magnitude of 100 Oe is applied at the region L_1 (source) to steadily excite SWs. Note that ω_f is determined externally regardless of the intrinsic magnetic properties of the magnetic system. The SW dynamics is described by the Landau-Lifshitz-Gilbert (LLG) equation including the spin torques,

$$\frac{\partial \hat{m}}{\partial t} = \gamma \hat{m} \times \vec{H}_{\text{eff}} + \alpha \hat{m} \times \frac{\partial \hat{m}}{\partial t} + \tau_s \tag{2}$$

where γ is the gyromagnetic ratio, and $\vec{H}_{\text{eff}}$ is the effective field [Eq. (4)]. For the numerical modeling, standard parameters of Permalloy are used for a magnetic nanowire: $\alpha = 0.01$, $M_S = 800$ emu/cm^3, the exchange stiffness constant $A = 1.3 \times 10^{-6}$ erg/cm, $P = 0.7$, and the unit cells of $4 \times 4 \times 6$ nm^3. For a legitimate comparison between theoretical and numerical results, reflected SWs from long edges of nanowire should be properly prevented. We have tested two different absorbing boundary conditions (ABCs) based on site-dependent damping constant; One uses smooth increase of α near edges [25] and the other uses abrupt increase of α at boundary cells [26]. We found both ABCs can successfully prevent the SW reflection when the final value of α at the boundary region is 1 (i.e., 100 times larger than the intrinsic α), which is consistent with [26].

Figure 1(b) shows a snapshot image of the SW propagation without the current injection, obtained from the micromagnetic simulation (nanowire width = 120 nm and thickness = 6 nm). The SW amplitude exponentially attenuates as the distance from the source increases. Consequently, we consider a small amplitude SW which decays exponentially in the perturbative limit,

$$\hat{m} = \hat{x} + \vec{m}_0 \exp[i(\omega t + kx)] \exp[-x/\Lambda] \tag{3}$$

where $|\vec{m}_0| \ll 1$, and Λ is the characteristic attenuation length. In a nanowire, a net effective field is given by

$$\vec{H}_{\text{eff}} = H_k m_x \hat{x} + D\nabla^2 \hat{m} - H_d m_z \hat{z} \tag{4}$$

where H_k and H_d are the easy axis and the hard axis anisotropy fields, respectively, and D is $2A/M_S$. Since we assume no crystalline anisotropy field, H_k and H_d are determined by the shape of nanowire and given by $(N_y - N_x)M_S$ and $(N_z - N_y)M_S$, respectively. Here, N_x, N_y, and N_z are demagnetization factors along the length, width, and thickness directions, respectively (i.e., for a infinitely long nanowire with cross section of 120×6 nm^2, $N_x = 0$, $N_y = 0.899$, $N_z = 11.667$ [27] in the cgs unit). By inserting Eqs. (3) and (4) into Eq. (2), and using $\text{Re}[\omega] = \omega_f$, $\text{Im}[\omega] = 0$, and $1/(k\Lambda)^2 \ll 1$, one finds

$$\begin{aligned} &-(\omega_f - u_0 k)^2 + (u_0/\Lambda)^2 + \gamma H_d \beta u_0/\Lambda \\ &+ [-\gamma(H_k + H_d + Dk^2) + \beta u_0/\Lambda]^2 \\ &- (-2\gamma Dk/\Lambda + \alpha\omega_f - \beta u_0 k)^2 \\ &\qquad - \gamma^2 H_d(H_k + H_d + Dk^2) = 0. \end{aligned} \tag{5}$$

Neglecting small terms proportional to α^2, β^2, and $\alpha\beta$, the dispersion relation with considering the attenuation is given as,

$$\begin{aligned} (\omega_f - u_0 k)^2 = \gamma^2(H_k + H_d + Dk^2)(H_k + Dk^2) \\ + 4\gamma Dk\Gamma/\Lambda. \end{aligned} \tag{6}$$

where $\Gamma = \alpha\omega_f - \beta u_0 k - \beta u_0(H_k + H_d/2 + Dk^2)/(2Dk) - \gamma Dk/\Lambda$. One can reproduce the SW Doppler shift in Ref. [9] when dropping the $4\gamma Dk\Gamma/\Lambda$ term.

From Eq. (5), the SW attenuation length Λ in the presence of the current is given by

$$\Lambda = \frac{2a}{b - \sqrt{b^2 - 4ac}}, \tag{7}$$

where $a = 2\gamma Dk\beta u_0$, $b = (\omega_f - u_0 k)u_0 + 2\gamma^2 Dk(H_k + H_d/2 + Dk^2) + \beta u_0(\alpha\omega_f - \beta u_0 k)$, and $c = \gamma(H_k + H_d/2 + Dk^2)(\alpha\omega_f - \beta u_0 k)$. Equation (7) is the main finding in this Letter.

To get the insight of underlying physics, we show an extreme case of $2\gamma^2 Dk(H_k + H_d/2 + Dk^2) \gg (\omega_f - u_0 k)u_0$ and $\beta u_0(\alpha\omega_f - \beta u_0 k)$. In this case, Eq. (7) is further simplified as

$$\Lambda = \frac{2\gamma Dk}{\alpha\omega_f - \beta u_0 k}. \tag{8}$$

From Eq. (8) [28], it is evident that when β is nonzero, the SW attenuation length Λ can increase or decrease depending on the relative sign of u_0 and k (i.e., the electron-flow direction and the SW propagation direction). In other words, for a nonzero $\beta(>0)$, the SW amplitude is less (more) attenuated when $u_0 k$ is positive (negative); i.e., SWs and electrons propagate toward the same (opposite) position. It indicates that the nonadiabatic spin torque compensates the intrinsic damping torque and thus can control the SW attenuation.

Another interesting point is that Λ becomes negative for one polarity of the current when $\beta u_0 k/\alpha$ exceeds ω_f, and thus SW amplitude exponentially grows, i.e., SW amplification. The threshold spin torque u_c for the SW amplification can be obtained from the dispersion relation [Eq. (6)] and the condition that the denominator of Eq. (7) becomes negative;

$$u_c = \frac{\alpha\omega_f}{\beta} \times \left[\frac{D}{-2\pi M_S + \sqrt{(H_d/2)^2 + (1-\alpha/\beta)^2(\omega_f/\gamma)^2}}\right]^{1/2}. \tag{9}$$

It is worthwhile comparing our result to the SW instability [14,22–24]. In the SW instability, the threshold spin current for nucleating multidomains from the uniformly magnetized state is proportional to $|1 - \beta/\alpha|^{-1}$ and thus goes to infinite when $\beta = \alpha$. It is because no external source of SW is in the system and in other words, ω is not ω_f but $u_0 k$. For the SWs excited externally and locally, however, it is nonzero when $\beta = \alpha$ because ω_f is externally determined. Instead, it goes to the infinite when $\beta = 0$ since there is no nonadiabatic torque and thus no antidamping effect in this situation.

In order to verify the theoretical prediction, we perform micromagnetic simulations. Figure 2(a) shows numerical results of the dispersion relation in the adiabatic limit ($\beta = 0$). When the frequency $f(= \omega_f/2\pi)$ is smaller than a bottom limit f_0 ($\sim$8 GHz), SWs cannot propagate along the nanowire. This bottom limit is consistent with the Kittel's ferromagnetic resonance frequency [29]. As was predicted, we observe the SW Doppler shift (Δk). Note that we get Δk instead of Δf since we assume a fixed ω_f in this work. Δk is positive (negative) when electrons flow in the opposite (same) direction with SWs. Figure 2(b) shows the spatial variation of M_y as a function of β. The wavelength of SWs does not change with β. It means that the SW Doppler shift is entirely associated with the adiabatic torque. However, the amplitude of SWs changes with β. For $u_0 = \pm 100$ m/s, the SW amplitude enhances (reduces) with increasing β, consistent with the theoretical prediction.

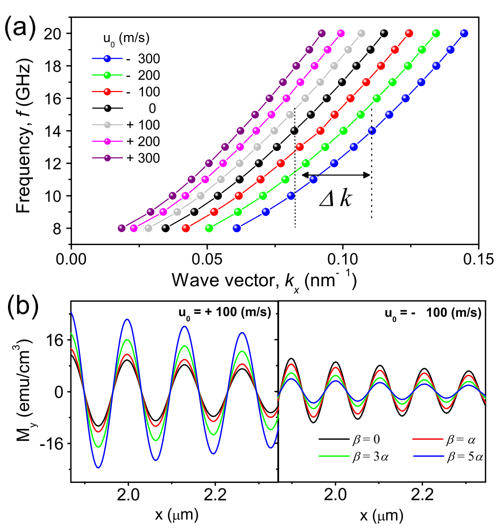

FIG. 2 (color online). (a) Dispersion relations (f versus k_x) at various current densities. (b) The spatial variation of y component of the magnetization for several cases of nonadiabatic spin torque terms. Left and right plots correspond to the cases of $u_0 = +100$ and -100 m/s, respectively.

In Figs. 3(a)–3(d), Λ is plotted as a function of the current at various frequencies and β-terms. The numerical values of Λ are in excellent agreement with the predicted values from Eq. (7). For $u_0 > 0$, Λ increases with increasing β. Remembering that the magnitude of β is still controversial, one can experimentally determine the β by comparing the current-dependent Λ to our theoretical prediction. When the nonadiabatic torque βu_0 is large enough, Λ becomes negative, and thus the SW amplitude exponentially grows as the distance from the source increases, i.e., SW amplification. A numerical example of the SW amplification is shown in Fig. 3(e). The SW amplification is maintained for the critical distance x_c where $M_y/M_S < 1$.

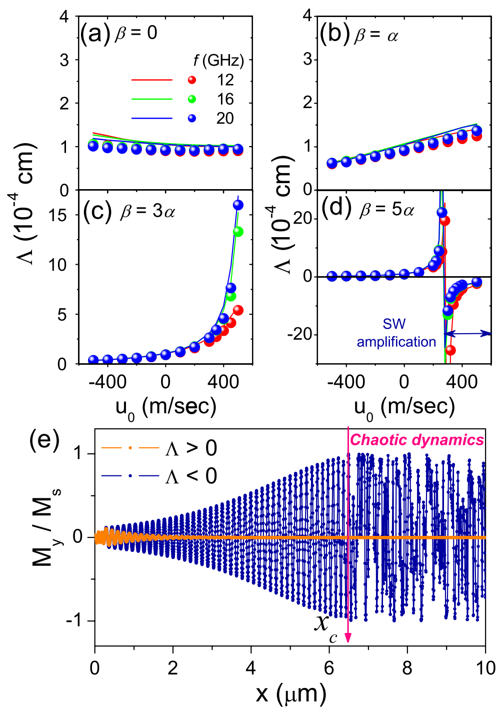

FIG. 3 (color online). (a)–(d) The attenuation length ($=\Lambda$) as a function of the current at various frequencies and β terms. The symbols are the simulation results. The lines are the theoretical values predicted from Eq. (7). (e) The spatial variation of the normalized M_y component is shown for $f = 10$ GHz, $|H_{ac}| = 100$ Oe, $u_0 = +200$ m/s, and the time step of $t = 20$ ns. Here, the cases of positive ($\Lambda > 0$) and negative ($\Lambda < 0$) attenuation lengths correspond to the cases of $\beta = 10\alpha$ and $\beta = 0$, respectively. For $\Lambda > 0$, a chaotic dynamics appears when M_y reaches M_S at beyond the critical position (x_c).

Beyond this limit, the magnetization dynamics becomes chaotic. Therefore, one should place the SW detector at the position shorter than x_c for applications of the SW amplification.

Finally, we show how high accuracy in measuring the SW amplitude is needed to experimentally estimate β. The SW amplitude $A^{\pm}$ at $u_0 = \pm u$ normalized by that at $u_0 = 0$ is given by $A^{\pm} = \exp[x/\Lambda_{u_0=\pm u} - x/\Lambda_{u_0=0}]$. Using the above equations and the Taylor's expansion, the difference of the normalized amplitudes depending on the current polarity, $\Delta A(= A^+ - A^-)$ is approximated by

$$\Delta A = \frac{x\beta u}{\gamma D}. \tag{10}$$

Note that Eq. (10) is independent of ω_f and k, and thus a universal relation regardless of the nanowire geometry. When $u = 10$ m/s and $x = 1$ μm, ΔA for a Permalloy nanowire is then 1.75β. Thus, if β is the order of 0.01, $\Delta A \times 100\%$ is 1.75%. Consequently, when the amplitude of SWs can be determined with enough accuracy which can distinguish the difference of about 1%, one can determine the absolute value of β from such kind of experiments. To achieve such accuracy, for instance, one can use tunnel magnetoresistance which was recently adopted to measure the domain wall motion in nanowires [30].

To conclude, the theoretical equation for the SW attenuation affected by the spin-transfer torque is presented. Some theories predicted that the magnitude of β can be controlled by changing the density of magnetic scatterer [12,14]. Our results open a way to estimate the β-term in various itinerant ferromagnets and will be potentially useful for applications of SW-active devices where the gain is required.

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