Spin-Wave Excitation in Magnetic Insulators by Spin-Transfer Torque

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(Received 17 January 2012; published 23 May 2012)

We study the excitation of spin waves in magnetic insulators by the current-induced spin-transfer torque. We predict preferential excitation of surface spin waves induced by an easy-axis surface anisotropy with critical current inversely proportional to the penetration depth and surface anisotropy. The surface modes strongly reduce the critical current and enhance the excitation power of the current-induced magnetization dynamics.

TABLE I. Parameters for YIG.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>YIG</th>
<th>Unit</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \gamma )</td>
<td>( 1.76 \times 10^{11} )</td>
<td>( 1/(\text{Ts}) )</td>
</tr>
<tr>
<td>( M_s )</td>
<td>( 1.56 \times 10^{5} )</td>
<td>( \mu_0 \text{m} )</td>
</tr>
<tr>
<td>( \omega_M = \gamma \mu_0 M_s )</td>
<td>( 34.5 )</td>
<td>( \text{GHz} )</td>
</tr>
<tr>
<td>( \Lambda_{ex} )</td>
<td>( 4.74 \times 10^{-6} )</td>
<td>( \text{m}^2/\text{s} )</td>
</tr>
<tr>
<td>( \alpha )</td>
<td>( 6.7 \times 10^{-5} )</td>
<td>( \cdots )</td>
</tr>
<tr>
<td>( \omega_0 = \gamma H_0 )</td>
<td>( 0.5 \omega_M )</td>
<td>( \text{GHz} )</td>
</tr>
<tr>
<td>( K_s )</td>
<td>( 5 \times 10^{-5} )</td>
<td>( 1/\text{m}^2 )</td>
</tr>
</tbody>
</table>

*Ref. [5].

\( K_s \) ranges 0.01–0.1 erg/cm\(^2\) or \( 10^{-5} \cdots 10^{-4} \text{ J/m}^2 \), Refs. [12,13].

Spintronics is all about manipulation and transport of the spin, the intrinsic angular momentum of the electron [1]. These two tasks are incompatible, since manipulation requires strong coupling of the spin with the outside world, which perturbs transport over long distances. In normal metals, spin can be injected and read out easily, but the spin information is lost over short distances [2]. In spin-based interconnects, transporting spins over longer distances is highly desirable [3].

The long-range transport of spin information can be achieved by encoding the information into spin waves that are known to propagate coherently over centimeters [4]. It has been demonstrated in Refs. [5–7] for the magnetic insulator Yttrium-Iron-Garnet (YIG) that spin waves can be actuated electrically by the spin-transfer torque [8,9] and detected by spin pumping [10] at a distant contact. In the experiment by Kajiwara [5], Pt was used as spin current injector and detector, making use of the (inverse) spin Hall effect [11]. In a current injector and detector, transporting spins over longer distances is strongly decreased in the presence of an easy-axis surface anisotropy (EASA). Simultaneously, EASA increases the power of the spin wave excitation by at least two orders of magnitude.

We study a structure as depicted in Fig. 1, where a nonmagnetic (N) metallic thin film of thickness \( t \) is in contact with a ferromagnetic insulator (FI), whose equilibrium magnetization is along the \( z \) direction. The spin current injected into the ferromagnetic insulator is polarized transverse to the magnetization \( \mathbf{J}_c = J_c \mathbf{m} \times \mathbf{z} \times \mathbf{m} \). The bulk magnetization is described by the Landau-Lifshitz-Gilbert (LLG) equation:

\[
\mathbf{m} = -\gamma \mathbf{m} \times [\mathbf{H}_0 + (A_{ex}/\gamma)\nabla^2 \mathbf{m} + \mathbf{h}] + \alpha \mathbf{m} \cdot \nabla \mathbf{m}.
\] (1)

where \( \mathbf{H}_0 \) includes the external and internal magnetic field, \( A_{ex} \) is the exchange constant, \( \mathbf{h} \) is the dipolar field that satisfies Maxwell’s equations, and \( \alpha \) is the total effective damping parameter. The damping caused by the 2-magnon processes in ultrathin ferromagnets [16] is assumed to be constant within the small frequency range of interest. In the quasistatic approximation, i.e., disregarding retardation in the electromagnetic waves, \( \nabla \times \mathbf{h} = 0 \) and \( \nabla \cdot \mathbf{b} = \nabla \cdot (\mathbf{h} + \mu_0 \mathbf{M} \cdot \mathbf{m}) = 0 \). All quantities are position- and...
time-dependent. In the absence of pinning, the total torque vanishes at the interface [17]:

$$A_{ex} \mathbf{m} \times \frac{\partial \mathbf{m}}{\partial n} - \frac{2 \gamma K_s}{M_s} (\mathbf{m} \cdot \mathbf{n}) \mathbf{m} \times \mathbf{n} + \gamma J_s \mathbf{m} \times \hat{z} \times \mathbf{m} = 0,$$

(2)

where \( \mathbf{n} \) is the outward normal as seen from the ferromagnet. The first term in Eq. (2) is the surface exchange torque, the second term the torque due to a perpendicular uniaxial surface anisotropy \( \mathbf{H}_a = \frac{2K_s}{M_s} (\mathbf{m} \cdot \mathbf{n}) \mathbf{n} \) and \( K_s = \int dx K_s \) across the surface, and the last term is the current-induced spin-transfer torque [18]. We parameterize the surface anisotropy and spin current as wave numbers \( k_x = 2\gamma K_s/A_{ex} M_s \) and \( k_y = \gamma J_s/A_{ex} M_s \). The dipolar fields \( h_y \) and \( h_x \) are continuous across the interface. Eqs. (1) and (2) in combination with Maxwell’s equations describe the low energy magnetization dynamics and can be transformed into a sixth-order differential equation for the scalar potential \( \psi \) with \( \mathbf{h} = -\nabla \psi \) [19,20].

The method described above extends previous studies by Hillebrands [20] and references therein by including the current-induced spin-transfer torque. We predict the critical conditions under which magnetization dynamics becomes amplified by the current-induced driving torque.

We start with the limiting case of \( d \to \infty \) (semi-infinite ferromagnet). After linearization and Fourier transformation in both time and space domains, Eq. (1) reduces to a fourth-order differential equation in \( \psi \). Focusing for simplicity first on the case of vanishing in-plane wave vector \( q = (q^x, q^y) = 0 \), the scalar potential can be written as

$$\psi(r) = \sum_{j=1} a_j e^{i\omega t} e^{i q r},$$

with

$$q^j_\omega = -i \left( \omega_0 + \frac{2}{2} \omega_M \pm \sqrt{\omega^2 + \frac{1}{4} \omega_M^2 + i \omega \omega} \right)^{1/2},$$

(3)

and \( |q_1| \gg |q_2| \) when \( \omega \sim \omega_0 \). Imposing the boundary condition in Eq. (2), up to the first order in \( k_j \):

$$0 = 2q_1 q_2 (q_1 + q_2) + ik_j \left[ (q_1 + q_2)^2 + \frac{\omega_M}{A_{ex}} \right] + 4k_j \omega.$$

(4)

The solutions of Eq. (4) are the complex eigenfrequencies \( \omega \), whose real part represents the energy and imaginary part, the inverse lifetime. To zeroth order in dissipation, i.e., with vanishing bulk damping \( (\alpha = 0) \) and spin current injection \( (k_j = 0) \), and using \( |q_1| \gg |q_2| \), Eq. (4) simplifies to

$$k_j = 2i q_2 (1 + \omega / \omega_M) (\lambda q_2^2),$$

which has no nontrivial solution for \( k_j \leq 0 \). The single real solution for \( k_j > 0 \) obeys \( \omega < \sqrt{\omega_0 (\omega_0 + \omega_M)} / A_{ex} \), such that both \( q_{1,2} \) are negative imaginary:

$$q_1 \approx -i \sqrt{2 \omega_0 + \omega_M} / A_{ex}, \quad q_2 \approx -ik_j \omega_0 / (2 \omega_0 + \omega_M + O(k_j^2)),$$

i.e., a surface spin wave induced by the easy-axis surface anisotropy. With the criteria \( \Im \omega < 0 \) and to leading order in \( 0 < k_j \ll q_1 \), Eq. (4) leads to the critical current:

$$k_j^c = \frac{-\alpha \left( \omega_0 + \omega_M / 2 \right)^2}{A_{ex} \omega_0} + \frac{\omega_0 + 2 \omega_M}{4 \omega_0} \sqrt{2 \omega_0 + \omega_M / A_{ex}}.$$

(5)

When there is no surface anisotropy \( (k_j \to 0) \), the critical current diverges because the macrospin mode cannot be excited in a semi-infinite film. Using the parameters given in Table I in Eq. (5), we estimate the critical current for exciting the EASA-induced surface wave \((q = 0)\) to be

$$k_j^c = -0.08 \frac{d}{k_c},$$

where \( k_c = \alpha \omega_0 + \omega_M / 2 d / A_{ex} \) is the critical current for bulk excitation in a YIG thin film of thickness \( d = 0.61 \mu m \) (used below).

EASA pulls down a surface spin wave for the following reason: when \( k_j = J_s = 0 \), the boundary condition in Eq. (2) requires cancellation between the exchange and surface anisotropy torques: \( \partial_x m_x - k_m = \partial_y m_y = 0 \). The exchange torque depends on the magnetization derivative in the normal direction and can only take one sign in the whole film, and \( m_{x,y} \to 0 \) as \( x \to -\infty \), therefore \((1 / m_s) \partial_x m_x > 0 \). Torque cancellation (for a nontrivial solution) is therefore possible only for \( k_j > 0 \). The surface spin wave induced by EASA \((k_j > 0)\) for the in-plane magnetized film \((m_r \sim 1)\) discussed in this Letter is analogous to the surface spin waves for the perpendicular magnetized film \((m_r \sim 1)\) induced by easy-plane surface anisotropy \((k_j < 0)\) studied before in YIG films [17,21–24]. For perpendicular magnetization, a different boundary condition, \( \partial_x m_x + k_r m_{y,z} = 0 \), results in a surface wave for \( k_j < 0 \).

We now include all ingredients: finite thickness \((d = 0.61 \mu m)\), surface anisotropy, intrinsic magnetic damping, spin current injection, exchange coupling, and dipolar fields. We calculate numerically the complex eigenfrequencies \( \omega(q, k_j) \) as a function of the in-plane wave vector \( q \) and the applied spin current at the surface \( k_j \). \Im \omega, the effective dissipation, can be either positive (damping) or negative (amplification) when driven by the spin-transfer torque.

First, we disregard the surface anisotropy: \( K_s = k_s = 0 \). With \( \theta \) the angle between \( \mathbf{q} \) and \( \mathbf{m} \), the results for \( \theta = 90^\circ \) are shown in Fig. 2. In the top left panels \( \Re \omega \), the magnetostatic surface wave (MSW) is seen to cross the flat bulk bands [19]. When no spin current is applied \((k_j = 0)\), the dissipative part \( \Im \omega \sim \alpha (\omega_0 + \omega_M / 2) > 0 \), as shown in the top middle panels. At a spin current that is 20% of that required for bulk excitation: \( k_j = 0.2 k_c \), the

FIG. 1 (color online). An electrically insulating magnetic film of thickness \( d \) with magnetization \( \mathbf{m} \) (\( \parallel \hat{z} \) at equilibrium) in contact with a normal metal. A spin current \( J_s \parallel \hat{z} \) is generated in the normal metal and absorbed by the ferromagnet.
FIG. 2 (color online). Spin wave band structure and magnetization profiles in YIG for \( d = 0.61 \, \mu m \) without surface anisotropy: \( k_j = 0 \) at \( \theta = \angle (\mathbf{m}, \mathbf{q}) = 90^\circ \). Top (from left to right): \( \text{Re}(\omega/\omega_M) \) vs. \( qd \), \( \text{Im}(\omega/\omega_M) \) at \( k_j = 0 \), \( \text{Im}(\omega/\omega_M) \) at \( k_j = -0.2k_c \). Bottom: \( n_\alpha \) of the same 6 modes for \( qd = 0.09 \) and 3.74 indicated by the dashed vertical lines in the top panels. The colors label different bands. The amplitude of the green mode (shaded [yellow] panel) is amplified.

We now turn on EASA: \( k_s = 25.0/\mu m \) (or \( K_s = 5 \times 10^{-5} \, J/m^2 \)) at the top surface \( (x = 0) \). Fig. 3 shows the results for \( \theta = 90^\circ \). The changes of \( \text{Re}\omega \) and \( \text{Im}\omega \) at \( k_j = 0 \) are modest (Fig. 2), but an additional band (black) appears, viz. the surface spin wave band induced by EASA. The spin-transfer torque strongly affects this mode because of its strong surface localization [25]. As seen in the top right panel, almost the whole band is strongly amplified by a spin current injection of \( k_j = 0.2k_c \). Inspecting the spin wave profiles at two different \( q \) values, we observe a surface spin wave near \( x = 0 \) for the black band at small \( q \) (shaded [yellow] panel in the middle row in Fig. 3). At larger \( q \), the 1st (black) band loses its surface wave features to the 5th (red) band (see top right panel in Fig. 3). The red band mode starts out as a magnetostatic surface spin wave, but the EASA enhances its surface localization by hybridization with the black mode to become strongly amplified by the spin current at higher \( q \). Also in the lower panel of Fig. 3, we observe that the red band has acquired the surface character.

We introduce an approximate power spectrum (Fig. 4) that summarizes all information about the mode-dependent current-induced amplification

\[
P(\omega) = \sum_n \int_{\text{Im}\omega_n < 0} |\text{Im}\omega_n(\mathbf{q})| \delta(\omega - \text{Re}\omega_n(\mathbf{q})) d\mathbf{q}
\]

with \( n \) the band index is the density of states at frequency \( \omega \) weighted by its amplification. \( P(\omega) \) represents the frequency dependence of the excitation power but not its absolute scale. Without surface anisotropy, only a few modes are excited even at a relatively large current \( (k_j = 0.2k_c) \). However, when \( k_s = 25/\mu m \), the excitation is strongly enhanced by more than two orders of magnitude due to the easily excitible surface spin wave modes. Furthermore, we observe broadband excitation over a

FIG. 3 (color online). Same as Fig. 2 but with \( k_s = 25/\mu m \).
much larger range of frequencies. This power spectrum is rather smooth, while the experiments by Kajiwara et al. [5] show a large number of closely spaced peaks. The latter fine structure is caused by size quantization of spin waves due to the finite lateral extension of the sample that has not been taken into account in our theory since it complicates the calculations without introducing new physics. The envelope of the experimental power spectrum compares favorably with the present model calculations.

The insets in Fig. 4 show the integrated power and allow the following conclusions: (i) the excitation power is enhanced by at least two orders of magnitude by the EASA; (ii) the critical current for magnetization dynamics allows the following conclusions: (i) the excitation power favorsably with the present model calculations.

envelope of the experimental power spectrum compares been taken into account in our theory since it complicates fine structure is caused by size quantization of spin waves show a large number of closely spaced peaks. The latter

anisotropy as well. As seen from Fig. 3, the surface anisotropies for capped as compared to free surfaces. A Pt Reference [12] reports an enhancement of the YIG surface would be further reduced (increased) by a larger EASA.

@!=@

mode (black band) has group velocity These values are calculated for a film thickness of 1

Compared to the original estimate $J_c$ for the excited spin waves are localized at the interface.

the threshold current required to excite the spin waves and dramatically increases the excitation power. Multiple spin wave modes can be excited simultaneously at different frequencies and wave vectors, thereby explaining recent experiments. Surface spin wave excitations could be useful in low-power future spintronics-magnonics hybrid circuits.

In conclusion, we predict that an easy-axis surface anisotropy as well as size quantization gives rise to a surface spin wave mode, which reduces the threshold current required to excite the spin waves and dramatically increases the excitation power. Multiple spin wave modes can be excited simultaneously at different frequencies and wave vectors, thereby explaining recent experiments. Surface spin wave excitations could be useful in low-power future spintronics-magnonics hybrid circuits.

This work was supported by the National Natural Science Foundation of China (Grant No. 11004036), the special funds for the Major State Basic Research Project of China (No. 2011CB925601), the FOM foundation, DFG Priority Program SpinCat, and EG-STREP MACALO.

J.X. acknowledges the hospitality of the G.B. Group at the Kavli Institute of NanoScience in Delft.