Current-induced spin torque resonance of magnetic insulators

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We formulate a theory of the AC spin Hall magnetoresistance (SMR) in a bilayer system consisting of a magnetic insulator such as yttrium iron garnet (YIG) and a heavy metal such as platinum (Pt). We derive expressions for the DC voltage generation based on the drift-diffusion spin model and quantum mechanical boundary condition at the interface that reveal a spin torque ferromagnetic resonance (ST-FMR). We predict that ST-FMR experiments will reveal valuable information on the current-induced magnetization dynamics of magnetic insulators and AC spin Hall effect.

Ferrimagnetic insulators such as yttrium iron garnet (YIG) with high critical temperatures and very low magnetization damping have been known for decades to be choice materials for in optical, microwave or data storage technologies $^1$. Near-dissipationless propagation of spin waves make YIG wires and circuits interesting for low power data transmission and logic devices. A crucial breakthrough was the discovery that the magnetization in YIG can be excited electrically by Pt contacts $^2$, thereby creating an interface between electronic/spintronic and magnonic circuits. However, the generation of coherent spin waves by the current-induced spin orbit torques in Pt is a strongly non-linear process and the low critical threshold currents found by the experiments $^2$ cannot yet be explained by theory $^2$. Here we suggest and model a simpler method to get to grips with the important current-magnetization interaction in the YIG/Pt system without a problematic threshold, viz. by employing the recently discovered magneto-resistance of YIG/Pt bilayers or Spin Hall Magnetoresistance (SMR) $^4$ to detect current-induced spin torque ferromagnetic resonance (ST-FMR).

The SMR is the dependence of the electrical resistance of the normal metal on the magnetization angle of a proximity insulator and is caused by a concerted action of the Spin Hall Effect (SHE) $^7$ and its inverse (ISHE). An alternative mechanism of the SMR phenomenology in terms of an equilibrium proximity magnetization close to the YIG interface has been proposed $^5$. However, this interpretation has been challenged by experiments $^4$. Moreover, while experiments of many groups are described quantitatively well by the SMR model with one set of parameters $^8$–$^{12}$, we are not aware of a transport theory that explains the observed magnetoresistance in terms of a monolayer-order magnetic Pt.

Current-induced spin torque ferromagnetic resonance (ST-FMR) has been demonstrated $^{13}$–$^{15}$ in bilayer thin films made from metallic ferromagnets (FM) and non-magnetic metals (N). In these experiments the SHE transforms an in-plane alternating current (AC) into an oscillating transverse spin current. The resultant spin transfer resonates with the magnetization at the FMR frequency. The effects induced simultaneously by the Oersted field can be distinguished by a different symmetry of the resonance on detuning. The magnetization dynamics leads to a time dependence of the bilayer resistance by the anisotropic magnetoresistance (AMR). Mixing the applied current and the oscillating resistance generates a DC voltage that is referred to as spin torque diode effect $^{16}$–$^{17}$.

The longitudinal spin Seebeck effect was found to be frequency independent up to 30 MHz $^{18}$. The DC ISHE induced by spin pumping has been observed by many groups, but detection of the AC spin Hall effect $^{23}$ has only recently been reported in metallic structures $^{19}$–$^{21}$ as well as in Pt/YIG under parametric microwave excitation $^{22}$. A DC voltage can be generated in Pt/YIG under FMR conditions by rectification of the AC spin Hall effect by means of the SMR, but this signal was found to be swamped by the DC spin Hall effect $^{23}$. A study of the spin Hall impedance concludes that the material constants of Pt/YIG bilayers do not depend on frequency up to 4 GHz $^{24}$.

In this paper we suggest to combine the principles sketched above to realize ST-FMR for bilayers of a ferromagnetic insulator (FI) such as YIG and a normal metal with spin orbit interaction (N) such as platinum $^{25}$ (see Fig. 1). We derive the magnetization dynamics and DC voltages generated by the SMR-induced spin torque diode effect as a function of the external magnetic field. Our theory should help to better understand the elusive current-induced magnetization dynamics of ferromagnetic insulators which should pave the way for low-power devices based on magnetic insulators $^1$.

The spin current through an F$|N$ interface is governed by the complex spin-mixing conductance $G_{↑↓}^f$ $^{27}$. The prediction of a large Re $G_{↑↓}^f$ for interfaces between YIG and simple metals by first-principle calculations $^{28}$ has been confirmed by recent experiments $^{24}$. The spin transport in N (spin Hall system) can be treated by spin-diffusion theory with quantum mechanical boundary conditions at the interface to the insulating ferromagnet $^{25}$. The AC current with frequency $\omega_a = 2\pi f_a$ induces a spin accumulation distribution $\mu_s(z,t)$ in N that obeys the spin-diffusion equation

$$\partial_t \mu_s = D \partial_z^2 \mu_s - \frac{\mu_s}{\tau_{sd}},$$ (1)
where $D$ is the charge diffusion constant and $\tau_{sf}$ spin-flip relaxation time in N. In position-frequency space the solution for the spatiotemporal dependence of the spin accumulation reads $\mu_s(z, \omega) = A e^{-\kappa(\omega)z} + B e^{\kappa(\omega)z}$, where $\kappa(\omega) = (1 + i \omega \tau_{sf}) / \chi^2$, $\chi = \sqrt{D / \tau_{sf}}$ is the spin-diffusion length, and the constant column vectors $A$ and $B$ are determined by the boundary conditions for the spin current density in the z-direction $J_{s,z}(z)$, where $J_{s,z} / |J_{s,z}|$ is the spin polarization vector, which is continuous at the interface to the ferromagnet at $z = 0$ and vanishes at the vacuum interface at $z = d_N$. For planar interfaces

$$J_{s,z}(z, \omega) = -\theta_{SH} J_c(\omega) \hat{y} - \sigma_2 \frac{\mu_s(z, \omega)}{2e},$$

where $\theta_{SH}$ is the spin Hall angle, $\sigma$ the electrical conductivity, and $J_c(\omega) = 2\pi J_c^0 \delta(\omega_{a0} - \omega)$ the currents not accounting for spin-orbit interaction. $J_{s,z}(d_N, \omega) = 0$ and $J_{s,z}(0, \omega) = \int_{-\infty}^{\infty} J_{s,z}(0, t) e^{-i \omega t} dt$, where $J_{s,z}(0, t) = J_{s}^T + J_{s}^P = J_{s}^{(F)}$ with

$$J_{s}^T = \frac{G_{s}}{e} \hat{M} \times (\hat{M} \times \mu_s(0)) + \frac{G_{t}}{e} \hat{M} \times \mu_t(0),$$

$$J_{s}^P = \frac{\hbar}{e} \left( G_{s} \hat{M} \times \partial_t \hat{M} + G_{t} \partial_t \hat{M} \right),$$

where $\hat{M}$ is the unit vector along the FI magnetization and $G_{s}^{1+} = G_s + i G_t$ the complex spin-mixing interface conductance per unit area of the FI/N interface. The imaginary part $G_t$ can be interpreted as an effective exchange field acting on the spin accumulation, which is usually much smaller than the real part. A positive $J_{s}^{(P)}$ corresponds here to up-spins flowing from FI into N. For Pt(10 nm)/MgO(2 nm)/Pt(10 nm) $J_{s}^{(P)}(\omega_{a}) = 1.5 \times 10^{-3}$ at the FMR frequency $f_a = 15.5$ GHz with $\tau_{sf}^{(Pt)}(\omega_{a}) = 0.01(1.15)$ ps, indicating that the condition $\omega_a \tau_{sf}^{(Pt)} \ll 1$ is fulfilled for these metals [23]. In this limit the frequency dependence of the spin diffusion length may be disregarded such that

$$\mu_s(z, t) \rightarrow -\hat{y} \mu_{s0}(t) \frac{\sinh \frac{2z - d_N}{2\lambda}}{\sinh \frac{d_N}{2\lambda}} + J_{s}^{(F)} \frac{2e \lambda \cosh \frac{z - d_N}{2\lambda}}{\sinh \frac{d_N}{2\lambda}},$$

where $\mu_{s0}(t) = (2e \lambda / \sigma) \theta_{SH} J_c(t) \tanh [d_N / (2\lambda)]$ with $J_c(t) = J_c^0(\omega) e^{i \omega t}$ and $T = \sigma G^{1+} / [\sigma + 2\lambda G^{1+} \coth (d_N / \lambda)]$. The ISHE drives a charge current in the $x$-$y$ plane by the diffusion spin current along $z$. The total charge current density reads

$$J_{c}(z, t) = J_{c}^{0}(t) \hat{x} + \sigma \theta_{SH} \left( \nabla \times \frac{\mu_s(z, t)}{2e} \right).$$

The averaged current density over the film thickness is

$$J_{c,z}(t) = \frac{d_{N}}{2} \int_{0}^{d_{N}} J_{c,z}(z, t) dz = J_{SMR}(t) + J_{SP}(t)$$

where $J_{SMR}(t)$ and $J_{SP}(t)$ are SRM rectification and spin pumping-induced charge currents, $\rho = \sigma^{-1}$ is the resistivity of the bulk normal metal layer and we recognize the conventional DC SMR with $\Delta \rho_0$ and $\theta_{SH}$ and spin pumping-induced charge currents, $\rho = \sigma^{-1}$ is the resistivity of the bulk normal metal layer and we recognize the conventional DC SMR with $\Delta \rho_0$ and $\theta_{SH}$.

$$J_{c}(t) = J_{c}^{0}(t) \left[ 1 - \frac{\Delta \rho_0}{\rho} - \frac{\Delta \rho_1}{\rho} \right] + \frac{\Delta \rho_0}{\rho} \hat{y} \theta_{SH}(\omega),$$

where $J_{SMR}(t)$ and $J_{SP}(t)$ are SMR rectification and sp}
We henceforth disregard the very low in-plane magnetocrystalline anisotropy field of $H_s \sim 3$ Oe reported [8]. The external magnetic field $H_{ex}$ is applied at a polar angle $\theta$ in the $x$-$y$ plane. It is convenient to consider the magnetization dynamics in the $XYZ$-coordinate system (Fig. 1) in which the magnetization is stabilized along the $X$-axis by a sufficiently strong external magnetic field. Denoting the transformation matrix as $R(\theta)$, the magnetization $\mathbf{M}_R(t) = R(\theta) \mathbf{M}(t)$ precesses across the $X$-axis, where $\mathbf{M}_R(t) = \mathbf{M}_r(t) + \mathbf{m}_r(t) = (M_x, m_y, m_z(t))$ as shown in Fig. 1. $\mathbf{M}_r(t)$ and $\mathbf{m}_r(t)$ are the static and the dynamic components of the magnetization, respectively. The LLG equation in the $XYZ$-system then becomes

$$\dot{\mathbf{M}}_R = -\gamma \mathbf{M}_R \times \mathbf{H}_{eff,R} + \frac{\alpha}{2} \mathbf{M}_R \times \dot{\mathbf{M}}_R$$

where the effective magnetic field in the $XYZ$-system is

$$\mathbf{H}_{eff,R} = \mathbf{H}_x \mathbf{x} + \mathbf{H}_y e^{i\omega_{ac} t} \mathbf{y} + (\mathbf{H}_z e^{i\omega_{ac} t} - 4\pi m_z(t)) \mathbf{z}$$

with $H_x = H_{ex}$, $H_y = (H_{ac} + H_i) \cos \theta$ and $H_z = H_r \cos \theta$ with

$$H_{r(i)} = \frac{\hbar}{2eM_d} \theta_{SH} J_{cR}^0 \Re(\text{Im}(\eta)), \quad (12)$$

a modulated damping $\alpha = \alpha_0 + \Delta \alpha$ and $g$-factor $\beta = 1 - \Delta \beta$ with $\Delta \alpha (\Delta \beta) = \gamma h^2/(2eM_d) \Re T(\text{Im} T)$.

For a small-angle precession around the equilibrium direction $\mathbf{M}_R = (0, \delta m_y e^{i\omega_{ac} t}, \delta m_z e^{i\omega_{ac} t}) (\Re[\delta m_y] \Re[\delta m_z] \ll M_s)$. Disregarding higher orders in $\delta m_y(z)$ in the $R$-transformed LLG equation we arrive at the (Kittel) relation between AC current frequency and resonant magnetic field $H_F = -2\pi M_s + \sqrt{(2\pi M_s)^2 + (\omega_a/\gamma)^2}$.

A DC voltage is generated by two different mechanisms, viz. the time-dependent oscillations of the SMR in $N$ (spin torque diode effect) and the ISHE generated by spin pumping. This is quite analogous to electrically detected FMR in which the magnetization is driven by microwaves in cavities or coplanar wave guides. In metal bilayers, the spin pumping signal due to the ISHE can be separated from effects of the magnetoresistance of the metallic ferromagnet by sample design and angular dependences [32, 33]. Here we focus on the current-induced magnetization dynamics that induces down-converted DC and second harmonic components in the normal metal. Indicating time-average by $\langle \cdots \rangle_t$ the open-circuit DC voltage is $V_{DC} = h p(\mathbf{J}_{ex}(t))_t = V_{SMR} + V_{SP}$, where $V_X = h p(\mathbf{J}_{ex}(t))_t$. The SMR rectification and spin pumping-induced DC voltage are

$$V_{SMR} = -\frac{h \Delta \rho_1 J_{cR}^0 F_S(H_{ex})}{4} \Delta C \left[ \frac{C (H_r + \alpha H_{ac}) + C_+ H_{ac} (H_{ex} - H_F)}{\Delta} \cos \theta \sin 2\theta, \right.$$

$$V_{SP} = \frac{h \rho J_{cR}^0 F_S(H_{ex})}{4} \Delta C \left[ \frac{C_+ H_r^2 + \alpha H_r H_{ac}}{\Delta} + \frac{C_+ H_{ac}^2 - \alpha H_r H_{ac}}{\Delta} \right] \cos \theta \sin 2\theta, \quad (14)$$

where $C = \omega_a / \sqrt{1 + \omega_a^2}$ and $C_+ = 1 + 1 / \sqrt{1 + \omega_a^2}$ with $\omega_a = \omega_{2\pi M_s}$. $F_S(H_{ex}) = \Delta^2 / [\Delta (H_{ex} - H_F)^2 + \Delta^2]$. $\Delta = \omega_a / \gamma$, the line width, $H_{ac} = 2\pi J_{cR}^0 d_N / c$ the Oersted field from the AC current determined by Ampère’s Law (in the limit of an extended film), and $c$ speed of light. Using the material parameters for YIG [8] and Pt [30, 33] shown in Tables I and II we compute the DC voltages in Eqs. (13) and (14). The calculated $V_{SMR}$ is plotted in Fig. 2 as a function of an external magnetic field and for different $d_{Fe}$, resolved in terms of the contributions to the FMR caused by the spin transfer torque (symmetric) and the Oersted magnetic field (asymmetric). In Fig. 3 we show the total DC voltage with both spin torque diode and spin pumping contributions. The DC voltage in F/Pt bilayers depends more sensitively on $d_{Fe}$ for $F = \text{YIG}$ than for $F = \text{Py}/\text{CoFeB}$ because spin pumping is more important when the Gilbert damping is small. ST-FMR measurements are carried out at relatively high current density, so Joule heating in Pt can cause observable effects, the most notable being the spin Seebeck effect (SSE), which adds a constant background DC voltage to the SMR rectification signal [34].

The ST-FMR spectra in Fig. 3 are enhanced for thicker

| \hline
| $T$ & $\text{Am}^{-1}$ & $\alpha$ & \\
| \hline
| 0$\text{YIG}$ & 1.76 & 1.56 & 6.7 \times 10^{-11} \\
| 0$\text{Pt}$ & 1.2 \times 10^{-11} & 1.56 \times 10^{-11} & 6.7 \times 10^{-11} \\
| 0$\text{Reference 2}$ & \\
| \hline
| $G_\alpha$ & $\Omega^{-1}$ & $\rho$ & $\text{nm}$ & $\theta_{SH}$ & \\
| \hline
| 0$\text{Pt}$ & 3.8 \times 10^{-11} & 11 & 5 & 0.12 \\
| 0$\text{Reference 3}$ & \\
| 0$\text{Reference 4}$ & \\
| \hline

TABLE I. Material parameters for the FI layer.

TABLE II. Material parameters for the N layer.
FIG. 2. (a) The ferromagnet thickness dependence of calculated SMR rectified voltage for YIG|Pt at $f_a = 9$ GHz with current density $J^0_c = 10^{10}$ A/m$^2$ and F(N) layer length and width $h = w = 30$ $\mu$m and $\theta = 45^\circ$. (b) $d_{F(N)} = 4(6)$ nm.

FIG. 3. Dependence of the ST-FMR spectra on $d_F$ at $f_a = 9$ GHz and $\theta = 45^\circ$. Inset: Contributions by SMR rectification and spin pumping for $d_F = 4$ nm.

FIG. 4. ST-FMR spectra dependence on $d_F$ in a trilayer setup to observe the spin torque induced DC voltages without artifacts of the Oersted field. ($f_a = 9$ GHz and $\theta = 45^\circ$)

In summary, we predict observable AC current-driven ST-FMR in bilayer systems consisting of a ferromagnetic insulator such as YIG and a normal metal with spin-orbit interaction such as Pt. Our main results are the DC voltages caused by an AC current as a function of in-plane external magnetic field and film thickness of a magnetic insulator. The DC voltages generated in YIG|Pt bilayers depend sensitively on the ferromagnet layer thickness because of the small bulk Gilbert damping. The predictions can be tested experimentally by ST-FMR-like experiments with a magnetic insulator that would yield important insights into the nature of the conduction electron spin-magnon exchange interaction and current-induced spin wave excitations at the interface of metals and magnetic insulators.

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