Magnetic dynamics with spin transfer torques near the Curie temperature.

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In this paper we explore the combined action of large thermal fluctuations and spin transfer torques on the behavior of magnetic layers in spin valves. We find that at temperatures near T_c, spin currents can measurably change the size of the magnetization (i.e. there is a *longitudinal* spin transfer effect), and we propose an experimental signature for this effect. We also find the temperature dependence of the applied-field applied-current phase diagram for a free magnetic layer, valid for temperatures near the Curie temperature T_c. In this study we employ both an atomistic stochastic Landau-Lifshitz-Slonczewski simulation at high temperatures, and also find the thermally averaged magnetic dynamics within mean field theory, yielding a Landau-Lifshitz-Bloch + Slonczewski equation. We show that this simplified equation describes the full stochastic model reasonably well. We also show how the Landau-Lifshitz-Gilbert equation describing a fixed magnetization size can be modified to describe the magnetic dynamics in this temperature regime.

I. INTRODUCTION

Spin transfer torque describes the interaction between the spin of itinerant, current-carrying electrons and the spins of the equilibrium electrons which comprise the magnetization of a ferromagnet. This torque results from the spin-dependent exchange-correlation electron-electron interaction, and leads to the mutual precession of equilibrium and non-equilibrium spin around the total spin. In spin valves with sufficiently high current density, spin transfer torque can excite a free ferromagnetic layer to irreversibly switch between two stable configurations (typically along an easy-axis, parallel or anti-parallel to an applied magnetic field), or to undergo microwave oscillations. Previous considerations of spin transfer torque mostly focus on the *transverse* response of the magnetization to spin currents. This is appropriate since the temperatures used in spin valve experiments are substantially below the Curie temperature T_c of the ferromagnets, so that longitudinal fluctuations can be ignored.

Even far from the Curie temperature, temperature plays an important role in quantitatively analyzing the dependence of the magnetic orientation on the applied field and applied current. Previous works on the effect of finite temperature on spin dynamics in the presence of spin transfer torque utilize the macrospin approximation, which assumes that the size of the magnetization is fixed. Refs. 1 and 2 add a Sloncewski torque to the Langevin equation describing the stochastic spin dynamics, while Ref. 3 solves the Focker-Planck equation for a single spin, with the spin transfer torque term added to the deterministic dynamics. Ref. 4 uses the Keldysh formalism to formally derive the stochastic equation of motion for the non-equilibrium (i.e., current-carrying) system, and also considers a single spin of fixed magnitude. These treatments have been successful in describing thermal characteristic of nanomagnets under the action of spin torques, such as dwell times and other details of thermally activated switching.

For materials like GaMnAs, experiments are done near T_c, so that the *size* of the magnetization is substantially reduced from its 0 temperature value, and undergoes sizeable fluctuations. In this case, the applicability of a macrospin model is not clear. For field-driven dymamics, there is theoretical work which accounts for longitudinal fluctuations near T_c [5]. This formal treatment culminates in the construction of the Landau-Lifshitz-Bloch equation (LLB), which is an extension of the familiar Landau-Lifshitz equation with an addition longitudinal degree of freedom. In this work, we consider temperatures near the Curie temperature and include both longitudinal fluctuations and the influence of spin transfer torque.

There are a number of open questions regarding magnetic dynamics near T_c, including the temperature dependence of the spin transfer torque itself, and the temperature dependence of more basic magnetic properties such as magnetic damping, magneto-crystalline anisotropy, etc. We want to consider only the influence of the spin torque (understood in its 0 temperature sense) on the thermally reduced magnetization. To do so, we use an atomistic approach for the stochastic dynamics of a local moment ferromagnet with the inclusion of spin transfer torque. Rather than attempt to deal with all of the high-temperature behaviors of magnetic properties, we make simple assumptions about the temperature dependence of the magnetic anisotropy, demagnetization field, and damping. Using this model we show that spin currents can change the *size* of the magnetization, and give an expression for this "spin-current longitudinal susceptibility", and propose an experimental scheme to measure this effect.

We then construct a Landau-Lifshitz-Bloch + Slonczewski (LLBS) equation to describe both longitudinal fluctuations and spin transfer torques. Following Ref. 6, we verify the applicability of the LLBS equation by comparing its results to the atomistic results. We then analyze the LLBS equation to find the applied field-applied current phase diagram for different temperatures. We find that critical switching currents are reduced by the same mechanism exploited in heat assisted magnetic recording, namely the temperature-induced reduction in the magnetic anisotropy [7]. We also find that regions of the phase diagram which have been experimentally unattainable become relevant at high temperatures. The dependence of critical currents versus temperature in these regions can provide quantitative details about the temperature dependence of spin transfer torque.

II. METHOD

To study the interplay between temperature and spin transfer torque, we consider a spin valve with a fixed layer magnetization in the $+\hat{z}$-direction with Curie temperature T_c^1, and a free layer with a smaller Curie temperature T_c^2 (see Fig. (1)). This allows for a nearly temperature independent spin current flux incident on the free layer. The details of how spin current propagates through a ferromagnet which is undergoing large thermal fluctuations is an interesting question in its own right, and is the subject of future work. For now we make the simplest assumption, namely that all of the incoming spin current is absorbed uniformly throughout the free layer magnetization. This is certainly an over-estimation, so that our results represent an upper bound on the temperature-dependent effects. But it is likely not a gross over-estimation: substantial spatial and temporal inhomogeneities in the magnetization should induce rather irregular spatial patterns in the spin currents carried by propagating states. This will lead to large dephasing effects, so that the total spin current should rapidly decay away from the interface as in the conventional picture of spin transfer torques. We also suppose that the spin transfer torque is uniformly distributed throughout the layer. It is known that the torque is generally localized near the interface [8]. However in this temperature regime, and for thin layers ($\sim$ 3 nm), magnetic non-uniformities in the direction transverse to current flow should be more substantial than non-uniformities *along* the current flow resulting from a localized spin transfer torque.

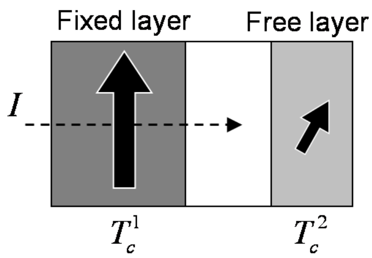

FIG. 1: Schematic of system, two ferromagnetic layers with different Curie temperatures. We suppose that $T_c^1 > T_c^2$.

A. Stochastic LL with spin transfer

We adopt three approaches to model the system. The first is an atomistic lattice model of normalized spins $\mathbf{S}$. We include nearest-neighbor Heisenberg coupling with exchange constant J, and an easy-axis anisotropy field of magnitude $H_{\rm an}$ in the $\hat{z}$-direction. To model the temperature dependence of the anisotropy, we make the simple ansatz that the magnitude of anisotropy is proportional to the reduced magnetization $m(T) = M_s(T)/M_s(T = 0)$:

$$H_{\rm an}(T) = H_{\rm an}(T = 0)m(T), \tag{1}$$

so that the anisotropy field on spin i is given by $H^i_{\rm an}(T) = H_{\rm an}(T = 0)|\langle\mathbf{S}\rangle|S_i^z$, where the brackets here indicate a spatial average. A hard-axis anisotropy field with magnitude $H_{\rm d}$ in the $\hat{y}$-direction is added to model the demagnetization field of the thin layer. We again make a simplified ansatz for the form of this field to avoid evaluating a non-local field, making the numerics more tractable. We take the demagnetization field to be uniform on all spins and given by $H^i_{\rm d}(T) = -H_{\rm d}(T = 0)\langle S^y\rangle\hat{y}$. This form of the hard-axis field ensures that $H_{\rm d} \propto M_s(T)$, and roughly captures the nonlocal nature of the field. Finally, we include an applied field $H_{\rm app}$ in the $\hat{z}$-direction. The Hamiltonian for spin i is then:

$$\begin{aligned} H_i \;=\; & J\sum_{j\in \rm n.n.} \mathbf{S}_i\cdot\mathbf{S}_j + \frac{H_{\rm an}(T=0)|\langle\mathbf{S}\rangle|}{2}\left(S_i^z\right)^2 \\ & - H_{\rm d}(T=0)S_i^y\langle S^y\rangle + H_{\rm app}S_i^z. \end{aligned} \tag{2}$$

To model nonzero temperatures, we add damping α and a stochastic field $\mathbf{H}_{\rm fl}$ to the equation of motion implied by Eq. (2), with the standard statistical properties:

$$\langle H_{\rm fl}(t)\rangle \;=\; 0, \tag{3}$$

$$\langle H^i_{\rm fl}(t)H^j_{\rm fl}(t')\rangle \;=\; \frac{\alpha}{1+\alpha^2}\frac{2k_BT}{\gamma z}\delta_{ij}\delta(t-t'). \tag{4}$$

where z is the spin on each lattice site (in units of μ_B, z is typically of order 1). We numerically integrate the equation of motion using a second-order Heun scheme, as described in Ref. 9. We add a Slonczewski-like spin transfer torque term to the equation of motion for the ith spin, which is given finally as:

$$\begin{aligned} \dot{\mathbf{S}}_i \;=\; & -\gamma\mathbf{S}_i\times(\mathbf{H}_{\rm eff}+\mathbf{H}_{\rm fl}) - \gamma\alpha\left(\mathbf{S}_i\times\mathbf{S}_i\times\mathbf{H}_{\rm eff}\right) \\ & + H_{\rm st}\left(\mathbf{S}_i\times\mathbf{S}_i\times\hat{z}\right). \end{aligned} \tag{5}$$

where $\mathbf{H}_{\rm eff} = H_{\rm app}\hat{z} + H_{\rm an}|\langle\mathbf{S}\rangle|S_i^z\hat{z} - H_{\rm d}\langle S^y\rangle\hat{y}$, γ is the gyromagnetic ratio, and we characterize the spin transfer torque with $H_{\rm st} = pJ^{\rm e}a^3/ez\ell\gamma$. Here $J^{\rm e}$ is the charge current density, p is the spin polarization of the current, a is the lattice constant, ℓ is the free layer thickness, and e is the electron charge. We use both a bulk geometry consisting of a $N = 48^3$ periodic array of spins in 3 dimensions, and a layer geometry with an array of 100 $\times$

100 × 15 spins. We employ the bulk geometry in comparing the stochastic model behavior with predictions from mean field theory, and the layer geometry for studying the effect of spin current on magnetization size.

B. Landau-Lifshitz Bloch + Slonczewski equation

In the second approach, we add a Slonczewski torque term to the LLB equation. To derive the LLB equation, a probability distribution for the spin orientation is assumed, which is used to find the ensemble average of Eq. (5). In addition the nearest neighbor exchange field is replaced by its mean-field value. The details of the derivation follow closely those in Ref. (5), so we omit them here. The final LLB+Slonczewski equation takes the form:

$$\begin{aligned}\dot{\mathbf{m}} &= -\gamma\left(\mathbf{m}\times\mathbf{H}_{\text{eff}}\right)+\frac{2\gamma\alpha T}{J_0 m^2}\mathbf{m}\cdot\left(\mathbf{H}_{\text{eff}}+\frac{H_{\text{st}}}{\alpha}\hat{z}\right)\mathbf{m}\\ &\quad -\frac{\gamma\alpha}{m^2}\left(1-\frac{T}{J_0}\right)\mathbf{m}\times\mathbf{m}\times\left(\mathbf{H}_{\text{eff}}+\frac{H_{\text{st}}}{\alpha}\hat{z}\right),\end{aligned} \quad (6)$$

with an effective field given by:

$$\begin{aligned}\mathbf{H}_{\text{eff}} &= H_{\text{app}}\hat{z}+H_{\text{an}}m^2 m_z\hat{z}-H_d m_y\hat{y}\\ &\quad -\frac{1}{2\chi}\left(\frac{m^2}{m_e^2}-1\right)\mathbf{m},\end{aligned} \quad (7)$$

where $m_e(T)$ is the zero field, zero current equilibrium magnetization: $m_e(T) = B(J_0/k_B T)$, and B is the Brouillon function. $\chi(T)$ is the longitudinal susceptibility: $\chi(T) = \partial m(T)/\partial H_{\text{app}}$. J_0 is the 0th component of the Fourier transformed exchange, and $\mathbf{m}$ is a vector with size between 0 and 1. The double cross product in Eq. (7) is the familiar Landau Lifshitz damping term, which describes the relaxation of the magnetization *direction* to the nearest energy minimum. The term longitudinal to $\mathbf{m}$ distinguishes the LLB equation from the LL equation. This longitudinal term describes the relaxation of the *size* of the magnetization to its steady state value, which is determined by the temperature, applied fields, and applied currents.

The detailed dependence of the magnetic anisotropy on temperature generally material specific. In our model, the anisotropy and demagnetization fields depend on temperature through their m dependence, and vary as $m(T)$ and $m(T)^3$, respectively. (Note the m^3 dependence of the anisotropy depends partially on our choice of uniaxial anisotropy $E_{\text{an}} = Km_z^2$. A similar choice of $E_{\text{an}} = -K(m_x^2 + m_y^2)$ would yield an anisotropy field dependence of m^2, see Ref. 10 for a further discussion of this point). The magnetic exchange J_0 can also depend on temperature. This dependence is stronger for ferromagnets with indirect exchange interactions (such as GaMnAs - where the magnetic interactions are mediated by hole carriers), and weaker for local moment systems with direct exchange (such as Fe). For simplicity we treat J_0 as temperature-independent.

Finally we consider the standard Landau-Lifshitz equation with a reduced saturation magnetization. We find in Sec. (III D) that an appropriately modified damping coefficient leads to qualitative agreement the more complicated models and the Landau-Lifshitz equation.

III. RESULTS

A. Longitudinal spin current susceptibility

We first consider the longitudinal spin transfer effect. Using the LLB+Slonczewski equation, it is straightforward to show that the size of the magnetization changes in the presence of spin current as:

$$\delta m(I,T) = H_{\text{st}}\frac{\chi(T)}{\alpha} \quad (8)$$

To verify the applicability of this expression, we considered the full stochastic simulation with 100×100×15 spins (initialized to thermal equilibrium) and find the response of the magnetization to applied fields and currents. Fig. (2) shows the longitudinal susceptibility to magnetic field and spin current, where the spin current susceptibility χ_I is defined as $\chi_I = \delta m/H_{\text{st}}$. χ is plotted in dimensionless form, and is scaled by the exchange field J/μ_B. We find that χ_H and $\chi_I\alpha$ correspond very well, demonstrating that Eq. (8) accurately describes the numerical stochastic model.

To get a feel for the magnitude of the longitudinal spin transfer, we rewrite the change in m in dimension-ful terms:

$$\delta m = \frac{pJ^e a^3}{ez\ell\gamma}\left(\frac{\mu_B}{J_0\alpha}\right)\chi(T) \quad (9)$$

Taking the exchange field J_0/μ_B =150 T (which corresponds to a T_c of 150 K in a cubic nearest neighbor Heisenberg model), a 0 temperature saturation magnetization of 10^6 A/m, $\chi = 7$, $J^e = 10^{11}$A/m^2, $p = 0.5$, $\alpha = 0.01$, $\ell = 3$nm results in a 5% change in m from its 0 temperature value, which should be measurable. (Using Fig. (2), we find that $\chi = 7$ corresponds to $T = 0.95\ T_c$.)

A notable aspect of this longitudinal spin transfer is that the size of the magnetization can either be increased or decreased according to the direction of current flow. For electron flow from fixed to free layer, the free layer moment *increases*, while for electron flow in the opposite direction *decreases* the free layer moment. This contrasts with current-induced Joule heating, which always decreases the magnetization.

This distinction can be exploited to probe the longitudinal spin transfer by using the experimental scheme showed in Fig. (3). Here we suppose that $T_c^1 \gg T > T_c^2$. We choose sign conventions such that a positive H_{app} aligns with the fixed layer, and a positive current represents electron flow from fixed to free layer. In the absence of a longitudinal spin transfer (χ_I=0, black line in Fig.

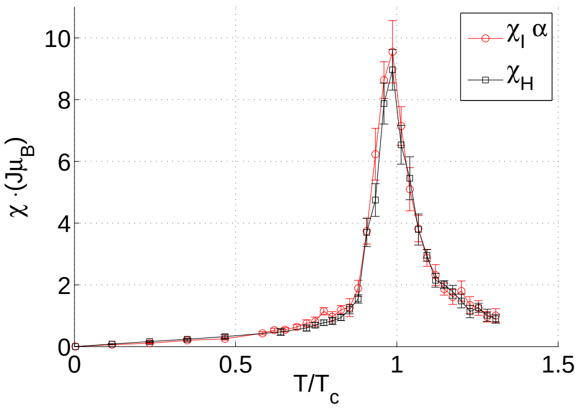

FIG. 2: The magnetic field and spin current susceptibility versus temperature for the stochastic landau lifshitz equation in the layer geometry. The spin current susceptibility is multiplied by α. The error bars indicate statistical uncertainty.

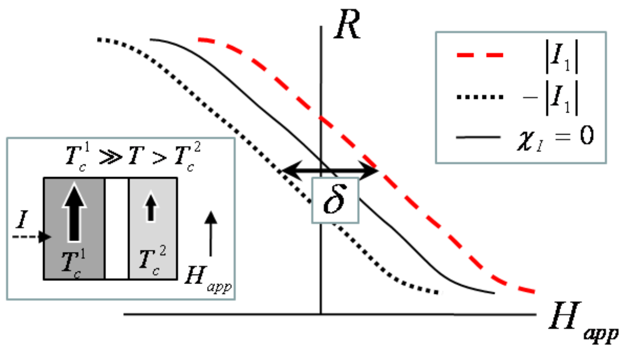

FIG. 3: The magnetic field and spin current susceptibility versus temperature for the stochastic landau lifshitz equation in the layer geometry. The spin current susceptibility is multiplied by α.

(3)), the application of a magnetic field will partially order the free layer to align or anti-align with the fixed layer. This should cause the resistance R of the device to change in some way, according to the GMR effect and magnetic order induced in the free layer (the detailed dependence of R on H_{app} is not important here). If a positive current I_1 is applied, then the longitudinal spin transfer induces partial ordering of the free layer, so that $m\,(H_{\mathrm{app}}=0)=+\chi_I I_1$. Then the curve of $m\,(H_{\mathrm{app}})$, and therefore the curve $R\,(H_{\mathrm{app}})$ is simply shifted by $+\chi_I I_1$ (the red dashed curve in Fig. (3)). If a negative current $-|I_1|$ is applied, then $m\,(H_{\mathrm{app}}=0)=-\chi_I I_1$ and the $m\,(H_{\mathrm{app}})$ and $R\,(H_{\mathrm{app}})$ curves are shifted by $-\chi_I I_1$ (black dotted curve in Fig. (3)). This shift represents a unique signature of longitudinal spin transfer.

Plugging in numbers as above, we estimate a total shift $\delta = 2\chi_I I_1$ between $R\,(H_{\mathrm{app}})$ for positive and negative current to be on the order of 1 T. Eq. (9) indicates that materials with small exchange field (or small T_c), and those that can support large current densities show the effect most strongly. This suggests that weak metallic ferromagnets such as Gd($T_c = 300K$), and Fe alloys such as FeS_2 and $FeBe_5$($T_c = 270K$) [11] may be good candidates for free layer material.

B. LLBS vs SLL

In this section we compare the results obtained from the full 3-dimensional stochastic LL+S equation with those obtained from the mean-field LLBS equation. In our numerics, we use dimensionless time $\tau = (\gamma J/\mu_B)t$, which rescales the magnetic fields H_{eff} by the exchange field $H_{\mathrm{ex}} = J/\mu_B$. Dimensionless fields are denoted by lowercase: $h_{\mathrm{app}} = H_{\mathrm{app}}\mu_B/J$, etc. We consider a current-induced magnetic excitation for the bulk lattice geometry at various temperatures. The magnetization is initialized at a 45° angle with respect to $+\hat{z}$-direction, and spin transfer torque is applied to excite the magnetization away from the $\hat{z}$-direction. The parameters used are an applied field of $h_{\mathrm{app}} = 0.0001$, a demagnetization field of $h_{\mathrm{d}} = 0.01$, a spin torque of $h_{\mathrm{st}} = -0.0002$, and damping of $\alpha = 0.1$ (the artificially high damping was chosen to allow the numerical simulation of the switching to be carried out in a reasonable time frame). As we vary temperature, we obtain trajectories of varying complexity. Fig. (4) compares the LLBS and several realizations of the stochastic LL equation. For this range of parameters, the magnetic dynamics evolves from steady oscillations to current induced switching as the temperature is increased. The trajectories for $t = 0.08$ indicate that a realization of stochastic dynamics can exhibit the crossover from precession to stable switching, whereas at this temperature the thermally averaged trajectory shows only oscillations. Generally, the level of correspondence between the two is qualitatively good, although not surprisingly varies. It would be necessary to model the switching event many times to obtain an average trajectory to compare to the LLBS, but computational constraints make this unfeasible. We can nevertheless conclude from this data that the LLBS equation qualitatively captures the features of the full stochastic simulations.

C. Applied field - applied current phase diagram

We now investigate the effect of the longitudinal degree of freedom on the applied field - applied current phase diagram of the free magnetic layer. Fig. (5) shows the generic topology for regions of stability for the parallel ("P", or $+\hat{z}$-direction) and antiparallel ("AP", or $-\hat{z}$-direction) fixed points. We focus on the stability

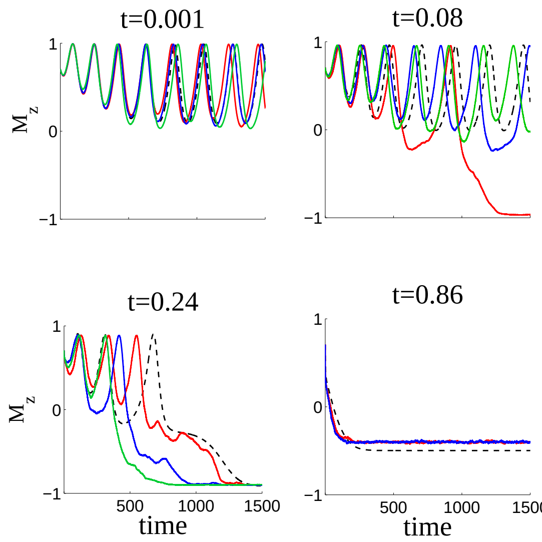

FIG. 4: Comparison of trajectories of current-induced excitations for the atomistic stochastic simulation and the LLB+Slonczewski equation for various reduced temperatures $t = T/T_c$. The dashed line gives the LLB+Slonczweski trajectory, while the solid lines show various realizations of the stochastic trajectory.

of the AP configuration for positive applied fields (the dashed boundary in the upper-half-plane of Fig. (5). We first briefly describe the main qualitative features before providing a mathematical description. For applied fields between $h_{\mathrm{an}}m^3$ and $h_{\mathrm{an}}m^3 + h_{\mathrm{d}}m$, the stability boundary is a horizontal parabola, while for other values of applied field, the stability boundary is linear with slope $1/\alpha$. For applied fields with magnitude less than $h_{\mathrm{an}}m$, there is hysteresis in the current switching. For T=0, this phase diagram reduces to the known form found experimentally [12]. As T increases, the size of the hysteretic region (and the switching current) decreases. Also the range of field with the parabolic boundary decreases, and the outer edge of the parabola gets pulled in closer to 0. For sufficiently high temperatures, this parabolic stability boundary should be experimentally accessible.

We now analyze these behaviors quantitatively. We determine the stability of fixed points using the standard method of linearizing Eq. (7) about a fixed point and finding parameter-dependent eigenvalues λ. A positive real part of λ indicates a loss of stability. This analysis leads to the following condition for instability of the antiparallel configuration (where it should be noted that m depends on h_{st} through $m = m_e + \chi\left(h + \frac{h_{\mathrm{st}}}{\alpha}\right)$):

$$\mathrm{Re}\left[h_{\mathrm{st}}^{\mathrm{crit}} + \alpha\left(h + h_{\mathrm{an}}m^3 + \frac{h_{\mathrm{d}}}{2}m\frac{1-T}{1-3T} - \frac{m}{2\chi}\left(1 - \frac{m^2}{m_e^2}\right)\frac{2T}{1-3T}\right) - \frac{m\sqrt{-\left(h + h_{\mathrm{an}}m^3\right)\left(h + h_{\mathrm{an}}m^3 + h_{\mathrm{d}}m\right)}}{1-3T}\right] = 0.$$

This leads to a cubic equation for $h_{\mathrm{st}}^{\mathrm{crit}}$. Assuming $m_e \gg \chi\left(h + \frac{h_{\mathrm{st}}}{\alpha}\right)$, and expanding to 0th order in χ leads to an approximate, closed form for $h_{\mathrm{st}}^{\mathrm{crit}}$. Again we distinguish between different regimes of applied field. For $h \notin [h_{\mathrm{an}}m^3, h_{\mathrm{an}}m^3 + h_{\mathrm{d}}m]$

$$h_{\mathrm{st}}^{\mathrm{crit}} = \alpha\left(h + \frac{h_{\mathrm{d}}}{2}m_e + h_{\mathrm{an}}m_e^3\frac{1-3T}{1-T}\right), \quad (10)$$

where again m_e is the equilibrium magnetization in the absence of applied field and applied current. Eq. (10) shows that the slope of the boundary is temperature independent, and is given by $1/\alpha$ (the intrinsic damping α is assumed to be temperature independent). The temperature independence of the slope follows from the fact that the spin transfer torque increases like $1/m(T)$, but the effective damping rate damping increases as $1/m(T)$. The intercepts of this boundary line are temperature dependent due to the temperature dependence of m. The contribution from the easy-axis anisotropy field has an additional temperature dependence, but the magnitude of this field is much smaller than the demagnetization field, so it does not play an important role. The critical current at zero field is reduced by $m(T)$ because of the reduction in the demagnetization field. This is important because the demagnetization field is usually larger than applied fields, and is therefore the primary impediment to current induced switching. Its reduction through increased temperature offers a route to reduced critical switching currents.

For $h \in [h_{\mathrm{an}}m^3, h_{\mathrm{an}}m^3 + h_{\mathrm{d}}m]$, a very large spin torque is required to stabilize the AP configuration. The values of current for which the AP configuration is stabilized are much higher than those attainable experimentally, so that for this range of fields the AP configuration is not seen [13]. The approximate critical current along the AP stability boundary is:

$$h_{\mathrm{st}}^{\mathrm{crit}} = \frac{m_e\sqrt{h(h_{\mathrm{d}}m_e - h)}}{1-T}. \quad (11)$$

The reduction in the outer boundary of the parabolic stability line is reduced at high temperature, and this reduction can also be traced back to the reduced magnetic anisotropy. For low temperatures, the application of spin transfer torques results in a elliptical precession mostly in the easy plane away about the $-\hat{z}$ fixed point. To stabilize the AP configuration in this regime, the spin transfer torque must overcome the *precessional* torque (usually, the spin transfer torque must overcome the much weaker *damping* torque). Assuming $h = h_{\mathrm{d}} m/2$ for definiteness, the precessional torque decreases with T as $h_{\mathrm{d}} m(T)$, while the STT increases like $1/m$. This implies a value for the maximum reach of the parabola of $I = m^2(T) h_{\mathrm{d}}/(2(1-T))$. Plugging in typical values for material parameters (the same used in Sec. (III A)) leads to a critical current of $10^{12}\mathrm{A/m}^2$ for $T = 0.95 T_c$. This is an order of magnitude smaller than the 0 temperature case. The behavior of this critical current versus temperature at a fixed applied field is shown in Fig. (6). (Solid line gives LLBS result). Note that the curve does not extend to $T = T_c$, because for a finite applied field, we necessarily move outside the region $h \in [h_{\mathrm{an}} m^3, h_{\mathrm{an}} m^3 + h_{\mathrm{d}} m]$ as $T \to T_c$. It should also be noted that the stochastic trajectories (shown in Fig. (4)) indicate that thermal fluctuations can effectively drive the system out of the precessional state and into the static antiparallel configuration.

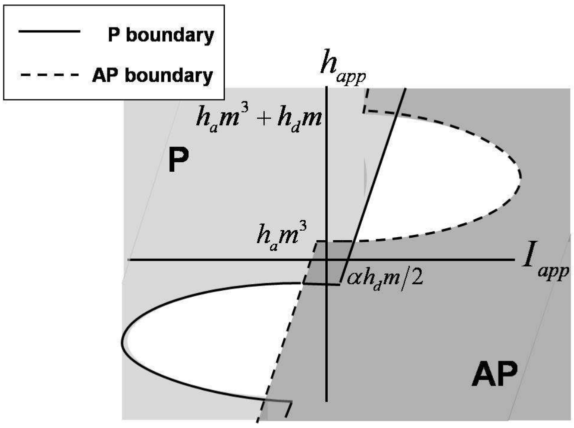

FIG. 5: Schematic of parallel/anti-parallel stability versus applied field and applied current. The hysteretic box near the origin and the fully unstable regions (white parabolic shapes) contract in size with increasing temperature.

D. Comparison with Landau Lifshitz

We finally turn to the differences between the LLBS and the Landau-Lifshitz-Slonczewski equation with reduced magnetization. Based on the qualitative behavior of the LLBS equation, a suitable form for a temperature dependent LLS equation for a nanomagnet of reduced size m and orientation $\hat{n}$ is:

$$\dot{\hat{n}} = -\hat{n} \times \mathbf{H}_{\mathrm{eff}} - \frac{\alpha}{m} \hat{n} \times \hat{n} \times \mathbf{H}_{\mathrm{eff}} - \frac{H_{\mathrm{st}}}{m} \hat{n} \times \hat{n} \times \hat{z} \quad (12)$$

where $\mathbf{H}_{\mathrm{eff}} = \mathbf{H}_{\mathrm{app}} - m H_{\mathrm{d}} n_y \hat{y} + m^3 H_{\mathrm{an}} n_z \hat{z}$, and the temperature dependence is contained entirely in $m(T)$. Clearly the divergence of the damping at $T = T_c$ is unphysical, however a more detailed treatment of damping near T_c is beyond the scope of this paper. The differences between this LL equation and the LLBS equation are quantitative (as opposed to qualitative) in nature. One difference is in the dependence of the critical current on temperature for $h \in [h_{\mathrm{an}} m^3, h_{\mathrm{an}} m^3 + h_{\mathrm{d}} m]$. Fig. (6) shows the prediction based on the LL equation.

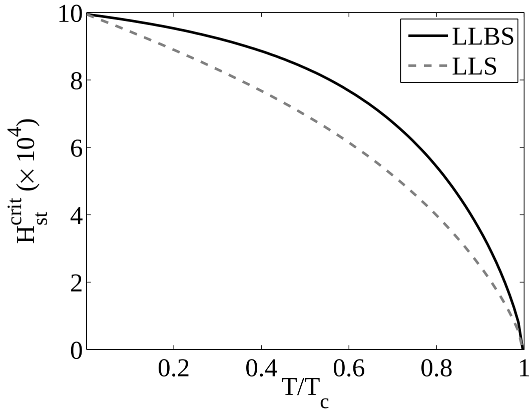

FIG. 6: Critical current versus temperature for LLBS and LL equations. The parameters are: $h_{\mathrm{app}} = -0.001$, $h_{\mathrm{d}} = 0.01, h_{\mathrm{an}} = 0.0001$. Recall that all fields are scaled by the exchange field.

Fig. (6) is indicative of the fact that for the applied-field applied current phase diagram, the spin-current longitudinal susceptibility plays a role that is secondary to the more pronounced effects of temperature reduced anisotropies.

IV. DISCUSSION

In this paper explored the consequences of the combined action of longitudinal fluctuations and spin transfer torques on the behavior of magnetic layers in spin valves. To do so we studied an atomistic, stochastic Landau-Lifshitz-Slonczewski simulation at large temperatures. We find that there is a longitudinal spin transfer effect, and estimate that at temperatures near T_c, spin currents can measurably change the size of the magnetization. We then supplemented the Landau Lifshitz Bloch equation with a Slonczewski torque term, and verified that this model captures the qualitative features of the stochastic simulations. We showed that the applied field-applied current phase diagram undergoes large changes in the presence of high temperatures, and that these changes may be useful for reducing critical switching currents and for studying the detailed behavior of the temperature dependence of the spin transfer torque.

The experimental system relevant for the effects we describe (shown schematically in Fig. (1)) should be relatively straightforward to fabricate. Jiang *et al.* considered a similar system [14], although this particular work dealt with other issues such as the ferrimagnet compensation point for magnetization and total angular momentum. By considering simpler ferromagnets with different Curie temperatures, the role of temperature may be more easily inferred. It is of course necessary to account for Joule heating in assessing the detailed temperature dependence of the spin transfer torque. However recent experimental on domain wall motion illustrates the feasibility of compensating for this effect [15]. On the other hand, experiments conducted at fixed current with varying ambient temperatures and applied fields may offer a more straightforward route to observing the longitudinal spin transfer effect.

Many experiments done with dilute magnetic semiconductors deal with domain wall motion, where thermal effects play an important role in even the qualitative aspects of the domain wall behavior[15]. There are additional challenges associated with extending this work from spin valves to continuous magnetic textures. Among these is the renormalization of the exchange interaction associated with the coarse graining of the magnetization, which becomes more important at higher temperatures [16]. In addition, the crucial role played by the demagnetization field in intrinsic domain wall pinning implies that the finite temperature treatment of the demagnetization field must also be handled more carefully. For these reasons the spin valve geometry may provide greater experimental control and admit a simpler theoretical description.

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