Enhanced Gilbert Damping in Thin Ferromagnetic Films

Yaroslav Tserkovnyak and Arne Brataas

Lyman Laboratory of Physics, Harvard University, Cambridge, Massachusetts 02138

Gerrit E. W. Bauer

Department of Applied Physics and DIMES, Delft University of Technology, 2628 CJ Delft, The Netherlands
(Received 11 October 2001; published 28 February 2002)

The precession of the magnetization of a ferromagnet is shown to transfer spins into adjacent normal metal layers. This "pumping" of spins slows down the precession corresponding to an enhanced Gilbert damping constant in the Landau-Lifshitz equation. The damping is expressed in terms of the scattering matrix of the ferromagnetic layer, which is accessible to model and first-principles calculations. Our estimates for permalloy thin films explain the trends observed in recent experiments.

The magnetization dynamics of a bulk ferromagnet is well described by the phenomenological Landau-Lifshitz-Gilbert (LLG) equation [1]

$$\frac{d\mathbf{m}}{dt} = -\gamma \mathbf{m} \times \mathbf{H}_{\text{eff}} + \alpha \mathbf{m} \times \frac{d\mathbf{m}}{dt}, \qquad (1)$$

where $\mathbf{m}$ is the magnetization direction, γ is the gyromagnetic ratio, and $\mathbf{H}_{\text{eff}}$ is the effective magnetic field including the external, demagnetization, and crystal anisotropy fields. The second term on the right-hand side of Eq. (1) was first introduced by Gilbert [1] and the dimensionless coefficient α is called the Gilbert damping constant. For a constant $\mathbf{H}_{\text{eff}}$ and $\alpha = 0$, $\mathbf{m}$ precesses around the field vector with frequency $\omega = \gamma H_{\text{eff}}$. When damping is switched on $\alpha > 0$, the precession spirals down to a time independent magnetization along the field direction on a time scale of $1/\alpha\omega$. The study of α in bulk metallic ferromagnets has drawn significant interest over several decades. Notwithstanding the large body of both experimental [2] and theoretical [3] work, the damping mechanism in bulk ferromagnets is not yet fully understood.

The magnetization dynamics in thin magnetic films and microstructures is technologically relevant for, e.g., magnetic recording applications at high bit densities. Recent interest by the basic physics community in this topic is motivated by the spin-current induced magnetization switching in layered structures [4–6]. The Gilbert damping constant was found to be $0.04 < \alpha < 0.22$ for Cu-Co and Pt-Co [5,7], which is considerably larger than the bulk value $\alpha_0 \approx 0.005$ in Co [6,8]. Previous attempts to explain the additional damping in magnetic multilayer systems involved an enhanced electron-magnon scattering near the interface [9] and other mechanisms [10], both in equilibrium and in the presence of a spin-polarized current.

In this Letter we propose a novel mechanism for the Gilbert damping in normal-metal-ferromagnet (N-F) hybrids. According to Eq. (1), the precession of the magnetization direction $\mathbf{m}$ is caused by the torque $\propto \mathbf{m} \times \mathbf{H}_{\text{eff}}$. This is physically equivalent to a volume injection of what we call a "spin current." The damping occurs when the spin current is allowed to leak into a normal metal in contact with the ferromagnet. Our mechanism is thus the inverse of the spin-current induced magnetization switching: A spin current can exert a finite torque on the ferromagnetic order parameter, and, *vice versa,* a moving magnetization vector loses torque by emitting a spin current. In other words, the magnetization precession acts as a spin pump which transfers angular momentum from the ferromagnet into the normal metal. This effect can be mathematically formulated in terms of the dependence of the scattering matrix of a ferromagnetic layer attached to normal metal leads on the precession of $\mathbf{m}$, analogous to the parametric charge pumping in nonmagnetic systems [11]. The damping contribution is found to obey the LLG phenomenology. Enhancement of the damping constant $\alpha' = \alpha - \alpha_0$ can be expressed in terms of the scattering matrix at the Fermi energy of a ferromagnetic film in contact with normal metal reservoirs, which can be readily obtained by model or first-principles calculations. Our numerical estimates of α' compare well with recent experimental results [12]. Earlier experiments reported in Ref. [13] can also be understood by our model [14].

We consider a ferromagnetic film sandwiched between two paramagnetic layers as shown in Fig. 1. Spin pumping is governed by the ferromagnetic film and the vicinity of the N-F interfaces. The normal metal layers are, therefore, interpreted as reservoirs attached to nonmagnetic leads. The quantity of interest is the 2×2 current matrix in spin space $\hat{I} = \hat{1}I_c/2 - \hat{\boldsymbol{\sigma}} \cdot \mathbf{I}_s e/\hbar$ for the charge (I_c) and spin flow ($\mathbf{I}_s$) from the magnetic film into adjacent normal metal leads, where $\hat{1}$ is the unit matrix and $\hat{\boldsymbol{\sigma}}$ is the vector of Pauli spin matrices.

When no voltages are applied and the external field is constant, the charge current vanishes. Two contributions to the spin current $\mathbf{I}_s$ on either side of the ferromagnet may be distinguished, viz., $\mathbf{I}_s^{\text{pump}}$ and $\mathbf{I}_s^{(0)}$. $\mathbf{I}_s^{\text{pump}}$ is the spin current pumped into the normal metal to be discussed below, whereas $\mathbf{I}_s^{(0)}$ is the current which flows back into the ferromagnet. The latter is driven by the accumulated spins
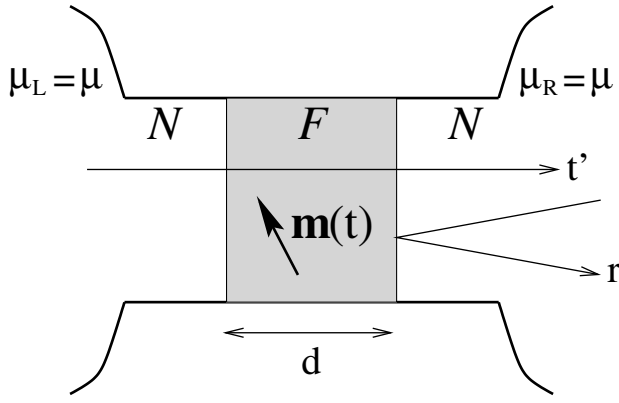
FIG. 1. Ferromagnetic film (F) adjacent to two normal metal layers (N). The latter are viewed as reservoirs in common thermal equilibrium. The reflection and transmission amplitudes r and t' shown here govern the spin current pumped into the right lead.

in the normal metal and, e.g., gives rise to the spin-current induced magnetization switching [4–6]. Here we model the normal metal as an ideal sink for the spin current, such that a spin accumulation does not build up. This approximation is valid when the spins injected by $\mathbf{I}_s^{\text{pump}}$ decay and/or leave the interface sufficiently fast, e.g., when the dimensionless conductance of the N-F interface is smaller than $h/\tau_{\text{sf}}\delta$ [14]. Here, τ_{sf} is the spin-flip relaxation time and δ is the energy level spacing at the Fermi energy of a normal metal film with a thickness which is the smaller one of the geometrical film thickness and the spin-flip diffusion length.

The current $\hat{I}(t)$ pumped by the precession of the magnetization into the right and left paramagnetic reservoirs, connected to the ferromagnet by normal metal leads (R) and (L), may be calculated in an adiabatic approximation since the period of precession $2\pi/\omega$ is typically much larger than the relaxation times of the electronic degrees of freedom of the system. The adiabatic charge-current response in nonmagnetic systems by a scattering matrix which evolves under a time-dependent system parameter $X(t)$ has been derived in [11,15]. Adopting Brouwer's notation [11], the generalization to the 2×2 matrix current (directed into the normal metal lead $l = R$ or L) reads

$$\hat{I}(t)^{\text{pump}} = e \, \frac{\partial \hat{n}(l)}{\partial X} \frac{dX(t)}{dt}, \qquad (2)$$

where the matrix emissivity into the lead l is

$$\frac{\partial \hat{n}(l)}{\partial X} = \frac{1}{4\pi i} \sum_{mnl'} \frac{\partial \hat{s}_{mn,ll'}}{\partial X} \hat{s}_{mn,ll'}^\dagger + \text{H.c.} \qquad (3)$$

and $\hat{s}$ is the 2×2 scattering matrix of the ferromagnetic insertion. m and n label the transverse modes at the Fermi energy in the normal metal leads and $l' = R, L$. Spin-flip scattering in the contact is disregarded. $\hat{s}$ depends on the magnetization direction $\mathbf{m}$ of the ferromagnet through the projection matrices $\hat{u}^\uparrow = (\hat{1} + \hat{\boldsymbol{\sigma}} \cdot \mathbf{m})/2$ and $\hat{u}^\downarrow = (\hat{1} - \hat{\boldsymbol{\sigma}} \cdot \mathbf{m})/2$ (Ref. [16]): $\hat{s}_{mn,ll'} = s_{mn,ll'}^\uparrow \hat{u}^\uparrow + s_{mn,ll'}^\downarrow \hat{u}^\downarrow$. The

spin current pumped by the magnetization precession is obtained by identifying $X(t) = \varphi(t)$, where φ is the azimuthal angle of the magnetization direction in the plane perpendicular to the precession axis. The resulting current is traceless, $\hat{I}^{\text{pump}} = -(e/\hbar)\hat{\boldsymbol{\sigma}} \cdot \mathbf{I}_s^{\text{pump}}$, i.e., the charge current indeed vanishes, and

$$\mathbf{I}_s^{\text{pump}} = \frac{\hbar}{4\pi} \left(A_r \mathbf{m} \times \frac{d\mathbf{m}}{dt} - A_i \frac{d\mathbf{m}}{dt} \right), \qquad (4)$$

where the interface parameters are

$$A_r = \frac{1}{2} \sum_{mn} \{ |r_{mn}^\uparrow - r_{mn}^\downarrow|^2 + |t_{mn}'^\uparrow - t_{mn}'^\downarrow|^2 \}, \qquad (5)$$

$$A_i = \text{Im} \sum_{mn} \{ r_{mn}^\uparrow (r_{mn}^\downarrow)^* + t_{mn}'^\uparrow (t_{mn}'^\downarrow)^* \}. \qquad (6)$$

Here, $r_{mn}^\uparrow$ ($r_{mn}^\downarrow$) is the reflection coefficient for spin-up (spin-down) electrons in the lth lead and $t_{mn}'^\uparrow$ ($t_{mn}'^\downarrow$) is the transmission coefficient for spin-up (spin-down) electrons into the lth lead. (See Fig. 1 for $l = R$.) Using unitarity of the scattering matrix for each spin direction, we can summarize Eqs. (5) and (6) by $A_r + iA_i = g^{\uparrow\downarrow} - t^{\uparrow\downarrow}$, where $g^{\sigma\sigma'} = \sum_{mn}\{\delta_{mn} - r_{mn}^\sigma (r_{mn}^{\sigma'})^*\}$ is the (dc) conductance matrix [16,17] and $t^{\uparrow\downarrow} = \sum_{nm} t_{mn}'^\uparrow (t_{mn}'^\downarrow)^*$. The spin current (4) trivially vanishes for the steady state, i.e., when $d\mathbf{m}/dt = 0$, and for unpolarized contacts $s_{mn,ll'}^\uparrow = s_{mn,ll'}^\downarrow$.

Per revolution, the precession pumps an angular momentum into an adjacent normal metal layer which is proportional to A_r, in the direction of the (averaged) applied magnetic field, and decaying in time. At first sight, it is astonishing that a pump can be operated by a single parameter varying in time, whereas the "peristaltic" pumping of a charge current requires at least two periodic parameters [11]. However, there are actually two periodic parameters (out of phase by $\pi/2$) hidden behind $\varphi(t)$, viz., the projections of the unit vector defined by φ in the plane perpendicular to the axis of precession.

By conservation of angular momentum, the spin torque on the ferromagnet resulting from the spin pumping into the nonmagnetic leads gives an additional term to the LLG equation (1). After including this term, Eq. (1) remains valid, but the gyromagnetic ratio and the damping constant are renormalized:

$$\frac{1}{\gamma} = \frac{1}{\gamma_0} \{ 1 + g_L [A_i^{(L)} + A_i^{(R)}]/4\pi M \}, \qquad (7)$$

$$\alpha = \frac{\gamma}{\gamma_0} \{ \alpha_0 + g_L [A_r^{(L)} + A_r^{(R)}]/4\pi M \}. \qquad (8)$$

Here, g_L is the Landé factor and M is the total magnetic moment (in units of μ_B) of the ferromagnetic film; the subscript 0 denotes the bulk values of γ and α; the superscripts (L) and (R) denote parameters evaluated on the left and right side of the F layer, respectively. Equations (7) and (8) are the central result of this Letter. A_r and A_i affect, e.g., ferromagnetic resonance experiments as a shift

of the resonance magnetic field via $A_i^{(L)} + A_i^{(R)}$, whereas $A_r^{(L)} + A_r^{(R)}$ increases the relative resonance linewidth.

From now on we focus on ferromagnetic films which are thicker than the coherence length $\lambda_{fc} = \pi/(k_\uparrow - k_\downarrow)$, where $k_{\uparrow\downarrow}$ are the spin-up and spin-down Fermi wave vectors, i.e., thicker than a few monolayers in the case of transition metals. In this regime, spin-up and spin-down electrons transmitted or scattered from one N-F interface interfere incoherently at the other interface, $t^{\uparrow\downarrow}$ vanishes, and the mixing conductance $g^{\uparrow\downarrow}$ is governed by the reflection coefficients of the isolated N-F interfaces.

$A_i = \mathrm{Im}\, g^{\uparrow\downarrow}$ vanishes for ballistic and diffusive contacts as well as nonmagnetic tunnel barriers [16]. First-principles calculations find very small A_i for Cu-Co and Fe-Cr [18]. It is, therefore, likely that A_i may be disregarded in many systems. If A_i does vanish on both sides of the ferromagnetic film, it follows from Eqs. (7) and (8) that the resonance frequency is not modified $\gamma = \gamma_0$ and the enhancement of the Gilbert damping is given by $\alpha' = g_L[A_r^{(L)} + A_r^{(R)}]/4\pi M$.

The coefficient A_r can be estimated by simple model calculations [16]. For ballistic (point) contacts, $A_r^B = (1 + p)g$ with the polarization $p = (g^{\uparrow\uparrow} - g^{\downarrow\downarrow})/(g^{\uparrow\uparrow} + g^{\downarrow\downarrow})$ and the average conductance $g = (g^{\uparrow\uparrow} + g^{\downarrow\downarrow})/2$. For diffusive N-F hybrids, $A_r^D = g_N$, the conductance of the normal metal part. A nonmagnetic tunneling barrier between F and N suppresses the spin current exponentially. The magnetization precession of a magnetic insulator can also emit a spin current into a normal metal, since $g^{\uparrow\downarrow}$ does not necessarily vanish because the phase shifts of reflected spin-up and spin-down electrons at the interface may differ [18].

Let us now estimate the damping coefficient α' for thin films of permalloy ($Ni_{80}Fe_{20}$, Py), a magnetically very soft material of great technological importance. Mizukami et al. [12] measured the ferromagnetic resonance linewidth of N-Py-N sandwiches and discovered systematic trends in the damping parameter as a function of Py layer thickness d for different normal metals. The spin polarization of electrons emitted by Py has been measured to be $p \approx 0.4$ in point contacts [19], the magnetization per atom is $f \approx 1.2$, and the Landé factor is $g_L \approx 2.1$ [12]. The interface conductance of metallic interfaces with Fe or Co is of the order of $10^{15}\ \Omega^{-1}\,m^{-2}$, with significant but not drastic dependences on interface morphology or material combination [20]. This corresponds to roughly one conducting channel per interface atom. Assuming the Fermi surface of the normal metal is isotropic, we arrive at the estimate $\alpha' \approx 1.1/d(\text{Å})$. The factor $1/d$ does not reflect an intrinsic effect; a reduced total magnetization is simply more sensitive to a given spin-current loss at the interface. Comparing with the intrinsic $\alpha_0 \approx 0.006$ of permalloy [12,21] the spin-current induced damping becomes important for ferromagnetic layers with thickness $d < 100$ Å. We can refine the estimate by including the significant film-thickness dependence of the magnetization measured

by the same group [12]. We, therefore, improve our above estimate as

$$\alpha'(d) \approx \kappa \times \frac{1.1}{d(\text{Å})} \times \frac{f_0}{f(d)}, \tag{9}$$

where f_0 and $f(d)$ are the atomic magnetization of the permalloy bulk and films. κ is an adjustable parameter representing the number of scattering channels in units of one channel per interface atom, which should be of the order of unity.

The experimental results for the damping factor α and the relative magnetization f/f_0 for N-Py-N sandwiches with $N =$ Pt, Pd, Ta, and Cu are shown in the insets of Fig. 2. Our estimate (9) appears to well explain the dependence of α on the permalloy film thickness d (see Fig. 2) for reasonable values of κ. First-principles calculations are called for to test these values.

The lack of a significant thickness dependence of damping parameter of the Cu-Py system requires additional attention. An opaque interface might be an explanation, but it appears more likely that due to long spin-flip relaxation times in Cu, the 5 nm thick buffer layers in [12] do not provide the ideal sink for the injected spins as assumed above. This means that a nonequilibrium spin accumulation on Cu opposes the pumped spin current and nullifies the additional damping when $h/\tau_{sf}\delta$ is comparable
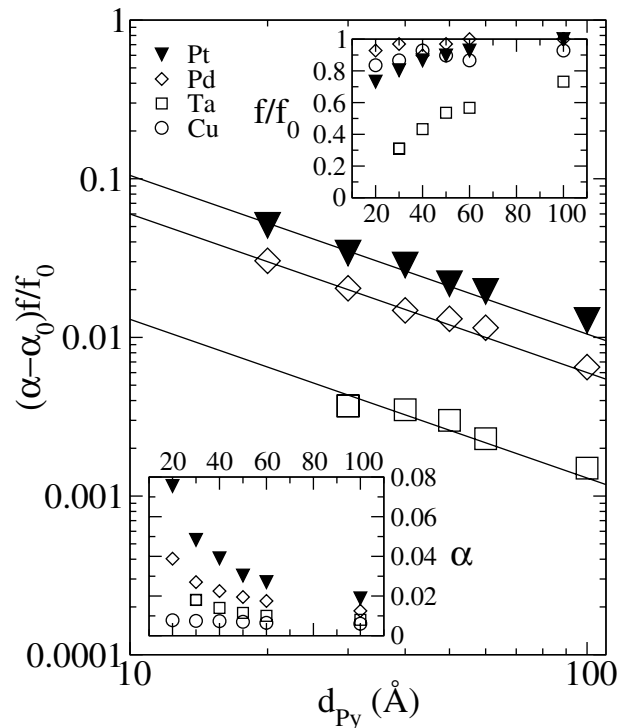

FIG. 2.   The lines show our theoretical result (9) with $\kappa = 1.0$, 0.6, and 0.1; the data points are derived from the measurements [12] shown in the two insets. Insets: Measured Gilbert damping constant α (lower inset) and the relative atomic magnetization f/f_0 (upper inset) in permalloy film of varied thickness d_{Py} in a trilayer structure N-Py-N.

or smaller than the conductance g. For 5 nm Cu buffers, $g\delta/h \sim 10^{13}$ s^{-1}, whereas $1/\tau_{sf} \sim 10^{12}$ s^{-1} [22]. It follows that Cu is indeed a poor sink for the injected spins and the Gilbert damping constant is not enhanced. On the other hand, Pt, Ta, and Pd are considerably heavier than Cu and, since $1/\tau_{sf}$ scales as Z^4 [23], where Z is the atomic number, have much larger spin-relaxation rates and our arguments hold.

A physical picture of the effect of magnetization precession in layered systems has been proposed earlier by Hurdequint et al. [24] in order to explain ferromagnetic and conduction electron spin resonance experiments. These authors realized that the precessing magnetization is a source of a nonequilibrium spin accumulation which diffuses out of the N-F interfaces into the adjacent normal metal layers where it can dissipate by spin-flip processes. Enhanced Gilbert damping in thin ferromagnetic films in contact with normal metal has also been discussed by Berger [9] for a ballistic N-F interface in a spin-valve configuration. His expression for the damping coefficient [Eq. (20) in Ref. [9]] scales like ours as a function of layer thickness, but differs as a function of material parameters. For example, in contrast to our result, Berger's expression does not vanish with vanishing exchange splitting.

In conclusion, we demonstrated that the Gilbert damping constant is enhanced in thin magnetic films with normal metal buffer layers by a spin-pump effect through the N-F contact. The damping is significant for transition metal films thinner than about 10 nm. Recent experiments on permalloy films [12] are well explained.

Note added.—After submission of this Letter, we became aware of new, exciting experimental results [25] on enhanced Gilbert damping in ultrathin iron films which support the spin-pumping mechanism proposed in this Letter [14].

[1] T. L. Gilbert, Phys. Rev. 100, 1243 (1955); L. D. Landau, E. M. Lifshitz, and L. P. Pitaevski, Statistical Physics, Part 2 (Pergamon, Oxford, 1980), 3rd ed.

[2] S. M. Bhagat and P. Lubitz, Phys. Rev. B 10, 179 (1974), and references therein.

[3] V. Korenman and R. E. Prange, Phys. Rev. B 6, 2769 (1972); V. S. Lutovinov and M. Y. Reizer, Zh. Eksp. Teor. Fiz. 77, 707 (1979) [Sov. Phys. JETP 50, 355 (1979)]; H. Suhl, IEEE Trans. Magn. 34, 1834 (1998); V. L. Safonov and H. N. Bertram, Phys. Rev. B 61, R14 893 (2000).

[4] J. C. Slonczewski, J. Magn. Magn. Mater. 159, L1 (1996); 195, L261 (1999).

[5] E. B. Myers, D. C. Ralph, J. A. Katine, R. N. Louie, and R. A. Buhrman, Science 285, 867 (1999).

[6] J. A. Katine, F. J. Albert, R. A. Buhrman, E. B. Myers, and D. C. Ralph, Phys. Rev. Lett. 84, 3149 (2000).

[7] C. H. Back, R. Allenspach, W. Weber, S. S. P. Parkin, D. Weller, E. L. Garwin, and H. C. Siegmann, Science 285, 864 (1999).

[8] F. Schreiber, J. Pflaum, Z. Frait, T. Mühge, and J. Pelzl, Solid State Commun. 93, 965 (1995).

[9] L. Berger, Phys. Rev. B 54, 9353 (1996).

[10] J.-E. Wegrowe, Phys. Rev. B 62, 1067 (2000); Y. B. Bazaliy, B. A. Jones, and S.-C. Zhang, Phys. Rev. B 57, R3213 (1998).

[11] P. W. Brouwer, Phys. Rev. B 58, R10 135 (1998).

[12] S. Mizukami, Y. Ando, and T. Miyazaki, Jpn. J. Appl. Phys. 40, 580 (2001); J. Magn. Magn. Mater. 226, 1640 (2001).

[13] B. Heinrich, K. B. Urquhart, A. S. Arrott, J. F. Cochran, K. Myrtle, and S. T. Purcell, Phys. Rev. Lett. 59, 1756 (1987).

[14] Y. Tserkovnyak, A. Brataas, and G. E. W. Bauer (to be published).

[15] M. Büttiker, H. Thomas, and A. Prêtre, Z. Phys. B 94, 133 (1994).

[16] A. Brataas, Y. V. Nazarov, and G. E. W. Bauer, Phys. Rev. Lett. 84, 2481 (2000); Eur. Phys. J. B 22, 99 (2001).

[17] X. Waintal, E. B. Myers, P. W. Brouwer, and D. C. Ralph, Phys. Rev. B 62, 12 317 (2000).

[18] K. Xia, P. J. Kelly, G. E. W. Bauer, A. Brataas, and I. Turek, cond-mat/0107589 (unpublished).

[19] R. J. Soulen, Jr. et al., Science 282, 85 (1998).

[20] K. Xia, P. J. Kelly, G. E. W. Bauer, I. Turek, J. Kudrnovsky, and V. Drchal, Phys. Rev. B 63, 064407 (2001).

[21] C. E. Patton, Z. Frait, and C. H. Wilts, J. Appl. Phys. 46, 5002 (1975).

[22] R. Meservey and P. M. Tedrow, Phys. Rev. Lett. 41, 805 (1978).

[23] A. A. Abrikosov and L. P. Gor'kov, Zh. Eksp. Teor. Fiz. 42, 1088 (1962) [Sov. Phys. JETP 15, 752 (1962)].

[24] P. Monod, H. Hurdequint, A. Janossy, J. Obert, and J. Chaumont, Phys. Rev. Lett. 29, 1327 (1972); H. Hurdequint and G. Dunifer, J. Phys. (Paris), Colloq. 49, C8-1717 (1988); H. Hurdequint, J. Magn. Magn. Mater. 93, 336 (1991).

[25] R. Urban, G. Woltersdorf, and B. Heinrich, Phys. Rev. Lett. 87, 217204 (2001).