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Currents and torques due to spin-dependent diffraction in ferromagnetic/spin spiral bilayers

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Abstract

Spin-dependent transport through the interface between a ferromagnet and a spin spiral is investigated using both ballistic and diffusive models. We find that spin-dependent interferences lead to a new type of diffraction called 'spin diffraction'. It is shown that this spin diffraction leads to local spin and electrical current along the interface, as well as spin transfer torque acting on the spin spiral. This study also emphasizes that in highly inhomogeneous magnetic configurations, diffracted electrons must be taken into account to properly describe the spin transport.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

In the last decade, considerable progress has been achieved in the investigation of magnetization control by a spin-polarized current [1]. Many experimental and theoretical efforts have focused on metallic magnetic multilayered nanostructures (SV) [2–4], magnetic tunnel junctions (MTJ) [5–7] and recently antiferromagnets [8]. It seems that, although important issues remain under investigation, most of the microscopic and macroscopic features of spin transfer torque within the macrospin approximation have now been understood [9] and spin torque is now seriously considered for applications [10].

The development of micromagnetic simulations [11] and the recent observation of current-induced domain wall motion [12] have underlined the question of spin transport in inhomogeneous systems.

The interaction between spin transport and inhomogeneous magnetization has been studied at the interface between a normal metal (N) and a ferromagnet (F), assuming smoothly varying magnetization at the N/F interface [13, 14]. It was found that, due to the inhomogeneous magnetization at the interface, the interfacial spin accumulation can possess components transverse to the local interfacial magnetization, leading to the so-called 'self-torque' effect [13, 14]. In other words, the interfacial magnetic inhomogeneities lead to an interfacial torque for a single magnetic layer sandwiched between non-magnetic electrodes.

The other topic that has attracted considerable attention is the phenomenon of current-induced domain wall motion (DWM). In such inhomogeneous systems, the Slonczewski torque [1] (*adiabatic* spin transfer torque) still exists but is not sufficient to interpret the experimental data [15, 16]. Besides the usual adiabatic spin transfer torque [1], another torque (usually called *non-adiabatic* torque) has been proposed by Berger [17] that arises from the non-adiabaticity of the magnetic system. Many studies have then investigated the nature of this torque. For example, Tatara *et al* [18] proposed a momentum-transfer torque, arising from the fast spin texture of the domain wall. Independently, Thiaville *et al* [19] and Barnes *et al* [20] proposed a modified Landau–Lifshitz–Gilbert equation taking into account the non-adiabatic torque. Zhang and Li [21] also showed that in the case of a slowly varying magnetic film, the spin relaxation of conduction electrons could lead to this non-adiabatic torque. Microscopic studies on the properties of this torque have been carried out either within a simple classical model [22], solving the Schrödinger equation [23–25], using Boltzmann formalism [26, 27] or Keldysh theory [28]. The non-adiabatic torque derived within these works proves its non-local origin. Nguyen *et al* [29] have shown that such non-adiabatic torque may be dramatically enhanced in dilute magnetic semiconductors, due to strong spin–orbit interaction. Recently, Tatara *et al* [30] proposed a microscopic derivation taking into account both non-locality and spin relaxation.

Finally, note that other mechanisms have been proposed to interpret experimental results without the need of non-adiabatic torque, such as thermal activation [31] or the use of Landau–Lifshitz damping instead of Gilbert damping [32].

The effective torque acting on a spatially varying magnetization can be written as [19–21, 25]:

$$\mathbf{T} = a\frac{\partial \mathbf{M}}{\partial t} + b\mathbf{M}\times\frac{\partial \mathbf{M}}{\partial t} - c_1\mathbf{M}\times[\mathbf{M}\times(\mathbf{j}_e\cdot\nabla)\mathbf{M}] - c_2\mathbf{M}\times(\mathbf{j}_e\cdot\nabla)\mathbf{M}. \quad (1)$$

The first two terms are, respectively, the renormalization factor of the gyromagnetic ratio and damping parameter, while the two following terms are proportional to the current density $\mathbf{j}_e$ and stand for spin transfer torque in spatially varying magnetic structures. The first term is called *adiabatic torque* while the second one is the *non-adiabatic torque*. For smooth enough inhomogeneities of the magnetic structure, the adiabatic approximation is usually assumed: the electron spin follows the local magnetization producing a small torque proportional to the spatial derivative of the magnetization [23, 26] (third term of equation (1)). In this case, it is usually accepted that the non-adiabatic term is small ($c_2/c_1 \approx 0.01$) but cannot be neglected in domain wall experiments [16, 19, 24]. Note that c_1 and c_2 are non-local coefficients [23–26, 28] which means that, at each point of the structure, one needs to consider the contribution of all electrons.

However, most of these studies addressed inhomogeneous systems in the adiabatic limit and the question of highly non-adiabatic systems is generally omitted: in such systems, the non-adiabatic torque should be of the same order as the adiabatic torque or even higher.

In this paper, we study the crossover between the adiabatic and non-adiabatic regimes in an 'academic' system using two different descriptions of the spin transport: ballistic and diffusive. The system under investigation consists of two adjacent layers—one with homogeneous magnetization (F) and the other with a spin spiral (SS). This system is actually justified by the recent observation of inhomogeneous magnetic states in CoFe-based metallic spin valves [33]. The authors observed that under the influence of the spin-polarized current, magnetic vortices with radius 5 nm appear. Our model uses a simplified picture of such systems in which a ferromagnetic layer has an inhomogeneous magnetization. It, however, allows us to predict new effects which take place in such systems and in particular a new phenomenon of spin diffraction.

In our system, spin-polarized electrons moving from F into SS are diffracted by the spin spiral and the diffraction pattern is spin-dependent. As a result, one can expect the current perpendicular to the F/SS interface to produce a non-zero torque acting on the magnetization of SS as well as a local spin current and even charge current (Hall effect and spin Hall effect [34]) along the F/SS interface.

This system is completely different from the one studied in self-torque investigations [13, 14]. As a matter of fact, self-torque arises from lateral spin diffusion due to magnetic inhomogeneities at the N/F interface (no spin-diffraction effect is considered), whereas our system gives rise to specific spin-dependent patterns. It is also different from the previous study of Xiao *et al* [23] on spin spirals since in our case the adiabatic approximation is no longer valid. Indeed, in our model the electrons moving through the interface keep the memory of their spin state over a finite length. This case has not been studied yet and strong differences from the adiabatic approximation can be expected. Moreover, contrary to most of the theoretical studies on domain wall motion, the electrical current is not injected along the domain wall direction, but transversely, yielding a spin-diffraction effect.

Note, finally, that such helicoidal spin structures exist in some compounds such as MnSi [35], oxide materials such as $SrFeO_3$, NaCuO [36] or rare-earth-based compounds [37]. This helicoidal structure can also be a simplified picture of narrow stripe domains with domain wall width comparable to domain width. The model system we study here is thus a first step towards domain wall motion and will serve as a reference for forthcoming studies on this topic.

This paper is organized as follows. In section 2, we present the quantum and diffusive models of spin transport at the F/SS interface. The results are described in section 3 and a discussion on experimental aspects is proposed in section 4. The conclusion is given in section 5.

2. Model and formalism

The structure under investigation is depicted in figure 1 and consists of two regions: the left one ($z < 0$) is a ferromagnet F, with a magnetization $\mathbf{P} = P\mathbf{z}$ and the right one ($z > 0$) is a spin spiral SS with a helicoidal magnetization $\mathbf{M} = M(\sin(\theta_0 + Qx)\mathbf{x} + \cos(\theta_0 + Qx)\mathbf{z})$ expanding in the (x, z) plane. Here, the interface is parallel to the (x, y) plane and z is perpendicular to it. We consider that the region F (respectively SS) is semi-infinite and connected to a ferromagnetic (respectively, spin spiral) reservoir. The bias voltage V is applied in the z direction across the interface.

2.1. Formulation of spin transport

To model the spin-dependent transport in such a system, we use two different approaches: ballistic and diffusive. This method gives information about the mechanisms that are model-dependent and those that are system-dependent. These formalisms are described in [4].

In the ballistic model, we use the Keldysh out-of-equilibrium technique [4, 38] that expresses the Keldysh Green functions $G^{-+}_{\sigma\sigma'}(\mathbf{rr}')$ as a function of the basis of wavefunctions

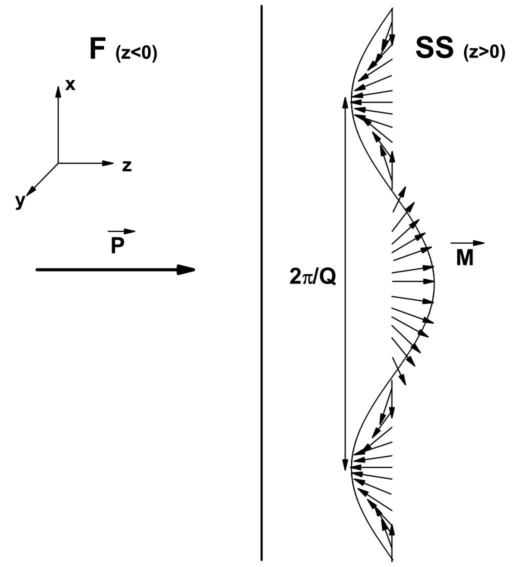

Figure 1. Cartoon of the bilayer structure. The left semi-infinite layer is a ferromagnet with a homogeneous magnetization and the right semi-infinite layer is a spin spiral with wavelength $2\pi/Q$.

$\Psi_{\mathrm{l(r)}\sigma}(\mathbf{r})$ for an electron with initial spin projection σ moving from the left (right) to the right (left):

$$G^{-+}_{\sigma\sigma'}(\mathbf{rr}') = f_{\mathrm{l}}(\mu_{\mathrm{l}})\Psi_{\mathrm{l}\sigma'}(\mathbf{r}')\Psi^{*}_{\mathrm{l}\sigma}(\mathbf{r}) + f_{\mathrm{r}}(\mu_{\mathrm{r}})\Psi_{\mathrm{r}\sigma'}(\mathbf{r}')\Psi^{*}_{\mathrm{r}\sigma}(\mathbf{r}) \quad (2)$$

where $f_{\mathrm{l(r)}}$ are the Fermi distribution functions in the left and right electrodes and $\mu_{\mathrm{l(r)}}$ are the chemical potentials in these electrodes so that $V = (\mu_{\mathrm{l}} - \mu_{\mathrm{r}})/e$. In this formalism, the spin-dependent conductivities and the spin density can be expressed as a function of the Keldysh Green functions, $G^{-+}_{\sigma\sigma'}(\mathbf{rr}')$ (see equations (6)–(9)).

In the diffusive approach, we use the formalism developed by Zhang *et al* [40, 41], where the electrical and spin currents are driven by the electrical field, the charge accumulation and the spin accumulation. Using this model, we assume that the s–d exchange coupling of SS is not too strong so that spin precession can be taken into account in the transport description [4].

The spin density (or spin accumulation in diffusive systems) exerts a torque on the local magnetization

$$\mathbf{T} = \frac{J_{\mathrm{sd}}}{\hbar M \mu_{\mathrm{B}}}\mathbf{M}\times\mathbf{m} = \frac{J_{\mathrm{sd}}}{\hbar\mu_{\mathrm{B}}}\begin{pmatrix} -m_y\cos\theta(x) \\ m_x\cos\theta(x) - m_z\sin\theta(x) \\ m_y\sin\theta(x)\end{pmatrix} \quad (3)$$

where γ is the gyromagnetic ratio, μ_{B} is the Bohr magneton and J_{sd} is the s–d exchange coupling. We can distinguish two components in the spin torque: one lies in the (x, z) plane and corresponds to the usual Slonczewski term [1] (in-plane torque—proportional to m_y), while the second one is perpendicular to the (x, z) plane and corresponds to the effective field term (out-of-plane torque, proportional to both m_x and m_z).

2.2. Ballistic description

In the ballistic model of the F/SS interface, the Hamiltonian of the system is:

$$H = \frac{p^2}{2m} - U - J_{\mathrm{sd}}(\vec{\sigma}\cdot\mathbf{M}) \quad (4)$$

where U is the potential profile, $\vec{\sigma}$ is the vector of Pauli matrices and $\mathbf{M}$ is the local magnetization. The wavefunctions are the eigenstates of the linear system:

$$H = \begin{pmatrix} \frac{-\hbar^2}{2m}\Delta - E & 0 \\ 0 & \frac{-\hbar^2}{2m}\Delta - E\end{pmatrix} - J_{\mathrm{sd}}\begin{pmatrix}\cos\theta & \sin\theta \\ \sin\theta & -\cos\theta\end{pmatrix} \quad (5)$$

where θ is non-zero only in the SS where $\theta = \theta_0 + Qx$. In the following, an electron state is determined by the Hartree–Fock spin-dependent wavefunction $(\Psi^{\uparrow}, \Psi^{\downarrow})$, where $\uparrow$ ($\downarrow$) refers to majority (minority) electron spin projection. Note that *majority* (*minority*) refers to the local magnetization of F.

Because the Hamiltonian depends on x through θ, there are two ways to solve this eigenvalue problem: one is to consider that Q is small enough so that the spatial variation of H can be neglected (adiabatic assumption) or, on the contrary, one has to properly consider this variation using a comprehensive method (non-adiabatic assumption).

2.2.1. Adiabatic assumption. We firstly build an adiabatic model in which the spin of the electron is assumed to follow the local magnetization. We consider that $\theta \approx \theta_0$ and solve the Schrödinger equation in such a simple system. To model the F/SS bilayer, we then take into account the spatial variation in the expression of the wavefunctions themselves setting $\theta = \theta_0 + Qx$. This way, the diffraction is ignored and at each point (x, y, z) of the structure, currents and torques are defined by the value of the angle θ at this point. We will use this oversimplified model as a reference in order to illustrate the role of spin diffraction.

2.2.2. Non-adiabatic model. The previous model is not correct when the inhomogeneities (given by the wavelength of SS, $2\pi/Q$) are very steep ($Q/2\pi$ is larger than k_{F}). In this case, we cannot neglect the spatial variation of θ and the Hamiltonian must be solved in the local coordinate system [42]. The linear system (H, Ψ) is then transformed, by rotation, into a new linear system $(\tilde{H}, \tilde{\Psi})$, in which $\tilde{\Psi}$ are plane waves. After obtaining the wavefunctions $\tilde{\Psi}$ in the local system of coordinates of SS, we go back to the initial system of coordinates of F to get Ψ. Note that in this case, the wavevectors for majority and minority spins depend upon the spin spiral wavelength $2\pi/Q$ [39]. Finally, the boundary conditions at the interface $z = 0$ lead to a linear system that can be solved using Fourier transformation, in order to change the x-dependence into a κ-dependence (κ is the in-plane component of the wavevector). We do not further enter into the details of the calculation (see [39]); the spin-dependent electrical current densities $\mathbf{J_e}$ and spin density $\mathbf{m}$ are calculated using the usual definition [4]:

$$m_x + \mathrm{i}m_y = 2\frac{J_{\mathrm{sd}}}{\mu_{\mathrm{B}}}\frac{a_0^3}{(2\pi)^2}\int G^{-+}_{\uparrow\downarrow}(\mathbf{r}, \mathbf{r}, \epsilon) \quad (6)$$

$$m_z = \frac{J_{\rm sd}}{\mu_{\rm B}}\frac{a_0^3}{(2\pi)^2}\int\left[G^{-+}_{\uparrow\uparrow}(\mathbf{r},\mathbf{r},\epsilon) - G^{-+}_{\downarrow\downarrow}(\mathbf{r},\mathbf{r},\epsilon)\right] \quad (7)$$

$$\mathbf{J}_{\uparrow(\downarrow)} = \frac{\hbar e}{4\pi m_{\rm e}}\int[\nabla_{\mathbf{r}} - \nabla_{\mathbf{r}'}]G^{-+}_{\uparrow\uparrow(\downarrow\downarrow)}(\mathbf{r},\mathbf{r}',\epsilon)|_{\mathbf{r}=\mathbf{r}'} \quad (8)$$

$$\mathbf{J} = \mathbf{J}_\uparrow + \mathbf{J}_\downarrow. \quad (9)$$

2.3. Diffusive description

Following the formalism developed by Zhang *et al* [40, 41], in a diffusive system, where spin diffusion is not negligible, the electric current is:

$$j_{\rm e}^z = 2C_0E^z - 2D_0\frac{\partial n^0}{\partial z} - 2\mathbf{D}\frac{\partial \mathbf{m}}{\partial z} \quad (10)$$

$$j_{\rm e}^x = -2D_0\frac{\partial n^0}{\partial x} - 2\mathbf{D}\frac{\partial \mathbf{m}}{\partial x} \quad (11)$$

where $\mathbf{m}$ and n^0 are spin and charge accumulation respectively, while C_0 and D_0 denote the conductivity and diffusion constant, respectively; they are related by the Einstein relation $C_0 = e^2 g(E_{\rm F})D_0$, where $g(E_{\rm F})$ is the density of states at the Fermi level. The spin polarization parameters are defined by the relations $\mathbf{C} = \beta\mathbf{M}C_0$ and $\mathbf{D} = \beta'\mathbf{M}C_0$. Equivalently, the spin current is given by:

$$\mathbf{j}_m^z = 2\mathbf{C}E^z - 2\mathbf{D}\frac{\partial n^0}{\partial z} - 2D_0\frac{\partial \mathbf{m}}{\partial z} \quad (12)$$

$$\mathbf{j}_m^x = -2\mathbf{D}\frac{\partial n^0}{\partial x} - 2D_0\frac{\partial \mathbf{m}}{\partial x}. \quad (13)$$

The equation of motion of the spin accumulation may be written:

$$\frac{\partial \mathbf{m}}{\partial t} + \nabla\cdot\mathbf{j}_m + \frac{J_{\rm sd}}{\hbar}[\mathbf{m}\times\mathbf{M}] = -\frac{\mathbf{m}}{\tau_{\rm sf}} \quad (14)$$

where the last term on the left hand side represents the precessional motion of the spin accumulation due to the s–d interaction. Note that equation (14) is valid for weakly ferromagnetic alloys with a strong spin-diffusion length, as discussed in [4]. The current-induced spin accumulation can be divided into two terms:

$$\mathbf{m}(x,z) = \alpha_m\mathbf{M} + \delta\mathbf{m} \quad (15)$$

where the first term on the right hand side is adiabatic and proportional to the local magnetization $\mathbf{M}$ while the other one describes the deviation from the adiabatic process. If we take into account that the adiabatic part does not give any contribution to the out-of-equilibrium spin current, we rewrite the equation of motion:

$$\frac{\partial \delta\mathbf{m}}{\partial t} + \nabla\cdot\mathbf{j}_m + \frac{J_{\rm sd}}{\hbar}[\delta\mathbf{m}\times\mathbf{M}] = -\frac{\delta\mathbf{m}}{\tau_{\rm sf}}. \quad (16)$$

Equivalently, the continuity equation for the charge accumulation and spin current reads:

$$\frac{\partial n^0}{\partial t} + \nabla\cdot\mathbf{j}_{\rm e} = 0. \quad (17)$$

It is instructive to define $\lambda_{\rm sf} = \sqrt{2D_0\tau_{\rm sf}}$ and $\lambda_J = \sqrt{2D_0\hbar/J_{\rm sd}}$, where $\tau_{\rm sf}$ is the spin-flip relaxation time, and $J_{\rm sd}$ is the exchange between the itinerant electrons and the magnetic background. Similarly to the ballistic model, we perform a rotation of the coordinate system which has the following matrix form:

$$\begin{pmatrix} n^0 \\ m_x \\ m_y \\ m_z \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos\phi\cos\theta & -\sin\phi & \cos\phi\sin\theta \\ 0 & \sin\phi\cos\theta & \cos\phi & \sin\phi\sin\theta \\ 0 & -\sin\phi & 0 & \cos\theta \end{pmatrix} \cdot \begin{pmatrix} n^{0\prime} \\ m'_x \\ m'_y \\ m'_z \end{pmatrix} \quad (18)$$

where $\theta = \theta_0 + Qx$ and $\phi = \pi/2$. After this transformation, we obtain a system of differential equations with constant coefficients for $n^{0\prime}$ and $\mathbf{m}' = (m'_x, m'_y, m'_z)$. The basic matrix equation for the SS layer in steady state is then:

$$\begin{pmatrix} -\Delta & \mathrm{i}\beta' M_d^0 Q\frac{\partial}{\partial x} & 0 & -\beta' M_d^0\Delta \\ -\mathrm{i}\beta' M_d^0 Q\frac{\partial}{\partial x} & -\Delta + Q^2 + \frac{1}{\lambda_{\rm sd}^2} & \frac{M_d^0}{\lambda_J^2} & -2Q\frac{\partial}{\partial x} \\ 0 & -\frac{M_d^0}{\lambda_J^2} & \Delta + \frac{1}{\lambda_{\rm sd}^2} & 0 \\ -\beta' M_d^0\Delta & 2Q\frac{\partial}{\partial x} & 0 & -\Delta + Q^2 + \frac{1}{\lambda_{\rm sd}^2} \end{pmatrix} \cdot \begin{pmatrix} n^{0\prime}(x,z) \\ m'_x(x,z) \\ m'_y(x,z) \\ m'_z(x,z) \end{pmatrix} = 0 \quad (19)$$

where Δ is Laplacian. The equation for the F layer is simpler because it does not contain any terms proportional to Q. Starting with the analytical solution in each of the two layers in the form:

$$f(x,z) = \sum_{n,i} C_{n,i}\mathrm{e}^{-q_{i,n}z}\mathrm{e}^{-\mathrm{i}k_n x} \quad (20)$$

where $k_n = nQ$ ($n = 0, 1, \ldots$) and $q_{i,n}$ are the roots of the secular equation corresponding to equation (19), we write down the boundary conditions between the layers. Similarly to the previous section, the continuity of both spin accumulation and currents at the interface gives recurrent formulae to determine the coefficients $C_{n,i}$.

3. Results and discussion

The diffusive transport equation proposed above is valid as long as $\lambda_{\rm sf} \ll 2\pi/(k_{\rm F}^\uparrow - k_{\rm F}^\downarrow)$ (see the discussion in [4]) and applies to weakly ferromagnetic alloy with strong scattering and low Curie temperature (such as NiFeCu alloys, for example). We then take $\lambda_J \approx 1$ nm [43] and $\lambda_{\rm sf} = 15$ nm in the diffusive model. In this model, the shortest spin spiral wavelength under consideration is $2\pi/Q = 6.3$ nm.

For the ballistic model, the Fermi wavevectors for majority and minority spins are respectively set to $k_{\rm F}^\uparrow = 1.1$ Å^{-1}, $k_{\rm F}^\downarrow = 0.6$ Å^{-1}; the shortest wavelength for the spin spiral is $2\pi/Q \approx 0.6$ nm in the ballistic model, which corresponds to a

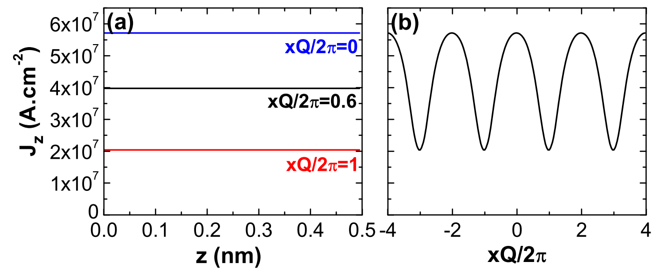

Figure 2. Intensity of the perpendicular current density J_e^z as a function of (a) z at $xQ/2\pi = 0, 0.6, 1$ and (b) x at $z = 0$ nm. $2\pi/Q = 0.6$ nm.

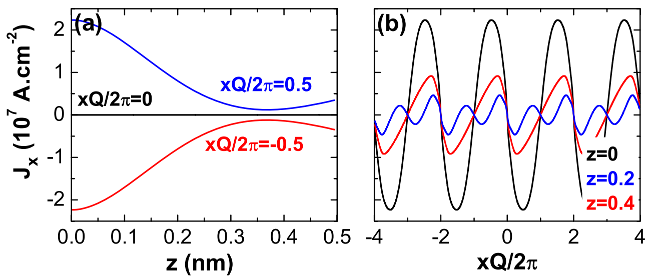

Figure 3. Intensity of the longitudinal current density J_e^x as a function of (a) z at $xQ/2\pi = -0.5, 0, 0.5$ and (b) x at $z = 0, 0.2$ and 0.4 nm. $2\pi/Q = 0.6$ nm.

highly inhomogeneous magnetic system. In this configuration, the spin spiral structure is close to that of an antiferromagnet. For simplicity, the Fermi wavevectors of SS and F are assumed to be the same.

Finally, we consider the linear approximation for the ballistic model: for a small enough bias voltage, only the electrons originating from the left reservoir with an energy located between μ_l and μ_r significantly contribute to the transport (current densities and torques). This approximation is justified for both the electrical currents and in-plane torque. However, electrons under the Fermi level are known to contribute to the non-equilibrium out-of-plane torque. We numerically verified that this contribution is small and does not alter the general behavior of this component (especially the interference scheme). Furthermore, this approximation allows a better comparison between the ballistic and diffusive models (in the latter, only Fermi electrons are taken into account). We set $\mu_l - \mu_r = 0.038$ μeV so that the amplitude of the electrical current is comparable in both ballistic and diffusive models.

3.1. Adiabatic model

Usually, in micromagnetic simulations and in weakly inhomogeneous ferromagnetic systems, it is convenient to assume that the spin transport (currents and torques) is defined through the adiabatic model presented in section 2.2.1. We will show, however, that in the case of high inhomogeneities, this assumption is not valid. Figures 2 and 3 display the perpendicular and longitudinal electrical currents J_e^x and J_e^z as a function of z and x in the adiabatic model (no spin diffraction). The perpendicular current J_e^z is unchanged in the z-direction (figure 2(a)) and oscillates in the x-direction (figure 2(b)), in agreement with the giant magnetoresistance (GMR) effect: because no non-locality has been introduced in the adiabatic model, the intensity of J_e^z only depends on the local angle between the magnetizations of the two layers.

Because $\partial J_e^z/\partial z = 0$, one expects that $\partial J_e^x/\partial x = 0$ in the stationary state ($\partial\rho/\partial t = 0$). However, the longitudinal current J_e^x also oscillates following x (figure 3(b)) and decreases following z (figure 3(a)) (the intensities of J_e^z and J_e^x do not depend on Q). As a consequence, the divergency of the current density $\nabla \cdot \mathbf{J_e}$ is non-zero in the adiabatic model, as long as $2\pi/Q$ is close to 0.6 nm (J_e^z is constant along z, whereas J_e^x oscillates as a function of x). This indicates that the adiabatic assumption is not valid for strong magnetic inhomogeneities. However, for very smooth magnetization variation ($2\pi/Q \gg 1$ nm), $\partial J_e^x/\partial x$ becomes very small and the divergency $\nabla \cdot \mathbf{J_e}$ goes to zero. In this extreme case, the adiabatic assumption is justified. However, for high inhomogeneities, one needs to consider the more comprehensive model proposed in sections 2.2.2 and 2.3.

3.2. Spin diffraction at F/SS interface

To understand the physics of the propagating waves in such a bilayer in a non-adiabatic regime, let us consider a plane wave originating from the left reservoir and initially in a majority spin state. After interaction (reflection transmission) with the interface, the electron is no longer in a pure spin state and can be described as the superposition of majority and minority spins (respectively described by $\Psi^{\uparrow}$ and $\Psi^{\downarrow}$—see [39] for example). Furthermore, because of the spatial variation of the magnetization of SS, we saw that the probability for an electron with an incident wavevector κ to propagate with a wavevector $\kappa + nQ$ ($\kappa + Q(n - 1/2)$) after reflection (transmission) by the interface is non-zero. This is illustrated by figure 4, where an incident wave initially in a pure spin state with incident wavevector κ diffracts forward and backward and changes to a mixed spin state.

The interferences between the diffracted waves may significantly affect the currents and torques, creating a strong non-locality in the system. As a matter of fact, in such inhomogeneous magnetic systems, spin torques and currents arise from the contribution of frontward and backward electrons diffracted from all over the structure.

3.3. Currents and torques in non-adiabatic regime

3.3.1. Current densities. We first analyze the influence of the diffraction on the spin-dependent electrical currents in this system. Figure 5 presents the intensity of the perpendicular current density J_e^z in the SS layer as a function of z (a) and x (b), in the ballistic model.

One can observe a drastic difference between figures 2 and 3. Within 0.5 nm near the interface, J_e^z oscillates with the SS magnetization, resembling figure 2(b), and the current is distributed along the interface, forming 'hot spots' (i.e. regions of high current density) corresponding to the parallel orientation of the local magnetizations of F and SS layers (figure 5(b)). But further from the interface,

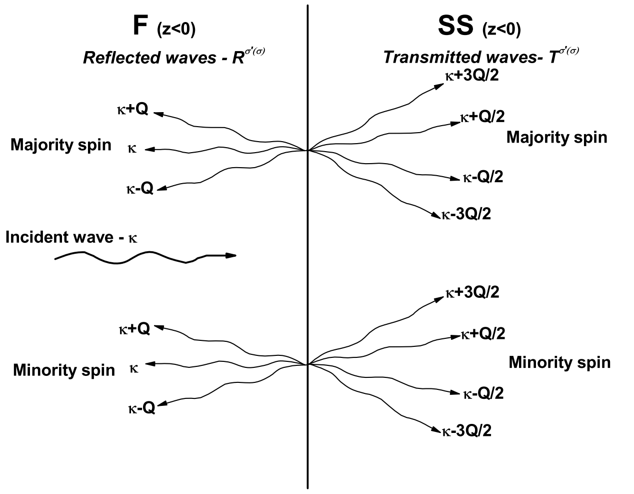

Figure 4. Schematics of the spin-dependent diffraction of a spin-polarized plane wave at the interface F/SS. The initially majority spin-polarized electron impinges with a wavevector κ on the interface and produces several reflected and transmitted waves with wavevectors $\kappa \pm Qn$ and $\kappa + Q(n \pm 1/2)$, respectively (n is an integer).

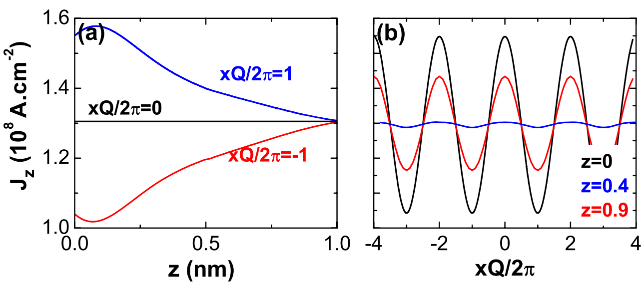

Figure 5. Intensity of the longitudinal current density J_e^x as a function of (a) z at $xQ/2\pi = -1, 0, 1$ and (b) x at $z = 0$, 0.4 and 0.9 nm. $2\pi/Q = 0.6$ nm.

the amplitude of the oscillations in J_e^z decreases and the current reaches a uniform value (figure 5(a)): the interferences between the diffracted waves become stronger (increasing with z) so that the spatial current modulation induced by the GMR effect is reduced, leading to an averaging of J_e^z.

This feature is independent of the model. Figure 6 displays the longitudinal and transverse electrical currents as a function of z in ballistic and diffusive regimes. In the ballistic regime, J_e^z (figure 6(a)) possesses several periods $T_1^{n,m} = 1/(k_{3(\kappa+Q(n-\frac{1}{2}))} - k_{4(\kappa+Q(m-\frac{1}{2}))})$, $T_2^{n,m} = 1/(k_{3(\kappa+Q(n-\frac{1}{2}))} - k_{3(\kappa+Q(m-\frac{1}{2}))})$ and $T_3^{n,m} = 1/(k_{4(\kappa+Q(n-\frac{1}{2}))} - k_{4(\kappa+Q(m-\frac{1}{2}))})$, depending on Q. The z-dependence of J_e^z in the diffusive model (figure 6(b)) is comparable to that in the ballistic model except that the oscillations do not exist. Note that the decay length is inversely proportional to Q: this means that the averaging effect due to interferences of spin-diffracted waves is stronger in a highly inhomogeneous system (large Q) than in a weakly inhomogeneous one (small Q). Then, the decrease in the longitudinal current and the damping of the current modulation due to the GMR effect occur on a longer characteristic length when Q is small.

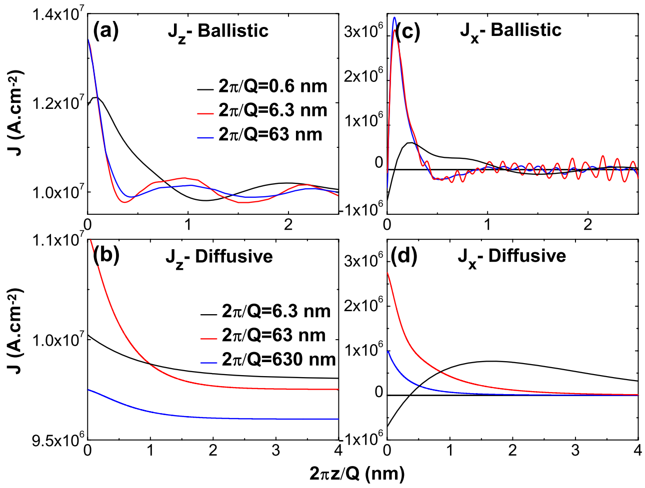

Figure 6. Intensity of the electrical current density as a function of z: J_e^z and J_e^x in ballistic ((a), (c)) and diffusive ((b), (d)) models.

The same structure is found for J_e^x (figures 6(c) and (d)), which is damped far from the interface: due to the spin diffraction, local spin and electrical currents arise along the F/SS interface. The maximum intensity of J_e^x may be increased by reducing Q (see the red solid lines): the averaging is strong enough to create a local transverse electrical current but not enough to damp it near the interface. Consequently, one may expect that local transverse current can be maintained over a rather long length in such systems, depending on Q. These currents are only local and eventually averaged out far from the interface (2–3 nm).

The angular dependence (or equivalently the x-dependence) is also strongly affected by SS wavelength. Figure 7 shows the longitudinal and transverse currents in ballistic and diffusive models, as a function of x. When $xQ = 2n\pi(n\pi)$ (n is an integer), the local magnetizations of F and SS are parallel (antiparallel), so that the longitudinal current is maximum (minimum) and the transverse current is zero. In highly inhomogeneous systems ($2\pi/Q = 0.6$ nm, black lines), when $x = \pi/2 + n\pi$, the transverse current is maximum and both J_e^z and J_e^x show a sine or cosine dependence on x.

However, when smoothing the inhomogeneities in SS, the angular dependence is reduced and diverges from the cosine or sine dependence. In particular, the current variation close to $x = n\pi$ (local antiparallel alignment) is sharpened. As a matter of fact, the contribution of electrons diffracted from regions where the magnetizations are parallel (or close to parallel) are reduced with increasing SS wavelength.

3.3.2. Torques. After the analysis of the current densities in the F/SS bilayer, one would expect the same kind of behavior for the in-plane and out-of-plane spin torques, namely an important spin torque close to the interface, which decreases far from the interface. Figure 8 displays the two components of spin torque calculated with ballistic and diffusive models, as a function of z. Similarly to the current densities, the in-plane and out-of-plane torques decrease rapidly to zero far from the interface. This decay is attributed to spin-diffraction-induced

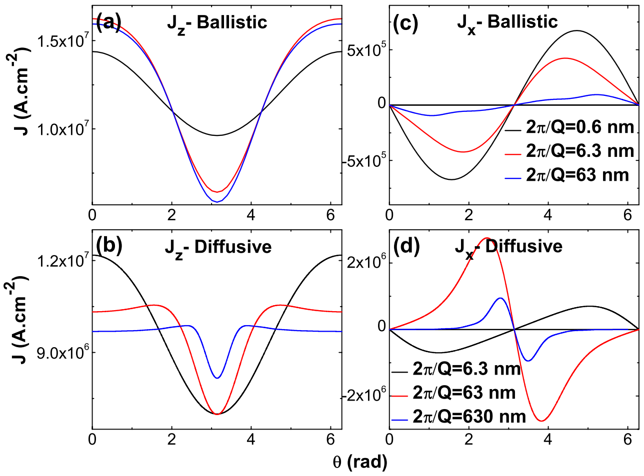

Figure 7. Intensity of the electrical current density as a function of x: J_e^z and J_e^x in ballistic ((a), (c)) and diffusive ((b), (d)) models.

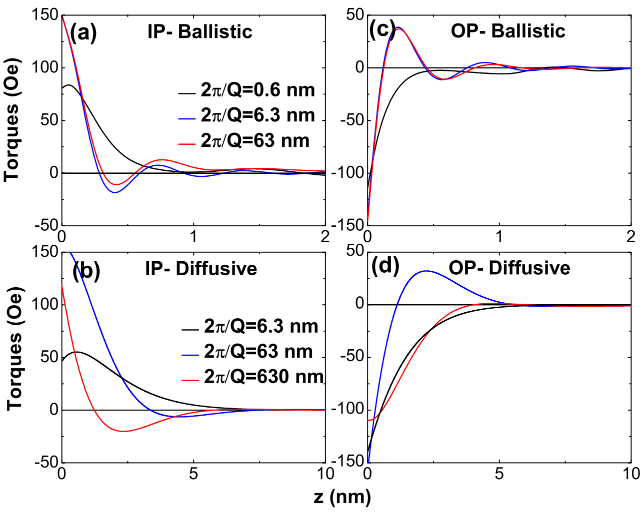

Figure 8. Intensity of spin transfer torques as a function of z: in-plane and out-of-plane components in ballistic ((a), (c)) and diffusive ((b), (d)) models.

averaging, as explained above (note that in the diffusive regime this decay follows an exponential oscillating function—see equation (20)).

However, the role of Q is not the same as for the current densities. Firstly, the decay length of the spin torque components does not depend on Q, contrary to the longitudinal and transverse currents. Secondly, the amplitudes of in-plane and out-of-plane components of spin torque are affected by Q in different ways.

Figure 9 displays the x-dependence of the in-plane and out-of-plane components of spin torque. The shape of the oscillation of both components becomes more asymmetric in the adiabatic regime. This asymmetry is even more important in the diffusive model; it indicates an increasing role of the longitudinal spin accumulation in the torques. As a matter of fact, in metallic spin valves, longitudinal spin accumulation is well known to affect the angular dependence of spin transfer torque, and makes it differ from a simple sine dependence [4].

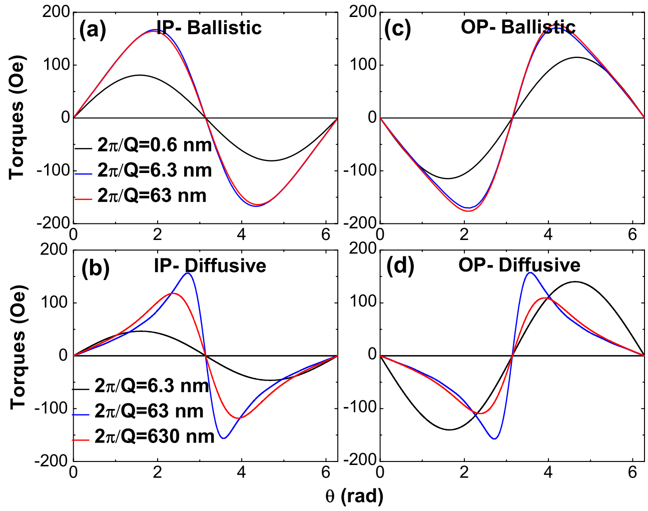

Figure 9. Intensity of spin transfer torques as a function of x: in-plane and out-of-plane components in ballistic ((a), (c)) and diffusive ((b), (d)) models.

When spin diffraction takes place, it reduces the influence of spin accumulation by averaging it out so that the angular dependence is found close to sine.

Furthermore, whereas the in-plane component decreases when increasing Q, the out-of-plane component is less dependent on Q. This interesting difference does not depend on the model (ballistic or diffusive). This characteristic may be explained as follows. Electron motion has a twofold character: non-adiabatic when it crosses the interface and adiabatic for its motion along the interface. It is known that the in-plane torque exists even for completely adiabatic motion but the out-of-plane torque appears only in the case of non-adiabatic motion. So there are two contributions into the in-plane torque (adiabatic and non-adiabatic) and only one (non-adiabatic) into the out-of-plane torque.

4. Discussion

In this section, we comment on a number of issues related to the experimental realization and observation of the phenomena described above.

The model presented in this article does not take into account the exchange interaction between the ferromagnet and the spin spiral. Indeed, such an interaction should locally modify the magnetic configuration across the interface, then reducing the effects predicted in this paper. However, usual systems would use a non-magnetic layer with large spin-diffusion length such as Cu or Ag to decouple the two magnetic layers without affecting the spin configuration of the itinerant electrons, which is close to the system described in this paper.

It was shown that currents and torques are modulated along the interface and that this modulation exists within some distance from the interface. This modulation may first lead to a specific dynamical behavior of the spin spiral structure: while the in-plane torque tends to align the spin spiral on the magnetization of the ferromagnet, the out-of-plane torque tends to orient the spin spiral out of the plane, then distorting the spin structure.

Furthermore, the current's and torque's modulations may lead to another interesting effect: if one uses a thin SS layer as a spacer between two ferromagnetic layers, the current-induced modulated spin torque may locally nucleate magnetization switching in the free layer. These 'hot spots' due to the SS layer should provide an efficient way to switch the magnetization of a ferromagnet at lower current density than in the case of a paramagnetic spacer.

Finally, our prediction that even smooth variation of magnetization leads to the existence of not only in-plane (adiabatic) but also out-of-plane (non-adiabatic) torque opens the possibility of an efficient mechanism for driving a domain wall by a current. Let us consider a sandwich consisting of a ferromagnet with a domain wall separated from a uniformly magnetized ferromagnetic layer by a paramagnetic spacer. Instead of injecting the current in the plane of the domain wall layer, we inject the current perpendicularly to the interfaces. This should lead to spin diffraction at the interface with the domain wall, as described in this paper, and to the combination between sizable in-plane and out-of-plane torques. As a consequence, the resulting distortion of the domain wall should lead to low critical current domain wall motion.

5. Conclusion

Ballistic and diffusive models of spin-dependent transport in a F/SS bilayer are proposed. We showed that due to the spatially varying local magnetization of the spin spiral, spin-dependent interferences occur leading to spin diffraction. This new type of diffraction is at the origin of the non-locality of the transport properties in the system and gives rise to complex physics inside the SS layer.

It was shown that currents and torques are modulated along the interface and that this modulation exists within some distance from the interface. Furthermore, in such systems, the amplitudes of in-plane and out-of-plane torque are expected to be of the same order, contrary to what is currently observed in metallic spin valves.

Finally, a number of experimental realizations are proposed in order to take advantage of both torque modulation and the amplitude of out-of-plane torque.

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References

[1] Slonczewski J C 1996 *J. Magn. Magn. Mater.* **159** L1–5
Berger L 1996 *Phys. Rev.* B **54** 9353

[2] Katine J A, Albert F J, Buhrman R A, Myers E B and Ralph D C 2000 *Phys. Rev. Lett.* **84** 3149
See also Fert A *et al* 2001 *Mater. Sci. Eng.* B **84** 1

[3] Stiles M D and Zangwill A 2002 *Phys. Rev.* B **66** 014407
Zhang S, Levy P M and Fert A 2002 *Phys. Rev. Lett.* **88** 236601

[4] Manchon A, Ryzhanova N, Strelkov N, Vedyayev A and Dieny B 2007 *J. Phys.: Condens. Matter* **19** 165212

[5] Slonczewski J C 2005 *Phys. Rev.* B **71** 024411
Theodonis I, Kioussis N, Kalitsov A, Chshiev M and Butler W H 2006 *Phys. Rev. Lett.* **97** 237205

[6] Manchon A, Ryzhanova N, Strelkov N, Vedyayev A, Chshiev M and Dieny B 2008 *J. Phys.: Condens. Matter* **20** 145208

[7] Ideka S, Hayakawa J, Lee Y M, Sasaki R, Meguro T, Matsukura F and Ohno H 2005 *Japan. J. Appl. Phys.* **44** L1442
Yuasa S, Fukushima A, Kubota H, Suzuki Y and Ando K 2006 *Appl. Phys. Lett.* **89** 042505

[8] Duine R A, Haney P M, Núñez A S and MacDonald A H 2007 *Phys. Rev.* B **75** 014433

[9] Ralph D C and Stiles M D 2008 *J. Magn. Magn. Mater.* **320** 1190

[10] Katine J A and Fullerton E E 2008 *J. Magn. Magn. Mater.* **320** 1217

[11] Stiles M D and Miltat J 2006 Spin transfer torque and dynamics *Spin Dynamics in Confined Magnetic Structures III* (*Springer Topics in Applied Physics* vol 101) ed B Hillebrands and A Thiaville (Berlin: Springer)

[12] Beach G S D, Tsoi M and Erskine J L 2008 *J. Magn. Magn. Mater.* **320** 1272

[13] Polianski M L and Brouwer P W 2004 *Phys. Rev. Lett.* **92** 026602

[14] Stiles M D, Xiao J and Zangwill A 2004 *Phys. Rev.* B **69** 054408

[15] Hayashi M, Thomas L, Bazaliy Ya B, Rettner C, Moriya R, Jiang X and Parkin S S P 2006 *Phys. Rev. Lett.* **96** 197207
Beach G S D, Knutson C, Nistor C, Tsoi M and Erskine J L 2006 *Phys. Rev. Lett.* **97** 057203

[16] Heyne L, Klaui M, Backes D, Moore T A, Krzyk S, Rudiger U, Heyderman L J, Rodriguez A F, Nolting F, Mentes T O, Nino M A, Locatelli A, Kirsch K and Mattheis R 2008 *Phys. Rev. Lett.* **100** 066603

[17] Berger L 1978 *J. Appl. Phys.* **49** 2156
Berger L 1984 *J. Appl. Phys.* **55** 1954

[18] Tatara G and Kohno H 2004 *Phys. Rev. Lett.* **92** 086601

[19] Thiaville A, Nakatani Y, Miltat J and Suzuki Y 2005 *Europhys. Lett.* **69** 990

[20] Barnes S E and Maekawa S 2005 *Phys. Rev. Lett.* **95** 107204

[21] Zhang S and Li Z 2004 *Phys. Rev. Lett.* **93** 127204

[22] Vanhaverbeke A and Viret M 2007 *Phys. Rev.* B **75** 024411

[23] Xiao J, Zangwill A and Stiles M D 2006 *Phys. Rev.* B **73** 054428

[24] Waintal X and Viret M 2004 *Europhys. Lett.* **65** 427

[25] Kohno H, Tatara G and Shibata J 2006 *J. Phys. Soc. Japan* **75** 113706

[26] Tserkovnyak Y, Skadsem H J, Brataas A and Bauer G 2006 *Phys. Rev.* B **74** 144405

[27] Piechon F and Thiaville A 2007 *Phys. Rev.* B **75** 174414

[28] Duine R A, Nunez A S, Sinova J and MacDonald A H 2007 *Phys. Rev.* B **75** 214420

[29] Nguyen A K, Skadsem H J and Brataas A 2007 *Phys. Rev. Lett.* **98** 146602

[30] Tatara G, Kohno H, Shibata J, Lemaho Y and Lee K-J 2007 *J. Phys. Soc. Japan* **76** 054707

[31] Duine R A, Nunez A S and MacDonald A H 2007 *Phys. Rev. Lett.* **98** 056605

[32] Stiles M D, Saslow W M, Donahue M J and Zangwill A 2007 *Phys. Rev.* B **75** 214423

[33] Acremann Y, Strachan J P, Chembrolu V, Andrews S D, Tyliszczak T, Katine J A, Carey M J, Clemens B M, Siegmann H C and Stohr J 2006 *Phys. Rev. Lett.* **96** 217202

[34] Nagaosa N 2006 *J. Phys. Soc. Japan* **75** 042001

[35] Panfilov A S 1999 *Low Temp. Phys.* **25** 432

[36] Capogna L, Mayr M, Horsch P, Raichle M, Kremer R K, Sofin M, Maljuk A, Jansen M and Keimer B 2005 *Phys. Rev.* B **71** 140402(R)

[37] Andrianov A V, Kosarev D I and Beskrovnyi A I 2000 *Phys. Rev.* B **62** 13844

[38] Keldysh L V 1965 *Sov. Phys.—JETP* **20** 1018

[39] Manchon A, Ryzhanova N, Vedyayev A and Dieny B 2008 *J. Appl. Phys.* **103** 07A721

[40] Zhang S, Levy P M and Fert A 2002 *Phys. Rev. Lett.* **88** 236601

[41] Shpiro A, Levy P M and Zhang S 2003 *Phys. Rev.* B **67** 104430

[42] Calvo M 1978 *Phys. Rev.* B **18** 5073

[43] Urazhdin S, Loloee R and Pratt W P Jr 2005 *Phys. Rev.* B **71** 100401(R)