Surface spin-transfer torque and spin-injection effective field in ferromagnetic junctions: Unified theory

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Abstract

We consider theoretically a current flowing perpendicular to interfaces of a spin-valve type ferromagnetic metallic junction. For the first time an effective approach is investigated to calculate a simultaneous action of the two current effects, namely, the nonequilibrium longitudinal spin injection and the transversal spin-transfer surface torque. Dispersion relation for fluctuations is derived and solved. Nonlinear problem is solved about steady state arising due to instability for a thick enough free layer.

1 Introduction

Great attention is attracted now to features of current flowing through ferromagnetic junctions of a spin-valve type, i.e. structures with contacting ferromagnetic thin layers, one of them having pinned and the other free spins. As experiments showed, current can influence substantially the free layer magnetic state of such junctions that leads to resistance jumps \cite{1}–\cite{3}, as well as microwave emission \cite{4}–\cite{6}.

Mechanisms of the current effect are not completely clear so far. A mechanism was proposed in \cite{7} of the current effect on the free ferromagnetic layer magnetization $\mathbf{M}$ due to injection of nonequilibrium longitudinal (i.e., collinear to $\mathbf{M}$) spins into the layer. Detailed theory of the mechanism was developed in \cite{8}–\cite{10}. The injection creates a nonequilibrium carrier spin polarization in the layer. This polarization, in its turn, contributes to the $sd$ exchange energy $U_{sd}(\mathbf{j})$ and to the corresponding $sd$ exchange effective field $\mathbf{H}_{sd}(\mathbf{j})$, which is dependent on the electric current density $\mathbf{j}$. For current densities exceeding some threshold, a reorientation first-order phase transition occurs and the magnetization vector $\mathbf{M}$ direction may change abruptly. Such a current induced

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magnetization reversal (or current induced switching) leads to resistance jumps, which corresponds to the experimental data \[1\]–\[3\].

On the other hand, another mechanism of current effect on the ferromagnetic layer state was proposed \[11, 12\] long before appearing of the experimental results \[1\]–\[6\]. According to the mechanism, the transverse (with respect to \(M\)) mobile electrons spin current is to vanish near the interface between two (pinned and free) contacting non-collinear ferromagnets. It occurs due to statistical velocity spread of the mobile electrons not because of the relaxation. The electrons interact with magnetic lattice (\(sd\) exchange interaction) and that is why a torque appears at the interface. This spin torque acts on the lattice and transfers the lost spin current to it. This spin-transfer torque may exceed the dissipation torque at large enough current densities. Then the initial state becomes unstable and the magnetization reversal occurs. As estimations show, such a switching mechanism can agree with experimental results \[1\]–\[6\] also.

In this connection, a very significant question arises: how are things going in the real experiments when both effects coexist, namely, longitudinal spin injection and corresponding current dependent effective field, and current dependent spin-transfer torque at the interface of the magnetic junction. Up to now, these effects were studied separately. Meanwhile, these effects do not only coexist, but influence each other also. Therefore, both effects are to be taken into account simultaneously, in scope of a unified theory, to understand better the experimental situation. Such a theory was developed for the first time in our previous preprint \[13\] and article \[14\]. The aim of the present paper is to elaborate a new effective approach for solving the problem. This approach allows not only get all the previously achievable results (see \[13, 14\]), but additionally to solve nonlinear problem about the steady state arising due to instability development. The approach based on the interesting feature of the effective field \(H_{sd}\) that was revealed in \[13, 14\]. This field is created due to the injection of nonequilibrium longitudinal mobile spins into the free layer volume and depends on the bulk parameters of the free layer. But a point of the field localization occurs near the interface. This feature suggests a possibility to include an effect of the field into boundary condition and shows the transversal spin flux becomes, in fact, discontinues at the interface. The simplification arises because of no singularity remains in the equations of motion.

2 Model

We will consider a spin-valve type magnetic junction with current flowing across the layer interfaces (CPP geometry). Ferromagnetic metal layer 1 has pinned orientations of the lattice and mobile electron spins. Another ferromagnetic metal layer 2 contacts with the layer 1 at a point \(x = 0\) and has free lattice\(^1\) magnetization \(M\) and mobile electron magnetization \(m\), so that the magnetizations direction can be changed by an external magnetic field \(H\) or spin-polarized current density \(j\). There is a very thin nonmagnetic spacer between the layers

\(^1\)We mean a lattice of magnetic ions considered in a continuum media approximation.
1 and 2, which will be considered as a geometrical plane \( x = 0 \). To close the electric circuit, a nonmagnetic metal layer 3 exists in the region \( x \geq L \).

External magnetic field \( \mathbf{H} \), lattice and mobile electron magnetizations of the layer 1, \( \mathbf{M}_1 \) and \( \mathbf{m}_1 \), as well as corresponding magnetizations of the layer 2, \( \mathbf{M} \) and \( \mathbf{m} \), are assumed lying in the junction plane \( x = 0 \). Vectors \( \mathbf{M}_1 \) and \( \mathbf{M} \) make an angle \( \chi \) of an arbitrary value under current \( j = 0 \). After the current is turned on electrons transfer from the layer 1 into the layer 2, that is \( j/e > 0 \). Then the electrons appear in a non-stationary quantum state in the layer 2 and “walk” between the spin subbands of this layer. It corresponds to precession of the mobile electron magnetization around the lattice magnetization \( \mathbf{M} \). We may introduce two components of the vector \( \mathbf{m} \) by an equality \( \mathbf{m} = \mathbf{m}_\parallel + \mathbf{m}_\perp \), where longitudinal component \( \mathbf{m}_\parallel \) is parallel to \( \mathbf{M} \) and transversal component \( \mathbf{m}_\perp \) is perpendicular to \( \mathbf{M} \). As it was first shown in [11, 12], angle of the precession decreases with coordinate \( x \) increasing and tends to zero at \( x \approx \lambda_F \), where \( \lambda_F \sim 1 \text{ nm} \) is an electron quantum wave length at the Fermi surface. This may be valid only because of electron velocity statistical spread and no relaxation processes are needed to provide such a behavior (more details may be found in the recent preprint [13] also). Therefore, both components \( \mathbf{m}_\parallel \) and \( \mathbf{m}_\perp \) exist in the region \( 0 \leq x \leq \lambda_F \), but only component \( \mathbf{m}_\parallel \) remains in the region \( x > \lambda_F \) (\( \mathbf{m}_\perp = 0 \)).

The region \( 0 \leq x \leq \lambda_F \) was introduced firstly by Slonczewski [11] and Berger [12] and further we will refer to it as “SB layer”. We follow the original works [11, 12] and assume ballistic regime of motion into the SB layer. The validity criteria for this assumption may be written as \( \lambda_F < l_p \), where momentum free pass length \( l_p \) may be estimated as \( l_p \sim 1–10 \text{ nm} \) for metals at room temperature.

3 Equations

We describe the motion of vector \( \mathbf{M} \) by means of the Landau–Lifshitz–Gilbert (LLG) equation [15]:

\[
\frac{\partial \mathbf{M}}{\partial t} + \gamma [\mathbf{M}, \mathbf{H}_{\text{eff}}] - \frac{\kappa}{M} [\mathbf{M}, \frac{\partial \mathbf{M}}{\partial t}] = 0,
\]  

(1)

where \( \gamma \) is the gyromagnetic ratio, \( \kappa \) is dimensionless damping constant (\( 0 < \kappa \ll 1 \)),

\[
\mathbf{H}_{\text{eff}} = \mathbf{H} + \mathbf{H}_a + A \frac{\partial^2 \mathbf{M}}{\partial x^2} + \mathbf{H}_d + \mathbf{H}_{sd}
\]  

(2)

is effective field, \( \mathbf{H}_a = \beta \mathbf{M} \) is anisotropy field, \( A \) is the intralattice inhomogeneous exchange constant, \( \mathbf{H}_d \) is demagnetization field, \( \mathbf{H}_{sd} \) is \( sd \) exchange effective field. The latter takes the form [15]

\[
\mathbf{H}_{sd}(x, t) = -\frac{\delta U_{sd}}{\delta \mathbf{M}(x, t)},
\]  

(3)
where $\delta(\ldots)/\delta \mathbf{M}(x, t)$ is a variational derivative, and $U_{sd}$ is sd exchange energy,

$$U_{sd} = -\alpha \int_0^L \mathbf{m}(x', t) \mathbf{M}(x', t) \, dx',$$  \hspace{1cm} (4)

parameter $\alpha$ being a dimensionless sd exchange constant (typical estimation is $\alpha \sim 10^4 - 10^6$ \[8\]). Due to the last term in (2), the motions of $\mathbf{M}$ and $\mathbf{m}$ vectors appear to be coupled.

We describe the motion of vector $\mathbf{m}$ by means of continuity equation (e.g., see [16, 17]):

$$\frac{\partial \mathbf{m}}{\partial t} + \frac{\partial J}{\partial x} + \gamma \alpha [\mathbf{m}, \mathbf{M}] + \frac{\Delta \mathbf{m}}{\tau} = 0,$$  \hspace{1cm} (5)

where $\tau$ is a time of relaxation to the local equilibrium value $\bar{\mathbf{m}} \equiv \bar{\mathbf{m}} \hat{\mathbf{M}}$, $\Delta \mathbf{m} \equiv \mathbf{m} - \bar{\mathbf{m}} = \Delta \mathbf{m} \hat{\mathbf{M}}$ and $\hat{\mathbf{M}} \equiv \mathbf{M}/M$ is the unit vector, $\mathbf{J}$ is mobile electron magnetization flux density.

4 Vector boundary conditions

The SB layer may influence as a special boundary condition if the thickness of the layer $2$ is large enough, namely, $L \gg \lambda_F$. Moreover, the value of $\lambda_F$ should be assumed the smallest one among the other lengths in the system (e.g., $\sqrt{A} \gg \lambda_F$ and $l \gg \lambda_F$). We may sum the Eqs. (1) and (5) and so consider the continuity equation for the total magnetization vector $\mathbf{M} + \mathbf{m}$. Let us integrate the sum equation over the region $0 \leq x \leq \lambda_F$, the integration being denoted as

$$\int_{-\varepsilon}^{\lambda_F + \varepsilon} (\ldots) \, dx \equiv \langle \ldots \rangle.$$  \hspace{1cm} (6)

Then the following three types of summands appear. The first type summands are proportional to small length $\lambda_F$ and may be neglected for actual values of parameters (fields: $H$, $H_a$, $H_d$; relaxation parameters: $\kappa$, $\tau$; and frequencies $\omega$ in a microwave region or lower). The second type of summands may be presented as derivatives with respect to coordinate $x$. These are spin fluxes of mobile electron and lattice spins. After integration in (6), they do not depend on $\lambda_F$ and remain finite in the limit $\lambda_F \to 0$. Finally, the third summand appears, which is proportional to a singular $\delta$-function inside the SB layer. This summand arises due to nonequilibrium longitudinal spin injection by current. After integration in (6), it does not depend on $\lambda_F$ also.

Going the way indicated, we obtain the following vector boundary condition:

$$\mathbf{J}(\lambda_F + \varepsilon) - \mathbf{J}(-\varepsilon) + \mathbf{J}_M(\lambda_F + \varepsilon) - \mathbf{J}_M(-\varepsilon) + \langle G_{sd} \rangle = 0.$$  \hspace{1cm} (7)

The mobile electron spin flux $\mathbf{J}$ in the condition (7) originates from the equation (5). The lattice magnetization flux $\mathbf{J}_M$ follows from the intralattice exchange effective field of Eq. (1). The following expression being valid:

$$\mathbf{J}_M = a \hat{\mathbf{M}} \frac{\partial \mathbf{M}}{\partial x}$$  \hspace{1cm} (8)
and \( a = \gamma AM \) is a lattice magnetization diffusion constant.

To calculate the last term \( \langle G_{sd} \rangle \) in (4), we take the sum equation for \( M + m \) and get
\[
\langle G_{sd} \rangle \equiv \gamma \langle [M, H_{sd}] \rangle + \gamma \alpha \langle [m, M] \rangle = \gamma \alpha \left( \langle M(x), \int_0^L (\delta \Delta m / \delta \hat{M}) \, dx \rangle \right),
\]
where formulae (3) and (4) are used under the calculations. The condition (7) remains valid at \( x = L \), if we replace: \( \lambda F + \varepsilon \rightarrow L + \varepsilon, \quad -\varepsilon \rightarrow L - \varepsilon \) and put \( \langle G_{sd} \rangle \rightarrow 0 \).

5 Nonequilibrium mobile electron spin flux and magnetization

To make the system of equations (1)–(5) and boundary condition (7) be well defined, we should express all the quantities involved via the vectors \( M \) and \( m \).

As it may be seen from (7), we should consider the spin flux \( J \) outside the SB layer only. In the region \( x > \lambda F \) electrons occupy subbands having spins parallel to \( M(\uparrow) \) and antiparallel to \( M(\downarrow) \). Therefore, the following representations are valid:
\[
m = \mu_B (n^\uparrow - n^\downarrow) \hat{M} \equiv m \hat{M}, \quad J = (\mu_B / e) (j^\uparrow - j^\downarrow) \hat{M},
\]
where \( \mu_B \) is the Bohr magneton, \( n^\uparrow, n^\downarrow \) and \( j^\uparrow, j^\downarrow \) are partial electron densities and current densities in the spin subbands, respectively. The total electron density \( n = n^\uparrow + n^\downarrow \) and \( j = j^\uparrow + j^\downarrow \) do not depend on \( x \) and \( t \) because of local neutrality conditions in metal and one-dimensional geometry of our model. We use the well-known ”diffusion-drift” formula for \( j^\uparrow, j^\downarrow \) to express the partial currents via \( j, n \) and \( m \).

The corresponding calculations are direct and were done completely in [8]. The final result for \( x > \lambda F \) may be written as follows
\[
J = \left( \frac{\mu_B}{e} Q j - \hat{M} \frac{\partial m}{\partial x} \right) \hat{M},
\]
where \( Q = (\sigma^\uparrow - \sigma^\downarrow) / (\sigma^\uparrow + \sigma^\downarrow) \) may be understood as an equilibrium conductivity spin polarization parameter with \( \sigma^\uparrow, \sigma^\downarrow \equiv \mu^\uparrow, \mu^\downarrow, \) and \( \hat{D} = (\sigma^\uparrow D^\uparrow + \sigma^\downarrow D^\downarrow) / (\sigma^\uparrow + \sigma^\downarrow) \) is the effective spin diffusion constant, quantities \( \mu^\uparrow, \mu^\downarrow \) and \( D^\uparrow, D^\downarrow \) being partial mobilities and diffusion constant, respectively. To obtain (10), an additional assumption should be made [8], namely, \( j / j_D \ll 1 \), where \( j_D = enl / \tau \) is a characteristic current density in the layer 2. With typical parameter values, \( n \sim 10^{22} \, \text{cm}^{-3}, \quad l \sim 3 \times 10^{-6} \, \text{cm}, \quad \tau \sim 3 \times 10^{-13} \, \text{s} \), we get \( j_D \sim 1.6 \times 10^{10} \, \text{A/cm}^2 \).

The lattice in the layer 1 is pinned. Therefore, the magnetization flux \( J(-\varepsilon) \) may be written similar to (10) but without the spatial derivative. We introduce \( \hat{M}_1 = M_1 / M_1 \) and have
\[
J(-\varepsilon) = \frac{\mu_B}{e} Q_1 j \hat{M}_1.
\]

The latter flux has longitudinal and transversal components, which are:
\[
J_{||}(-\varepsilon) = \frac{\mu_B}{e} Q_1 j \left( \hat{M}_1 \hat{M}(\varepsilon) \right) \hat{M}(\varepsilon)
\]
and
\[ J_\perp (-\varepsilon) = \frac{\mu_B}{e} Q_1 j \left[ \mathbf{\hat{M}} (\varepsilon), \left[ \mathbf{\hat{M}}_1, \mathbf{\hat{M}} (\varepsilon) \right] \right]. \] (12)

We should calculate now the nonequilibrium magnetization for \( x > \lambda_F \). In the region an effective frequency of the motion \( \omega \) is determined by the precession in relatively small fields: \( \omega \sim \gamma H, \gamma H_a, \gamma H_a \ll \gamma \alpha M \equiv \omega_{sd} \). We assume the conditions \( \omega \tau \ll 1 \) and \( \omega_{sd} \tau \gg 1 \) are valid for typical values \( \tau \sim 3 \times 10^{-13} \) s, \( \alpha \sim 2 \times 10^4, M \sim 10^3 \) G. It allows neglect the time derivative in (5) and substitute the formula (10) for flux \( J \).

Then the equation reduces to
\[ \frac{\partial^2 m}{\partial x^2} - \frac{\Delta m}{l^2} = 0, \] (13)

where \( l = \sqrt{D\tau} \) is the spin diffusion length. Analogous relations may be written, of course, for the nonmagnetic layer 3 with the following modifications: \( J_3 = 0 \) and \( Q_3 = 0 \).

The next step is to find the solution of (13) that satisfies the conditions: 1) \( J_\parallel (-\varepsilon) = J (\lambda_F + \varepsilon) \), 2) \( J (L - \varepsilon) = J (L + \varepsilon) \), 3) continuity of subband chemical potential difference at the interface point \( x = L \) (see [13, 14] for details). This solution takes the form
\[ \Delta m(x) = \frac{j}{jD} \frac{\mu_B n}{\sinh \lambda + \nu \cosh \lambda} \{ Q \cosh \xi + [Q_1 (\mathbf{\hat{M}}_1 \mathbf{\hat{M}} (\varepsilon)) - Q] \times \cosh(\lambda - \xi) + \nu \sinh(\lambda - \xi) \}, \] (14)

where \( \lambda = L/l, \xi = x/l \) and parameter \( \nu \) characterizes the influence of the layer 3 (typically \( \nu \sim 1 \), see [13, 14] for details). We substitute (14) into expression (10) and use the variational derivative \( \delta (\mathbf{\hat{M}}_1 \mathbf{M}(\varepsilon))/\delta \mathbf{M}(x) = \mathbf{\hat{M}}_1 \delta (x - \varepsilon) \). Then we get
\[ \langle G_{sd} \rangle = \gamma \alpha \mu_B n Q_1 l jD \left( 1 - \frac{\nu}{\sinh \lambda + \nu \cosh \lambda} \right) \left[ \mathbf{M}(\varepsilon), \mathbf{\hat{M}}_1 \right], \] (15)

because of \( \langle \delta (x - \varepsilon) \rangle = 1 \).

Turning back to condition (7), we see all the terms are now explicitly dependent on the vectors \( \mathbf{M} \) and \( \mathbf{m} \). The latter vectors satisfy the equation of motion (11–15). Therefore, the way opens now for solving our problem directly.

### 6 Initial steady state and fluctuation instability

To illustrate the possibilities of solution, we assume simple situation where field \( H \) is applied along the positive direction of \( z \) axis, \( H_a \) is parallel to this axis too and the angle \( \chi = \pi \). In the situation all the equations and boundary conditions are satisfied for the following steady state magnetization vector \( \mathbf{\hat{M}} \): \( \mathbf{\hat{M}}_x = \mathbf{\hat{M}}_y = 0, \mathbf{\hat{M}}_z = 1 \), and \( \mathbf{\hat{M}}_1 = -\mathbf{\hat{z}} \).
We introduce the fluctuations $\hat{\Delta M}$ by the equality: $\hat{\Delta M} = \hat{z} + \Delta \hat{M}$ with $\hat{\Delta M} \hat{\Delta M} = 0$ and $\Delta \hat{M} \equiv |\Delta \hat{M}| \ll 1$. We linearize the equations LLG and condition (7) with respect to $\Delta \hat{M}$ and lay $H_d = -4\pi M \Delta M \hat{x}$. Then we get the equations for fluctuations $\Delta \hat{M}_{x,y} \sim \exp(iqx - i\omega t)$:

\[
\frac{\partial^2 \Delta \hat{M}_x}{\partial x^2} - \left(\frac{\Omega_x - i\kappa \omega}{a}\right) \Delta \hat{M}_x - \frac{i\omega}{a} \Delta \hat{M}_y = 0,
\]

\[
\frac{\partial^2 \Delta \hat{M}_y}{\partial x^2} - \left(\frac{\Omega_y - i\kappa \omega}{a}\right) \Delta \hat{M}_y + \frac{i\omega}{a} \Delta \hat{M}_x = 0,
\]

(16)

where two characteristic frequencies appear: $\Omega_x = \gamma (H + H_a + 4\pi M)$, $\Omega_y = \gamma (H + H_a)$ and transversal part of the vector boundary condition (7) for fluctuations at $x = 0$ takes the form:

\[
\frac{\partial \Delta \hat{M}_x}{\partial x} = k \Delta \hat{M}_y - p \Delta \hat{M}_x, \tag{17}
\]

\[
\frac{\partial \Delta \hat{M}_y}{\partial x} = -k \Delta \hat{M}_x - p \Delta \hat{M}_y.
\]

Parameters $k = (\mu_B j Q_1)/(e\gamma M^2 A)$ and $p = \alpha \mu_B j Q_1 \tau |1 - \nu/(\sinh \lambda + \nu \cosh \lambda)|/(eAM)$ are proportional to current density $j$ and describe respectively the effects of spin-transfer torque and spin injection.

Transversal terms of the boundary condition (7) at $x = L$ become:

\[
\frac{\partial \Delta \hat{M}_x}{\partial x} = 0, \quad \frac{\partial \Delta \hat{M}_y}{\partial x} = 0. \tag{18}
\]

Solving the standard boundary problem (16)–(18), we come to the characteristic wave number

\[
q^2 = \frac{1}{2a} \left[\Omega_x + \Omega_y - 2i\kappa \omega \pm \sqrt{(\Omega_x - \Omega_y)^2 + 4\omega^2}\right] \tag{19}
\]

and dispersion relation

\[
qL \tan qL = -pL \pm ikL \tag{20}
\]

for determining the spin-wave spectrum and decrement $\omega(q)$.

7 Steady state arising due to instability development

The dispersion relation (20) is the same as in our previous works [13, 14] and therefore all the results about the instability thresholds and current dependent spectrum remain valid. But in our new approach, contrary to [13, 14], we have much more simple equations of motion, which do not contain any singular terms.
It allows solve more complicated problems now. One of such problem is: what will be a steady state the instability development leads to?

We confine ourselves to the case of large enough thickness \( L > l \), \( \sqrt{A} \) corresponding to the injection mechanism dominated. In such case, the only polar angle \( \theta(x) \) exists that determines the stationary magnetization \( \hat{M} \) orientation and the problem reduces to solving of nonlinear equation

\[
\delta^2 \frac{d^2 \theta}{dx^2} - \sin \theta \cos \theta - \frac{H}{H_a} \sin \theta = 0,
\]

where \( \delta = \sqrt{A/\beta} \). Transversal vector boundary conditions in the polar coordinates are

\[
\frac{d\theta}{dx} \bigg|_{x=0} = 0 \quad \text{and} \quad \frac{d\theta}{dx} \bigg|_{x=L} = 0.
\]

Equation (21) describes a "classical" particle moving in a potential field. As it is well known, such type of equation may be integrated easily. At \( p\delta < (1 + H/H_a)^{1/2} \), there is no instability and only one solution \( \theta = 0 \) exists. Under the opposite condition, the instability occurs and a nonzero solution appears which describes the new steady state orientation angle \( \theta(x) \) in the layer 2:

\[
\theta(x) = 2 \arccos \left\{ \left( \frac{h j}{j_0} \tanh(x/h\delta) + 1 \right)(1 - h^2)^{1/2} \times \left[ \left( \frac{h j}{j_0} + \tanh(x/h\delta) \right)^2 - h^2 \left( \frac{h j}{j_0} \tanh(x/h\delta) + 1 \right)^2 \right]^{-1/2} \right\},
\]

where \( h = (1 + H/H_a)^{-1/2} \), \( j_0 = e\delta H_a/\alpha \mu_B \tau Q_1 \).

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