Possibility of exchange switching ferromagnet-antiferromagnet junctions

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Abstract

Current flowing is studied in magnetic junctions consisting of a ferromagnetic metal (FM), antiferromagnetic conductor (AFM) and a nonmagnetic metal closing the electric circuit. The FM layer with high anisotropy and pinned spins of the magnetic atoms in the lattice acts as a spin injector relative to the AFM layer. To obtain resulting magnetization in the AFM layer, magnetic field is applied, which may be varied to control the magnetization. The spin-polarized current from the FM layer creates a torque and makes it possible to switch the magnetization. A possibility is shown to lower the threshold current density by the orders of magnitude by means of the magnetic field.

1 Introduction

Antiferromagnetic layers are used conventionally in spintronics to introduce additional magnetic anisotropy and pin the magnetization vector direction in a magnetic junction. Such an effect is based on the unidirectional anisotropy phenomenon [1]. Recently, interest has been revived to studying FM-AFM junctions with special attention to the current effect on the interface processes [2][3][4]. Besides, experimental works have been revealed [5][6] in which the magnetoresistive effect has been observed in FM-AFM structures with a point contact between the layers under high current densities (up to $10^9–10^{10}$ A/cm$^2$); the effect is due, apparently, to the $sd$ exchange interaction. The similar effect is observed in FM junctions with broad applications [7][8]. Suppositions have been voiced that such an effect in FM-AFM or AFM-AFM structures will have interesting features because of absence of demagnetization, in particular. Finally, opinions have been stated that the $sd$ exchange may lead to the spin transfer from the conduction electrons to the lattice, as in FM junctions, and cause

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instability with the magnetization switching. It has been supposed that nonlinear effects and the current-driven electron spin polarization may promote such an effect [9, 10, 11, 12].

Thus, the possibility of the current-driven exchange switching FM-AFM structures remains, apparently, the most incomprehensible in theory and experiment. The idea of the electron spin transfer to the lattice by the spin-polarized current has been put forward in famous works [13, 14]. The idea turned out to be very fruitful, so that it is interesting to know whether it is applicable to the FM-AFM structures. An attempt to clear the matter up is the goal of the present paper.

When we say about the spin transfer “to the lattice”, it is necessary to explain which of two (at least!) AFM sublattices spins transfer to. The basic works [13, 14], naturally, do not answer the question. A physical picture of the spin transfer, according to [14], consists in precession of the current polarization vector \( p \) around the lattice magnetization vector \( M \) and gradual decrease in the precession angle with motion along the current because of the statistical spreading of the electron velocities. Such a picture looks general enough and must remain for any model of the magnetic sublattices, including AFM. If such a picture is applied, then a torque \( T \) similar to that in [13] is to be added to the right-hand side of the Landau-Lifshitz (LL) equation for the AFM magnetization. Note that the torque \( T \) reveals in LL equation for the total magnetization \( M \), rather than the each of the sublattices \( M_1 \), \( M_2 \).

A formal derivation of the expression for the \( T \) moment was proposed by us in Refs. [15, 16, 17]. It should have in mind, however, that the torque originates in a very thin layer \( \lambda_F \sim 1-2 \) nm thick at the FM-AFM interface; such a length coincides with the precession damping lengths of the \( M \) and \( p \) vectors [13, 14]. So \( T \) moment exists only in that thin layer, not the whole range of the current interaction with the magnetic lattice. In accordance with Refs. [15, 16, 17], we take the \( T \) effect into account by means of the boundary conditions.

2 The structure and equations of motion

2.1 Conduction electrons in the AFM layer

The structure in study is shown schematically in Fig. 1. The sizes along 0\( y \) and 0\( z \) axes are large, so that demagnetizaion is absent in that dimensions. For simplicity, we assume the FM layer to be pinned. This layer injects spin polarized along \( p = \hat{M}_\text{FM} \) (here and below the hat over the vector indicates a unit vector). The main processes occur in the AFM layer. Let us consider the equations of motion for that layer.

The electron magnetization \( \mathbf{m}(x, t) \) obeys the continuity equation

\[
\frac{\partial \mathbf{m}_i}{\partial t} + \nabla_k J_{ik} + g \alpha_{sd} [\mathbf{m} \times \mathbf{M}]_i - \frac{\mathbf{m}_i - \bar{\mathbf{m}}_i}{\tau} = 0,
\]

where \( g \) is the gyromagnetic ratio, \( \alpha_{sd} \) is the \( sd \) exchange constant,

\[
J_{ik} \equiv \mathbf{M}_i \frac{\mu_B}{e} (j_k^+ - j_k^-) = \mathbf{M}_i \frac{\mu_B}{e} \left\{ P_{jk} - enD \frac{\partial P}{\partial x_k} \right\}
\]
Figure 1: The structure in study. The FM layer is a ferromagnetic metal with $M_{FM}$ magnetization pinned along $z$ axis; the AFM layer with $L$ thickness is an easy-plane-type antiferromagnetic conductor with sublattice magnetization vectors canted by $\theta$ angle in an applied field $H$ parallel to the anisotropy axis $n||z$: $|M_1| = |M_2| = M_0$. The electric current is closed by a nonmagnetic conductor NM, so that electron flux $j/e$ flows through all the layers.

is the electron spin current, $j^\pm_k$ are the partial electric current densities in the spin subbands, $j_k = j^+_k + j^-_k$ is the total current density, $D$ is the spin diffusion constant (it is assumed to be the same in both subbands), $m_i = m\tilde{M}_i$, $m = \mu_B nP$, $n$ is the electron density, $P$ is the spin polarization, $\bar{m}_i = \mu_B \bar{n} \bar{P} \tilde{M}_i$, $\bar{P} = \frac{\epsilon_F^{3/2} - (\epsilon_F - \epsilon_{sd})^{3/2}}{\epsilon_F^{3/2} + (\epsilon_F - \epsilon_{sd})^{3/2}}$ (3) is the equilibrium spin polarization, $\epsilon_{sd}$ is the spin subband energy shift, $\epsilon_F$ is the Fermi energy relative to the bottom of the lower subband.

The ways of the Eq. (1) transformation and simplification were discussed in detail in Refs. [15, 16, 17], so that we present only some necessary results here.

We consider only slow enough processes due to the vector $\mathbf{M}$ precession with $\omega \ll 1/\tau$ frequencies. A typical spin relaxation time in AFM is $\tau \sim 10^{-12}$ s. Therefore, the electrons have time to follow the precession, so that the time derivative may be omitted. In derivation of Eq. (2), the current density may be assumed to be low, namely, $j/j_D \ll 1$, where $j_D = enl/\tau \sim 2 \times 10^{10}$ A/cm$^2$. This condition is assumed to be fulfilled. As a result, Eq. (1) is reduced (see [15, 16, 17]) to

$$\frac{\partial^2 m}{\partial x^2} + \frac{m - \bar{m}}{\tilde{t}^2} = 0.$$ (4)

At the layer boundaries $x = 0$ and $x = L$ the solution must obey the continuity conditions for the spin currents and the differences of the Fermi
quasilevels \((\epsilon_F^+ - \epsilon_F^-)\). Then the solution of Eq. (4) takes the form (see [17])

\[
\Delta m \equiv \bar{m} - m = \mu_B n Q_{FM} \frac{j}{j_D} \frac{\nu \cos \chi}{\nu^* + \cos^2 \chi},
\]

where \(Q_{FM} = (j_{FM}^+ - j_{FM}^-)/j\) is the current polarization in the FM layer,
\(\cos \chi = M_{FM} \cdot \hat{M}(0), \nu = Z_{FM}/Z_{AFM}, \nu^* = Z_{FM}/Z_{NM} + \lambda Z_{FM}/Z_{AFM}, \lambda = L/l \ll 1, Z_q = \frac{l_q \rho_q}{1 - Q_q^2}\) are the spin resistances of the layers, \(l_q, \rho_q, Q_q\) being the spin diffusion length, the resistivity and the current polarization, respectively, in the \(q = FM, AFM, NM\) layers. Equation (5) describes distribution of the electron magnetization across the AFM layer.

2.2 Magnetic sublattices in the AFM layer

We derive the equations of motion for the \(M_1\) and \(M_2\) sublattice magnetizations by means of the procedure described in Ref. [18]. The desired equations take the following general form:

\[
\frac{\partial M_i}{\partial t} = g_i [M_i \times \hat{H}_i] + \mathbf{R}_i \quad (i = 1, 2),
\]

where the effective fields are determined with variational derivatives of the AFM energy \(W\):

\[
\hat{H}_i(x, t) = -\frac{\partial W}{\partial M_i(x, t)} \quad (i = 1, 2);
\]

the energy is \(W = \int w(M_1, M_2) d^3x\), the energy density is \(w = w_{ex} + w_{sd} + w_a + w_H\),

\[
w_{ex} = \frac{1}{4} \delta (M^2 - L^2) + \frac{1}{4} (\alpha + \alpha') \left( \frac{\partial M}{\partial x} \right)^2 + \frac{1}{4} (\alpha - \alpha') \left( \frac{\partial L}{\partial x} \right)^2
\]

is the sublattice exchange interaction energy density, \(w_{sd} = -\alpha_{sd} \mathbf{n} \cdot \mathbf{M}\) is the \(sd\) exchange interaction energy density,

\[
w_a = -\frac{\beta}{4} ((\mathbf{n} \cdot \mathbf{M})^2 + (\mathbf{n} \cdot \mathbf{L})^2) - \frac{\beta'}{4} ((\mathbf{n} \cdot \mathbf{M})^2 - (\mathbf{n} \cdot \mathbf{L})^2)
\]

is the anisotropy energy density (the anisotropy axis \(\mathbf{n}\) is shown in Fig. 1), \(w_H = -\mathbf{H} \cdot \mathbf{M}\) is the Zeeman energy density in the external field.

The relaxation terms may be taken in the Gilbert form [11]

\[
\mathbf{R}_i = \frac{\alpha G}{M_0} \left[ \mathbf{M}_i \times \frac{\partial \mathbf{M}_i}{\partial t} \right] \quad (i = 1, 2),
\]

\(|M_1| = |M_2| = M_0\).

It is convenient to introduce total magnetization \(\mathbf{M} = M_1 + M_2\) and antiferromagnetism vector \(\mathbf{L} = M_1 - M_2\) instead of sublattice vectors \(M_1, M_2\). We assume sublattices to be equivalent. Therefore, a nonzero magnetization \(\mathbf{M} \neq 0\) appears only under applied field \(\mathbf{H}\) parallel to \(\mathbf{n}\) axis, as shown in Fig. 1. The \(\theta\) angle between the sublattices and \(\mathbf{n}\) axis
is determined by $H$ field with $\cos \theta = H/H_{ex}$, $H \leq H_{ex} = 2\delta M_0$. The $\theta$ angle is related with the squares of the vector lengths:

$$M^2 - L^2 = 4M_0^2 \cos \theta, \quad M^2 + L^2 = 4M_0^2.$$  \hspace{1cm} (8)

It is seen that the angle is conserved with motion if the vector lengths conserve. The magnetization is assumed to be low enough, $M \equiv |M| \ll M_0$.

For simplicity, we neglect demagnetization and weak relativistic non-collinear Dzyaloshinskii interaction [18].

The equations for $M$ and $L$ vectors take the form

$$\frac{\partial M}{\partial t} = g \left[ M \times H_{FM}^{\text{eff}} \right] + g \left[ L \times H_{AFM}^{\text{eff}} \right] + R_1 + R_2,$$  \hspace{1cm} (9)

$$\frac{\partial L}{\partial t} = g\delta \left[ M \times L \right] + g \left[ M \times H_{AFM}^{\text{eff}} \right] + g \left[ L \times H_{FM}^{\text{eff}} \right] + R_1 - R_2.$$  \hspace{1cm} (10)

New effective fields appear here:

$$H_{FM}^{\text{eff}} = \frac{1}{2} (\alpha + \alpha') \frac{\partial^2 M}{\partial x^2} + H + H_{FM}^{\text{FM}} + H_{sd},$$  \hspace{1cm} (11)

$$H_{AFM}^{\text{eff}} = -\frac{1}{2} (\alpha - \alpha') \frac{\partial^2 L}{\partial x^2} + H_{AFM}^{AFM},$$  \hspace{1cm} (12)

with anisotropy fields

$$H_{a}^{FM} = \frac{1}{2} (\beta + \beta') (n \cdot M) n, \quad H_{a}^{AFM} = \frac{1}{2} (\beta - \beta') (n \cdot L) n.$$  \hspace{1cm} (13)

The relaxation terms contain vector products $[M \times \partial M/\partial t]$, $[L \times \partial L/\partial t]$, and $[M \times \partial L/\partial t]$, $[L \times \partial M/\partial t]$. By transforming these products by means of Eqs. (9) and (10) with Eqs. (11) and (12) taking into account, we find that the products are of the same order of magnitude, so that we may put $R_1 \pm R_2 \equiv R = \frac{2\alpha'G}{M_0} \left[ M \times \frac{\partial M}{\partial t} \right]$ and use the Gilbert formula with $\alpha' \sim \alpha_G$. The $sd$ exchange effective field is

$$H_{sd}(x, t) = -\frac{\partial w_{sd}}{\partial M(x, t)} = H_{eq}(M)\dot{M} + \Delta H_{sd}.$$  \hspace{1cm} (13)

The first term in the right-hand side of Eq. (13) is the equilibrium contribution that directed along $M$ vector and vanishes from Eq. (9). The second term is due to the current. It determines with the nonequilibrium part $\Delta m$ of the electron magnetization. This term coincides with the one calculated in Refs. [16, 17]. So we may use those results here. We have

$$\Delta H_{sd} = \dot{M} \cdot H_{sd} \delta(x - 0),$$  \hspace{1cm} (14)

$$h_{sd} = (\alpha + \alpha')\mu_B n Q_{FM} \lambda \frac{j}{j_D} (\nu^* - \cos^2 \chi) (\nu^* + \cos^2 \chi)^2.$$  \hspace{1cm} (15)

According to Eqs. (11) and (12), the processes in the AFM layer differ substantially from those in the FM layer. In particular, even the lengths of the $M$ and $L$ vectors do not conserve in motion. The possibility of the electron spin transfer to the lattice similar to that in Refs. [13, 14] is to be discussed anew. At the same time, the electron interaction with the
lattice is described with $H_{sd}$ field and Eqs. (14)–(12), as in the FM layer. Therefore, it is necessary only, that two magnetic sublattices precess in synchronism as in ferromagnets.

Such a synchronism is attained easier if the sublattices are the same. Let the following conditions are fulfilled:

$$\alpha + \alpha' \gg |\alpha - \alpha'| \to 0 \quad \text{and} \quad \beta + \beta' \gg |\beta - \beta'| \to 0.$$  

The uniform exchange dominates in the equations because of the following estimates [18]:

$$\delta \sim \alpha/a^2 \sim 10^4 \gg 1, \quad \alpha \sim 10^{-12} \text{cm}^2, \quad a \sim 10^{-8} \text{cm}.$$  

As a result, Eqs. (9) and (10) are reduced to

$$\frac{\partial M}{\partial t} = g[M \times H_{sd}^{FM}] + R, \quad \frac{\partial L}{\partial t} = g\delta [M \times L]. \quad (16)$$

These equations describe precession of the magnetization vector and the precession-driven forced oscillation of the antiferromagnetism vector. Note that the vector squares $M^2$ and $L^2$ conserve in accordance with (16). So it follows from (8) that the sublattice canting angle $\theta$ conserves also in motion that is substantial in our model.

3 Spin currents and boundary conditions

Let us derive expressions for the spin currents in the lattice. According to Refs. [15, 16, 17], these expressions follow from the equations of motion. However, the second of Eqs. (16) does not contain nonlocal or singular terms. Therefore, the currents may be derived from the first equation only. We obtain from (16)

$$g\alpha \left[ M \times \frac{\partial^2 M}{\partial x^2} \right] = a \frac{\partial}{\partial x} \left[ M \times \frac{\partial M}{\partial x} \right] \equiv \frac{\partial J_M}{\partial x}, \quad (17)$$

with a magnetization flux

$$J_M = a \left[ M \times \frac{\partial M}{\partial x} \right], \quad (18)$$

where $a = g\alpha M$ has the meaning of a magnetization diffusion constant.

Let us consider the singular term due to $sd$ exchange in the first of Eqs. (16), namely, $g[M \times H_{sd}]$. By substituting (14) and (15) to it, we have

$$g[M \times H_{sd}] = \frac{\partial J_{sd}}{\partial x}, \quad (19)$$

where a $sd$ exchange flux appears

$$J_{sd} = gh_{sd} \left[ M(0) \times \bar{M}_{FM} \right] \theta(x - 0), \quad (20)$$

$\theta(x)$ being the Heaviside step function.

Now let us return to the electron spin current $J_{ik}$ (see Eq. (2)). Here $k = x$, and the current exists in two magnetic layers, FM and AFM. However, the spatial gradient of the spin polarization $P$ is absent in the FM layer, so that only the first term proportional to the current density $j$
remains in Eq. (2). The spin current itself contains spins collinear to the FM layer magnetization, i.e.,
\[ J_{\text{FM}} \equiv J(-0) = \hat{M}_{\text{FM}} \mu_B e P_{\text{FM}}. \] (21)

The spin current continuity conditions at \( x = 0 \) and \( x = L \) interfaces give the boundary conditions necessary to solve Eqs. (16). As the lattice is pinned in the FM layer, we have \( J_{\text{FM}} \equiv J_M(-0) = 0 \) and \( J_{\text{sd}} \equiv J_{\text{sd}}(-0) = 0 \). Then we obtain for \( x = 0 \)
\[ J(+0) - J(-0) + J_M(+0) + J_{\text{sd}}(+0) = 0. \] (22)

Let us project Eq. (22) to the \( \hat{M} \) direction and to plane perpendicular to that direction. Two equations are obtained:
\[ \hat{M} \left( \hat{M} \cdot J(-0) \right) = J(+0), \]
\[ [\hat{M} \times [J(-0) \times \hat{M}]] = J_M(+0) + J_{\text{sd}}(+0), \] (23)
which are the desired boundary conditions. At \( x = L \) boundary, the continuity condition gives instead of (22)
\[ J(L + 0) - J(L - 0) - J_M(L - 0) = 0. \] (24)

After projecting to \( \hat{M} \), two conditions are obtained:
\[ J(L + 0) = J(L - 0), \quad J_M(L - 0) = 0. \] (25)

To express (23) and (25) conditions explicitly as requirements imposed on the desired magnetization, we multiply the conditions by \( \hat{M} \) vectorially. Then we obtain finally
\[ \frac{\partial \hat{M}(x)}{\partial x} \bigg|_{x=0} = k \left[ \hat{M}_{\text{FM}} \times \hat{M}(0) \right], \] (26)
\[ \frac{\partial \hat{M}(x)}{\partial x} \bigg|_{x=L} = 0 \] (27)
with a parameter
\[ k = \frac{\mu_B Q_{\text{FM}}}{a M} e \frac{\nu^*}{e \nu^* + \cos^2 \chi}. \] (28)
which characterizes the torque value.

### 4 The macrospin approximation

The problem is simplified considerably if the AFM layer is thin enough, so that \( \lambda \ll 1, L \ll \sqrt{\alpha/\beta} \) conditions are fulfilled and the magnetization is almost constant in that layer. Then the following expansion is valid:
\[ M(x) = M(0) + M'(0)x + \frac{1}{2} M''(0)x^2 + \ldots, \] (29)
where the primes mean derivatives with respect to \( x \) coordinate. Thus, several functions of a single variable (time \( t \)) are introduced instead of a
single $\hat{M}(x, t)$ function of two variables. It means physically, that a single large domain ("macrospin") is placed in the AFM layer [13]. Such an approximation often corresponds to experiments and simplifies calculations.

By differentiating (29) with respect to $x$, substituting $x = L$ and using boundary conditions (26) and (27), we obtain

$$\hat{M}''(0) = -L^{-1}\hat{M}'(0) = -\frac{k}{L} [\hat{M}^{\text{FM}} \times \hat{M}].$$

(30)

The only term in Eq. (16) that contains the second derivative may be transformed with (30) taking into account:

$$a[\hat{M}(0) \times \hat{M}''(0)] = k[\hat{M}(0) \times [\hat{M}(0) \times \hat{M}^{\text{FM}}]].$$

(31)

By substituting (31) into Eq. (16), we have

$$\frac{d\hat{M}(0)}{dt} + g[\hat{M}(0) \times H'] + \frac{ak}{L}[\hat{M}(0) \times [\hat{M}(0) \times \hat{M}^{\text{FM}}]] -$$

$$-\alpha G \left[ \hat{M}(0) \times \frac{d\hat{M}(0)}{dt} \right] = 0,$$

(32)

where $H' = H + H_{\text{FM}}^{\text{FM}}$.

At first sight, Eq. (32) coincides practically with that in [13] that was reproduced in many publications. There is an important difference, however. In Eq. (2) an exchange interaction is absent because of fulfilling condition (14), so that the equation is valid, in principle, at any distance from $x = 0$ plane where torque $T$ originates. At $x > \lambda_F$, vectors $\hat{M}_{\text{FM}}$ and $\hat{M}$ become collinear, so the torque vanishes. However, the spin current created in the lattice by the torque remains.

5 Exchange instability conditions

Let us discuss solution of the linearized dynamical equation (32) for small harmonic fluctuations $\Delta \hat{M}_x, \Delta \hat{M}_y \sim \exp(-i\omega t)$ near the stationary state $\hat{M} = \hat{z}$. The dispersion equation takes the form

$$\omega^2 + 2i\nu \omega - w = 0,$$

(33)

where the following notations are used:

$$w = \frac{1}{1 + \alpha G} \left[ \Omega_x \Omega_y + \left( \frac{ak}{L} \right)^2 \right],$$

(34)

$$\nu = \frac{\alpha G}{1 + \alpha G} \left[ \frac{1}{2} (\Omega_x + \Omega_y) + \frac{ak}{L\alpha G} \left( \hat{M} \cdot n \right) \right],$$

(35)

$$\Omega_x = g((H \cdot \hat{M}) + H_{\text{FM}}^\text{FM} + 4\pi M), \quad \Omega_y = g((H \cdot \hat{M}) + H_{\text{FM}}^\text{FM}).$$

The general instability condition $\Im \omega > 0$ requires fulfilling inequality $\nu < 0$. Put, for example, the following typical parameter values: $\Omega_x \approx$
$4\pi gM \sim 10^{11} \text{ c}^{-1}$, $\Omega_y = g(H + H_{AFM}^2) \sim 10^9 \text{ c}^{-1}$, $\alpha'_G \approx 10^{-1} \ll 1$. Then $\nu < 0$ condition is approximately reduces to the following inequality:

$$1 + \frac{2\alpha k}{\alpha'_G L \Omega x} (\bar{M} \cdot \mathbf{n}) < 0.$$  \hspace{1cm} (36)

By solving the inequality (36) with respect to the current density and taking the definition (28) into account, we obtain the following instability condition:

$$- \left(\bar{M} \cdot \mathbf{n}\right) \frac{j}{\mathcal{E}} > \lambda \alpha'_G \frac{2\pi gM^2}{\mu_B Q_{FM}} \left(1 + \frac{1}{\nu'}\right).$$  \hspace{1cm} (37)

Note that we discuss only one current direction when the electrons flow from the FM layer towards to the AFM layer. It is seen from (37) that the instability occurs only at $\left(\bar{M} \cdot \mathbf{n}\right) < 0$ under such a current direction.

As $\mathbf{n}$ vector is co-directed with $\bar{M}_{FM}$, it is clear that two magnetizations $\bar{M}_{FM}$ and $\bar{M}$ should be antiparallel to create instability.

It follows from (37) that the threshold current density is proportional to the dissipation constant and squared magnetization. The latter is determined with the field and can be decreased. For example, if $M \sim 10$ G is taken, then the threshold current density decreases by two orders of magnitude in comparison with the typical value for the FM junctions.

6 Conclusions

The present work has preliminary character, however, it allows to make some conclusions:

- FM–AFM structure may be canted in a magnetic field, i.e., a resulting magnetization can be induced. It allows to consider an exchange switching by means of a current-driven torque similar to that in a FM junction.
- To decrease the threshold current, the AFM layer must have minimal dissipation and minimal value of the field-induced magnetization.
- The instability is possible under antiparallel relative orientation of the FM and AFM layers.

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