Journal of Physics: Condensed Matter

Alternating current driven instability in magnetic junctions

To cite this article: E M Epshtein and P E Zilberman 2009 *J. Phys.: Condens. Matter* **21** 146003

IOP Publishing Journal of Physics: Condensed Matter
J. Phys.: Condens. Matter **21** (2009) 146003 (4pp) doi:10.1088/0953-8984/21/14/146003

Alternating current driven instability in magnetic junctions

E M Epshtein and P E Zilberman

V A Kotelnikov Institute of Radio Engineering and Electronics of the Russian Academy of Sciences, Fryazino 141190, Russia

E-mail: epshtein36@mail.ru

Received 27 October 2008, in final form 19 January 2009
Published 13 March 2009
Online at stacks.iop.org/JPhysCM/21/146003

Abstract
An effect is considered of alternating (high-frequency) current on the spin-valve-type magnetic junction configuration. The stability with respect to small fluctuations is investigated in the macrospin approximation. When the current frequency is close to the eigenfrequency (precession frequency) of the free layer, parametric resonance occurs. Both collinear configurations, antiparallel and parallel, can become unstable under resonance conditions. The antiparallel configuration can also become unstable under non-resonant conditions. The threshold current density amplitude is of the order of the dc current density for switching of the magnetic junction.

1. Introduction

There is continuing attention on the behaviour of magnetic junctions under spin-polarized current flow. This is not surprising because of a number of interesting phenomena which have been observed, such as magnetic configuration switching [1], spin wave generation [2], etc. The effects can occur on a nanosize scale, since their characteristic lengths are the exchange and spin diffusion ones with typical values of the order of 10 nm [3]. This allows us to use such effects for high-density information recording by electric current, unattainable for switching magnetic elements by magnetic field alone.

The current driven switching of magnetic junctions is often accompanied with magnetization oscillations and other high-frequency effects (see, e.g., [2, 4–6]). In this connection, an interesting problem arises, namely, the effect of spin-polarized *alternating* current on magnetic junctions.

In this work, we consider an effect of alternating (high-frequency) current on the magnetic junction configuration. When the parametric resonance conditions are fulfilled, both collinear configurations, parallel and antiparallel, can become unstable. It should be noted that the parametric resonance in magnetics has been studied in many works (see, e.g., [7]). However, the parametric resonance considered there was excited by a high-frequency external magnetic field, i.e. the nonlocal Ampere field. In our case, the spin-polarized current interaction with the magnetic lattice is of exchange nature, so that it is localized in the above-mentioned small range. As will be shown below, the current-induced parametric resonance has additional features. Thus, the main effect takes place at the precession frequency, not the doubled one. The instability is possible, too, beyond the parametric resonance conditions. It appears that only the spin-injection mechanism [8, 9], not the spin-torque transfer (STT) mechanism [10, 11], contributes to the effects under consideration.

2. Model considered and main equations

We consider a conventional spin-valve model consisting of a pinned ferromagnetic layer (layer 1), thin spacer layer, ferromagnetic free layer (layer 2) and nonmagnetic layer (layer 3) closing an electric circuit. The alternating current flows perpendicular to the layer planes (CPP mode). We investigate the stability of collinear (parallel or antiparallel) relative orientations of the pinned and free layers against small magnetization fluctuations under an alternating current flowing with density

$$j(t) = j_0 \cos \Omega t. \tag{1}$$

The free layer is assumed to be thin compared to the spin diffusion length and domain wall thickness, so that the macrospin approximation is applicable [12]. In this approximation, the fluctuations are described by the modified Landau–Lifshitz–Gilbert (LLG) equation [12]:

$$\frac{\mathrm{d}\hat{\boldsymbol{M}}}{\mathrm{d}t} - \kappa\left[\hat{\boldsymbol{M}} \times \frac{\mathrm{d}\hat{\boldsymbol{M}}}{\mathrm{d}t}\right] + \gamma[\hat{\boldsymbol{M}} \times \boldsymbol{H}] + \gamma H_{\mathrm{a}}(\hat{\boldsymbol{M}}\cdot\boldsymbol{n})[\hat{\boldsymbol{M}} \times \boldsymbol{n}] + \gamma[\hat{\boldsymbol{M}} \times \boldsymbol{H}_{\mathrm{d}}] + \frac{ap}{L}[\hat{\boldsymbol{M}} \times \hat{\boldsymbol{M}}_1] + \frac{ak}{L}[\hat{\boldsymbol{M}} \times [\hat{\boldsymbol{M}} \times \hat{\boldsymbol{M}}_1]] = 0. \tag{2}$$

Here the following notations are used: $\hat{\boldsymbol{M}} = \boldsymbol{M}/|\boldsymbol{M}|$ is the unit vector along the free layer magnetization $\boldsymbol{M}$, $\hat{\boldsymbol{M}}_1$ is the same for the pinned layer, $\boldsymbol{H}$ is the external magnetic field, H_a is the anisotropy field, $\boldsymbol{n}$ is the unit vector along the anisotropy axis, $\boldsymbol{H}_d$ is the demagnetization field, L is the free layer thickness, a is the magnetization diffusion constant, γ is the gyromagnetic ratio and κ is the Gilbert damping constant. The parameters p and k describe the spin-polarized current effect on the free layer magnetic lattice due to the spin-injection mechanism [8, 9] and STT mechanism [10, 11], respectively. In the collinear configuration with $(\hat{\boldsymbol{M}}_1 \cdot \hat{\boldsymbol{M}}) = \pm 1$, these parameters take the following forms:

$$p = \frac{\mu_B \gamma \alpha \tau Q_1}{ea} \frac{\lambda \nu(\nu^* - 1) \pm 2b\nu^*}{(\nu^* + 1)^2} |j| \equiv p^{(\pm)}, \tag{3}$$

$$k = \frac{\mu_B Q_1}{eaM} \frac{\nu^*}{\nu^* + 1} j, \tag{4}$$

where e is the electron charge, μ_B is the Bohr magneton, Q_1 is the conductivity spin polarization in layer 1, $\lambda = L/l \ll 1$, l being the spin diffusion length in the free layer, τ is the spin relaxation time in the free layer, α is the dimensionless sd exchange interaction constant in the free layer, the $b = (\alpha_1 M_1 \tau_1)/(\alpha M \tau)$ ratio describes the pinned layer contribution (see [13] for more details), $\nu = Z_1/Z_2$, $\nu^* = \lambda\nu + (Z_1/Z_3)$; Z_i $(i = 1, 2, 3)$ being the spin resistances [14]

$$Z_i = \frac{l_i}{\sigma_i(1 - Q_i^2)} \tag{5}$$

and σ_i is the conductivity of the ith layer. The upper and lower signs in (3) refer to the parallel and antiparallel configurations, respectively.

The formulae (2)–(4) have been derived in the assumption of direct current flowing. However, at $\Omega\tau \ll 1$ (this is a typical situation) the conduction electron spins follow the magnetization oscillations without any appreciable inertia, so that all the conditions are fulfilled under which the formulae are valid [13]. So the electron contributions to LLG equation (2) described by the terms with quantities p and k contain an instantaneous value of the current density. Therefore, we may substitute the alternating current density (1) with frequency Ω for j in formulae (3) and (4). Then the parameters p and k in (2) will have a time dependence of the form $p^{(\pm)}(t) = p_0^{(\pm)}|\cos\Omega t|$ and $k(t) = k_0 \cos\Omega t$, respectively. As to the free layer magnetization, it does not follow the current modulation in an inertialess way, as is seen from (2) containing the time derivative of the magnetization.

In contrast to the contribution of the STT mechanism, the contribution of the spin-injection mechanism due to the sd effective field is proportional to the absolute value of the current density, so that it is the same for forward ($1 \to 2 \to 3$) and backward ($3 \to 2 \to 1$) currents. This leads to different spectra of $p^{(\pm)}(t)$ and $k(t)$ functions:

$$p^{(\pm)}(t) = p_0^{(\pm)}|\cos\Omega t| = p_0^{(\pm)}\left(\frac{2}{\pi} + \frac{4}{3\pi}\cos 2\Omega t - \frac{4}{15\pi}\cos 4\Omega t + \cdots\right), \tag{6}$$

i.e. only the even harmonics of Ω frequency (including dc component) are presented in the $p^{(\pm)}(t)$ spectrum, while the $k(t)$ function spectrum consists of the single frequency Ω. The presence of the dc component in the spectrum of $p^{(\pm)}$ in spite of the absence of such a component in the current is due to the 'rectifying' character of the injection mechanism, where the effective field is the same for both signs of the current density [13]. Therefore, the two mechanisms do not interact in the lowest-order resonance phenomena and they may be considered separately. We begin with the spin-injection mechanism.

Let the x axis be directed along the current, the yz plane be parallel to the layer planes, the free layer occupy the range $0 \leqslant x \leqslant L$ and the vectors $\boldsymbol{H}, \boldsymbol{n}$ and $\hat{\boldsymbol{M}}_1$ have the following components: $\boldsymbol{H} = \{0, 0, H_z\}$, $\boldsymbol{n} = \{0, 0, 1\}$, $\hat{\boldsymbol{M}}_1 = \{0, 0, 1\}$. We investigate the stability of equilibrium collinear configurations $\bar{\hat{M}}_z = \pm 1$ against the free layer magnetization small fluctuations $\hat{M}_x, \hat{M}_y$. The LLG equation components linearized in the fluctuations, taking the spin-injection mechanism into account only, take the form

$$\frac{d\hat{M}_x}{dt} + \kappa\bar{\hat{M}}_z \frac{d\hat{M}_y}{dt} + \gamma(H_z + H_a\bar{\hat{M}}_z)\hat{M}_y + \frac{ap^{(\pm)}(t)}{L}\hat{M}_y = 0, \tag{7}$$

$$\frac{d\hat{M}_y}{dt} - \kappa\bar{\hat{M}}_z \frac{d\hat{M}_x}{dt} - \gamma(H_z + H_a\bar{\hat{M}}_z + 4\pi M\bar{\hat{M}}_z)\hat{M}_x - \frac{ap^{(\pm)}(t)}{L}\hat{M}_x = 0. \tag{8}$$

The periodic time dependence of the coefficients of the last terms on the left-hand side of the equations leads to the possibility of parametric resonance.

3. Parametric resonance

It is well known [15] that parametric resonance occurs when the parameter modulation frequency is close to $2\omega_0/n$, where ω_0 is the eigenfrequency of the oscillatory system, $n = 1, 2, 3, \ldots$. If the modulation coefficient is small, the parametric instability range narrows and the instability threshold rises with increasing resonance order n.

In accordance with (6), we consider parametric excitation at a frequency 2Ω, with the first two terms taken into account on the right-hand side of (6). We assume the damping constant κ to be small and neglect it for the time being. Taking the Fourier transforms of (7) and (8) with respect to time, we have

$$-i\omega\hat{M}_x(\omega) + \gamma\left(H_z + H_a + \frac{3\epsilon^{(\pm)}}{\gamma}\right)\hat{M}_y(\omega) = -\epsilon^{(\pm)}[\hat{M}_y(\omega + 2\Omega) + \hat{M}_y(\omega - 2\Omega)], \tag{9}$$

$$-i\omega\hat{M}_y(\omega) - \gamma\left(H_z + H_a + 4\pi M\bar{\hat{M}}_z + \frac{3\epsilon^{(\pm)}}{\gamma}\right)\hat{M}_x(\omega) = \epsilon^{(\pm)}[\hat{M}_x(\omega + 2\Omega) + \hat{M}_x(\omega - 2\Omega)]; \tag{10}$$

$$\epsilon^{(\pm)} = \frac{2}{3\pi} \frac{ap_0^{(\pm)}}{L} \tag{11}$$

which is a quantity proportional to the current density amplitude with the dimension of frequency.

At typical values of the parameters ($M \sim 10^3$ Oe, $H_a \sim \epsilon^{(\pm)}/\gamma \sim 10^2$ Oe) a condition is fulfilled:

$$4\pi M \gg |H_z|, H_a, 3\epsilon^{(\pm)}/\gamma, \tag{12}$$

that is assumed below.

If we make the substitution $\omega \to \omega \pm 2\Omega$ in (10), (11), the equations are obtained where terms with $\hat{M}_{x,y}(\omega \pm 2\Omega)$ stand on the left-hand sides and the terms with $\hat{M}_{x,y}(\omega)$ and $\hat{M}_{x,y}(\omega \pm 4\Omega)$ with the $\epsilon^{(\pm)}$ coefficient on the right-hand sides.

If frequency $\epsilon^{(\pm)}$ (the coupling parameter) is small compared to the other characteristic frequencies (ω_0, Ω), the coupling with $\hat{M}_{x,y}(\omega \pm 4\Omega)$ can be neglected, because it leads to higher-order corrections in $\epsilon^{(\pm)}$. As a result, we obtain a closed system of equations for $\hat{M}_{x,y}(\omega)$ and $\hat{M}_{x,y}(\omega \pm 2\Omega)$. Equating the determinant of the system to zero, we get the dispersion equation

$$\begin{aligned}&\left\{\omega + (\epsilon^{(\pm)})^2\left[\frac{\omega + 2\Omega}{\Delta(\omega + 2\Omega)} + \frac{\omega - 2\Omega}{\Delta(\omega - 2\Omega)}\right]\right\}^2 - \omega_0^2 \\ &\quad + (\epsilon^{(\pm)})^2(4\pi\gamma M)^2\left[\frac{1}{\Delta(\omega + 2\Omega)} + \frac{1}{\Delta(\omega - 2\Omega)}\right] \\ &\quad - (\epsilon^{(\pm)})^4\omega_0^2\left[\frac{1}{\Delta(\omega + 2\Omega)} + \frac{1}{\Delta(\omega - 2\Omega)}\right]^2 = 0, \end{aligned} \tag{13}$$

where

$$\omega_0 = 2\gamma\left[\pi M\left(H_z\bar{\hat{M}}_z + H_a + \frac{3\epsilon^{(\pm)}}{\gamma}\bar{\hat{M}}_z\right)\right]^{1/2} \tag{14}$$

is the eigenfrequency neglecting the terms quadratic in $\epsilon^{(\pm)}$:

$$\Delta(\omega) = \omega_0^2 - \omega^2. \tag{15}$$

Note that the dc component of $p^{(\pm)}(t)$ (see (6)) renormalizes the system eigenfrequency.

Near the parametric resonance, we have $\omega \approx \omega_0 \approx \Omega$, so that we may retain only the summands with resonant denominator $\Delta(\omega - 2\Omega)$ in the terms with $\epsilon^{(\pm)}$ and replace ω and Ω with ω_0 everywhere but that denominator. This leads to the parametric resonance equation

$$\Delta(\omega)\Delta(\omega - 2\Omega) = (4\pi\gamma M\epsilon^{(\pm)})^2. \tag{16}$$

Let $\omega = \omega_0 + \nu$, $\Omega = \omega_0 + \delta$, where $|\nu|, |\delta| \ll \omega_0$. The equation for ν takes the form

$$\nu^2 - 2\delta\nu + \left(\frac{2\pi\gamma M\epsilon^{(\pm)}}{\omega_0}\right)^2 = 0, \tag{17}$$

which gives

$$\nu = \delta \pm \left[\delta^2 - \left(\frac{2\pi\gamma M}{\omega_0}\right)^2(\epsilon^{(\pm)})^2\right]^{1/2}. \tag{18}$$

It is seen from (18) that the parametric instability takes place at

$$\epsilon^{(\pm)} > \frac{\omega_0\delta}{2\pi\gamma M} \tag{19}$$

with the increment

$$(\mathrm{Im}\,\omega)_p = \left[\left(\frac{2\pi\gamma M\epsilon^{(\pm)}}{\omega_0}\right)^2 - \delta^2\right]^{1/2}. \tag{20}$$

In the presence of dissipation ($\kappa \neq 0$), damping takes place with decrement (see, e.g., [7])

$$|(\mathrm{Im}\,\omega)_d| = 2\pi\kappa\gamma M. \tag{21}$$

If $\kappa \ll 1$, the parametric instability threshold may be estimated from the condition

$$(\mathrm{Im}\,\omega)_p > |(\mathrm{Im}\,\omega)_d| \tag{22}$$

at zero resonance detuning ($\delta = 0$).

Equations (20)–(22) give the following condition for the parametric instability:

$$\epsilon^{(\pm)} > \kappa\omega_0. \tag{23}$$

The right-hand side of this inequality contains $\epsilon^{(\pm)}$ too. Therefore, the inequality is to be resolved with respect to $\epsilon^{(\pm)}$. As a result, the following instability threshold is obtained for the collinear configurations:

$$\epsilon_{th}^{(\pm)} = 6\pi\kappa^2\gamma M\left[\left(\frac{H_a \pm H_z}{9\pi\kappa^2 M} + 1\right)^{1/2} \pm 1\right]. \tag{24}$$

In contrast to the case of direct current, both collinear configurations can become unstable, but the corresponding thresholds are different. At $\epsilon^{(-)} > \epsilon_{th}^{(-)}$, $\epsilon^{(+)} < \epsilon_{th}^{(+)}$ the switching from the unstable antiparallel configuration to the stable parallel one is possible, while at $\epsilon^{(+)} > \epsilon_{th}^{(+)}$ both configurations are unstable. In the latter case, the free layer magnetization, apparently, will oscillate between two collinear configurations. To obtain a more precise answer, a nonlinear problem should be solved that requires separate consideration. We hope to return to that problem later.

At $H_a \pm H_z \gg 9\pi\kappa^2 M$ the instability threshold for both configurations takes the form

$$\epsilon_{th}^{(\pm)} = 2\kappa\gamma[\pi M(H_a \pm H_z)]^{1/2} = \kappa\omega_0^{(0)}, \tag{25}$$

where $\omega_0^{(0)} = 2\gamma[\pi M(H_a \pm H_z)]^{1/2}$ is the eigenfrequency in the absence of the electric current.

If $H_a - H_z \ll 9\pi\kappa^2 M$, we have

$$\epsilon_{th}^{(-)} = \tfrac{1}{3}\gamma(H_a - H_z) \tag{26}$$

for the antiparallel configuration. It is seen from (26) that the instability threshold can be lowered considerably with external magnetic field close to, but slightly lower than, the anisotropy field. Note that participation of the magnetic field does not prevent locality of the effect, because the magnetic field cannot do switching alone, without the current.

Let us compare the alternating current density amplitude j_{th} corresponding to the parametric instability threshold with the direct current density j_{dc} leading to the switching antiparallel orientation to the parallel one in the absence of the external magnetic field. The dc threshold corresponds to the condition [13] $ap^{(-)}/L > \gamma H_a$, where $p^{(-)}$ is determined by (3). In the parametric resonance case, $ap_{th}^{(-)}/L = 3\pi\kappa\gamma(\pi M H_a)^{1/2}$, so that

$$\frac{j_{th}}{j_{dc}} = 3\pi\sqrt{\pi}\kappa\left(\frac{M}{H_a}\right)^{1/2}. \quad (27)$$

At typical values of the parameters ($M/H_a \sim 10, \kappa = 3 \times 10^{-2}$), this ratio is of the order of 1. At lower damping constant, the parametric instability threshold will be smaller than the dc threshold.

4. Non-resonance instability

The instability of the antiparallel configuration under alternating current flow is also possible when the parametric resonance condition $\Omega \approx \omega_0$ is not fulfilled. It follows from (14) that the eigenfrequency ω_0 becomes imaginary at $\epsilon^{(-)} > \gamma(H_a - H_z)/3$ because of the contribution of the dc component in the spectrum of the $p^{(\pm)}(t)$ function, i.e. such a component of the *sd* exchange field. The cancelling of the eigenfrequency corresponds to an orientational phase transition similar to that under dc injection current [13]. The threshold amplitude of the alternating current is $\pi/2$ times as much as the corresponding dc threshold. Note that the threshold is higher then the parametric instability threshold $\epsilon_{th}^{(-)}$, so that the parametric instability develops first under fulfilled parametric resonance conditions.

5. Is the parametric instability possible due to the spin-torque transfer mechanism?

Since the $k(t)$ function describing the STT mechanism contribution has a single-mode spectrum, the lowest order of the parametric resonance corresponds to the $\Omega \approx 2\omega_0$ condition. However, if all the previous calculations are carried out for $p^{(\pm)}(t) = 0, k(t) \neq 0$, the contributions from $k(t)$ mutually cancel in resonance approximation ($\Omega \approx 2\omega \approx 2\omega_0$). Therefore, a negative answer should be given to the question stated above. This is due to the fact that the STT, in contrast with the spin injection, modifies the damping, not the eigenfrequency. As is known, the parametric resonance is not possible in such a situation. Therefore, observation of switching a magnetic junction by a high-frequency current means that the spin-injection mechanism takes place. The role of the STT mechanism is reduced to modifying the damping that determines the threshold of the parametric instability.

6. Conclusion

The analysis carried out shows that the collinear configurations of a magnetic junction can become unstable under parametric resonance conditions. The instability threshold is of the same order of magnitude as the corresponding threshold under direct current flowing through the magnetic junction. Depending on the alternating current density amplitude, the antiparallel configuration can be switched to a parallel one or both collinear configurations can be unstable with growing magnetic fluctuations. To elucidate the resulting state which the system considered comes to, further investigations, both theoretical and experimental, are needed.

Acknowledgments

The authors are grateful to Drs S G Chigarev and A I Krikunov for helpful discussion.

The work was supported by the Russian Foundation for Basic Research, grant nos. 06-02-16197 and 08-07-00290.

References

[1] Katine J A, Albert F J, Buhrman R A, Myers E B and Ralph D C 2000 *Phys. Rev. Lett.* **84** 3149
[2] Tsoi M, Jansen A G M, Bass J, Chiang W-C, Seck M, Tsoi V and Wyder P 1999 *Phys. Rev. Lett.* **80** 4281
[3] Bass J and Pratt W P Jr 2007 *J. Phys.: Condens. Matter* **19** 183201
[4] Krivorotov I N, Emley N C, Sankey J C, Kiselev S I, Ralph D C and Buhrman R A 2005 *Science* **307** 228
[5] Ralph D C and Stiles M D 2008 *J. Magn. Magn. Mater.* **320** 1190
[6] Xiao Z H, Ma X Q, Wu P P, Zhang J X, Chen L Q and Shi S Q 2007 *J. Appl. Phys.* **102** 093907
[7] Gurevich A G and Melkov G A 1996 *Magnetization Oscillations and Waves* (Boca Raton, FL: CRC Press)
[8] Heide C, Zilberman P E and Elliott R J 2001 *Phys. Rev.* B **64** 064424
[9] Gulyaev Yu V, Zilberman P E, Epshtein E M and Elliott R J 2002 *JETP Lett.* **76** 155
[10] Slonczewski J C 1996 *J. Magn. Magn. Mater.* **159** L1
[11] Berger L 1996 *Phys. Rev.* B **54** 9353
[12] Gulyaev Yu V, Zilberman P E, Krikunov A I, Panas A I and Epshtein E M 2007 *JETP Lett.* **86** 328
[13] Epshtein E M, Gulyaev Yu V and Zilberman P E 2006 Phenomenological theory of current driven exchange switching in ferromagnetic nanojunctions arXiv:cond-mat/0606102
[14] Epshtein E M, Gulyaev Yu V and Zilberman P E 2007 *J. Magn. Magn. Mater.* **312** 200
[15] Landau L D and Lifshitz E M 1976 *Mechanics* (London: Pergamon)