Spin-transfer torque induced vortex dynamics in Fe/Ag/Fe nanopillars

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Received 6 April 2011, in final form 10 June 2011
Published 8 September 2011
Online at stacks.iop.org/JPhysD/44/384002

Abstract

We report on the experimental and analytical work on spin-transfer torque induced vortex dynamics in metallic nanopillars with in-plane magnetized layers. We study nanopillars with a diameter of 150 nm, containing two Fe layers with a thickness of 15 nm and 30 nm, respectively, separated by a 6 nm Ag spacer. The sample geometry is such that it allows for the formation of magnetic vortices in the Fe discs. As confirmed by micromagnetic simulations, we are able to prepare states where one magnetic layer is homogeneously magnetized while the other contains a vortex. We experimentally show that in this configuration spin-transfer torque can excite vortex dynamics and analyse their dependence on a magnetic field applied in the sample plane. The centre of gyration is continuously dislocated from the disc centre, and the potential changes its shape with field strength. The latter is reflected in the field dependence of the excitation frequency. In the second part we propose a novel mechanism for the excitation of the gyrotropic mode in nanopillars with a perfectly homogeneously magnetized in-plane polarizing layer. We analytically show that in this configuration the vortex can absorb energy from the spin-polarized electric current if the angular spin-transfer efficiency function is asymmetric. This effect is supported by micromagnetic simulations.

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

1. Introduction

Since the theoretical prediction [1, 2] and experimental demonstration [3, 4] of spin-transfer torque induced effects, many studies focused on phenomena such as current-induced magnetization dynamics and switching. Spin-transfer devices are promising candidates for future information technology. Current-induced switching between discrete magnetization states may be employed in nonvolatile memories such as the magnetic random access memory. As spin-polarized currents can propel steady spin precession [5], spin-transfer torque devices are also envisaged to be used as integrated microwave sources. Therefore, finding highly tunable spin-transfer torque nano oscillators (STNOs) is a matter of great current interest. After the discovery that spin-transfer torque can also drive the oscillatory motion of a magnetic vortex [6–9], gyrotropic vortex motion is considered a promising STNO candidate. The magnetic vortex structure appears at certain dimensions as the ground state of ferromagnetic discs. The magnetization circulates around a core with magnetization perpendicular to the plane of the disc. The orientation of the magnetization in the core is called the polarization and the sense of rotation of the in-plane circulation constitutes the vortex chirality. This configuration is a consequence of the competition between the dipolar and exchange energies.

One of the first experimental studies focusing on the interaction of the magnetic vortex with spin-polarized electric currents was published by Kasai et al. in 2006 [10]. In their study, a Permalloy (Py, Fe19Ni81) disc 40 nm in thickness and several hundred nm in radius was connected to two Au electrodes. The electrodes contact the disc edge on opposite sides resulting in an in-plane current, i.e. the electrons traverse the Py disc laterally when the electrodes lie on unequal
electrical potentials. In [10], three such samples with different radii (410, 530 and 700 nm) are investigated. The discs exhibit a vortex structure as a remanent state of magnetization. Upon application of an ac current, vortex gyrotropic motion is detected at specific frequencies of the excitation by measuring the time-averaged anisotropic magnetoresistance of the configuration. The resonance frequency of the respective disc is found to be specific to its radius, which is consistent with [11].

The resonant excitation of a vortex by an in-plane ac current can lead to a switching of the vortex core polarity. Yamada et al [12] study a sample similar to the ones just mentioned, namely a Py disc of thickness 50 nm and diameter 500 nm. The magnetic configuration is observed by magnetic force microscopy, which allows distinction of the vortex core polarities. Contact to the disc is made by Au electrodes analogous to [10]. An ac current tuned to the characteristic frequency of 290 MHz is applied for about 10 s. The core polarity after the current interval is found to be opposite the one before in about half of the experiments, indicating that the core switched multiple times during the application of the current. This agrees with results from micromagnetic simulations. The switching probability in the experiment has a resonant character, and the height of its maximum decreases with decreasing current amplitude.

Since in the experiments reported above, the spin currents are due to charge transport, these are unavoidably accompanied by Oersted fields. Time-dependent magnetic fields have been shown to give rise to excitation of the vortex state, and thus in the presence of oscillating spin-polarized charge currents, one expects a combined effect of spin-transfer torque and fields. This combined action was investigated in detail by Bolte et al [13]. Here, the geometry of a ferromagnetic Py square (Ni$_{80}$Fe$_{20}$) of area $2 \times 2 \, \mu \text{m}^2$ and 20 nm thickness is studied. The square is contacted by two 40 nm thick Au electrodes at opposite sides. The overlap of the Py structure with each electrode is 150 nm. The Permalloy nanomagnet stabilizes a Landau pattern at low magnetic fields. Thus, the system is comparable to the ones in [10, 12]. Additional Py squares are located directly underneath the electrodes. In these structures, the influence of spin-transfer torque is negligible, so that the occurring vortex motion is almost entirely caused by the oscillating Oersted field generated by the electric current through the Au strip lines. In this way, the gyration due to a pure field or a mixture of field and spin-polarized current can be compared. By scanning transmission x-ray microscopy at the Ni L$_3$ absorption edge, the vortices are imaged at discrete phases of their motions. It is shown that vortices exhibiting the same polarities but opposite chiralities gyrate at different phases with respect to the ac current. This phase difference is traced back to the coupling of the vortex motion to the ac magnetic field. The experiment shows a phase splitting also for the motion of the spin-current-driven vortices, revealing the influence of the Oersted field in this geometry. Micromagnetic simulations indicate that the fields causing the splitting originate from the perpendicular currents in the regions around the Au/Py interfaces and from the inhomogeneous current distribution in the nanomagnet.

Spin-transfer torque driven vortex dynamics is not limited to current-in-plane geometries, as shown in [6]. Priibag et al investigate nanopillars exhibiting a spin-valve geometry. The samples contain two Py discs separated by a Cu spacer. In equilibrium, the thick nanomagnet stabilizes a vortex while due to shape anisotropy, the thin disc is found in a quasi-homogeneous configuration, as confirmed by micromagnetic simulations. For a dc current along the cylinder axis of the nanopillar, where the electron flow is directed from the thin towards the thick layer, the authors observe microwave signals due to periodic variation of the magnetoresistance. The dependence of dynamics on the sample current and the external magnetic field is investigated and compared with micromagnetic simulations, confirming that the signals arise from vortex gyration.

In [14], a similar geometry is studied in the Fe/Ag system. In this case, the thin Fe layer is an extended film, resulting in a stability of the homogeneously magnetized state increased by the absence of side surfaces in the vicinity of the nanocontact. The thick Fe nanomagnet, which is a disc 230 nm in diameter and 20 nm in thickness, can stabilize both a quasi-homogeneous state and a vortex while the thin layer magnetization is uniform. In contrast to [6], for this system dc current-induced vortex gyration is found for an electron flow directed from the vortex into the thin layer. The dynamical properties of the vortex and quasi-homogeneous state are compared with each other.

In nanopillars, the vortex gyration is possible due to the magnetostatic potential arising from the lateral confinement. The potential leads to a restoring force counteracting the gyrotropic term in the Thiele equation [15]. However, in [16] it has been shown that spin-transfer torque induced vortex gyration is also possible in extended films contacted by point-like electrodes 200 nm in diameter, where the spatial confinement of the vortex is absent and the magnetostatic potential arises from the circumferential Oersted field induced by the electrical current. For this case micromagnetic simulations show that the vortex core is driven to a stable circular motion around the contact centre on an orbit outside the contact area. The excitations are observed for dc currents of both polarities and external magnetic fields applied perpendicularly to the film planes. Within the framework of the rigid vortex model, the authors show that only the spin-polarization component perpendicular to the film planes causes a torque resulting in a precession of the vortex.

Recently, Ruotolo et al [17] investigated systems of four point contacts in a 15 nm thick Co layer with intercontact distances down to 500 nm, resulting in a lattice of neighbouring vortices and antivortices. The authors demonstrate that at a sufficiently small intercontact distance and above critical currents, the spin-transfer torque driven vortex dynamics is coupled. The coupling is shown to originate from the exchange interaction between neighbouring vortex–antivortex pairs, resulting in a phase-locked gyration of the vortices. This leads to the observation of a single high-frequency peak with an integrated power that scales with the square of the number of individual oscillators as expected. The experiment shows a route to a phase-locking of large numbers.
of vortex oscillators with possible application in tunable on-chip microwave sources.

Here, we study vortex dynamics in nanopillars, patterned from a molecular beam epitaxy grown system that contains Fe as ferromagnetic and Ag as spacer material (Fe/Ag(0 0 1)). This material combination has interesting and advantageous properties. First, it has been calculated [18] that the Fe/Ag(0 0 1) interface exhibits a resistance that strongly depends on the spin direction thus providing for a high spin polarization. Additionally, it has been shown [19] that an asymmetry in (spin channel averaged) ferromagnetic element and spacer resistances leads to an angular asymmetry in giant magnetoresistance (GMR) and also in spin torque efficiency. Experimental evidence for the appearance of this feature in Fe/Ag(001) spin valves was provided by Lehndorff et al [20].

In contrast to a vortex residing in an extended ferromagnetic film, due to the lateral confinement in a nanopillar we expect a more pronounced dependence of the magnetostatic energy on the vortex position, i.e. the coordinates of its core. This renders the pillar geometry especially interesting for studying the dependence of the vortex dynamics on external in-plane-aligned magnetic fields. Moreover, the finite lateral extension makes micromagnetic simulations of the system more feasible. As will be shown the measurements agree very well with the calculations. After the experimental part we will address the question of how the vortex motion in a pillar with in-plane magnetized polarizer is excited and propose a novel mechanism. We show by analytical calculation and micromagnetic simulations that the angular asymmetry in spin-transfer torque efficiency can lead to energy absorption of the magnetic vortex from the spin-polarized electric current.

### 2. Sample fabrication

After cleaning and annealing the GaAs(100) substrate, the sequence of metallic layers Fe (1 nm)/Ag (150 nm)/Fe (30 nm)/Ag (6 nm)/Fe (15 nm)/Au (50 nm) is deposited by molecular beam epitaxy. The 1 nm Fe layer serves as a seed layer for the adjacent Ag buffer layer, which will later on be the bottom electrode. It follows the spin-valve layer system with an Au capping layer on top, which will also provide a part of the top electrode. All layers grow in a single crystalline manner [21]. Both Fe layers have bcc structure and exhibit cubic magneto-crystalline anisotropy. Next the areas of the prospective bottom electrodes are covered with resist by optical lithography. The uncovered material is then removed with ion beam etching (IBE). Afterwards resist dots of 150 nm diameter are placed into the 15 × 15 µm² contact area located in the middle of the previously defined structures by means of electron beam lithography. We employ a hydrogen silsesquioxane (HSQ) electron beam resist (FOX-12). Successively, the layers are milled down practically onto the bottom electrode by IBE resulting in pillars with diameters of about 150 nm. We again use HSQ to insulate the nanopillars and planarize the sample. Additional protection and insulation is provided by a 50 nm thick Si₃N₄ layer which is deposited by means of plasma-enhanced chemical vapour deposition (PECVD). In the next step the sample is covered with optical resist apart from 10 × 10 µm² contact windows above the nanopillars, which are opened by IBE so that the pillar tops are cleared from insulating material. A short reactive ion etching (RIE) step is additionally performed to selectively remove the remaining HSQ. Finally, the top electrode is deposited by a negative optical lithography step and subsequent lift-off.

### 3. Vortex dynamics: experiment and simulation

All measurements are carried out at room temperature. To characterize the dc behaviour of the sample we apply a constant current while sweeping the magnetic field and measuring the resistance (two-point measurements). For measurements of the RF voltage we split the ac and the dc part of the voltage by a bias-T (frequency range: 45 MHz–26.5 GHz). The signal is then amplified by a 30 dB amplifier (0.5–26.5 GHz) and fed into a 50 GHz spectrum analyser. The sample holder is homemade and makes contact with the sample leads via spring pins with a diameter of 0.3 mm. The angles between the in-plane easy axes of the cubic magneto-crystalline anisotropy and the long axes of the top and bottom leads are 45° (insets in figure 1). (The two in-plane hard axes of the cubic anisotropy are parallel to the top and bottom leads, respectively.) In a previous work on Fe/Ag nanocontacts fabricated in a similar way [22], the cubic magneto-crystalline anisotropy constant was estimated to be about 45 kJ m⁻³. Figure 1 shows the measurements of the dc resistance at low currents for both easy axes (first axis: black, second axis: red). The difference of the two cases is likely to be caused by sample imperfections (e.g. electrode contacting, deviations from perfect cylindrical shape of the nanopillar) which can lead to nonsymmetric current distributions and additional contributions to anisotropy. At small applied fields, most interestingly, the axis-2 measurement exhibits states with resistances as low as the saturation resistance. Instead we...
would expect intermediate up to maximum resistance values due to the dipolar antiferromagnetic coupling between the layers. To investigate the sample characteristics under high current densities we apply a current of $-17 \, \text{mA}$ corresponding to a current density of $-9.62 \times 10^7 \, \text{A cm}^{-2}$ (figure 2(a)) (electron flow from bottom to top of the pillar). We expect increasing influence of spin torque effects and the Oersted field on the disc magnetizations. The field is applied parallel to easy axis 2 and swept from negative to positive (sweep 1, black) and back to negative values (sweep 2, red). Both sweeps show an increase in the GMR as the applied field is driven from the saturation range towards zero. Once a sufficiently low external field is reached ($-51 \, \text{mT}$ for sweep 1, $86 \, \text{mT}$ for sweep 2) the sample enters a state of low resistance which persists up to large opposite field values ($274 \, \text{mT}$ for sweep 1), where for both sweeps we see abrupt steps in resistance. Afterwards the resistance decreases as the sample is driven into saturation.

Micromagnetic simulations have been performed with the finite element code called TetraMag [23]. The magnetization configurations in each disc for the different applied field values are provided by the micromagnetic simulations. Thus a clear physical understanding of the GMR signal is possible. Typical material parameters of iron, $\mu_0 M_s = 2.15 \, \text{T}$ (saturation magnetization), exchange constant $A = 2.1 \times 10^{-11} \, \text{J m}^{-1}$ and anisotropy constant $K_c = 48 \, \text{kJ m}^{-3}$ are used. The sample volume is discretized into irregular tetrahedrons with cell size of about $1 \, \text{nm}$. The magnetic field is applied along the second easy axis and linearly increased/decreased over a 40 ns simulation time between the two extreme values. The damping factor is set to $\alpha = 0.1$. The simulations take into account the Oersted field, which is calculated from the current-density distribution obtained for the nanopillar contact geometry including the leads. The Oersted field distributed in the disc does not have a circular symmetry. This is due to the asymmetric input-output arrangement of the technical current (see figure 1). According to the simulations, the maximum field value at the perimeter of the disc is about $41 \, \text{mT}$ on one side and $33 \, \text{mT}$ on the opposite side. Figure 3 shows the dependence of the simulated GMR on the applied magnetic field for both sweep directions. The observed dependence is in good qualitative agreement with the experimentally measured GMR curve for sweep 2 although spin-transfer torque is not included by the simulation. From this we can infer that the spin-transfer torque has negligible influence on the time-averaged magnetoresistance. Since the GMR curves are symmetric in the direction of the field sweep we only discuss one sweep direction (from positive to negative). For field values larger than 200 mT a low resistance region is found, in which the magnetization in both discs is mostly aligned with the external field. At about 190 mT, a vortex nucleates in
the bottom disc, leading to a jump in the resistance which then further increases linearly with a decrease in the external field. The chirality of this vortex is determined by the direction of the Oersted field. When the external field approaches zero a vortex with the same chirality nucleates in the top disc, which is reflected in the low resistance value in this regime. With a further decrease in the external field both vortices are moved towards the perimeter of the discs. The vortex of the upper disc is expelled at about $-320\,\text{mT}$ and consequently the resistance increases. The resistance reaches its minimum above $-400\,\text{mT}$ when the vortex of the bottom disc is expelled as well. The parallel orientation of the two magnetizations is restored. The ‘wiggles’ in the resistance accompanying the nucleation of the first vortex are caused by excess energy. In the experiment, this energy is dissipated almost instantly and thus they are not observed.

During the measurement depicted in figure 2(a) in addition to the dc voltage we measured a spectrum at each field value. The star symbols in the graph mark the presence of peaks in the spectrum representing excitations of the magnetization. Obviously, the single vortex state is excited by the current. Figure 2(b) displays the spectral power density at each field point. In the first field range from 202 to 182\,\text{mT} the spectra are broad (750–1680\,\text{MHz}) and show multiple peaks of different amplitudes. As the field is lowered, the frequencies mostly red shift. At the next field value all secondary peaks have vanished. Starting from 177\,\text{mT} we measure single peak spectra with a peak at initially 1.275\,\text{GHz}, which continues the red shift observed until 136\,\text{mT}, where the frequency reaches its minimum of 1.218\,\text{GHz} with a FWHM of 3.1\,\text{MHz}. Afterwards the peak frequency increases until 106\,\text{mT} to a value of 1.356\,\text{GHz}. If we reverse the current sign the resistance profile indicates that we again have a vortex in a similar field range, while the other layer has homogeneous magnetization. But the rich dynamics observed with the previous current setting do not occur. In order to rule out thermal activation of the vortex motion we prepare the single vortex state by applying $-17\,\text{mA}$ and sweeping the applied field from positive saturation to 151\,\text{mT}. After reducing the current stepwise to zero while keeping the field constant we measure a current loop at this field from 0 to $-20\,\text{mA}$, then to $+20\,\text{mA}$ and finally back to 0\,\text{mA}. Only for negative current do we see clear excitations. Therefore, thermal activation cannot explain the excitations, which must be caused by the spin-transfer torque. These results lead us to the following conclusions: since both discs have a geometry, which at low effective fields favours the vortex state, and due to the match of the resistance profiles obtained by micromagnetic simulations to the measured ones we have assured ourselves that the investigated state comprises one vortex and one homogeneously magnetized disc. The homogeneous disc is aligned with the field, therefore its dipolar field reduces the applied field to sufficiently small effective values for the vortex to occur in the other disc. No significant in-plane stray field components are expected to arise from the vortex, thus the homogeneous disc is subject to the full applied field. If the observed excitations took place in the homogeneously magnetized disc, we would expect significantly higher frequencies than measured. Another very interesting aspect is the V-shaped dependence of the frequency on the field (figure 2(c)). The most redshifted peak measured occurs at an applied field of 136\,\text{mT}. This value is close to the micromagnetically computed dipolar field that the spacer facing part of the vortex disc is exposed to. If we assume that at this point the dipolar field and the applied field cancel each other, the measured frequency behaviour can most easily be explained: when the vortex enters the disc the effective field is roughly 60\,\text{mT}. The magnetostatic energy minimum, which is the centre of the vortex core motion, is shifted towards the disc rim. As the field is lowered the centre of motion moves towards the disc centre where the frequency reaches its minimum value. On further decreasing the applied field the average vortex position is shifted towards the opposite rim as the effective field tunes down to about $-50\,\text{mT}$. (At lower fields the sample enters the double vortex state.) The shifting of the average vortex position is connected to a shape change of the potential in which the vortex moves. This change in potential shape causes the gyration frequency to first decrease on the way to the disc centre and afterwards increase again while the vortex is shifted towards the rim. We, therefore, conclude that the vortex is excited and this excitation is caused by the spin torque. The vortex moves around an average position, which shifts through the disc as the magnetic field is swept. The curvature of the potential at its minimum is connected to the restoring force. With changing field, that curvature also changes. This is reflected in the observed V-shaped dependence of the frequency on the field.

4. Vortex gyration under influence of a homogeneously polarized current

The excitation of the gyrotropic motion of a vortex by a spin-polarized current is conventionally described by extending the Thiele equation with a spin torque term [9, 16]. However, this approach denies the possibility of exciting a vortex with a current that is homogeneously in-plane polarized [24]. Khvalkovskiy et al [24] pointed out that an inhomogeneous polarizer can supply the vortex with energy. Here, we will propose another mechanism by showing that even a perfectly homogeneous polarizer can pump energy into the gyrotropic motion of a vortex via spin torque, if the angular spin torque efficiency function is asymmetric. We start with the Landau–Lifschitz–Gilbert equation with spin torque

$$\frac{\partial m}{\partial t} = -\gamma m \times H_{\text{eff}} + \alpha m \times \frac{\partial m}{\partial t} + \sigma j g(\Lambda, \beta) m \times (m \times p)$$

(1)

where $m$ and $p$ are vector fields of unit length denoting the normalized magnetization and the polarization, respectively, $\alpha$ is the Gilbert damping parameter, $\gamma$ the gyromagnetic ratio and $j$ the current density. The effective field is defined as

$$H_{\text{eff}} = -\frac{1}{\mu_0 M_0} \frac{\delta W}{\delta m}$$

(2)

where $W$ is the energy density of the disc and $M_0$ is the saturation magnetization. The spin torque coefficients are...
absorbed by the disc: Let

\[ g(\Lambda, \beta) = \frac{P \Lambda}{(\Lambda^2 / 2) + \Lambda^{-1} \sin^2(\beta / 2)} \]  

(3)

with \( P \in [0, 1] \) being the degree of current polarization. It is related to the difference in spin dependent resistances \( R^+ \), \( R^- \) as \( P = (R^+ - R^-)/(R^+ + R^-) \). The parameter \( \Lambda \) is defined according to \( \Lambda = \sqrt{\Lambda G(R^+ + R^-) / 2} \), where \( \Lambda \) is the contact cross section and \( G \) is the spacer conductance as defined in [19]. Thus, \( \Lambda \) describes the mismatch between spacer and ferromagnet resistance and is therefore related to spin accumulation at the nonmagnet–ferromagnet interfaces. Let \( p \) from now on be homogeneous and constant in time. By combining equations (1) and (2) we get the power density absorbed by the disc:

\[ \frac{\partial W}{\partial t} = -\frac{\alpha \mu_0 M_s}{\gamma} \left| \frac{\partial m}{\partial t} \right|^2 - \frac{\mu_0 M_s}{\gamma} \sigma_{jg} \int \frac{\partial m}{\partial t} \times m \cdot p. \]

(4)

The total power is obtained by an integration of equation (4) over the disc volume. Let the cylinder axis of the disc be the \( z \)-axis while the disc is parallel to the \( x-y \)-plane. The direction of \( m \) is then described by the two angles \( \psi(x, t) \) and \( \theta(x, t) \) where \( \psi \) is the azimuthal angle (\( \psi = 0 \) is the \( x \)-direction) and \( \theta \) is the angle between \( m \) and the \( z \)-axis:

\[ m = \sin(\theta) \cos(\psi) \hat{e}_x + \sin(\theta) \sin(\psi) \hat{e}_y + \cos(\theta) \hat{e}_z. \]

(5)

We switch to cylindrical disc coordinates (radius \( \rho \), azimuth \( \chi \), height \( z \)). In order to carry out the integration we make use of the following relations describing the behaviour of the magnetization under a rotation by an angle \( \chi \) around the \( z \)-axis:

\[ \begin{align*}
\varphi(p, \chi; a, \chi_v) &= \varphi(p, \chi - \chi_v; a, 0) + \chi_v, \\
\frac{\partial \varphi}{\partial \chi_v} (p, \chi; a, \chi_v) &= \frac{\partial \varphi}{\partial \chi_v} (p, \chi - \chi_v; a, 0), \\
\theta(p, \chi; a, \chi_v) &= \theta(p, \chi - \chi_v; a, 0), \\
\frac{\partial \theta}{\partial \chi_v} (p, \chi; a, \chi_v) &= \frac{\partial \theta}{\partial \chi_v} (p, \chi - \chi_v; a, 0).
\end{align*} \]

(6)

Here \( a \) and \( \chi_v \) are the polar coordinates of the vortex core, respectively. It is assumed that the magnetization distribution is independent of the \( z \)-coordinate. In order to obtain an expression for the energy the vortex absorbs due to spin torque during one rotation when moving on a constant orbit radius \( a \) at a constant frequency \( \omega \) one has to evaluate the expression

\[ E_{\text{STT}} = -\int_0^T \int_0^{2\pi} \int_0^\Lambda \frac{\partial m}{\partial t} \times m \cdot \sigma_{jg} \frac{\partial m}{\partial t} \hat{R}_\chi, m) \]

(7)

Choosing \( p^T = (1, 0, 0) \), equation (7) can be transformed to

\[ \begin{align*}
E_{\text{STT}} &= -L \mu_0 \frac{\mu_0 M_s}{\gamma} \sigma_{jg} \int_0^{2\pi} \int_0^\Lambda d^2 x \tilde{g}(\Lambda, \chi_v, m) \\
&\quad \times \hat{R}_\chi, (\frac{\partial m}{\partial t} \times m).
\end{align*} \]

(8)

with \( g(\Lambda, \beta) := \tilde{g}(\Lambda, \cos \beta) = \tilde{g}(\Lambda, m_s) \). The operator \( \hat{R}_\chi \) returns the \( x \)-component of a vector \( f \) rotated around the \( z \)-axis by the angle \( \chi_v \), i.e., \( \hat{R}_\chi f = f_x \cos \chi_v - f_y \sin \chi_v \). To proceed further it is useful to express \( g \) in terms of the parameter \( \xi := \Lambda^2 - 1 \). \( \xi = 0 \) now corresponds to the symmetric case where we perform a Taylor expansion in \( \xi \):

\[ g(\xi, \beta) = \frac{P}{2} + 1 - \cos \beta - \frac{1}{4} P\xi \sum_{n=0}^{\infty} \left( \frac{1}{2} (\cos \beta + 1)^n \right)^n. \]

(9)

This expression converges for all angles only if \( -1 < \xi < 1 \). Keeping the first-order (\( n = 0 \)) term in \( \xi \) and inserting it into equation (8) yields:

\[ \frac{dE_{\text{STT}}}{d\xi} \bigg|_{\xi=0} = \frac{\pi}{8} e \int_\Lambda d^2 x \left[ m \times \hat{R}_\chi m \right]. \]

(10)

The formula clearly shows that only the core region contributes to the energy uptake. The sign of the integral is independent of core polarity and chirality because vortices of opposite polarity revolve at different sense of rotation. The integrand in equation (10) is sensitive to the core shape. Approximations for the bell-shaped profile of the stationary core have been given, for example, in [25, 26]. Since moving vortices are accompanied by a dip in the \( m_z \)-magnetization [27–29], we take the distributions \( m(x) \) and \( \partial / \partial t (m(x)) \) of a moving vortex obtained from an OOMMF simulation in order to compute expression (10) as well as the energy \( E_{\text{damp}} \) dissipated:

\[ E_{\text{damp}} = \frac{\alpha \mu_0 M_s}{\gamma} \int_0^T \int_0^{2\pi} \int_0^\Lambda d^2 x \left| \frac{\partial m}{\partial t} \right|^2. \]

(11)

Thus, an estimate of the critical current density necessary to obtain a balance between the spin-transfer torque induced energy uptake and dissipation can be given by imposing

\[ E_{\text{damp}} \approx \left. \frac{dE_{\text{STT}}}{d\xi} \right|_{\xi=0} \eta. \]

(12)

From equation (12), the critical current density for steady gyration is estimated to about \( 3.36 \times 10^7 \) A cm\(^{-2} \) (for \( P = 0.85, \Lambda = 1.41 \) and \( L = 30 \) nm). The vortex absorbs energy for negative \( j \) (meaning that electrons flow from the polarizer into the vortex) and positive \( \eta \) corresponding to \( \Lambda > 1 \). If either the sign of \( j \) or \( \eta \) is flipped, the term will lead to an additional damping force. The current density estimated above is one order of magnitude higher than usually applied currents. This suggests that the effect is not the only cause for steady gyration. However, simulations clearly confirm its existence (figure 4). We perform four simulations (by OOMMF, lateral cell size 2 nm, one \( z \)-node, disc diameter 150 nm, disc thickness 30 nm, \( \mu_0 M_s = 2.14 \) T, \( K_1 = 48 \) kJ m\(^{-3} \)) with the same initial magnetization distribution which is a vortex moving on an orbit with an intermediate radius. The homogeneous polarizer points into the \( x \)-direction. A current density of \( 1.5 \times 10^8 \) A cm\(^{-2} \) is employed. In the first simulation (black lines) the current sign is chosen such that the parallel alignment is favoured, while \( \Lambda = 1.41 \) corresponding to \( \eta = 0.988 \). According to equation (10) in first order in \( \xi \) the
spin current must lower the effective damping in comparison with simulation three (blue), where we have the same current but \( \xi = 0 \). Simulation two (red) is the same as simulation one, but this time the electron flow is reversed so that the current now favours the antiparallel configuration. In that case we expect the strongest damping. Simulation four has the same parameters as three, except that the current sign is negative this time. The damping should be the same as in simulation three. As is shown in figure 4 all cases qualitatively reproduce the analytical findings. From a fit of an exponential decay function to the data, damping parameters are extracted. Indeed, simulation one exhibits the smallest damping corresponding to the largest time constant \( \tau \approx 2.426 \) ns, simulation two has the largest damping (\( \tau \approx 2.033 \) ns) and simulation three has an intermediate value as expected (\( \tau \approx 2.224 \) ns). Simulation four yields \( \tau \approx 2.209 \) ns which is very close to the result of three as it should be. In realistic cases \( \Lambda \) can be much larger. For the Fe/Ag/Fe(0 0 1) system at low temperatures we have an experimental value of \( \Lambda = 3.4 \) obtained from spin torque measurements leading to \( \xi = 10.6 \) ( [20] and references therein). Equation (9) converges uniformly for all \( \beta \in [0, \pi] \) only if \( \xi |\beta| < 1 \), but was well suited to calculate the derivative (10). However, it fails if one wants to compute the energy contribution for larger \( |\xi| \). In this case one can expand in \( \cos \beta \):

\[
g(\Lambda, \beta) = \frac{P\Lambda^2}{1 + \Lambda^2} \sum_{n=0}^{\infty} \left( \frac{1 - \Lambda^2}{1 + \Lambda^2} \cos \beta \right)^n.
\]  

This series uniformly converges on \( \beta \in [0, \pi] \) for every \( \Lambda \).

5. Summary and discussion

In the first part we presented measurements of excited states of a nanopillar containing two ferromagnetic discs. One of the discs is in a vortex state while the other is homogeneously magnetized. We deduced that the measured peaks correspond to spin-torque-driven vortex motion. Several micromagnetic simulations indicate that the vortex resides in the thicker bottom disc. We see the excitations for electron flow from the bottom to the top disc. This would mean that the excitation of the vortex is caused by electrons which are reflected by the top layer. It is not clear why a direct electron flow from the top disc into the vortex cannot excite it. The experiment in [6] which also has a vortex and a homogeneous polarizer shows excitation only for electron flow into the vortex. On the other hand, the previously reported experiment which also involves the Fe/Ag(001) system [14] is in accordance with the presented data. The V-shaped frequency versus field dependence of the vortex excitations has been explained by the in-plane motion of the vortex potential minimum in the disc and the accompanying shape change of the potential under the action of the field.

In the second part we have shown that the angular asymmetry in the spin-transfer torque efficiency function enables the vortex to gain energy even when the polarizer is perfectly homogeneous. The same effect can also lead to an additional damping depending on the signs of current and asymmetry parameter \( \xi \). Our analytical method relies on the rigid vortex approximation, that is, the vortex degrees of freedom are reduced to rotational motion around the disc centre. The critical current estimated to maintain a stationary motion is an order of magnitude above experimental currents. We employ the OOMMF code to investigate the effect by means of micromagnetic simulations. Although the simulations performed so far are too simple to claim capturing the complex behaviour of a thick magnetic disc they clearly show the expected qualitative behaviour. Their results also confirm the small size of the effect. More advanced simulations will have to be performed in order to gain insight into, for example, the influence of a possible \( \gamma \)-dependence of the vortex magnetization distribution on the dynamical properties of the system [6, 30].

The novel mechanism of the \( \xi \)-dependent driving/damping of vortex motion is expected to be influential in the Fe/Ag/Fe(0 0 1) system, especially in the experiment reported by Lehndorff et al [14] for two reasons. (i) The Fe/Ag/Fe(0 0 1) system exhibits a rather strong \( \xi = 10.6 \) and (ii) the polarizing layer is laterally extended and thus much less susceptible for nonhomogeneous magnetization distributions.

Acknowledgment

A M D acknowledges financial support from the EU project STraDy (MOIF-CT-2006-039772).

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