Spin-torque switching and control using chirped AC currents

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Abstract

We propose to use oscillating spin currents with slowly varying frequency (chirp) to manipulate and control the magnetization dynamics in a nanomagnet. By recasting the Landau–Lifshitz–Slonczewski equation in a quantum-like two-level formalism, we show that a chirped spin current polarized in the direction normal to the anisotropy axis can induce a stable precession of the magnetic moment at any angle (up to 90°) with respect to the anisotropy axis. The drive current can be modest (10^6 \text{ A cm}^{-2} or lower) provided the chirp rate is sufficiently slow. The induced precession is stable against thermal noise, even for small nano-objects at room temperature. Complete reversal of the magnetization can be achieved by adding a small external magnetic field antiparallel to the easy axis. Alternatively, a combination of chirped ac and dc currents with different polarization directions can also be used to trigger the reversal.

Keywords: magnetization dynamics, spin torque transfer, spin-torque nano-oscillators, autoresonance

(Some figures may appear in colour only in the online journal)

1. Introduction

Many technological applications of magnetic nano-objects (nanomagnets) require to accurately control their magnetization dynamics [1–5]. This can be achieved in several ways, including static or oscillating magnetic fields, thermal effects, and spin-torque transfer (STT). The latter technique consists in injecting a spin-polarized current into a nanomagnet; the electron spins transfer some of their angular momentum to the magnetic material by applying a torque on its magnetic moment and thus inducing the switch. This technique was first proposed theoretically by Slonczewski [6] and Berger [7] and later realized experimentally and further developed by many others [8–11]. In the last decade, STT has given rise to new technological developments such as STT-based random-access memory [12] and spin-torque nano-oscillators (STNOs) [13]. Still more recent investigations in this field have been focusing on spin-Hall effects [14].

Achieving optimal switching of the magnetization is a compromise between Joule heating of the sample and reversal time. Although dc currents are the most widespread method to achieve fast switching [15, 16], recent theoretical and experimental work has shown that an ac current tuned at the resonant precession frequency could be even more efficient [17–19]. Various combinations of ac and dc currents and microwave magnetic fields were implemented to improve the efficiency of the switching [20–23]. A spin current excitation can also be used to induce persistent precession of the magnetic moment, thus enabling magnetic nanostructures to behave as tunable radiofrequency oscillators [24, 25]. Analyzing the tunability and stability of such devices in the presence of intrinsic effects (damping, magnetic anisotropy, thermal fluctuations) is therefore of utmost importance.

In this work, we will demonstrate that an oscillating spin current with slowly variable frequency (chirp) is a very efficient tool for manipulating the magnetization dynamics in a nanomagnet. We will focus on two important effects: (i) the fast switching of the magnetic moment and (ii) the precise control of its precession frequency.

A classical nonlinear oscillator can be excited and controlled by a chirped oscillating force using a well-known...
effect called autoresonance, which has been exploited for very diverse applications ranging from plasma [26] and atomic [27] physics to semiconductor nanostructures [28]. Autoresonant excitation occurs when a nonlinear oscillator starting in equilibrium is driven by a force \( F(t) = \epsilon \cos \left[ \frac{1}{2} \omega_d(t) t \right] \), with a time-dependent frequency \( \omega_d(t) \) that slowly passes through the linear frequency \( \omega_0 \) of the oscillator. It can be shown that, for the driving amplitude \( \epsilon \) above a certain threshold \( \epsilon_{th} \) (scaling as \( \epsilon_{th} \sim \alpha^{3/4} \), where \( \alpha = d\omega_d/dt \) is the chirp rate), the oscillator frequency ‘locks’ to the driving frequency continuously, so that the resonance condition is preserved for a long time. In that case, the amplitude of the oscillations grows without saturation, until of course some other effects kick in.

In two earlier studies [29, 30], we made use of the autoresonance mechanism to control the magnetization switching of a magnetic nanoparticle using a chirped microwave field. This technique was shown to reduce the static switching field and a magnetic nanoparticle using a chirped microwave field. This technique was shown to reduce the static switching field and

\[
\text{Equation (1) can be rewritten as: } \dot{\mathbf{m}} = \mathbf{H} \times \mathbf{m}, \quad \mathbf{H} = \mathbf{H}_{\text{eff}} + \mathbf{H}_{\text{an}}. \quad (1)
\]

We shall adopt an approach due to Feynman [31], which exploits the analogy between the magnetization dynamics and a two-level quantum-like system and was used earlier to study the autoresonant control of the magnetization dynamics [30]. The LLS equation is equivalent to a system of two coupled equations for the complex quantities \( A_1 \) and \( A_2 \):

\[
i A_1 = \frac{\kappa_0}{2} A_1 + \kappa A_2, \quad (6)
\]

\[
i A_2 = \frac{\kappa_0}{2} A_2 + \kappa^* A_1, \quad (7)
\]

where \( \kappa_0 = \tilde{H}_c, \kappa = \frac{1}{2} (\tilde{H}_c - i\tilde{H}_s) \), and \( \mathbf{m} \) is related to \( A_{1,2} \) through the expressions:

\[
m_x = A_1 A_2^* + A_1^* A_2, \quad (8)
\]

\[
m_y = i (A_1 A_2^* - A_1^* A_2), \quad (8)
\]

\[
m_z = |A_1|^2 - |A_2|^2, \quad (8)
\]

which shows that, in the Feynman representation, the system is fully described by the real amplitudes \( B_{1,2} \) and the phase difference \( \phi \) is \( \varphi_2 - \varphi_1 \).

In order to illustrate the autoresonant technique, we first consider the simple case where damping is neglected (\( \lambda = 0 \)) and the frequency varies linearly with time, \( \omega_d(t) = \omega_0 - \alpha t \). Other effects—including damping, thermal noise, and an external magnetic field—will be added in sections 3 and 4.

We focus on the case of an ac spin current of constant amplitude, polarized orthogonally to the axis of easy magnetization, i.e., \( \gamma \mathbf{I} = J_\perp(t) \mathbf{e}_z \), with \( J_\perp(t) = 2 \cos \varphi_d \) and \( \omega_d(t) = \varphi_d(t) \) is the chirped driving frequency. In this case, it follows from equation (5) that \( \mathbf{H} = (\omega m_x - J_\perp m_y) \mathbf{e}_z + J_\perp \mathbf{m} \mathbf{e}_y \), where \( \omega_r = 2 \gamma K \mathbf{m} / \mu_s \) is the resonant precession frequency. The autoresonance mechanism requires that the time-dependent drive frequency crosses the resonant frequency from above, so we set the initial driving frequency \( \omega_0 > \omega_r \).

We seek solutions to equations (6) and (7) under the initial conditions \( A_1 = 1 \) and \( A_2 = 0 \), i.e., \( \mathbf{m} = \mathbf{e}_z \). Using equations (9), we obtain:

\[
\dot{B}_1 = \epsilon (B_1^2 - B_2^2) B_2 \cos \varphi_d \cos \Delta \varphi \quad (10)
\]

\[
\dot{B}_2 = -\epsilon (B_1^2 - B_2^2) B_1 \cos \varphi_d \cos \Delta \varphi \quad (11)
\]

\[
\Delta \varphi = \omega_r (B_1^2 - B_2^2) + \frac{\epsilon}{B_1 B_2} \cos \varphi_d \sin \Delta \varphi. \quad (12)
\]
We then define $\phi = \Delta \varphi - \varphi_d - \pi/2$, and use the rotating wave approximation (neglecting the high frequencies) to derive the equations for the coupled variables $B_2$ and $\phi$:

$$B_2 = (\epsilon/2)(1 - 2B_2^2)B_1 \sin \phi$$  \hspace{1cm} (13)

$$\dot{\phi} = \omega_r - \omega_d - 2\omega_i B_2^2 + \epsilon/(2B_1 B_2) \cos \phi,$$  \hspace{1cm} (14)

where we recall that $B_1 = \sqrt{1 - B_2^2}$. Focussing on the weakly nonlinear regime ($B_1 \approx 1$ and $B_2 \ll 1$), we obtain:

$$\dot{B}_2 = (\epsilon/2) \sin \phi$$  \hspace{1cm} (15)

$$\dot{\phi} = \omega_r - \omega_d - 2\omega_i B_2^2 + (\epsilon/2B_1) \cos \phi.$$  \hspace{1cm} (16)

The above equations are typical of systems that can be driven into autoresonance [30]. Previous work [32] showed that the system is captured into autoresonance when the excitation amplitude $\varepsilon$ exceeds a certain threshold

$$\varepsilon > \varepsilon_{th} = 0.82(2\omega_r)^{-1/2}\alpha^{3/4}.\hspace{1cm} (17)$$

When the above condition is satisfied, the chirped spin current stays locked with the precession oscillations, and drives the magnetic moment away from the anisotropy axis even in the nonlinear regime. These theoretical results are in agreement with numerical simulations of the full LLS equation, carried out for a nanomagnet with volume $V = 2 \times 10^{-24} m^3$ (20 nm x 20 nm x 5 nm), anisotropy constant $K = 2.2 \times 10^5 J m^{-3}$, and magnetic moment $\mu = 3.35 \times 10^{-18} J T^{-1}$ (see figure 1). For these parameters, the resonant frequency is $\omega_r/2\pi = 7.36 \text{ GHz}$. In this and all subsequent numerical results, the chirped current was applied for the entire duration of the simulation. The numerical solutions were obtained using a standard second-order predictor-corrector method (Heun’s scheme).

Note that, according to equation (13), the time derivative of $B_2$ vanishes when $B_1 = B_2$. Thus, when $m_z = 0$, it is impossible to further populate the level $B_2$. This implies that one cannot fully reverse the magnetization (i.e. reach $m_z = -1$) using such spin current. The largest precession angle attainable with this technique is $\theta = 90^\circ$ (where $\theta$ is the angle between the magnetic moment and the $z$ axis) as can be seen from figure 1. In the absence of damping and thermal noise, the magnetic moment will precess indefinitely perpendicular to the anisotropy axis $e_z$.

However, we will show in section 4.1 that, by adding a small ($\approx 10 \text{ mT}$) external magnetic field antiparallel to the anisotropy axis, it is possible to fully reverse the magnetization using the autoresonant technique described above. A second reversal technique, based on the combination of two spin currents, parallel and perpendicular to $e_z$, will be illustrated in section 4.2.

3. Autoresonant control of the precession

We now show that the autoresonant technique can be used to bring the magnetic moment to rotate around the anisotropy axis at a certain target angle and precession frequency. This is an important feature that allows to convert an electric current into high-frequency magnetic rotation, with potential applications to nanoscale devices such as STNOs. In particular, we want to study the stability of the forced precession regime using a spin current, including the effect of the Gilbert damping term $\Gamma_G$ and thermal fluctuations.

To this end, we enforce a fixed precession angle by chirping the excitation frequency exponentially, from the initial value $\omega_0$ towards the asymptotic value $\omega_f$:

$$\omega_d = \omega_d = \omega_f + (\omega_0 - \omega_f)e^{-t/\tau}.$$  \hspace{1cm} (18)

However, we emphasize that the particular form of the function $\omega_d(t)$ is not important—the autoresonant mechanism works in any case as long as the frequency variation is sufficiently slow. The required slowness is determined by equation (17), which can be recast as

$$\alpha^{3/4} < 1.22(2\omega_f)^{1/2} \varepsilon = 0.86 \gamma I_s \omega_f^{1/2},$$

where we recall that $\alpha = d\omega_d/dt$ is the chirp rate and $I_s = 2\varepsilon/\gamma$. Thus, the slowness of the chirp is related to both the precession frequency and the current intensity.

3.1. Gilbert damping and stability properties

We proceed from equations (13)–(14), where we add a small dissipative term $(\lambda \omega_f/\varepsilon \ll 1)$. Assuming that, for $\omega_d < \omega_f$,
precession being achieved in either

dynamics, so that the magnetic moment reaches its final pre-

As the last term in the square root is small and

tees stability. Thus, the autoresonant regime is always stable,

\( \phi = \phi_0 + \delta \phi e^{i \omega t} \). Equations (19) and (20) lead to the

definition with respect to \( x \), yielding two characteristic

\( \nu_{\pm}^2 = 2 \pi \lambda \omega f_0^2 + 2 \pi \omega F_0^2 = 0 \), where

\( f_0 \equiv G_0 - F_0^2/(G_0/F_0) = (1 - 2x_0^2)/2 \) (the prime denotes

\( \nu_{\pm} = i \omega f_0 \pm i \sqrt{2 \pi \lambda \omega f_0^2 + (\lambda \omega f_0)^2} \).

As the last term in the square root is small and \( f_0 \) is positive,

both roots \( \nu_{\pm} \) have a positive imaginary part, which guarant-

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which is a standard value for STNOS [33]. For this current, the threshold chirp time $\alpha^{-1/2}$ is of the order of 0.5 ns (the actual time to reach the asymptotic precession angle will be a multiple of this time), as can be deduced from the inset of figure 1. But since the threshold current decreases almost linearly with decreasing $\alpha$, using a slower chirp can reduce the required current by a significant factor. For instance, decreasing $\alpha$ by a factor of 10, cuts the threshold current by a factor $10^{3/2} \approx 5.6$, while it increases the time to induce the precession by a factor $10^{1/2} \approx 3.2$. Since the energy is proportional to the current, the autoresonant procedure can be helpful to reduce the energy required to achieve complete magnetization switching. Of course, there is a trade-off to be made between the rapidity of the overall process and the intensity of the required current (or energy), but it is clear that competitively low currents can be achieved if one accepts to lengthen the time to induce the precession.

3.2. Thermal effects

In the results reported above, temperature effects were neglected. However, previous theoretical [29, 34] and experimental [35] studies showed that the autoresonant mechanism is rather robust against thermal noise. In order to check that the same conclusion holds in the present case, we introduced thermal fluctuations in our model. As is usually done [29], thermal fluctuations at temperature $T$ are represented as a random magnetic field $\mathbf{b}(t)$ with zero mean and autocorrelation function given by:

$$
\langle \mathbf{b}_i(t) \mathbf{b}_j(t') \rangle = \frac{2 \lambda k_B T}{(1 + \lambda^2) \gamma \mu_s} \delta_{ij} \delta(t - t'),
$$

where $i,j$ denote the cartesian components $(x,y,z)$, $\delta_{ij}$ is the Kronecker symbol (meaning that the spatial components of the random field are uncorrelated), and $\delta(t - t')$ is the Dirac delta function, implying that the autocorrelation time of $\mathbf{b}$ is much shorter than the response time of the system. The temperature is thus proportional to the autocorrelation function of the fluctuating field.

In figure 4, we plot results at room temperature ($T = 300 \, K$) for a 25 nm-diameter nanoparticle (blocking temperature $\sim 5000 \, K$) with damping $\lambda = 0.01$ and $\omega_f/2\pi = 4 \, GHz$. There is no external magnetic field. The three curves correspond to different values of the oscillating spin current amplitude. The amplitude $I_5 = 6.3 \, mT$ is just above the autoresonant threshold in the absence of thermal fluctuations and can thus control the precession in a stable way, as was done in figure 2 (black curve). However, this is no longer true at finite temperature (figure 4), where thermal noise drives the magnetic moment back to the $z$ axis. In order to induce a stable precession, the current needs to be increased slightly, up to 8 mT or higher.

The above phenomenon is consistent with what was observed in the past for finite-temperature systems that are excited autoresonantly [29, 34, 35]. In particular, the ability to hold the precession for increasing driving amplitude $I_5$ (figure 4) can be explained as follows. The autoresonant system is formally equivalent to a quasiparticle trapped in an effective potential well of height $V_0$ proportional to $I_5$ [34]. The noise drives the quasiparticle out of the well, on a time scale proportional to $\exp(V_0/k_B T)$ if the quasiparticle is initially deeply trapped in the well [36]. Therefore, increasing $I_5$ (and thus $V_0$) amounts to reducing the effect of the thermal noise, in accordance with figure 4. In addition, thermal fluctuations also modify the threshold phenomenon. At zero temperature, there exists a sharp threshold for the excitation amplitude $I_5$ above which the system is always captured into the autoresonant regime. In the present work the existence of such a threshold, which depends on the chirp rate $\alpha$, was confirmed in figure 1 (see inset). At finite temperature, the threshold is no longer sharp, but instead displays a certain width that is proportional to the square root of the temperature [29]. All these effects were observed in our numerical simulations in full agreement with the general autoresonance theory.

The above results show that the autoresonant technique is very stable against thermal fluctuations. Such stability properties are of great importance in real STT devices [37], where phase fluctuations due to the presence of thermal noise can have a disruptive effect. Here, we showed thermal fluctuations do not disrupt the autoresonant drive of the precession, provided the spin current is increased slightly above the nominal (zero-temperature) threshold. In addition, the autoresonant excitation is not sensitive to the precise temporal profile of the chirped current frequency, the only requirement being that the frequency varies slowly in time.

We also note that many simulations of STNOS were performed at zero [38] or very small [39] temperature, or they involved large nano-objects [40] (diameter $> 100 \, nm$) for which the blocking temperature is very high and therefore the effect of thermal noise is minor even at $T = 300 \, K$. The present autoresonant technique has proven to preserve the stability of the oscillations even for much smaller nano-objects (25 nm) at room temperature. It may therefore be more advantageous for such ultrasmall nano-oscillators.
3.3. Phase locking

The first procedure is based on a dc spin current, which counteracts the Gilbert damping term, thus preventing the magnetic moment to relax back to easy axis [38, 39, 41, 42]. Although a dc current may be easier to implement, our approach has some specific advantages. First, it is possible (by modulating the frequency variation) to control precisely the trajectory of the magnetic moment towards the desired precession angle. Second, the method is rather stable against damping and thermal fluctuations, as was shown in the preceding paragraphs.

Now, we show that the autoresonant technique is also useful to induce phase locking between the external signal and the response of the STNO. Usually, phase locking (or injection locking) is achieved by combining an external dc current with an ac drive signal [43]. When the ac drive is close enough to the natural frequency of the STNO, then the latter starts oscillating in phase at the same frequency of the drive. For a given dc current, phase locking is achieved only for a narrow range of drive frequencies.

Using our approach, it was possible to phase-lock the drive (chirped ac current) to the STNO precession response, without any external dc currents and for a wide range of precession frequencies. Indeed, the autoresonant technique was originally devised exactly for such a purpose: to bring a system to oscillate at a specified nonlinear frequency by slowly sweeping the frequency of the drive. This should work for any target frequency, provided the threshold condition, equation (17), is satisfied. Importantly, the threshold condition also tells us that the driving ac current can have a very small amplitude, provided the frequency variation rate is slow enough.

To demonstrate phase locking between the drive and the STNO precession, we plot in figure 5 (top) the difference between the drive frequency $\omega_d(t)$ and the instantaneous precession frequency $\omega_{mz}(t)$. As expected for an autoresonant process, the two frequencies remain close together for all times after the system has been captured in autoresonance. The bottom panel of figure 5 shows the phase difference between the drive and the precession, which remains remarkably constant after about 8 ns. Importantly, the phase locking appears to be robust against thermal fluctuations, as these simulations were performed for the case corresponding to room temperature conditions. Such robustness and flexibility should make the proposed technique competitive with respect to other approaches.

4. Magnetization reversal

As a further application, we propose two procedures to completely switch the magnetic moment from parallel to antiparallel to the anisotropy axis $e_z$. The first procedure is based on an external static magnetic field antiparallel to the anisotropy axis, combined with the autoresonant spin current described in the preceding sections. The second method relies on the combination of two types of spin currents (ac and dc) polarized in different directions.

4.1. External magnetic field

The presence of an external magnetic field $\mathbf{H}_0 = H_0 e_z$ affects the magnetization dynamics in two ways, through the torques $\Gamma_{1L}$ and $\Gamma_G$. As to $\Gamma_{1L}$, its primary effect is to move the peak of the energy barrier (the point where the instantaneous precession frequency vanishes) away from $\theta = 90^\circ$ (i.e. $m_z = 0$), towards values $\theta < 90^\circ$ ($m_z > 0$) for an external field antiparallel to $e_z$, and $\theta > 90^\circ$ ($m_z < 0$) for a field parallel to $e_z$ (figure 6). As to $\Gamma_G$, the part of the Gilbert torque that is due to the external field can be written:

$$\Gamma_G^{ext} = -\gamma \mu_0 \lambda \mathbf{m} \times (\mathbf{m} \times \mathbf{H}_0) = -\gamma \mu_0 \lambda H_0 (m_z \mathbf{m} - e_z).$$

Therefore, the $z$ component of the magnetic moment evolves under the action of $\Gamma_G^{ext}$ as follows:

$$\dot{m}_z = \gamma \mu_0 \lambda H_0 (1 - m_z^2) + \ldots \quad (23)$$

Of course, many other terms (notably the spin current) also affect the evolution of $m_z$. From the above expression, we see that the effect of $\Gamma_G^{ext}$ is to drive the magnetic moment towards $m_z = -1$ when $H_0 < 0$ and towards $m_z = 1$ when $H_0 > 0$.

However, according to equation (13), the autoresonant condition is always lost at $\theta = 90^\circ$ (when $B_1 = B_2$, or $m_z = 0$), irrespective of the external field. Thus, we have two possible scenarios, depending on the orientation of the external field (see figure 6):

1. If $H_0 < 0$ (antiparallel) the peak of the energy barrier is situated at a position $1 > m_z^2 > 0$. Starting from
drives the magnetic moment towards

denotes the peak of the barrier in either case. The point \( m_z = 0 \) cannot be crossed through autoresonant excitation. The magnetic moment starts at \( m_z = +1 \) and evolves right to left.

\[ m_z = 1, \] the autoresonant excitation induces precession with decreasing \( m_z \) and can bring the magnetic moment to overcome the energy barrier. Subsequently, the autoresonant phase locking is lost and the external-field Gilbert torque \( \Gamma_ez \) drives the magnetic moment towards \( m_z = -1 \).

2. If \( H_0 > 0 \) (parallel) the peak of the energy barrier is situated at a position \( m_z^* < 0 \). The autoresonant excitation can never bring the magnetic moment to cross the \( m_z = 0 \) plane and thus it can never overcome the barrier. In this case, \( \Gamma_ez \) brings the magnetic moment back to its initial value \( m_z = 1 \) [see equation (23)].

In figure 7, we present some numerical results that confirm the above scenarios. We consider an external field of intensity \( H_0 = \pm 50 \) mT, oriented either parallel or antiparallel to the anisotropy axis \( z \). Other parameters are identical to those corresponding to the red curve on figure 2. When the magnetic field is antiparallel to \( e_z \), the magnetic moment first starts precessing at decreasing azimuthal angle until it crosses the barrier, which is located around \( \theta = 79^\circ \) \( (m_z^* = 0.19) \), visible on figure 7 as the point where the autoresonant phase locking is lost). Subsequently, the magnetic moment relaxes towards \( m_z = -1 \) under the action of the external-field torque. In contrast, when \( H_0 \) is parallel to \( e_z \), the magnetic moment goes back to its original position \( m_z = +1 \), in agreement with the second scenario of our analysis.

For \( H_0 < 0 \) we were able to reverse the magnetic moment, in contrast to the case with no external field, for which the plane \( m_z = 0 \) could not be crossed. Thus, adding a small antiparallel magnetic field seems to be a good strategy to obtain complete reversal of the magnetization on a nanosecond timescale using the proposed autoresonant technique. Note however that the switching is triggered by the chirped AC current and not by the static field, which is far too small to induce alone the magnetization reversal. For instance, complete reversal can be achieved for \( H_0 = -10 \) mT, for which the energy barrier is situated at \( \theta = 88^\circ \) (not shown here). The role of the magnetic field is just to break the symmetry that places the maximum of the energy barrier at \( \theta = 90^\circ \) in the absence of an external field.

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4.2. Parallel spin current

The procedure is again based on the autoresonance technique and requires two spin currents polarized in the parallel and perpendicular directions with respect to \( e_z \). Let us first consider a purely parallel spin current: \( \gamma I_y = -J_y(t)e_z \). The effective field is then given by (we neglect damping for simplicity):

\[
\hat{H} = \omega_m e_z - J_\parallel (m_z e_z - m_z e_z),
\]

Using the two-level formalism described above, one can derive a closed-form solution for the real amplitude \( B_2 \):

\[
B_2(t) = \frac{B_2^0(0)e^{2t\Gamma}}{B_1^0(0) + B_2^0(0)e^{2t\Gamma}},
\]

where \( \Gamma(t) = \int_0^t J_y(t)dt \). Thus, for sufficiently large times, one obtains that \( B_2 \to 1 \), i.e. complete reversal of the magnetization by means of a dc spin current collinear with the anisotropy axis. From equation (24), it appears that the magnetic moment must be tilted away from the anisotropy axis at the initial time, i.e. \( B_2(0) \neq 0 \), in order for the reversal process to work. This suggests a way to combine two types of ac and dc spin currents in order to shorten the reversal time. Starting with a magnetic moment oriented along \( e_z \), a chirped current polarized along \( e_z \), first tilts the moment of a certain angle with respect to the anisotropy axis (this is the technique described earlier in this work); next, a dc current polarized along \( e_z \), completes the reversal according to equation (24).

Numerical simulations confirm this scenario (figure 8). Here, we show three cases where the \( J_\perp \) and \( J_\parallel \) currents are applied either separately or together: \( J_\perp \) alone can tilt the magnetic moment only up to \( 90^\circ \) \( (m_z = 0) \); \( J_\perp \) alone (3 mT in this case, with an initial tilt of \( 1\degree \)) can reverse the magnetization completely in about 15 ns; finally, when both currents are combined, the switching time is reduced to 8 ns. In the combined case, we used an ac spin current of magnitude 6 mT, although the theoretical threshold amplitude is close to 9 mT. This shows that the simultaneous use of the two types of excitations leads to a reduction of both the switching time and the autoresonance threshold for the \( J_\perp \) component.
Figure 8. Time evolution of $m_z$ for different types of spin currents: dc spin current of intensity $I_k = 3 \, \text{mT}$ parallel to the anisotropy axis $\mathbf{e}_z$ (red curve); ac chirped spin current perpendicular to $\mathbf{e}_z$ with $I_k = 9 \, \text{mT}$, $\alpha = 2 \, \text{GHz \, ns}$, and $\omega_0/2\pi = 20 \, \text{GHz}$ (blue curve); and the combination of both parallel and perpendicular currents (black curve). All cases include Gilbert damping $\lambda = 0.01$, but no thermal fluctuations.

5. Conclusion

In this work we explored the potential use of chirped spin currents to manipulate and control the magnetization dynamics. Such chirped currents could be produced by means of commercially available arbitrary waveforms generators, which can now reach the desired frequency range $^4$. We have shown that a chirped spin current polarized in the direction normal to the anisotropy axis can capture the magnetic moment into autoresonance and drive its precession to a direction normal to the anisotropy axis, it is possible to fully reverse the magnetization using a chirped spin current polarized in the direction perpendicular to the anisotropy axis. A second method to switch the magnetization relies on the combination of different types of spin currents. Different scenarios that combine chirped microwave fields with ac or dc spin currents could also be envisaged [21, 23] in the future.

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$^4$ See for instance: www.tek.com/signal-generator/awg5000-arbitrary-waveform-generator

References

[1] Hillebrands B and Fassbender J 2002 Nature 418 493
[2] Back C H et al 1999 Science 285 864
[3] Gerrits T, van den Berg H A M, Hohlfeld J, Bär L and Rasing T 2002 Nature 418 509
[4] Schumacher H W, Chappert C, Crozat P, Sousa R C, Freitas P P, Militat J, Fassbender J and Hillebrands B 2003 Phys. Rev. Lett. 90 017201
[5] Seki T, Usumiyi K, Nozaki Y, Yamamura H and Takahashi K 2013 Nat. Commun. 4 1726
[6] Slonczewski J C 1996 J. Magn. Magn. Mater. 159 L1
[7] Berger L 1996 Phys. Rev. B 54 9353
[8] Zhang S, Levy P M and Fert A 2002 Phys. Rev. Lett. 88 236601
[9] Stiles M D and Zangwill A 2002 Phys. Rev. B 66 014407
[10] Myers E B, Ralph D C, Katine J A, Louie R N and Buhrman R A 1999 Science 285 867
[11] Katine J A, Albert F J, Buhrman R A, Myers E B and Ralph D C 2000 Phys. Rev. Lett. 84 3149
[12] Wang K L, Alzate J G and Khalili Amiri P 2013 J. Phys. D: Appl. Phys. 46 074003
[13] Kiselev S I, Sankey J C, Krivorotov I N, Emley N C, Schoelkopf R J, Buhrman R A and Ralph D C 2003 Nature 425 380
[14] Sinova J, Valenzuela S O, Wunderlich J, Back C H and Jungwirth T 2015 Rev. Mod. Phys. 87 1213
[15] Bedau D, Liu H, Bouzaglou J-J, Kent D, Sun J Z, Katine J A, Fullerton E E and Mangin S 2010 Appl. Phys. Lett. 96 022514
[16] Swiebodzinski J, Chudnovskiy A, Dunn T and Kamenev A 2010 Phys. Rev. B 82 144404
[17] Cui Y T, Sankey J C, Wang C, Thadani K V, Li Z P, Myers E B, Ralph D C and Katine J A 1999 Science 285 864
[18] Fajans J, Gilson E and Friedland L 1999 Phys. Rev. B 60 014430
[19] Houssameddine D et al 2007 Nat. Mater. 6 447
[20] Bertotti G, Serpico C, Mayergoyz I D, Magni A, d'Aquino M and Bonin R 2005 Phys. Rev. Lett. 94 127206
[21] Fajans J, Gilson E and Friedland L 1999 Phys. Rev. Lett. 82 4444
[22] Meerson B and Friedland L 1990 Phys. Rev. A 41 5233–6
[23] Manfredi G and Hervieux P-A 2007 Appl. Phys. Lett. 91 061108
[24] Klughertz G, Hervieux P-A and Manfredi G 2014 J. Phys. D: Appl. Phys. 47 345004
[25] Klughertz G, Friedland L, Hervieux P-A and Manfredi G 2015 Phys. Rev. B 91 104433
[26] Feynman R, Vernon F L and Hellwarth R W 1957 J. Appl. Phys. 28 49
[27] Fajans J and Friedland L 2001 Am. J. Phys. 69 1096
[28] Zeng Z, Finocchio G and Jiang H 2013 Nanoscale 5 2219
[29] Barth I, Friedland L, Sarid E and Shagalov A G 2009 Phys. Rev. Lett. 103 155001
[30] Shalibo Y, Rofe Y, Barth I, Friedland L, Bialczack R, Martinis J M and Katz N 2012 Phys. Rev. Lett. 108 037701
[31] Dykman M I, Schwartz I B and Shapiro M 2005 Phys. Rev. E 72 021102
[37] Kim J-V 2006 Phys. Rev. B 73 174412
[38] Taniguchi T, Tsunegi S, Kubota H and Imamura H 2014 Appl. Phys. Lett. 104 152411
[39] Rippard W H, Deac A M, Pufall M R, Shaw J M, Keller M W, Russek S E, Bauer G E W and Serpico C 2010 Phys. Rev. B 81 014426
[40] Kubota H et al 2013 Appl. Phys. Express 6 103003
[41] Zeng Z et al 2012 ACS Nano 6 6115
[42] Slavin A and Tiberkevich V 2009 III Trans. Magn. 45 1875
[43] Rippard W H, Pufall M R, Kaka S, Silva T J, Russek S E and Katine J A 2005 Phys. Rev. Lett. 95 067203