Magnetic Moment Manipulation by a Josephson Current

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We consider a Josephson junction where the weak link is formed by a noncentrosymmetric ferromagnet. In such a junction, the superconducting current acts as a direct driving force on the magnetic moment. We show that the ac Josephson effect generates a magnetic precession providing then a feedback to the current. Magnetic dynamics result in several anomalies of current-phase relations (second harmonic, dissipative current) which are strongly enhanced near the ferromagnetic resonance frequency.

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Many interesting phenomena have been observed recently in the field of spintronics: the spin-dependent electric current and inversely the current-dependent magnetization orientation (see, for example, [1,2]). Moreover, it is well known that spin-orbit interaction may be of primary importance for spintronics, namely, for systems using a two-dimensional electron gas [3]. In the superconductor/ferromagnet/superconductor (S/F/S) Josephson junctions, the spin-orbit interaction in a ferromagnet without inversion symmetry provides a mechanism for a direct (linear) coupling between the magnetic moment and the superconducting current [4]. Similar anomalous properties have been predicted for Josephson junctions with a spin-polarized quantum point contact in a two-dimensional electron gas [5]. S/F/S junctions are known to reveal a transition to the π phase, where the superconducting phase difference φ in the ground state is equal to π [6]. However, the current-phase relation (CPR) in such a π-junction has a usual sinusoidal form $I = I_c \sin\varphi$, where the critical current I_c depends in a damped oscillatory manner on the modulus of the ferromagnet exchange field. In a noncentrosymmetric ferromagnetic junction, called hereafter φ_0-junction, the time reversal symmetry is broken, and the CPR becomes $I = I_c \sin(\varphi - \varphi_0)$, where the phase shift φ_0 is proportional to the magnetic moment perpendicular to the gradient of the asymmetric spin-orbit potential [4]. Therefore, manipulation of the internal magnetic moment can be achieved via the superconducting phase difference (i.e., by Josephson current).

In the present work, we study theoretically the spin dynamics associated with such φ_0-junctions. Though there is a lot of experimental progress in studying the static properties of S/F/S junctions, little is known about the spin-dynamics in S/F systems. Note here the pioneering work [7] where a sharpening of the ferromagnetic resonance was observed below the superconducting transition in $Nb/Ni_{80}Fe_{20}$ system. Theoretically, the single spin dynamics interplay with a Josephson effect has been studied in [8–11]. More recently, the dynamically induced triplet proximity effect in S/F/S junctions was studied in [12,13], while the junctions with composite regions (including several F regions with different magnetization) were discussed in [14,15]. Here, we consider a simple S/F/S φ_0-junction in a low frequency regime $\hbar\omega_J \ll T_c$ ($\omega_J = 2eV/\hbar$ being the Josephson angular frequency [16]), which allows us to use the quasistatic approach to treat the superconducting subsystem in contrast with the case analyzed in [12,13]. We demonstrate that a dc superconducting current may produce a strong orientation effect on the F layer magnetic moment. More interestingly, the ac Josephson effect, i.e., applying a dc voltage V to the φ_0-junction, would produce current oscillations and consequently magnetic precession. This precession may be monitored by the appearance of higher harmonics in CPR as well as a dc component of the current. In particular regimes, a total reversal of the magnetization could be observed. In the case of strong coupling between magnetic and superconducting subsystems, complicated nonlinear dynamic regimes emerge.

To demonstrate the unusual properties of the φ_0-junction, we consider the case of an easy-axis magnetic anisotropy of the F material (see Fig. 1). Both the easy axis and gradient of the asymmetric spin-orbit potential $\mathbf{n}$ are along the z-axis. Note that suitable candidates for the F interlayer may be MnSi or FeGe. In these systems, the lack

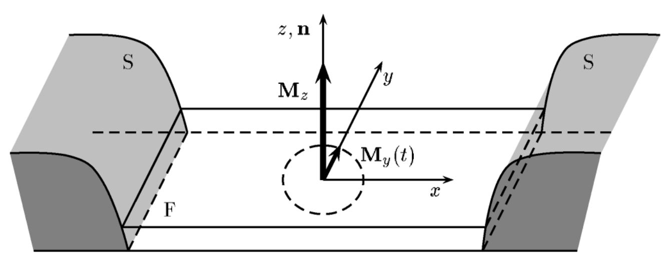

FIG. 1. Geometry of the considered φ_0-junction. The ferromagnetic easy-axis is directed along the z-axis, which is also the direction $\mathbf{n}$ of the gradient of the spin-orbit potential. The magnetization component $\mathbf{M}_y$ is coupled with Josephson current through the phase shift term $\varphi_0 \propto \mathbf{n}.(\mathbf{M} \wedge \nabla\Psi)$, where Ψ is the superconducting order parameter ($\nabla\Psi$ is along the x-axis in the system considered here).

of inversion center comes from the crystalline structure, but the origin of broken-inversion symmetry may also be extrinsic, like in a situation near the surface of a thin F film. In the following, we completely disregard the magnetic induction. Indeed, the magnetic induction in the (xy) plane is negligible for the thin F layer considered in this Letter, whereas the demagnetization factor cancels the internal induction along the z-axis ($N = 1$). The coupling between F and S subsystems due to the orbital effect has been studied in [17], and it appears to be very weak, and quadratic over magnetic moment $\mathbf{M}$ for the case when the flux of $\mathbf{M}$ through the F layer is small in comparison with flux quantum $\Phi_0 = h/2e$.

The superconducting part of the energy of a φ_0-junction is

$$E_s(\varphi, \varphi_0) = E_J[1 - \cos(\varphi - \varphi_0)], \tag{1}$$

where $E_J = \Phi_0 I_c/2\pi$ is the Josephson energy, I_c is the critical current, and φ_0 is proportional to the M_y component of the magnetic moment (see Fig. 1). Therefore, when the magnetic moment is oriented along the z-axis, we have the usual Josephson junction with $\varphi_0 = 0$. Assuming the ballistic limit we may estimate the characteristic Josephson energy as [6] $\Phi_0 I_c/S \sim T_c k_F^2 \sin\ell/\ell$ with $\ell = 4hL/\hbar v_F$, where S, L and h are the section, the length and the exchange field of the F layer, respectively. The phase shift is

$$\varphi_0 = \ell \frac{v_{\text{so}}}{v_F} \frac{M_y}{M_0} \tag{2}$$

where the parameter v_{so}/v_F characterizes the relative strength of the spin-orbit interaction [4]. Further on, we assume that $v_{\text{so}}/v_F \sim 0.1$. If the temperature is well below the Curie temperature, $M_0 = \|\mathbf{M}\|$ can be considered as a constant equal to the saturation magnetization of the F layer. The magnetic energy contribution is reduced to the anisotropy energy

$$E_M = -\frac{K\mathcal{V}}{2}\left(\frac{M_z}{M_0}\right)^2, \tag{3}$$

where K is an anisotropy constant and $\mathcal{V}$ is the volume of the F layer.

Naturally, we may expect that the most interesting situation corresponds to the case when the magnetic anisotropy energy does not exceed too much the Josephson energy. From the measurements [18] on permalloy with very weak anisotropy, we may estimate $K \sim 4 \times 10^{-5}$ K $\cdot$ Å^{-3}. On the other hand, typical value of L in S/F/S junction is $L \sim$ 10 nm and $\sin\ell/\ell \sim 1$. Then, the ratio of the Josephson over magnetic energy would be $E_J/E_M \sim 100$ for $T_c \sim$ 10 K. Naturally, in the more realistic case of stronger anisotropy, this ratio would be smaller, but it is plausible to expect a great variety of regimes from $E_J/E_M \ll 1$ to $E_J/E_M \gg 1$.

Let us now consider the case when a constant current $I < I_c$ is applied to the φ_0-junction. The total energy is (see, e.g., [16]):

$$E_{\text{tot}} = -\frac{\Phi_0}{2\pi}\varphi I + E_s(\varphi, \varphi_0) + E_M(\varphi_0), \tag{4}$$

and both the superconducting phase shift difference φ and the rotation of the magnetic moment $M_y = M_0 \sin\theta$ (where θ is the angle between the z-axis and the direction of $\mathbf{M}$) are determined from the energy minimum conditions $\partial_\varphi E_{\text{tot}} = \partial_{\varphi_0} E_{\text{tot}} = 0$. It results in

$$\sin\theta = \frac{I}{I_c}\Gamma \quad \text{with} \quad \Gamma = \frac{E_J}{K\mathcal{V}} \ell \frac{v_{\text{so}}}{v_F}, \tag{5}$$

which signifies that a superconducting current provokes the rotation of the magnetic moment M_y in the (yz) plane. Therefore, for small values of the rotation, $\theta(I)$ dependence is linear. In principle, the parameter Γ can be larger than one. In that case, when the condition $I/I_c \geq 1/\Gamma$ is fulfilled, the magnetic moment will be oriented along the y-axis. Therefore, applying a dc superconducting current switches the direction of the magnetization, whereas applying an ac current on a φ_0-junction could generate the precession of the magnetic moment.

We briefly comment on the situation when the direction of the gradient of the spin-orbit potential is perpendicular (along y) to the easy axis z. To consider this case, we simply need to take $\varphi_0 = \ell(v_{\text{so}}/v_F)\cos\theta$. The total energy (4) has two minima $\theta = (0, \pi)$, while applying the current removes the degeneracy between them. However, the energy barrier exists for the switch from one minimum into another. This barrier may disappear if $\Gamma > 1$ and the current is large enough $I > I_c/\Gamma$. In this regime, the superconducting current would provoke the switching of the magnetization between one stable configuration $\theta = 0$ and another $\theta = \pi$. This corresponds to the transitions of the junction between $+\varphi_0$ and $-\varphi_0$ states. The readout of the state of the φ_0-junction may be easily performed if it is a part of some SQUID-like circuit (the φ_0-junction induces a shift of the diffraction pattern by φ_0).

In fact, the voltage-biased Josephson junction, and thus the ac Josephson effect, provides an ideal tool to study magnetic dynamics in a φ_0-junction. In such a case, the superconducting phase varies with time like $\varphi(t) = \omega_J t$ [19]. If $\hbar\omega_J \ll T_c$, one can use the static value for the energy of the junction (4) considering $\varphi(t)$ as an external potential. The magnetization dynamics are described by the Landau-Lifshitz-Gilbert equation (LLG) [20]

$$\frac{d\mathbf{M}}{dt} = \gamma\mathbf{M} \times \mathbf{H}_{\text{eff}} + \frac{\alpha}{M_0}\left(\mathbf{M} \times \frac{d\mathbf{M}}{dt}\right), \tag{6}$$

where $\mathbf{H}_{\text{eff}} = -\delta F/\mathcal{V}\delta\mathbf{M}$ is the effective magnetic field applied to the compound, γ the gyromagnetic ratio, and α a phenomenological damping constant. The corresponding free energy $F = E_s + E_M$ yields

$$\mathbf{H}_{\text{eff}} = \frac{K}{M_0}\left[\Gamma \sin\left(\omega_J t - r\frac{M_y}{M_0}\right)\hat{\mathbf{y}} + \frac{M_z}{M_0}\hat{\mathbf{z}}\right], \qquad (7)$$

where $r = \ell v_{\text{so}}/v_F$. Introducing $m_i = M_i/M_0$, $\tau = \omega_F t$ ($\omega_F = \gamma K/M_0^2$ is the frequency of the ferromagnetic resonance) in LLG Eq. (6) leads to

$$\dot{m}_x = m_z(\tau)m_y(\tau) - \Gamma m_z(\tau)\sin(\omega\tau - rm_y)$$
$$\dot{m}_y = -m_z(\tau)m_x(\tau) \qquad \dot{m}_z = \Gamma m_x(\tau)\sin(\omega\tau - rm_y), \qquad (8)$$

where $\omega = \omega_J/\omega_F$. The generalization of Eq. (8) for $\alpha \neq 0$ is straightforward. One considers first the "weak coupling" regime $\Gamma \ll 1$ when the Josephson energy E_J is small in comparison with the magnetic energy E_M. In this case, the magnetic moment precess around the z-axis. If the other components verify $(m_x, m_y) \ll 1$, then the Eqs. (8) may be linearized, and the corresponding solutions are

$$m_x(t) = \frac{\Gamma\omega\cos\omega_J t}{1-\omega^2} \quad \text{and} \quad m_y(t) = -\frac{\Gamma\sin\omega_J t}{1-\omega^2}. \qquad (9)$$

Near the resonance $\omega_J \approx \omega_F$, the conditions of linearization are violated, and it is necessary to take the damping into account. The precessing magnetic moment influences the current through the φ_0-junction like

$$\frac{I}{I_c} = \sin\omega_J t + \frac{\Gamma r}{2}\frac{1}{\omega^2 - 1}\sin 2\omega_J t + \ldots; \qquad (10)$$

i.e., in addition to the first harmonic oscillations, the current reveals higher harmonics contributions. The amplitude of the harmonics increases near the resonance and changes its sign when $\omega_J = \omega_F$. Thus, monitoring the second harmonic oscillations of the current would reveal the dynamics of the magnetic system.

The damping plays an important role in the dynamics of the considered system. It results in a dc contribution to the Josephson current. Indeed, the corresponding expression for $m_y(t)$ in the presence of damping becomes

$$m_y(t) = \frac{\omega_+ - \omega_-}{r}\sin\omega_J t + \frac{\alpha_- - \alpha_+}{r}\cos\omega_J t, \qquad (11)$$

where

$$\omega_\pm = \frac{\Gamma r}{2}\frac{\omega \pm 1}{\Omega_\pm} \quad \text{and} \quad \alpha_\pm = \frac{\Gamma r}{2}\frac{\alpha}{\Omega_\pm}, \qquad (12)$$

with $\Omega_\pm = (\omega \pm 1)^2 + \alpha^2$. It thus exhibits a damped resonance as the Josephson frequency is tuned to the ferromagnetic one $\omega \to 1$. Moreover, the damping leads to the appearance of out of phase oscillations of $m_y(t)$ [term proportional to $\cos\omega_J t$ in Eq. (11)]. In the result, the current

$$I(t) \approx I_c\sin\omega_J t + I_c\frac{\omega_+ - \omega_-}{2}\sin 2\omega_J t + I_c\frac{\alpha_- - \alpha_+}{2}\cos 2\omega_J t + I_0(\alpha) \qquad (13)$$

acquires a dc component

$$I_0(\alpha) = \frac{\alpha\Gamma r}{4}\left(\frac{1}{\Omega_-} - \frac{1}{\Omega_+}\right). \qquad (14)$$

This dc current in the presence of a constant voltage V applied to the junction means a dissipative regime which can be easily detected. In some aspect, the peak of dc current near the resonance is reminiscent of the Shapiro steps effect in Josephson junctions under external r.f. fields. Note that the presence of the second harmonic in $I(t)$ Eq. (13) should also lead to half-integer Shapiro steps in φ_0-junctions [21].

The limit of the "strong coupling" $\Gamma \gg 1$ (but $r \ll 1$) can also be treated analytically. In this case, $m_y \approx 0$ and solutions of Eq. (8) yields

$$m_x(t) = \sin\left[\frac{\Gamma}{\omega}(1 - \cos\omega_J t)\right]$$
$$m_z(t) = \cos\left[\frac{\Gamma}{\omega}(1 - \cos\omega_J t)\right], \qquad (15)$$

which are the equations of the magnetization reversal, a complete reversal being induced by $\Gamma/\omega > \pi/2$. Strictly speaking, these solutions are not exact oscillatory functions in the sense that $m_z(t)$ turns around the sphere center counterclockwise before reversing its rotation, and returns to the position $m_z(t = 0) = 1$ clockwise, like a pendulum in a spherical potential [see Fig. 2(c)].

Finally, we have performed numerical studies of the nonlinear LLG Eq. (6) for some choices of the parameters when the analytical approaches fail. To check the consistency of our numerical and analytical approaches,

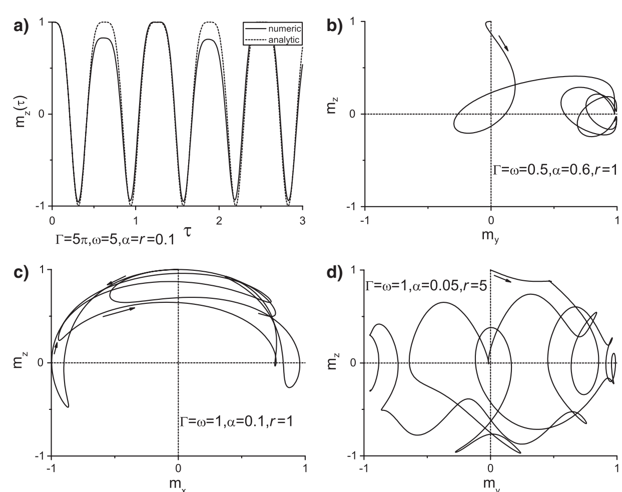

FIG. 2. Results of numerical analysis of the magnetic moment dynamics of the φ_0-junction. (a) Reversal of m_z from analytical expression Eq. (15) (dashed curve) and numerical one (normal curve). The other curves are related to the **M** trajectory: (b) in strong damping case (c) and (d) in the strong coupling regime revealing incomplete and complete magnetic moment reversal, respectively.

we present in Fig. 2(a) the corresponding $m_z(t)$ dependences for low-damping regimes. They clearly demonstrate the possibility of the magnetization reversal. In Figs. 2(b)–2(d), some trajectories of the magnetization vectors are presented for general coupling regimes. These results demonstrate that the magnetic dynamics of S/F/S φ_0-junction may be pretty complicated and strongly nonharmonic.

If the φ_0-junction is exposed to a microwave radiation at angular frequency ω_1, the physics that emerge are very rich. First, in addition to the Shapiro steps at $\omega_J = n\omega_1$, half-integer steps will appear. Second, the microwave magnetic field may also generate an additional magnetic precession with ω_1 frequency. Depending on the parameters of φ_0-junction and the amplitude of the microwave radiation, the main precession mechanism may be related either to the Josephson current or the microwave radiation. In the last case, the magnetic spin-orbit coupling may substantially contribute to the amplitude of the Shapiro steps. Therefore, we could expect a dramatic increase of this amplitude at frequencies near the ferromagnetic resonance. When the influence of the microwave radiation and Josephson current on the precession is comparable, a very complicated regime may be observed.

In the present work, we considered the case of the easy-axis magnetic anisotropy. If the ferromagnet presents an easy-plane anisotropy, then qualitatively the main conclusions of our article remain the same because the coupling between magnetism and superconductivity depends only on the M_y component. However, the detailed dynamics would be strongly affected by weak in-plane anisotropy.

To summarize, we have demonstrated that S/F/S φ_0-junctions provide the possibility to generate magnetic moment precession via Josephson current. In the regime of strong coupling between magnetization and current, magnetic reversal may also occur. These effects have been studied analytically and numerically. We believe that the discussed properties of the φ_0-junctions could open interesting perspectives for its applications in spintronics.

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