Quantization of Electric Flux in Spin Josephson Effect and Spin-FET

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We study spin Josephson effect in the presence of external magnetic and electric field. The magnetic field can be used to maintain spin supercurrent between two antiferromagnetically coupled atomic spins, or to overcome the dipole-dipole interaction in a junction made by ferromagnetic metal. The Josephson spin current couples to electric field via spin-orbit coupling or Aharonov-Casher effect, which implies a \(\phi\)-junction that can be controlled by electric field, and quantization of electric flux defined by \(\Phi_E = EL\), where \(L\) is the trajectory that electric field \(E\) is applied. This quantization also manifests itself in spintronic devices such as the current-voltage characteristics of spin-FET, and the persistent spin current due to Aharonov-Casher effect in a ring.

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Introduction.- The basic principle of spintronics is to utilize spin current, instead of charge current, to transport information. Various mechanisms such as spin field effect transistor[spin-FET][1], spin-transfer torque[2,3], and spin pumping[4] have been proposed to generate and control the spin current, which made building practical devices possible. The spin current generated by these means is analogous to charge current generated from a battery, which means a certain way of energy delivery is necessary to maintain the spin current. On the other hand, the discovery of spin Josephson effect[5–7,10] indicates the possibility to build devices with dissipationless spin current. By weakly coupling two systems that break \(SO(3)\) symmetry in spin space, a spontaneous spin current flows across the junction. Speaking in general terms, spin Josephson effect also covers the spin supercurrent induced between two ferromagnetically or antiferromagnetically coupled atomic spins, if the two spins show certain noncollinearity.[11,13]

A natural consequence of the spin current, in combination with angular momentum conservation, is the precession of order parameter.[8,10,14] in such a way that one Goldstone mode becomes dynamical while the whole junction remains in equilibrium. This fascinating phenomenon of spontaneous precession, often overlooked in the literature, serves as a new mechanism for a \(SO(3)\)-breaking order parameter to precess. However, such precession decays rapidly within the scale of nanoseconds due to Gilbert damping,[15] which also dissipates the energy stored in the junction and diminishes the spin supercurrent. Moreover, as demonstrated later, the direct dipole-dipole interaction between two sides of the junction is much larger than the Josephson effect, and diminishes the spin supercurrent within a time scale much shorter than the precession period. Hence to build a practical device, it is necessary to use certain handles, such as external fields, to overcome these mechanisms and maintain the spin supercurrent in a junction or between two atomic spins.

In this Letter, we study the effect of static external magnetic and electric field on the spin Josephson effect, and demonstrate their importance in building practical devices. There are two classes of spin Josephson effect in a junction according to the condensate that transports the Josephson spin current. In the first class, the spin supercurrent is carried out by spinful Cooper pairs, for instance the combination \(\langle c_\uparrow c^\dagger_\downarrow \rangle - \langle c^\dagger_\uparrow c_\downarrow \rangle\) in the triplet SC/insulator/triplet SC junction.[2] In the second class, the spinful particle-hole pair \(\langle c^\dagger_\downarrow c^\dagger_\uparrow \rangle\) does the job, as in ferromagnetic metal/insulator/ferromagnetic metal(FMM/I/FMM) junction.[6] We focus on the second class because of the following reasons: (a) The order parameter of FMM couples to external magnetic field directly and affects the Josephson spin current. (b) The charge neutral but spinful \(\langle c^\dagger_\downarrow c^\dagger_\uparrow \rangle\) couples to external electric field via spin-orbit interaction(SOI) or Aharonov-Casher(AC) effect. (c) There is no Josephson charge current in this problem. We show that magnetic field can be used to overcome the direct dipole-dipole interaction in such junction, as well as to maintain the spin supercurrent between two antiferromagnetically coupled atomic spins. The spin supercurrent oscillates periodically with external electric field, which leads to quantization of electric flux defined by electric field multiplied by the trajectory. As the magnetic flux quantum serves as a reference for precise measurement of magnetic field, we anticipate that electric flux quantum can have important applications in precise measurement of electric field. Such quantization also manifests itself in other spintronic devices, such as the current-voltage characteristics of spin-FET and persistent spin current in a ring due to AC effect.

FMM/I/FMM junction.- We start from the FMM/I/FMM junction without any external field. We follow the analysis in Ref. [3] but with a different mean field treatment of Hubbard interaction

\[
U n_\uparrow(r)n_\downarrow(r) \rightarrow -c^\dagger_\uparrow(r)c_\downarrow(r)\Delta(r) - c^\dagger_\downarrow(r)c_\uparrow(r)\Delta(r)^\dagger
\]

(1)

to describe bulk FMM, where \(\Delta^\dagger = U\langle S^+\rangle = U\langle S^x + iS^y\rangle = U\langle c^\dagger_\downarrow c^\dagger_\uparrow\rangle\). The reason to use this mean field scheme...
is that our magnetization lies in the \( S^x S^y \)-plane with an angle \( \theta \) that is also the phase of \( \Delta = |\Delta| e^{i\theta} \); hence the dynamics of magnetization is directly related to the uncertainty relation \( [\theta, S^z] = i \), as demonstrated below.

We can directly utilize the formalism in Ref. 3 to calculate the spin current by setting \( m = (2U/3)(S^z) = 0 \) therein. Considering the case \(|\Delta_L| = |\Delta_R| = |\Delta|\), and a phase difference \( \theta = \theta_L - \theta_R \). The spin supercurrent due to coherent tunneling of \( \langle c_L^+ c_R \rangle \) is

\[
J_s = \frac{1}{2} \langle N_{L\uparrow} - N_{L\downarrow} \rangle = \frac{|T|^2}{2} S(0, |\Delta|) \sin \theta
\]

\[
= J_s^0 \sin \theta = N \langle \hat{S}_L^z \rangle = -N \langle \hat{S}_R^z \rangle
\]

where \(|T|^2\) represents the tunneling amplitude, and \( N \) is total number of sites on either side of the junction. The function \( S(a, b) \) satisfies \( S(a, 0) = 0 \) and \( S(0, b) \neq 0 \). Eq. 2 implies that magnetization on each of the junctions rotates out of \( S^x S^y \)-plane but in opposite direction, hence the precession of the magnetization is realized. Notice that this precession does not affect our mean field treatment as long as \( \Delta \) is defined on a rotating reference frame.

Eq. 2 implies that the dynamics of the magnetization can be mapped into an effective two-spin model \( H_{2\text{-spin}} = J_{eff} S_L \cdot S_R \), where \( J_{eff} = J_s^0/N S^2 \), and \( S \) is the magnitude of magnetization per site. 1 13 Within this two spin model (which ignores the crucial direct dipole-dipole interaction of the two spin reservoirs, as addressed later), if the two spins are deviated from their collinear ground state by \( \theta \) in the \( S^x S^y \)-plane, then the spins precess due to uncertainty relation \( [\theta, S^z] = i \), as demonstrated in Fig. 1(a). This is the same principle that causes precession and spin supercurrent in the FMM/I/FMM junction. 1 13 The current is related to bound state energy stored in the two spins \( E_b = J_{eff} S^2 \cos \theta \) by

\[
\langle \hat{S}_L^z \rangle = -\partial E_b/\partial \theta
\]

which is analogous to the relation between Josephson current and Andreev bound state \( I = -\partial E_b/\partial \phi \).

Precession of the magnetization implies that the spin supercurrent described by Eq. 2 has its direction of polarization constantly change with time, due to rotation of \( S^x S^y \)-plane. Thus the spin supercurrent is an ac current from a lab observer’s point of view, even though the precession is analogous to dc Josephson effect. Defining the direction of magnetization as \( \langle S^z \rangle = S \), the spin supercurrent from a lab observer’s point of view is

\[
J_{lab}^x = -a J_{eff} S^2 \sin \theta \sin \omega t
\]

\[
J_{lab}^y = a J_{eff} S^2 \sin \theta \cos \omega t
\]

where \( a \) is the lattice constant, and \( \omega = 2 J_{eff} S \cos(\theta/2) \).

**Effect of static magnetic field.** We now discuss the effect of magnetic field by assuming that direct dipole-dipole interaction between two FMM’s can be ignored, and address the validity of this assumption later. The effect of static magnetic field can be described by adding the potential \(-\mu_B B_L S_L - \mu_B B_R S_R\) into the effective two-spin model where, in general, \( B_L \) and \( B_R \) do not have to be equal. Its effect is twofold. Firstly, for a uniform field of the scale of exchange coupling \( \mu_B B_R \sim J_{eff} \), the total energy is

\[
E_{total} = E_b + E_B = J_{eff} S^2 \cos \theta - 2 \mu_B B S \cos \frac{\theta}{2}
\]

The sign of \( J_{eff} \) depends on \( S(0, |\Delta|) \), which is determined by several bulk parameters. We emphasize that \( J_{eff} > 0 \) and \( J_{eff} < 0 \) junctions behave differently under magnetic field. For ferromagnetic coupling \( J_{eff} < 0 \), Lamour precession and Josephson precession are along the same direction, so magnetic field speeds up the precession. The precession then decays due to Gilbert damping. For antiferromagnetic coupling \( J_{eff} > 0 \), Eq. 5 implies an angle

\[
\theta_0 = 2 \arccos \left( \frac{\mu_B B}{2 J_{eff} S} \right)
\]

that minimizes \( E_{total} \), at which no precession occurs. This is because Lamour precession and Josephson precession are in opposite direction. Whatever is the initial angle, the system releases to \( \theta_0 \) through Gilbert damping, and then the two spins remain static in this energy minimum. However, the spin supercurrent persists due to Eq. 3. This spin current has a fixed direction of polarization that depends on the orientation of the two static spins. Therefore the application of magnetic field stops the precession at \( \theta = \theta_0 \) and creates a persistent, field-adjustable dc spin supercurrent for the \( J_{eff} > 0 \) case. This serves as a novel mechanism to maintain a spin supercurrent, up to a time scale that practical usage, such as transporting information, is feasible.
Although this magnetic field induced spin supercurrent is likely to be hindered by the direct dipole-dipole interaction in a realistic FMM/I/FMM junction, as addressed later, it can be realized in two atomic spins coupled antiferromagnetically via superexchange mechanism. [11] In such case, the spin current due to back and forth hopping of electrons is $J_{eff}S_R \times S_L$, analogous to Eqs. (2) and (3), which causes two spins to precess. The dipole-dipole interaction is typically several orders smaller than $J_{eff}$. Applying magnetic field controls the angle of two spins via Eq. (8), stops the precession and hence no Gilbert damping, and yields a static interatomic spin current via Eq. (9). Applying this mechanism to a bulk antiferromagnet is a more delicate issue as it involves the quantum fluctuations in the canted antiferromagnetic state, [16–21] which will be addressed elsewhere.

Secondly, if $g\mu_B B_R \neq g\mu_B B_L$, but the two fields are of the same scale as $J_{eff}$ (assuming dipole-dipole interaction can be ignored), the spins rotate with time and the relative angle $\theta(t)$ changes with time, analogous to ac Josephson effect. [10] In such case, the equation of motion is nonlinear and there is no guarantee that $\theta(t)$ would be periodic in time, unlike in SC Josephson effect where the voltage across the junction determines the periodicity of phase difference. Note that SC Josephson effect has only one energy scale $\mu$ that determines the dynamics of the $U(1)$ phase, whereas in spin Josephson effect there are two energy scales $\{\mu_B (B_R - B_L), J_{eff}\}$ that determine the dynamics of $\theta(t)$. Lastly, if $g\mu_B B_i \gg J_{eff}$, then the two spins are pinned along the direction of $B_i$ with negligible precession or vibration.

Effect of static electric field. We now address the effect of electric field on the FMM/I/FMM junction. Applying a gate voltage on the thin insulating interface has two effects. Firstly, for a 3D junction, it changes the potential barrier at the interface, thus changing the tunneling amplitude $|T|^2$. Therefore one should replace $J_0^s \rightarrow J_0^s(E)$ and $J_{eff} \rightarrow J_{eff}(E)$. Secondly, for an interface with Rashba SOI, the propagation of spinful $\langle c_{1}^{\dagger}c_{j}\rangle$ picks up a phase that depends linearly on the electric field. For an interface without SOI, the AC effect gives a phase of the same form. This can be understood by constructing a Ginzburg-Landau(GL) free energy for the FMM order parameter that contains AC phase

$$f = f_{\text{m}} + \alpha|\psi|^2 + \frac{\beta}{2}|\psi|^4 - \frac{\hbar^2}{2m}(\partial_x - ik_{0x})\psi|^2$$

(7)

where $k_{0x} = \mu \times (\gamma + \lambda E)$ and $k_{0x}$ is its component along the junction(not to be confused with the rotating reference frame of magnetization), and $m$ represents the effective mass. $\gamma$ represents the intrinsic Rashba SOI and $\lambda$ characterizes the field dependence, and $\lambda = 1/\hbar c^2$ if only AC effect is present. Defining the dimensionless quantity

$$g(x) = \psi(x)/\psi_{\infty} = v_x/(-\alpha/\beta)$$

its value in the interface $0 < x < L$ where electric field is applied is

$$g(x) = 1 - \frac{x}{L} e^{ik_{0x}} + \frac{2}{L} e^{-ik_{0x}(L-x)+i\theta}$$

(8)

such that it satisfies the boundary condition $g(0) = 1$ and $g(L) = e^{i\theta}$, and the Laplace equation $(\partial_x - ik_{0x})^2 g = 0$. [22] The GL free energy integrated over the interface is $\Delta F = A\hbar^2 \psi_2^2/Lm [1 - \cos(\theta - k_{0x}L)]$, where $A$ is the cross section area. Therefore the AC phase changes the bound state energy, and by Eq. (4) or by expression of current from Eq. (7) one gets

$$J_s = J_s^0(E)\sin(\theta - \varphi_{AC})$$

(9)

where $\varphi_{AC} = \mu \times (\gamma + \lambda E) \cdot \hat{x}$. If direction of the two spins are fixed by anisotropy or strong magnetic field (see below) then $\mu$ is fixed and Eq. (9) implies a $\varphi$-junction that the phase shift can be arbitrarily controlled by the gate voltage. In this case, one certainly has to disregard the angular momentum conservation and the precession of magnetization described by Eq. (2). However, as long as the two magnetizations point at different directions, the spin supercurrent still flows with polarization $\mu \parallel S_R \times S_L$. Therefore controlling the spin supercurrent by electric field, as described by Eq. (9), is still feasible. This situation is very similar to the spin-FET,[1] in the sense that transmission of spin current across two FMM’s can be controlled by gate voltage, except we are in the quantum regime and the spin current is spontaneous.

Consider the case that the magnetization of the two FMM’s are fixed by strong magnetic field or anisotropy, so $\mu$ is fixed and $E \times \mu \parallel \hat{x}$. The periodicity in the argument of Eq. (9) implies quantization of one dimensional electric flux

$$\Phi_E = EL = n \frac{2\pi}{\lambda \mu} = n\Phi_E^0.$$  

(10)

If only AC effect is present and assuming $\langle c_{1}^{\dagger}c_{j}\rangle$ has magnetic moment $g\mu_B$ with $g \approx 2$, then $\Phi_E^0 = \hbar c^2/g\mu_B = 3.215 \times 10^6 \sqrt{V}$. The spin supercurrent oscillates with periodicity $\Phi_E^0$, and in principle can be measured by detecting the induced electric dipole field. [23, 24] Although the flux quantum $\Phi_E^0$ has been mentioned schematically in Ref. [23], it was not until our proposal does this flux quantum have a concrete meaning, i.e. from the periodicity of Josephson spin current. In reality, the SOI is important to observe this quantization. Because for AC effect alone in a thin interface of size $\sim \mu m$, it requires an unreasonably large field to even observe one flux quantum. This reflects the fact that AC effect is extremely small compare to other interference effects such as Aharonov-Bohm(AB) effect. [25]

In reality, the oscillation with flux quantum may be smeread out due to the field dependent critical current $J_s^s(E)$. This can be overcomed by considering a planar junction that sandwiches a 2D electron gas between two
2D FMM’s (2DFMM/2DEG/2DFMM). Then the gate voltage does not change the tunneling barrier because the motion of \langle c_{1\uparrow}c_{1\downarrow}\rangle is confined in 2D, so \( J^0_s(E) \rightarrow J^0_{seff} \). In addition, the strong SOI in 2DEG may reduce the value of flux quantum to a more observable region. For a 2D junction of the size \( \mu m^2 \), one can estimate \( J^0_{seff} \) by assuming \( |T| \sim meV, \Delta \sim 10^{-1}eV, \) and density of states \( N_F \sim 1/eV, \) and we adopt the formula of critical current for estimation \( I_{c} = g_{\mu B} \pi^2 N^2 \Delta |T|^2 \gamma /h \sim 10^5 g_{\mu B}/s \), \( \gamma \sim 10^9 \) is number of cross section channels by assuming lattice constant \( a \sim nm \). The induced dipolar field \( \mathbf{D} \) gives a voltage drop \( \sim 10^{-3}V \) for two points that are \( r \sim 10^{-5} \) away from the junction, which is within the precision of current technique. This yields \( J_{seff} \sim 10^{-15}eV \) and the precession period \( \sim s \).

However, in contrast to two coupled atomic spins, the direct dipole-dipole energy in a FMM/1/FMM junction is not negligible. For one spin on the left, the dipole-dipole energy from all the \( 10^6 \) spins on the right is \( U_{dip} \sim 10^{-10}eV \), which is much larger than \( J_{seff} \) and tends to align the magnetizations along the direction of the junction. This means even if the two spins are prepared to be noncollinear, their dynamics is not determined by the spin Josephson effect but the dipole-dipole interaction, so the precession of spins and the ac spin current described by Eq. 1 are unlikely to be observed in a realistic junction. Moreover, the spin current decays within the scale of \( h/U_{dip} \sim 10^{-8} \). This motivates us to propose the device in Fig. 1(b) that uses magnetic field to overcome the dipole-dipole interaction. A current carrying wire is placed above and along the direction of interface, which produces a magnetic field that is larger than the dipole-dipole energy but smaller than the Hubbard interaction in Eq. 1. \( U_{dip} \ll g_{\mu B}B \ll U \).

The magnetic field pins the direction of magnetization, so no precession occurs, but the spin current and the effect of electric field are well described by Eq. 1. Ideally, the spin current persists as long as the magnetic field is present.

Eq. 1 demonstrates a significant difference between spin Josephson effect and SC Josephson effect. In SC Josephson effect, although propagation of Cooper pairs is coupled to the gauge field via AB effect, the effect of gauge field does not show up unless the junction is made into a closed trajectory such that the phase integrated along the path \((2e/h) \oint A \cdot dl = 2\pi \Phi_B / \Phi_0 \) is gauge invariant. Superconducting quantum interference devices (SQUID) are designed according to this principle. In spin Josephson effect, the propagation of \langle c_{1\uparrow}c_{1\downarrow}\rangle couples to \( \mathbf{E} \times \mathbf{\mu} \) which is not a gauge field, therefore one can apply an electric field to control the Josephson spin current even in an open trajectory like Fig. 1(b).

The quantization of electric flux should also show up in spintronic devices that use polarized electrons to transport the spin current, as long as the electrons remain phase coherent along the trajectory. For instance, the original proposal of spin-FET in Ref. 1 already implies this quantization. In an ideal 1D spin-FET, the spin current is proportional to \( \cos^2(\Delta \theta/2) \), where \( \Delta \theta = (k_1 - k_1)L = 2mnL/h^2 \). The SOI parameter \( \eta \) corresponds to \( h^2 \mu e/\gamma \lambda E/m \) from Hamiltonian of the form of the last term in Eq. 7. The period of oscillation indicates \( \Delta \theta = 2\pi n = 2\eta \mu eL = 2\pi \Phi_E / \Phi_0 \). Therefore the electric flux, once again defined as field multiplied by the trajectory, is quantized with the same flux quantum (if one assumes \( \mu = 2\mu_e \) in Eq. 10). Since the field is controlled by the gate voltage, this quantization should show up as the periodicity of current-voltage characteristics. Note that the intrinsic SOI is \( \eta \sim 3.9 \times 10^{-12}eVm \) in 2DEG, which gives \( L(\Delta \theta = \pi) = 0.67\mu m \). Assuming \langle c_{1\uparrow}c_{1\downarrow}\rangle in the 2DFMM/2DEG/2DFMM junction experiences the same order of \( \eta \) and \( L(\Delta \theta = \pi) \), then Eq. 1 implies one can control the phase shift of \( \chi \)-junction by changing the width of 2DEG in \( \mu m \) scale, which is within the reach of current techniques.

The quantization in spin-FET is analogous to the persistent charge current in a metallic ring, where the quantization of magnetic flux is seen as long as the electrons remain phase coherent. Whether the charge current is carried by electrons or Cooper pairs is not essential to this quantization. Along this line of argument, the persistent spin current in a ring due to AC effect also manifests this quantization. Following Ref. 20, this can be seen from the spin-dependent phase shift of eigen energies \( E_n(\Phi_{AB} + \sigma_z \Phi_{AC}) \). The total energy \( E = \sum_n E_n \) summing over occupied states changes periodically with \( \Phi_E / \Phi_0 \), and so is the spin current calculated by \( J_s \propto \gamma \sigma_z \partial E / \partial \Phi_{AC} \).

Conclusions.- In summary, we show that magnetic field can be used to maintain the spin supercurrent between two antiferromagnetically coupled atomic spins, or to overcome dipole-dipole interaction in a FMM/1/FMM junction. The spin current couples to the electric field via SOI or AC effect, which implies a \( \chi \)-junction that can be controlled by the electric field or width of the interface. The spin current oscillate periodically with electric field, which gives a concrete meaning to the quantization of electric flux. In the case of pure AC effect, the flux quantum \( \Phi_E^0 = h c^2 / g_{\mu B} \) is determined only by fundamental constants.

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