Controllable π junction with magnetic nanostructures

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We propose a Josephson device in which 0 and π states are controlled by an electrical current. In this system, the π state appears in a superconductor/normal metal/superconductor junction due to the nonlocal spin accumulation in the normal metal which is induced by spin injection from a ferromagnetic electrode. Our proposal offers not only possibilities for the application of superconducting spin-electronic devices but also the in-depth understanding of the spin-dependent phenomena in magnetic nanostructures.

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Nowadays spin-electronics is one of the central topics in condensed matter physics.[1–3] There has been considerable interest in the spin injection, accumulation, transport, and detection in ferromagnet/normal metal (F/N) hybrid structures.[4–10] Twenty years ago, Johnson and Silsbee demonstrated the spin injection and detection in a F/N/F structure for the first time.[4] Recently, spin accumulation has been observed at room temperature in all-metallic spin-valve geometry consisting of a F/N/F junction by Jedema *et al.*[5] In their system, the spin-polarized bias current is applied at one F/N junction, and the voltage is measured at another F/N interface, for the parallel (P) and antiparallel (AP) alignments of the Fs magnetizations. They observed the difference of the nonlocal voltages between the P and AP alignments due to spin accumulation in N. Also in a F/I/N/I/F (I indicates an insulator) structure, clear evidence of spin accumulation in N has been shown.[6] In hybrid structures consisting of a ferromagnet and a superconductor (S), a suppression of the superconductivity due to spin accumulation in S has been studied theoretically and experimentally.[11–13]

Furthermore, ferromagnetic Josephson (S/F/S) junctions have been studied actively in recent years.[14–19] In the S/F/S junctions, the pair potential oscillates spatially due to the exchange interaction in F.[14,15] When the pair potentials in two Ss take different sign, the direction of the Josephson current is reversed compared to that in ordinary Josephson junctions. This state is called the π state in contrast with the 0 state in ordinary Josephson junctions because the current-phase relation of the π state is shifted by "π" compared to that of the 0 state. The observations of the π state have been reported in various systems experimentally.[16–19] The applications of the π state to quantum computing also have been proposed.[20–22] Another system to realize the π state is a S/N/S junction with a voltage-control channel.[23,24] In the system, the nonequilibrium electron distribution in N induced by the bias voltage plays an important role, and the sign reversal of the Josephson critical current as a function of the control voltage has been demonstrated.[23,24]

In this paper, we propose a Josephson device in which the 0 and π states are controlled electrically. In this device, spin accumulation is generated in a nonmagnetic metal by the spin-polarized bias current flowing into the nonmagnetic metal from a ferromagnet. In a metallic Josephson junction consisting of the spin accumulated nonmagnetic metal sandwiched by two superconductors, the π state appears due to the spin split of the electrochemical potential in the nonmagnetic metal. The magnitude of the spin accumulation is proportional to the value of the spin-polarized bias current, and therefore the state of the Josephson junction is controlled by the current. Our proposal leads to an in-depth understanding of the spin-dependent phenomena in magnetic nanostructures as well as possibilities for the application of superconducting spin-electronic devices.

We consider a magnetic nanostructure with two superconductors as shown in Fig. 1. The device consists of a nonmagnetic metal N (the width $w_{\rm N}$, the thickness $d_{\rm N}$) which is connected to a ferromagnetic metal F (the width $w_{\rm F}$, the thickness $d_{\rm F}$) at $x=0$ and sandwiched by two superconductors S1, S2 located at $x=L$. In this device, the electrode F plays a role as a spin-injector to the electrode N, and the S1/N/S2 junction is a metallic Josephson junction. The spin-diffusion length $\lambda_{\rm N}$ in N is much longer than the length $\lambda_{\rm F}$ in F,[4–8] and we consider the structure with dimensions of $\lambda_{\rm F} \ll (w_{\rm N(F)}, d_{\rm N(F)}) \ll \lambda_{\rm N}$ which is a realistic geometry.[5,6]

In the electrodes N and F, the electrical current with spin σ is expressed as

$$\mathbf{j}_\sigma = -(\sigma_\sigma/e)\boldsymbol{\nabla}\mu_\sigma, \tag{1}$$

where σ_σ and μ_σ are the electrical conductivity and the electrochemical potential (ECP) for spin σ, respectively. Here ECP is defined as $\mu_\sigma = \epsilon_\sigma + e\phi$, where ϵ_σ is the chemical potential of electrons with spin σ and ϕ is the electric potential. From the continuity equation for charge, $\boldsymbol{\nabla}\cdot(\mathbf{j}_\uparrow + \mathbf{j}_\downarrow) = 0$, and that for spin, $\boldsymbol{\nabla}\cdot(\mathbf{j}_\uparrow - \mathbf{j}_\downarrow) = e\,\partial(n_\uparrow - n_\downarrow)/\partial t$ (n_σ is the carrier density for spin σ), we obtain[8,10]

$$\boldsymbol{\nabla}^2(\sigma_\uparrow\mu_\uparrow + \sigma_\downarrow\mu_\downarrow) = 0, \tag{2}$$

$$\boldsymbol{\nabla}^2(\mu_\uparrow - \mu_\downarrow) = (\mu_\uparrow - \mu_\downarrow)/\lambda^2, \tag{3}$$

where $\lambda = \sqrt{D\tau_{sf}}$ is the spin diffusion length with the diffusion constant $D = (N_\uparrow + N_\downarrow)/(N_\uparrow D_\downarrow^{-1} + N_\downarrow D_\uparrow^{-1})$ (N_σ and D_σ are the density of states and the diffusion constant for spin σ, respectively) and the scattering time of an electron $\tau_{sf} = 2/(\tau_{\uparrow\downarrow}^{-1} + \tau_{\downarrow\uparrow}^{-1})$ ($\tau_{\sigma\bar{\sigma}}$ is the scattering time of an electron from spin σ to $\bar{\sigma}$). In order to derive Eqs. (2) and (3), we take the relaxation-time approximation for the carrier density,

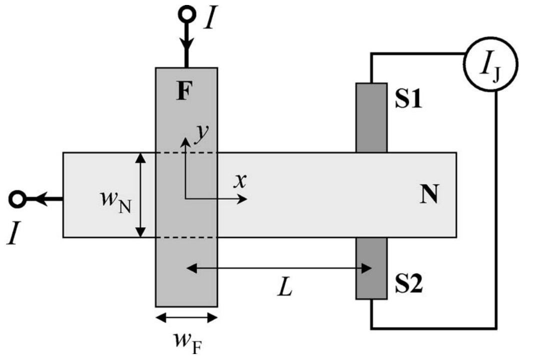

FIG. 1. Structure of a controllable π junction with magnetic nanostructures. The bias current I flows from a ferromagnet (F) to the left side of a normal metal (N). The Josephson current I_J flows in a superconductor/normal metal/superconductor (S1/N/S2) junction located at $x=L$.

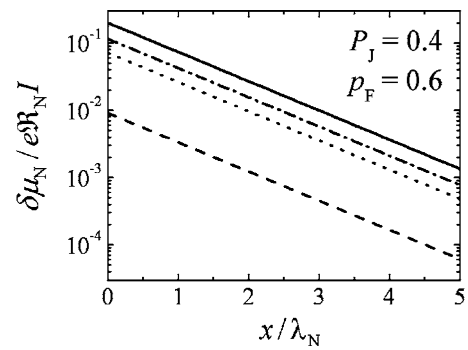

FIG. 2. Spatial variation of the split of the electrochemical potential in N. The solid line is for the tunnel-limit case ($R \gg \Re_N, \Re_F$), the dashed, dotted, and dot-dashed lines are for the metallic-limit cases ($R=0$) with $r=\Re_F/\Re_N=0.01$, 0.1, and 0.2, respectively.

$\partial n_\sigma / \partial t = -\delta n_\sigma / \tau_{\sigma\bar{\sigma}}$, and use the relations $\sigma_\sigma = e^2 N_\sigma D_\sigma$ and $\delta n_\sigma = N_\sigma \delta\epsilon_\sigma$, where δn_σ and $\delta\epsilon_\sigma$ are the carrier density deviation from equilibrium and the shift in the chemical potential from its equilibrium value for spin σ, respectively. In addition, the detailed balance equation $N_\uparrow \tau_{\uparrow\downarrow}^{-1} = N_\downarrow \tau_{\downarrow\uparrow}^{-1}$ is also used. We use the notations $\sigma_N = 2\sigma_N^\uparrow = 2\sigma_N^\downarrow$ in N and $\sigma_F = \sigma_F^\uparrow + \sigma_F^\downarrow$ ($\sigma_F^\uparrow \neq \sigma_F^\downarrow$) in F hereafter.

At the interface between N and F, the interfacial current I_σ flows due to the difference of ECPs in N and F: $I_\sigma = (G_\sigma/e)(\mu_F^\sigma|_{z=0^+} - \mu_N^\sigma|_{z=0^-})$, where G_σ is the spin-dependent interfacial conductance. We define the interfacial charge and spin currents as $I = I_\uparrow + I_\downarrow$ and $I_{spin} = I_\uparrow - I_\downarrow$, respectively. The spin-flip effect at the interface is neglected for simplicity. In the electrode N with the thickness and the contact dimensions being much smaller than the spin-diffusion length ($d_N, w_N, w_F \ll \lambda_N$), μ_N^σ varies only in the x direction.[8] The charge and spin current densities in N, $j = j_\uparrow + j_\downarrow$ and $j_{spin} = j_\uparrow - j_\downarrow$, are derived from Eqs. (1)–(3), and satisfy the continuity conditions at the interface: $j = I/A_N$ and $j_{spin} = I_{spin}/A_N$, where $A_N = w_N d_N$ is the cross-sectional area of N. From these conditions, we obtain ECP in N, $\mu_N^\sigma(x) = \bar{\mu}_N + \sigma\delta\mu_N$, where $\bar{\mu}_N = (eI/\sigma_N A_N)x$ for $x<0$, $\bar{\mu}_N = 0$ for $x>0$, and $\delta\mu_N = (e\lambda_N I_{spin}/2\sigma_N A_N)e^{-|x|/\lambda_N}$. In the electrode F, the spin split of ECP, $\delta\mu_F^\sigma$, decays in the z-direction because the thickness of F and the dimension of the interface are much larger than the spin-diffusion length in F ($d_F, w_N, w_F \gg \lambda_F$).[8] In a similar way to the case of N, ECP in F is obtained from the continuity conditions for charge and spin currents. ECP in F is expressed as $\mu_F^\sigma(z) = \bar{\mu}_F + \sigma\delta\mu_F^\sigma$, where $\bar{\mu}_F = (eI/\sigma_F A_J)z + eV$ and $\delta\mu_F^\sigma = [e\lambda_F(p_F I - I_{spin})/2\sigma_F^\sigma A_J]e^{-z/\lambda_F}$ with the contact area $A_J = w_N w_F$, the voltage drop at the interface $V = (\bar{\mu}_F - \bar{\mu}_N)/e$, and the polarization of the current in F, $p_F = (\sigma_F^\uparrow - \sigma_F^\downarrow)/\sigma_F$. The influence of the electrodes S1 and S2 on ECP in N may be neglected. When the superconducting gap in S1 and S2 is much larger than the spin split $\delta\mu_N$ at $x=L$, almost no quasiparticle is excited above the gap at low temperature. Therefore, the spin current does not flow into S1 and S2, and the behavior of ECP in N is not modified by the connection to the electrodes S1 and S2.

In order to obtain the relation between the bias current I and the shift of ECP, $\delta\mu_N$, at the right side in N ($x>0$), we substitute the obtained μ_N^σ and μ_F^σ for the expressions of I and I_{spin}, and eliminate V. As a result, we obtain the relation between I and I_{spin}, and finally we get the relation between I and $\delta\mu_N$ as follows:

$$\delta\mu_N(x) = e\Re_N I \frac{\dfrac{P_J}{1-P_J^2}\left(\dfrac{R}{\Re_N}\right) + \dfrac{p_F}{1-p_F^2}\left(\dfrac{\Re_F}{\Re_N}\right)}{1 + \dfrac{2}{1-P_J^2}\left(\dfrac{R}{\Re_N}\right) + \dfrac{2}{1-p_F^2}\left(\dfrac{\Re_F}{\Re_N}\right)} e^{-x/\lambda_N}, \tag{4}$$

where $\Re_N = \lambda_N/(\sigma_N A_N)$ and $\Re_F = \lambda_F/(\sigma_F A_J)$ indicate the nonequilibrium resistances of N and F, respectively, $R = G^{-1} = (G_\uparrow + G_\downarrow)^{-1}$ is the interfacial resistance, and $P_J = (G_\uparrow - G_\downarrow)/G$ is the polarization of the interfacial current. When the F/N interface is the tunnel junction ($R \gg \Re_N, \Re_F$), Eq. (4) reduces to a simple form $\delta\mu_N(x) = (e\Re_N I P_J/2)e^{-x/\lambda_N}$. On the other hand, when the F/N junction is of a metallic contact ($R=0$), Eq. (4) becomes $\delta\mu_N(x) = e\Re_N I p_F r e^{-x/\lambda_N}/(2r + (1-p_F^2))$, where $r = \Re_F/\Re_N$ is a mismatch factor of the resistances in F and N. Figure 2 shows the spacial variation of $\delta\mu_N(x)$ both for the tunnel- and metallic-limit cases with $P_J = 0.4$ and $p_F = 0.6$ (Refs. 1 and 25). As shown in this figure, in the case of the metallic contact, $\delta\mu_N$ becomes larger with decreasing the resistance mismatch.[8]

Next we consider how spin accumulation affects the Josephson current I_J flowing through the S1/N/S2 junction located at $x=L$ (Fig. 1). In the metallic Josephson junction, the Andreev bound state plays a key role for the Josephson effect.[18,26] The Andreev bound state is formed by a multiple Andreev reflection of an electron with the wave number $k_e = (\sqrt{2m}/\hbar)\sqrt{E_F + E}$ and a hole with $k_h = (\sqrt{2m}/\hbar)\sqrt{E_F - E}$, respectively, where E is the energy of the electron and hole measured from the Fermi energy E_F. As shown in Fig. 3, when there is the spin split $\delta\mu_N$ in N, a spin-up (-down) electron with the energy $E \approx \delta\mu_N$ ($-\delta\mu_N$) is injected into S from N at low temperatures. The injected electron captures another electron with the energy $E \approx -\delta\mu_N$ ($\delta\mu_N$) from the opposite spin band in order to form a Cooper pair in S. Therefore, a spin-up (-down) hole with the energy $E \approx \delta\mu_N$

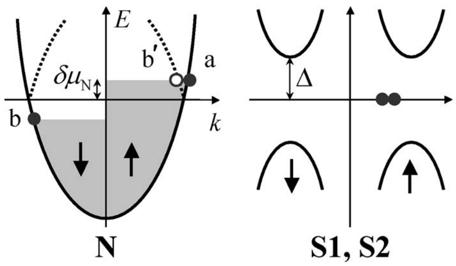

FIG. 3. Schematic diagram of energy vs momentum in the Andreev reflection when there is spin accumulation in N. The filled and open circles represent an electron and a hole, respectively. In N, the solid and dashed lines denote electron and hole bands, respectively, the shaded area indicates an occupation by electrons. In the Andreev reflection, a spin-up electron (a) injected into Ss captures another electron with spin down (b), and a spin-up hole (b′) is reflected back to N.

($-\delta\mu_N$) is reflected back to N (Andreev reflection).[26] In other words, the spin-up (-down) electron with $k_e \approx (\sqrt{2m}/\hbar)\sqrt{E_F+(-)\delta\mu_N}$ and the spin-up (-down) hole with $k_h \approx (\sqrt{2m}/\hbar)\sqrt{E_F-(+)\delta\mu_N}$ mainly contribute to the formation of a Cooper pair. Note that the values of the wave numbers k_e and k_h differ due to the spin split $\delta\mu_N$ in contrast with the case of a no spin split ($\delta\mu_N=0$) in which $k_e \approx k_h$.

The split $\delta\mu_N$ corresponds to the exchange energy E_{ex} of a ferromagnet in a superconductor/ferromagnet/superconductor (S/F/S) Josephson junction as follows:[14–19] In the S/F/S systems, Cooper pairs are formed by the Andreev reflection of spin-σ electrons with the wave number $k^F_{e,\sigma} \approx (\sqrt{2m}/\hbar)\sqrt{E_F+\sigma E_{ex}}$ and holes with $k^F_{h,\sigma} \approx (\sqrt{2m}/\hbar)\sqrt{E_F-\sigma E_{ex}}$ at the energy $E\approx 0$. In the case that the exchange interaction is much weaker than the Fermi energy ($E_{ex} \ll E_F$), the stable state (0 or π) in the system depends on the dimensionless parameter $\alpha_F=(E_{ex}/E_F)(k_F d_F)$, where d_F is the thickness of F and k_F is the Fermi wave number.[18] At $\alpha_F=0$ the system is in the 0 state, and the first 0-π transition occurs at $\alpha_F=\pi/2$, and then the system is in the π state at $\alpha_F=\pi$.[18] Because the value of E_{ex} is fixed in the S/F/S system, the 0 and π states change periodically with the period $2\pi(E_F/E_{ex})$ as a function of d_F. As a result, the d_F dependence of the Josephson critical current shows a cusp structure and the critical current becomes minimum at the 0-π transition.[17,18]

In analogy with the case of the S/F/S junction discussed above, when there is spin accumulation in N as shown in Fig. 3, the 0 or π state is realized in the S1/N/S2 junction depending on the parameter $\alpha=(\delta\mu_N/E_F)(k_F w_N)$. In this case, the width w_N is fixed, and the 0 and π states are controlled through the value of $\delta\mu_N$ which is proportional to the bias current I [see Eq. (4)]. The N part of the system is in the nonequilibrium state by the spin current in contrast with F in the equilibrium state of the S/F/S junction. However, one can discuss the critical current in the nonequilibrium S1/N/S2 junction in the same way as the equilibrium S/F/S junction because the critical current is dominated by the energy of the quasiparticles in N, not by the flow of the current.[23,24] Although we discuss the single-channel Josephson junction for simplicity so far, in the multichannel case the Josephson coupling varies in the x direction because the spin split of ECP decays in the x direction. The total Josephson coupling is expressed as an integral of the local Josephson coupling for the x direction, and therefore the ground state in the Josephson junction shows a similar dependence on the bias current to that in the single-channel case.

From the point of view of a more detailed description, the free energy in the system is obtained by the summation of the energy of the Andreev bound states.[20] The bound state energy is calculated from the Bogoliubov–de Gennes equation,[27] and the free energy is minimum for the phase difference 0 (π) for the 0 (π) state. In the S1/N/S2 junction with no spin accumulation in N ($\delta\mu_N=0$), the bound states with the energy $E>0$ contribute to the free energy. On the other hand, when the spin accumulation exists in N, the spin-up (-down) bound states with the energy $E>\delta\mu_N$ ($-\delta\mu_N$) contribute to the free energy because ECP is shifted by $\delta\mu_N$ ($-\delta\mu_N$) in N. The 0-π transition occurs due to the shift of the energy region of the Andreev bound states which contribute to the free energy.

As an example, we consider the case that the F/N interface consists of a tunnel junction. The material parameters $P_J=0.4$, $\rho_N=\sigma_N^{-1}=2\ \mu\Omega$ cm, $\lambda_N=1\ \mu$m, $w_N=800$ nm, and $d_N=10$ nm, which lead to $\Re_N=2.5\ \Omega$, are taken. The distance between F and Ss is taken to be $L=500$ nm. When no bias current is applied between F and N ($I=0$), the S1/N/S2 junction is in the ordinary 0 state because there is no spin split of ECP ($\delta\mu_N=0$). With increasing the bias current, the magnitude of the Josephson critical current decreases because the parameter α increases due to the increase of the spin split. When the bias current reaches the value $I=I_0 \approx 3$ mA which induces the spin split $\delta\mu_N \approx 1$ meV at $x=500$ nm, the parameter $\alpha \approx \pi/2$ and the first transition to the π state from the 0 state occurs (the values of $E_F=5$ eV and $k_F=1$ Å^{-1} are taken).[28] As a result, the magnitude of the Josephson critical current takes its minimum at $I=I_0$, and increases with the increasing bias current $I>I_0$. When the bias current attains $I=2I_0$, the magnitude of the Josephson critical current becomes maximum because of $\alpha\approx\pi$, and decreases with the increasing bias current $I>2I_0$. For $I=3I_0$ corresponding to $\alpha \approx 3\pi/2$, the second transition to the 0 state from the π state occurs.

Here we discuss the effect of the spin accumulation on the superconducting gap.[11] The spin split $\delta\mu_N$ at $x=L$ in N causes the split of ECP of Ss by $\delta\mu_N$ near the S/N interfaces. The spin split in Ss decreases exponentially with the spin-diffusion length λ_S from the interface. In the superconductors, the superconducting gap is not suppressed by spin accumulation until $\delta\mu_N$ exceeds the critical value of the spin split $\delta\mu_{Nc}$.[11] At low temperatures much lower than the superconducting critical temperature ($T \ll T_c$), the critical value of the spin split is obtained as $\delta\mu_{Nc} \lesssim \Delta_0$ by solving the gap equation,[11] where Δ_0 is the superconducting gap for $\delta\mu_N=0$ at $T=0$. In the case discussed above, $\delta\mu_N \approx 1$ meV at the first 0-π transition ($\alpha\approx\pi/2$). For example, $\Delta_0 \approx 1.5$ meV for niobium,[29] and therefore the superconducting gap is almost not affected by the spin accumulation at the first 0-π transition. When superconductors with the higher value of T_c, e.g., MgB_2 ($T_c \approx 39$ K)[30] or high-T_c materials (T_c is several 10 Ks),[29] are used as the electrodes S1 and S2, the superconductivity is robust even at the second ($\delta\mu_N \approx 3$ meV, $\alpha \approx 3\pi/2$) and higher 0-π transitions.

In summary, we have proposed the Josephson device in which the 0 and π states are controlled electrically. The spin split of the electrochemical potential is induced in the electrode N by the spin-polarized bias current flowing from F to N. The π state appears in the S1/N/S2 junction due to the nonlocal spin accumulation in N. Because the magnitude of the spin accumulation is proportional to the value of the spin-polarized bias current, the 0 and π states of the Josephson junction are controlled by the current. Our proposal provides not only possibilities for the application of superconducting spin-electronic devices but also a deeper understanding of the spin-dependent phenomena in the magnetic nanostructures.

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