Anisotropy of p-wave Josephson junction*

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The anisotropy of the dc Josephson current in the superconducting junctions of the p-wave equal spin pairing symmetry is theoretically investigated by the Furusaki–Tsukada-like formula. The current phase relations exhibit different oscillation periods and different phase shifts for the current along different directions, respectively.

Keywords: tunnel junction, Josephson current, p-wave superconductor

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1. Introduction

The recently discovered superconductivity in the weak ferromagnetic metals, such as UGe_2,[1] URhGe,[2] and $ZrZn_2$,[3] has aroused the considerable interest in the coexistence between the superconductivity and the ferromagnetism. In those ferromagnetic superconductors (FSs), the superconductivity and the ferromagnetism are both attributed to the conduction electrons. In the ferromagnets the electron spins are aligned in parallel, and in the conventional superconductors the electrons are condensed into Cooper pairs where the electron spins are aligned in antiparallel. Therefore, the spin-singlet pairing is somewhat inconsistent with the ferromagnetism. Although there were some suggestions of the *s*-wave pairing for the FSs,[4,5] the electrons in the FSs are more likely to form the spin-triplet Cooper pairs, which can have uniform magnetic moments.[6,7] For example, the p-wave equal spin pairing was proposed for $ZrZn_2$.[8,9] The low energy tunneling spectra and the Josephson effects in the FS tunnel junctions are controlled by the Andreev reflection process[10] and can provide the important information to reveal the pairing symmetry of the superconducting gap. Thus, the tunneling spectra and the Josephson effects in the equal spin pairing FS junctions have received much attention.[11−14]

In the Josephson junction, there is a supercurrent driven by the macroscopic phase difference φ between the two superconducting leads. The current phase relation $I(\varphi)$ is one of the key features in the dc Josephson effect and can often be well approximated by $I = I_0 \sin\varphi$ with I_0 being the critical current. In some cases the critical current is achieved at $\varphi_{\max} < \pi/2$ or $\varphi_{\max} > \pi/2$ resulting in a deviated sinusoidal current phase relation. In the Josephson junction with the ferromagnetic interface layer, the sign of the critical current can be changed by tuning the temperature or the thickness of the ferromagnet layer.[15−21] The Josephson junction with a negative critical current is the so-called π junction.[22,23] The more general φ_0 junction, where $I = I_0 \sin(\varphi + \varphi_0)$, and the $\cos\varphi$ current phase relation were also reported.[24,25] Besides the phase shift, the oscillation period of the current phase relation can also be changed in the junction consisting of two superconductors with different pairing symmetries. The Josephson current between s wave and $d_{x^2-y^2}$ wave superconductors is a π periodic function of the phase difference when the junction is formed along the *c* axis.[26] There are many different kinds of orbital pairing symmetries for the p-wave superconductors which lead to some novel current phase relations. For example, when the orbital pairing symmetries of the two superconductors are different, the current phase relation in the equal spin pairing p-wave Josephson junction can be the $\pi/2$ periodic function $\sin 4\varphi$.[27]

In this paper, the anisotropy of the dc Josephson effect in the junctions consisting of two equal spin pairing p-wave FSs with different orbital pairing symmetries is investigated theoretically. The Josephson

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current is calculated from the Andreev reflection coefficients by a Furusaki–Tsukada-like formula.[28−30] The Andreev reflection coefficient is obtained from the Bogoliubov-de Gennes equations.[31,32] It is shown that the internal phases of the two superconductors play an important role in the current phase relation. The internal phases are different for the Josephson current along different directions, resulting in different oscillation periods and phase shifts in the current phase relations.

2. Model and formulation

The system under consideration is an FS–Insulator–FS junction. First, we assume that the current is along the x direction and the thin insulator layer at $x = 0$ is modeled by a δ-type barrier potential with strength V_0. The FS layer is considered as an equal spin pairing p-wave superconductor, which is the case proposed for $ZrZn_2$.[8,9] There is only one of the exchange-splitting bands, namely the spin-up band, which is superconducting. When the length of the junction is less than the spin-flipping length and the magnetizations of the two FS layers are in a parallel configuration, the finite Josephson current can only exist in the spin-up channel and the junction is actually a normal-metal junction without supercurrent for the spin-down channel. There are many possible orbital pairing symmetries for the p-wave superconductors. In the reciprocal space, we assume that the wave vector dependences of the pair potentials take the forms

$$\begin{aligned}\Delta_{\rm L} &= \Delta_{^1E} \\ &= {\rm i}\Delta_0 \sin\frac{k_z a}{2}\left(\cos\frac{k_x a}{2} - \cos\frac{k_y a}{2}\right), \end{aligned} \quad (1)$$

and

$$\begin{aligned}\Delta_{\rm R} &= \Delta_A \\ &= {\rm i}\Delta_0 \cos\frac{k_z a}{2}\left(\sin\frac{k_y a}{2} + {\rm i}\sin\frac{k_x a}{2}\right) \end{aligned} \quad (2)$$

for the left and the right FS layers, respectively. Here, Δ_0 is the strength of the pair potential and a is the lattice constant. The pair potentials defined in Eqs. (1) and (2) are the two possible choices in the superconducting states proposed for $ZrZn_2$, and the z direction is along the [001] direction of the lattice.[8] The subscripts ^1E and A denote the two irreducible corepresentations of the magnetic point group $\boldsymbol{D}_4(\boldsymbol{C}_4)$, respectively.

We consider the Josephson junction without the impurity scattering as a simplest example. The wave functions of the quasiparticles in the spin-up channel in the left or right FS layer $\begin{pmatrix} u & v \end{pmatrix}$ satisfy the Bogoliubov-de Gennes equation[11,31]

$$\begin{pmatrix} H_0 & \Delta_{\rm L,R} \\ \Delta^*_{\rm L,R} & -H_0 \end{pmatrix}\begin{pmatrix} u \\ v \end{pmatrix} = E\begin{pmatrix} u \\ v \end{pmatrix}, \quad (3)$$

where H_0 is the single particle Hamiltonian, and the energy of the quasi-particle, E, is measured from Fermi energy $E_{\rm F}$. Since the Cooper pair consists of two electrons with the same spins, the exchange energy in FS layer has been removed by redefining the Fermi energy in Eq. (3). The solution of the Bogoliubov-de Gennes equation (3) is written as

$$\begin{aligned}\Psi_{\rm L} =& \begin{pmatrix} u_{\rm L} \\ v_{\rm L}\,{\rm e}^{-{\rm i}\gamma_{\rm L}^+}\,{\rm e}^{-{\rm i}\varphi_{\rm L}} \end{pmatrix}{\rm e}^{{\rm i}k_x x} \\ &+ a\begin{pmatrix} v_{\rm L} \\ u_{\rm L}\,{\rm e}^{-{\rm i}\gamma_{\rm L}^+}\,{\rm e}^{-{\rm i}\varphi_{\rm L}} \end{pmatrix}{\rm e}^{{\rm i}k_x x} \\ &+ b\begin{pmatrix} u_{\rm L} \\ v_{\rm L}\,{\rm e}^{-{\rm i}\gamma_{\rm L}^-}\,{\rm e}^{-{\rm i}\varphi_{\rm L}} \end{pmatrix}{\rm e}^{-{\rm i}k_x x} \end{aligned} \quad (4)$$

in the left FS layer and

$$\begin{aligned}\Psi_{\rm R} =& c\begin{pmatrix} u_{\rm R} \\ v_{\rm R}\,{\rm e}^{-{\rm i}\gamma_{\rm R}^+}\,{\rm e}^{-{\rm i}\varphi_{\rm R}} \end{pmatrix}{\rm e}^{{\rm i}k_x x} \\ &+ d\begin{pmatrix} v_{\rm R} \\ u_{\rm R}\,{\rm e}^{-{\rm i}\gamma_{\rm R}^-}\,{\rm e}^{-{\rm i}\varphi_{\rm R}} \end{pmatrix}{\rm e}^{-{\rm i}k_x x} \end{aligned} \quad (5)$$

in the right FS layer. In Eq. (4), the first term on the right-hand side is the incoming electron-like quasiparticle (ELQ), the second one is the Andreev reflected hole-like quasi-particle (HLQ) with amplitude a,[10] and the last one is the normal reflected ELQ with amplitude b. In Eq. (5), the first term on the right-hand side is the transmitted ELQ with amplitude c and the second one is transmitted HLQ with amplitude d. The pair potential and the quasi-particle energy are much smaller than the Fermi energy. Consequently, the wave vectors of the quasi-particles in Eqs. (4) and (5) are approximated as $k_x = k_{\rm F}\cos\theta$ where $k_{\rm F}$ is the Fermi wave vector and θ is the incident angle. The Bogoliubov amplitudes read

$$u_{\rm L,R} = \frac{1}{\sqrt{2}}\left(1 + \frac{\sqrt{E^2 - |\Delta_{\rm L,R}|^2}}{E}\right)^{1/2},$$

$$v_{\mathrm{L,R}} = \frac{1}{\sqrt{2}}\left(1 - \frac{\sqrt{E^2 - |\Delta_{\mathrm{L,R}}|^2}}{E}\right)^{1/2}. \quad (6)$$

In Eqs. (4) and (5), φ_{L} and φ_{R} are the macroscopic phases of the Cooper pairs, and $\gamma_{\mathrm{L}}^{\pm}$ and $\gamma_{\mathrm{R}}^{\pm}$ are the internal phases. The internal phases occur in the anisotropic superconductors where the quasi-particles moving along different directions have different wave vectors and thus feel different pair potentials.[28] The internal phases affect the tunneling spectra and the Josephson current significantly[27,28,33,34] and are obtained as

$$\mathrm{e}^{\mathrm{i}\gamma_{\mathrm{L,R}}^{\pm}} = \frac{\Delta_{\mathrm{L,R}}(\pm k_x)}{|\Delta_{\mathrm{L,R}}(\pm k_x)|}. \quad (7)$$

According to the orbital pairing symmetries defined in Eqs. (1) and (2), the internal phases along the positive and the negative x directions have the relation

$$\gamma_{\mathrm{L}}^{-} = \gamma_{\mathrm{L}}^{+}, \quad \text{and} \quad \gamma_{\mathrm{R}}^{-} = \pi - \gamma_{\mathrm{R}}^{+}. \quad (8)$$

The Andreev reflection amplitude $a(E)$ can be obtained by matching the boundary conditions at $x = 0$[32]

$$\Psi_{\mathrm{L}}|_{x=0} = \Psi_{\mathrm{R}}|_{x=0}, \text{ and}$$

$$\frac{\partial \Psi_{\mathrm{R}}}{\partial x}\Big|_{x=0} - \frac{\partial \Psi_{\mathrm{L}}}{\partial x}\Big|_{x=0} = 2Zk_{\mathrm{F}}\Psi_{\mathrm{L}}\Big|_{x=0}, \quad (9)$$

where the dimensionless barrier strength $Z = mV_0/(\hbar^2 k_{\mathrm{F}})$ and m is the effective mass of the electron. The Josephson current at temperature T is determined by the Andreev reflection process and can be obtained by the Furusaki–Tsukada-like formula[28–30]

$$I = \frac{e}{\hbar}\sum_{k_y,k_z}\sum_{\omega_n}\frac{k_{\mathrm{B}}T|\Delta_{\mathrm{L}}|}{\Omega_{n\mathrm{L}}}\left[a(\mathrm{i}\omega_n) - \widetilde{a}(\mathrm{i}\omega_n)\right], \quad (10)$$

where $\Omega_{n\mathrm{L}} = \mathrm{sgn}(\omega_n)\sqrt{\omega_n^2 + |\Delta_{\mathrm{L}}|^2}$, and $\omega_n = (2n+1)\pi k_{\mathrm{B}}T$. The Andreev reflection amplitude $a(\mathrm{i}\omega_n)$ can be obtained from $a(E)$ by the analytic continuation $E \to \mathrm{i}\omega_n$. The Andreev reflection amplitude $\widetilde{a}(\mathrm{i}\omega_n)$ describes the process where an incoming HLQ from the left superconductor is reflected as an ELQ and can be obtained by the symmetry consideration $\widetilde{a}(k_x, E) = a^*(-k_x, E)$. With the help of internal phase relation (8), the dc Josephson current along the x direction can be finally written as

$$I_x = \frac{ek_{\mathrm{B}}T}{\hbar}\sum_{k_y,k_z}\sum_{\omega_n}\frac{\Xi_x|\Delta_{\mathrm{L}}||\Delta_{\mathrm{R}}|\sin\varphi\sin\gamma_{\mathrm{L}}^x}{\Xi_x^2 + \left[\left(1 + 2\dfrac{Z^2}{\cos^2\theta}\right)\Omega_{n\mathrm{L}} + \Omega_{n\mathrm{R}}\right]^2\omega_n^2\cos^2\gamma_{\mathrm{R}}^x}, \quad (11)$$

where

$$\Xi_x = \left[\omega_n^2 + \left(1 + 2\frac{Z^2}{\cos^2\theta}\right)\Omega_{n\mathrm{L}}\Omega_{n\mathrm{R}}\right]\sin\gamma_{\mathrm{R}}^x + |\Delta_{\mathrm{L}}||\Delta_{\mathrm{R}}|\cos\varphi\sin\gamma_{\mathrm{L}}^x, \quad (12)$$

the macroscopic phase difference $\varphi = \varphi_{\mathrm{L}} - \varphi_{\mathrm{R}}$, and $\gamma_{\mathrm{L,R}}^x$ denotes the internal phase along the positive x direction.

For the junctions prepared along the y and the z directions, one can derive the current formula in a similar way. The main difference comes from the internal phases. For the current along the y direction, the internal phases for the particles moving along the positive and the negative y directions have the relation

$$\gamma_{\mathrm{L}}(-k_y) = \gamma_{\mathrm{L}}(+k_y), \text{ and } \gamma_{\mathrm{R}}(-k_y) = -\gamma_{\mathrm{R}}(+k_y). \quad (13)$$

The dc Josephson current along the y direction is given by

$$I_y = \frac{ek_{\mathrm{B}}T}{\hbar}\sum_{k_x,k_z}\sum_{\omega_n}\frac{\Xi_y|\Delta_{\mathrm{L}}||\Delta_{\mathrm{R}}|\cos\varphi\sin\gamma_{\mathrm{L}}^y}{\Xi_y^2 + \left[\left(1 + 2\dfrac{Z^2}{\cos^2\theta}\right)\Omega_{n\mathrm{L}} + \Omega_{n\mathrm{R}}\right]^2\omega_n^2\sin^2\gamma_{\mathrm{R}}^y}, \quad (14)$$

where

$$\Xi_y = \left[\omega_n^2 + \left(1 + 2\frac{Z^2}{\cos^2\theta}\right)\Omega_{n\mathrm{L}}\Omega_{n\mathrm{R}}\right]\cos\gamma_{\mathrm{R}}^y - |\Delta_{\mathrm{L}}||\Delta_{\mathrm{R}}|\sin\varphi\sin\gamma_{\mathrm{L}}^y, \quad (15)$$

and $\gamma_{\mathrm{L,R}}^y$ denotes the internal phase along the positive y direction. For the current along the z direction, the internal phases for the particles to move along the positive and the negative z directions have the following relations:

$$\gamma_{\mathrm{L}}(-k_z) = \pi + \gamma_{\mathrm{L}}(+k_z), \quad \text{and} \quad \gamma_{\mathrm{R}}(-k_z) = \gamma_{\mathrm{R}}(+k_z). \quad (16)$$

The dc Josephson current along the z direction is given by

$$I_z = -\frac{ek_{\rm B}T}{\hbar}\sum_{k_x,k_y}\sum_{\omega_n}\frac{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2\sin\eta\cos\eta}{\omega_n^2\left[\Omega_{n{\rm L}}+\left(1+2\dfrac{Z^2}{\cos^2\theta}\right)\Omega_{n{\rm R}}\right]^2+|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2\sin^2\eta}, \tag{17}$$

where the total phase difference $\eta=\varphi+\gamma_{\rm L}^z-\gamma_{\rm R}^z$ including both the macroscopic phase difference and the internal phase difference with $\gamma_{\rm L,R}^z$ representing the internal phase along the positive z direction.

3. Results and discussion

With the help of Eqs. (11), (14) and (17), the current can be calculated numerically so that one can study the anisotropy of the dc Josephson effect. We take the Fermi surface parameter $k_{\rm F}a\approx 1.6$,[35] the dimensionless barrier strength $Z=0.5$, and the temperature $k_{\rm B}T=0.01\Delta_0$. The current is normalized by the resistance of the junction in a normal state, which is given by[26]

$$R_N^{-1}=\frac{e^2}{h}\int_0^{\pi/2}\frac{\cos^2\theta}{Z^2+\cos^2\theta}\sin\theta{\rm d}\theta. \tag{18}$$

The current phase relation is shown in Fig. 1. The currents along different directions exhibit different oscillation periods and different phase shifts. For the current along the z direction, it is a $\sin 4\varphi$-like function of the macroscopic phase difference with a period of $\pi/2$. For the current along the x direction, it is a $\sin 2\varphi$-like function of the macroscopic phase difference with a period of π. For the current along the y direction, it is also a π periodic function but has a π phase shift relative to the current along the x direction. These results can be understood by investigating the symmetries of the internal phases. The $\sin 4\varphi$ relation for the current along the z direction has already been reported[27] and can be explained by the four-fold rotation symmetry of the internal phase difference in the k_x–k_y plane. Josephson current (17) can be expanded into power series[26,27]

$$I_z=-\frac{ek_{\rm B}T}{\hbar}\sum_{k_x,k_y}\sum_{\omega_n}\sum_{m=0}^{\infty}\zeta^{m+1}\cos^m 2\eta\sin 2\eta, \tag{19}$$

where

$$\zeta=\frac{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2}{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2+2\omega_n^2\left[\Omega_{\rm R}+\left(1+2\dfrac{Z^2}{\cos^2\theta}\right)\Omega_{\rm L}\right]^2}$$

$$<1. \tag{20}$$

Under a $\pi/2$ rotation in the k_x–k_y plane, $\Delta_{\rm L}$ becomes $-\Delta_{\rm L}$ and $\Delta_{\rm R}$ becomes $-{\rm i}\Delta_{\rm R}$, which results in a π phase shift and a $-\pi/2$ phase shift for internal phases $\gamma_{\rm L}^z$ and $\gamma_{\rm R}^z$, respectively. Therefore, total phase difference η gains a $-\pi/2$ phase shift and both $\sin 2\eta$ and $\cos 2\eta$ are antisymmetric under the rotation operation. After the summation over k_x and k_y, only the terms with odd m in the power series (19) can survive. The first nonzero term comes from $m=1$ and is proportional to $\sin 4\varphi$. The higher order terms can only give the higher order harmonics. Therefore, the current phase relation is a $\pi/2$ periodic function.

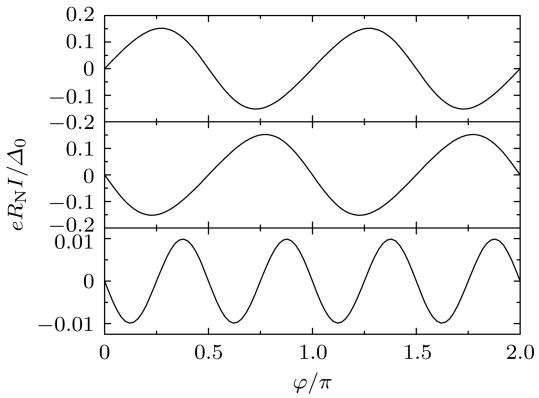


Fig. 1. Josephson current I versus macroscopic phase difference φ, where current I is along the x direction for the upper panel, the y direction for the middle panel, and the z direction for the lower panel, respectively.

For the currents along the x and y directions, there is no simple rotation symmetry of the pair potentials in the transverse plane. To understand the current phase relation, we focus on the case of the small angle incidence from which the main contribution of the current arises. In this case, the internal phases along the positive x direction satisfy the relations $\sin\gamma_{\rm L}^x=\pm 1$, $\cos\gamma_{\rm R}^x\approx\pm 1$, and $\sin\gamma_{\rm R}^x\approx 0$. The dominant contribution of the current along the x direction reduces to

$$I_x \sim \frac{ek_{\rm B}T}{\hbar}\sum_{\omega_n}\frac{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2\sin 2\varphi}{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2(1+\cos 2\varphi)+2[(1+2Z^2)\Omega_{n{\rm L}}+\Omega_{n{\rm R}}]^2\omega_n^2}. \tag{21}$$

One can find that the lowest harmonics in Eq. (21) is $\sin 2\varphi$ and the current is a π periodic function of macroscopic phase difference. In the region of the small angle incidence, the internal phases along the positive y direction satisfy the relations $\sin\gamma_{\rm L}^y = \pm 1$, $\cos\gamma_{\rm R}^y \approx 0$, and $\sin\gamma_{\rm R}^y \approx \pm 1$. The dominant contribution of the current along the y direction reduces to

$$I_y \sim -\frac{ek_{\rm B}T}{\hbar}\sum_{\omega_n}\frac{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2\sin 2\varphi}{|\Delta_{\rm L}|^2|\Delta_{\rm R}|^2(1-\cos 2\varphi)+2[(1+2Z^2)\Omega_{n{\rm L}}+\Omega_{n{\rm R}}]^2\omega_n^2}. \tag{22}$$

The lowest harmonics in Eq. (22) is also $\sin 2\varphi$ and the current is a π periodic function of the macroscopic phase difference but the current along the y direction has a π phase shift relative to the current along the x direction.

4. Summary

In this paper, the anisotropy of the dc Josephson effect in the junctions consisting of two equal spin pairing p-wave FSs with different orbital pairing symmetries is investigated theoretically. The dc Josephson current is obtained by the Furusaki–Tsukada-like formula. The pair potential of the p-wave superconductor is anisotropic. It is shown that the internal phase determined by the pairing symmetry in the $\boldsymbol{k}$ space has a great influence on the current phase relation of the Josephson junction. We take the two pair potentials proposed for $ZrZn_2$ as example. The current phase relation along the z direction is a $\sin 4\varphi$-like function with a period of $\pi/2$, the one along the x direction is a $\sin 2\varphi$-like function with a period of π, and the one along the y direction is also a π periodic function but has a π phase shift relative to that along the x direction. The results can be explained by the symmetry consideration of the internal phases resulting from the orbital pairing symmetries of the pair potentials.

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