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Anomalous Josephson Effect in p -Wave Dirty Junctions

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The Josephson effect in p -wave superconductor/diffusive normal metal/ p -wave superconductor junctions is studied theoretically. Amplitudes of Josephson currents are several orders of magnitude larger than those in s -wave junctions. Current-phase (J - φ) relations in low temperatures are close to those in ballistic junctions such as $J \propto \sin(\varphi/2)$ and $J \propto \varphi$ even in the presence of random impurity potentials. A cooperative effect between the midgap Andreev resonant states and the proximity effect causes such anomalous properties and is a character of the spin-triplet superconductor junctions.

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The *internal* π -phase shift (sign change) of pair potentials is essential for unconventional superconductivity and is the source of the midgap Andreev resonant state (MARS) [1–3]. It is now known that the MARS is responsible for anomalous transport properties in superconducting junctions [4]. In normal metal/superconductor junctions, transport properties are affected also by the proximity effect which is interpreted in terms of diffusion of Cooper pairs into normal metals. In what follows, we assume that normal metals are in the diffusive transport regime due to impurity scatterings. Recent theoretical studies have revealed sensitivity of the proximity effect to the internal phase of pair potentials [5,6]. In normal metals attached to unconventional superconductors, Cooper pairs have a sign degree of freedom reflecting the π -phase shift of pair potentials. Suppression of the proximity effect is usually expected because the wave function of a Cooper pair originated from the positive part of pair potentials cancel that originated from the negative part [5,6]. Two of us, however, discussed anomalous enhancement of the zero-bias tunneling conductance due to the proximity effect in a presence of the MARS [7,8].

In superconductor/normal metal/superconductor junctions, another phase degree of freedom affects quantum transport. Namely, the *external* phase difference across the junctions φ drives Josephson currents. An importance of studying the Josephson effect is growing these days because quantum interference devices consisting of Josephson junctions can be basis of future technologies. In fact, a recent experiment has tried to apply high- T_c superconductors to coherent devices [9]. In unconventional junctions, the MARS is considered to have the phase degree of freedom. When MARSs are formed at the two junction interfaces, the external phase may modify interference effects between the two MARSs and Josephson currents. The research in this direction can shed new light on quantum transport in unconventional superconductors.

In this Letter, we theoretically study Josephson currents between two p -wave superconductors through normal metals by solving the Bogoliubov–de Gennes equation using the recursive Green function method [10,11]. We show that amplitudes of Josephson currents in the p -wave junctions are much larger than those in the s -wave junctions when transmission probabilities of junction interfaces are small. The local density of states in normal metals has a zero-energy peak reflecting anomalous diffusion of the MARSs into a normal metal and that spatial profiles of the zero-energy peak depend strongly on φ . As a consequence, current-phase (J - φ) relations remarkably deviate from the sinusoidal function in low

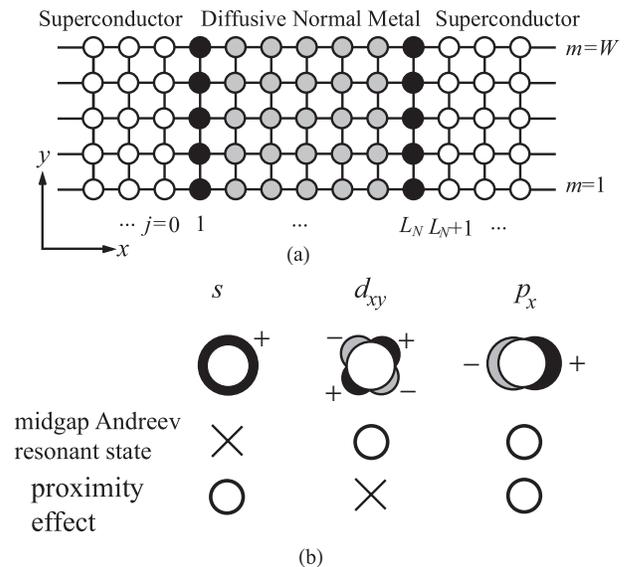


FIG. 1. A schematic figure of a Josephson junction on the tight-binding lattice is shown in (a). In (b), we illustrate the pair potentials in momentum space, where open circles represent the Fermi surface. The pair potentials are classified into three groups by the presence or absence of the two interference effects.

temperatures and are close to those in ballistic junctions such as $J \propto \sin(\varphi/2)$ and $J \propto \varphi$ [12–15]. The resonant tunneling through the MARS in normal metals is responsible for such unusual Josephson effect. The obtained results imply high potentials of spin-triplet superconducting junctions as coherent devices.

We consider three pairing symmetries on two-dimensional superconductors: (i) $\Delta_{\mathbf{k}} = \Delta_0$ for s -wave, (ii) $\Delta_0 2\bar{k}_x \bar{k}_y$ for d_{xy} -wave, and (iii) $\Delta_0 \bar{k}_x$ for p_x -wave symmetries. Here Δ_0 is the maximum amplitude of pair potentials at zero temperature, and $\bar{k}_x = k_x/k_F$ and $\bar{k}_y = k_y/k_F$ are normalized wave numbers on the Fermi surface in the x and y directions, respectively. Josephson currents are parallel to the x direction and junction interfaces are parallel to the y direction as shown in Fig. 1(a). The pair potentials in momentum space are illustrated in Fig. 1(b). An interference of a quasiparticle enables formation of the MARS at a junction interface when $\Delta_{k_x, k_y} \Delta_{-k_x, k_y} < 0$ [4,16]. The pair potentials in the d_{xy} and p_x symmetries satisfy the relation for all wave numbers. The absence of the proximity effect in normal metals is described by a relation [5,6] $\Delta_{k_x, k_y} = -\Delta_{k_x, -k_y}$. The pair potential in the d_{xy} symmetry satisfies the relation. Thus the proximity effect is expected in both the s - and p_x -wave symmetries. In Fig. 1(b), we classify the pairing symmetries into three groups by the presence (○) or absence (×) of the two interference effects [7,8]. Within p -wave symmetries, we pay special attention to the p_x -wave symmetry because the proximity effect and MARS are present at the same time. On the other hand, in the p_y -wave symmetry, neither is present [7].

Let us consider Josephson junctions on the two-dimensional tight-binding model as shown in Fig. 1(a). A vector $\mathbf{r} = j\mathbf{x} + m\mathbf{y}$ points to a lattice site, where \mathbf{x} and \mathbf{y} are the unit vectors in the x and y directions, respectively. The junction consists of three regions: a normal metal (i.e., $1 \leq j \leq L_N$) and two superconductors (i.e., $-\infty \leq j \leq 0$ and $L_N + 1 \leq j \leq \infty$). In the y direction, the number of lattice sites is W and we assume the periodic boundary condition. Electronic states in superconducting junctions are described by the mean-field Hamiltonian

$$H_{\text{BCS}} = \frac{1}{2} \sum_{\mathbf{r}, \mathbf{r}'} \left[\tilde{c}_{\mathbf{r}}^\dagger h_{\mathbf{r}, \mathbf{r}'} \tilde{c}_{\mathbf{r}'} - \bar{\tilde{c}}_{\mathbf{r}} h_{\mathbf{r}, \mathbf{r}'}^* \bar{\tilde{c}}_{\mathbf{r}'}^\dagger \right] + \frac{1}{2} \sum_{\mathbf{r}, \mathbf{r}' \in \text{S}} \left[\tilde{c}_{\mathbf{r}}^\dagger \hat{\Delta}_{\mathbf{r}, \mathbf{r}'} \bar{\tilde{c}}_{\mathbf{r}'}^\dagger - \bar{\tilde{c}}_{\mathbf{r}} \hat{\Delta}_{\mathbf{r}, \mathbf{r}'}^* \tilde{c}_{\mathbf{r}'} \right], \quad (1)$$

$$h_{\mathbf{r}, \mathbf{r}'} = -t \delta_{|\mathbf{r}-\mathbf{r}'|, 1} + (\epsilon_{\mathbf{r}} - \mu + 4t) \delta_{\mathbf{r}, \mathbf{r}'}, \quad (2)$$

$$\hat{\Delta}_{\mathbf{r}, \mathbf{r}'} = \begin{cases} i \mathbf{d}_{\mathbf{r}, \mathbf{r}'} \cdot \hat{\boldsymbol{\sigma}}_2 & : \text{triplet,} \\ i d_{\mathbf{r}, \mathbf{r}'} \hat{\sigma}_2 & : \text{singlet,} \end{cases} \quad (3)$$

$$\tilde{c}_{\mathbf{r}} = \begin{pmatrix} c_{\mathbf{r}, \uparrow} \\ c_{\mathbf{r}, \downarrow} \end{pmatrix}, \quad (4)$$

where $c_{\mathbf{r}, \sigma}^\dagger$ ($c_{\mathbf{r}, \sigma}$) is the creation (annihilation) operator of an electron at \mathbf{r} with spin $\sigma = (\uparrow \text{ or } \downarrow)$ and $\bar{\tilde{c}}$ means the transpose of \tilde{c} . The hopping integral t is considered among nearest neighbor sites. We assume that t and the Fermi energy μ are common in superconductors and a normal metal. In a normal metal, on-site potentials are given randomly in the range of $-V_I/2 \leq \epsilon_{\mathbf{r}} \leq V_I/2$. We introduce insulating barriers at $j = 1$ and L_N , where $\epsilon_{\mathbf{r}}$ is given by V_B . Two superconductors in which $\epsilon_{\mathbf{r}}$ are taken to be zero are identical to each other. In the p_x -wave symmetry, a spin vector of Cooper pairs \mathbf{d} points to the z direction. The arguments below do not depend on the directions of \mathbf{d} . The Hamiltonian is diagonalized by the Bogoliubov transformation and the Bogoliubov–de Gennes equation is numerically solved by the recursive Green function method [10,11]. Josephson currents are given by

$$J = -ieT \sum_{\omega_n} \text{Tr} [\hat{G}_{\omega_n}(\mathbf{r}', \mathbf{r}) - \hat{G}_{\omega_n}(\mathbf{r}, \mathbf{r}')] \quad (5)$$

with $\mathbf{r}' = \mathbf{r} + \mathbf{x}$, where \hat{G}_{ω_n} is the Green function and $\omega_n = (2n + 1)\pi T$ is the Matsubara frequency with n and T being an integer and a temperature, respectively. In Eq. (5), Tr means the trace in the Nambu space and the summation over m . In this Letter, the unit of $\hbar = k_B = 1$ is used, with k_B being the Boltzmann constant. The local density of states is also calculated from $N(E, j) = -\text{Im} \text{Tr} \hat{G}_{E+i\gamma}(\mathbf{r}, \mathbf{r})/\pi$, where E is measured from the Fermi energy and γ is a small imaginary part. Throughout this Letter, we fix parameters as $L_N = 70$, $W = 25$, $\mu = 2t$, and $V_I = 2t$. Under these parameters, normal metals are in the diffusive transport regime, where the mean free path in normal metals is estimated to be about $\ell \sim 6$ lattice constants and the Thouless energy E_{th} is calculated to be $1.6 \times 10^{-3}t$. Results discussed below are qualitatively insensitive to these parameters.

At first we show that the maximum amplitudes of Josephson currents in the p_x -wave symmetry $J_c(p_x)$ become much larger than those in the s -wave $J_c(s)$. In Fig. 2(a), ratios $J_c(p_x)/J_c(s)$ are plotted as a function of

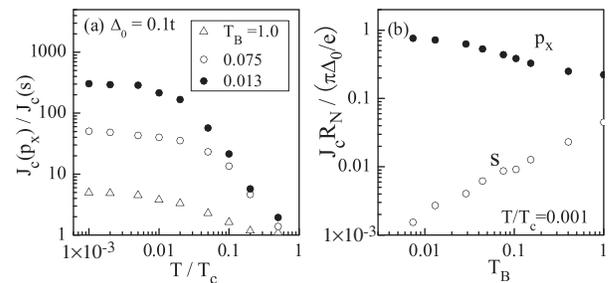


FIG. 2. The maximum amplitudes of Josephson currents in the p_x -wave symmetry $J_c(p_x)$ are compared with those in the s -wave symmetry $J_c(s)$ in (a), where T_B is the transmission probability of potential barriers in the normal states. In (b), $J_c(p_x)$ and $J_c(s)$ are plotted as a function of T_B at $T = 0.001T_c$.

temperatures for $\Delta_0 = 0.1t$. Here we choose several values of the barrier potentials V_B at $j = 1$ and L_N . The resulting normal transmission probabilities of the barrier T_B are 1.0, 0.075, and 0.013 for $V_B/t = 0, 6$, and 15, respectively. The ratios $J_c(p_x)/J_c(s)$ increase with decreasing T and amazingly become more than 100 in low temperatures for small T_B . The amplitudes of Josephson currents in the p_x -wave junctions are much larger than those in the s -wave junctions. In Fig. 2(b), $J_c R_N$ normalized by $\pi\Delta_0/e$ is plotted as a function of T_B at $T = 0.001T_c$, where R_N is the normal resistance of junctions. The results show that $J_c R_N$ in the s wave decreases with decreasing T_B , whereas that in the p_x wave increases.

We next focus on the current-phase relations of the Josephson effect. In Fig. 3, Josephson currents are plotted as a function of φ for the p_x -wave symmetries at $V_B = 0$. Parameters are chosen as $\Delta_0 = 0.01t$ and $0.0001t$ in (a) and (b), respectively. The current-phase relations are almost sinusoidal function in a high temperature at $T = 0.5T_c$. At $T = 0.001T_c$, however, the current-phase relations are close to $J \propto \varphi$ and $J \propto \sin(\varphi/2)$ in (a) and (b), respectively. These are characteristic current-phase relations in ballistic Josephson junctions in the s -wave symmetry [12–14]. We have confirmed that these current-phase relations remain even in the presence of potential barriers (i.e., $V_B \neq 0$).

The results imply large contributions of the multiple Andreev reflection in low temperatures. In general, Josephson currents can be decomposed into a series of $J = \sum_{n=1}^{\infty} J_n \sin(n\varphi)$, where J_n for $n \geq 2$ represent contributions of the multiple Andreev reflection. Roughly speaking, J_n is proportional to $\{T_N\}^n$ with T_N being the transmission probability of a quasiparticle from the left superconductor to the right superconductor through the normal segment (including two barriers and a normal metal). Thus the multiple Andreev reflection is negligible (i.e., $J_1 \gg J_2 \gg J_3 \gg \dots$) for $T_N \ll 1$. On the other hand, in the case of $T_N = 1$, the multiple Andreev reflection leads to the deviation of current-phase relations from

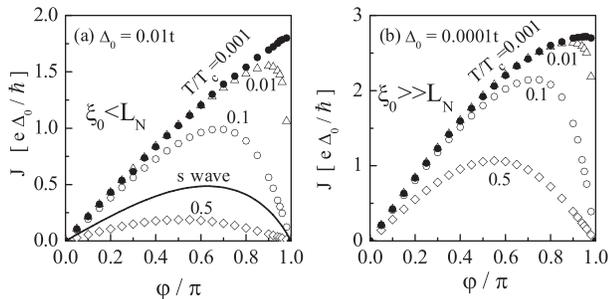


FIG. 3. Current-phase relations for the p_x -wave symmetries are shown for several temperatures at $V_B = 0$, where $\Delta_0 = 0.01t$ in (a) and $\Delta_0 = 0.0001t$ in (b). For comparison, results in the s -wave junctions at $T = 0.001T_c$ are shown with a solid line in (a), where the amplitude of Josephson current is multiplied by 5.

the sinusoidal function. It is noted at $T_N = 1$ that we obtain $J \propto \varphi$ and $J \propto \sin(\varphi/2)$ at the zero temperature for $L_N \gg \xi_0$ and $L_N \ll \xi_0$, respectively [12–14].

In Fig. 3(a), we also show the current-phase relations in the s -wave symmetry at $T = 0.001T_c$ with a solid line. The current-phase relation in the s wave is described almost by the sinusoidal function [17] because impurity potentials in normal metal suppress T_N and therefore the multiple Andreev reflection. In the p_x -wave junctions, the coherence lengths ξ_0 are estimated to be about 50 lattice constants in (a) and 5000 in (b). Thus $L_N > \xi_0$ and $L_N \ll \xi_0$ are satisfied in (a) and (b), respectively. The current-phase relations such as $J \propto \sin(\varphi/2)$ in (b) and $J \propto \varphi$ in (a) are universal properties of the p_x -wave junctions in low temperatures because they are independent of the strength of barrier potentials and the degree of disorder in normal metals. The calculated results in Fig. 3 indicate $T_N = 1$ even in the presence of impurity potentials. The large amplitudes of the Josephson current in Fig. 2 are also explained by $T_N = 1$.

The calculated results in Figs. 2 and 3 show the specific properties of Josephson currents in the p_x -wave junctions. In what follows, we analyze quasiparticle states in normal metals to understand the origin of the anomalous Josephson effect. In Fig. 4, we show the local density of states in normal metals for the s - and p_x -wave symmetries, where $\Delta_0 = 0.005t$, $\gamma = 0.05\Delta_0$, and N_0 denotes the normal density of states. At $\varphi = 0$ in the s -wave junctions in (a), the local density of states for $E < E_{th} \sim 0.3\Delta_0$ is suppressed because of the proximity effect. The suppression of the local density of states indicates the conversion of quasiparticles to Cooper pairs in normal metals. At $\varphi = \pi$ in (b), the local density of states recovers its amplitude for $E < E_{th}$. The wave function of Cooper pairs from the left superconductor and that from the right one cancel each other around $\varphi \sim \pi$, as schematically illustrated in a picture below the calculated results.

The local density of states is drastically changed in the p_x -wave symmetry as shown in (c) and (d). Zero-energy peaks whose width is determined by γ can be seen, which means formation of the MARS in normal metals. Although the MARS originally localizes at junction interfaces [4], the MARS penetrates into normal metals in the presence of the proximity effect. Spatial profiles of the local density of states depend remarkably on the external phase difference as shown in (c) and (d). At $\varphi = 0$, the zero-energy peak disappears at the center of normal metals ($j \sim 35$) because the wave function of the MARS from the left superconductors cancels out that from the right one, as shown schematically in a lower panel in (c). On the other hand, in (d), wave functions of the MARS in the two superconductors have the same sign with each other. Thus the two MARSs can penetrate deeply into normal metals and the zero-energy peak can be seen everywhere. We note that the penetration of the MARS is possible only when the proximity effect is present in normal metals. In fact, we

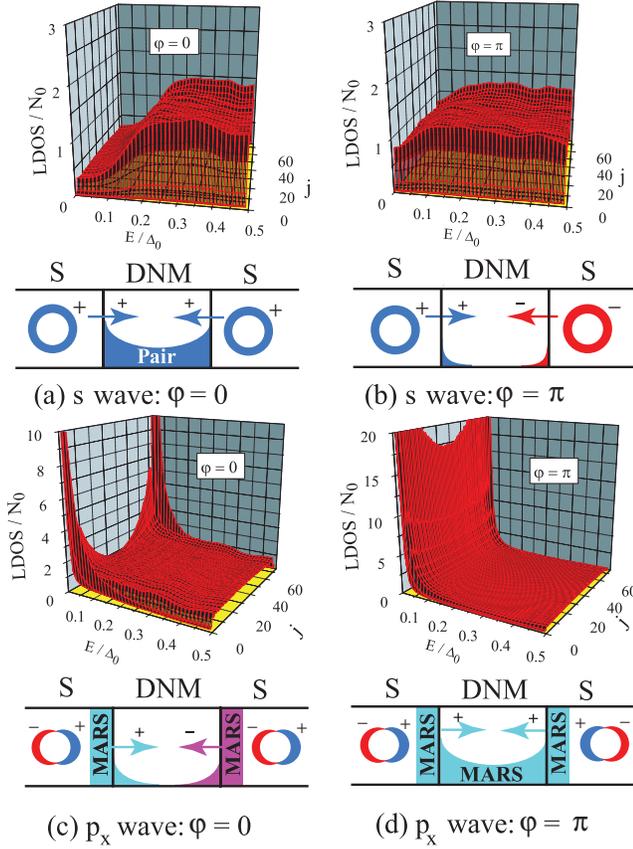


FIG. 4 (color). Local density of states (LDOS) in normal metals ($1 \leq j \leq L_N = 70$) is shown for the s -wave and p_x -wave symmetries. The left and right superconductors are attached at $j = 0$ and $j = 71$, respectively. Note that E_{th} is about $0.3\Delta_0$. In the schematic pictures, DNM and S denote a diffusive normal metal and a superconductor, respectively. The local density of states shown here is calculated in the absence of Josephson currents. We have confirmed that the results at $\varphi = 0.99\pi$ qualitatively show the same behavior as those at $\varphi = \pi$.

have confirmed that no zero-energy peak is found in normal metals in the d_{xy} -wave symmetry (results are not shown) and that the ensemble average of Josephson currents vanishes because the proximity effect is absent in normal metals [5]. Figure 4 indicates that the proximity effect bridges the two MARSs in the two superconductors. Thus $T_N = 1$ holds because of the resonant transmission through the MARS in normal metals. The Josephson effect specific to the p_x -wave symmetry discussed in Figs. 2 and 3 are a consequence of the diffusion of the MARS into normal metals.

In summary, we found anomalous behaviors of Josephson currents in superconductor/normal metal/superconductor junctions in the p_x -wave symmetry. The maxi-

imum amplitudes of Josephson currents J_c in the p_x -wave junctions become much larger than those in the s -wave junctions. It is known that large values of J_c are desired in device applications because $J_c R_N$ limits operation speeds of Josephson devices. Current-phase relations in low temperatures are close to those in ballistic junctions such as $J \propto \sin(\varphi/2)$ and $J \propto \varphi$ independent of the strength of potential barriers at interfaces and the degree of disorder in normal metals. The two the midgap Andreev resonant states penetrate deeply into normal metals, which causes the unusual Josephson effect in p_x -wave superconducting junctions. The anomalous Josephson effect is a novel feature of phase-sensitive transport in spin-triplet superconducting junctions.

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