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Coherence effects in S/I/N/I/FS tunnel junctions*

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The dc Josephson effect in superconductor / insulator / normal metal / insulator /ferromagnetic superconductor junctions has been studied. We calculate the dc Josephson current based on the Bogoliubov de Gennes equation. The Josephson current is derived as a function of exchange field in ferromagnetic superconductor, normal metal thickness and insulating barrier strength. It is found that there exists an oscillation relation between the critical Josephson current and the normal metal thickness. The oscillation amplitude decreases as the thickness of the normal metal increases or the exchange field augments.

Keywords: ferromagnetic superconductor, Josephson current, current oscillation
PACC: 7450, 7470H

Josephson effect is a well-known manifestation of macroscopic phase coherence. The mechanism of the effect is the phase coherent transport of quasiparticles from one superconductor to another in the case in which the two superconductors are weakly coupled with each other.[1,2] Early studies of weak links mainly focused on the superconductor / insulator /superconductor junctions.[3] A new era in the research of Josephson effect started from the discovery of a peculiar scattering mechanism at the superconductor /normal metal (S/N) boundary, known as the Andreev reflection.[4] In the superconductor/normal metal /superconductor (S/N/S) junctions, the Andreev reflection indicates that the dc Josephson effect arises due to the phase coherent transport in the normal region. The mechanism may be the transport of Cooper pairs from one superconducting electrode to another.[5−8]

Recently, the coexistence of superconductivity and ferromagnetism has been a very interesting subject.[9,10] It was predicted by Fulde and Ferrel,[11] and Larkin and Ovchinnikov[12] (FFLO) that pairing can exist when electron momenta at Fermi energy are different in two spin directions. Experimentally, $CeCoIn_5$ was proposed as a candidate for the observation of the FFLO state.[13] Unlike the conventional Cooper pair, in which two electrons have opposite spins and momenta $(K\uparrow, -K\downarrow)$, the FFLO pair in the presence of an exchange field has a finite centre-of-mass momentum $Q = 2h_0/\hbar v_{\rm F}$, where $2h_0$ is the exchange energy corresponding to the difference in energy between the spin-up and spin-down bands, and $v_{\rm F}$ is the Fermi velocity. The FFLO state with $[(K + Q/2)\uparrow, (-K + Q/2)\downarrow]$ was never observed in bulk materials. It stems from the fact that in a bulk ferromagnet (F), h_0 is at least two orders of magnitude larger than the energy gap Δ_0 of a bulk superconductor (S), while the FFLO state can appear only in the region where $h_0 < 0.754\Delta_0$ for the s-wave S[11,12] or $h_0 < 1.06\Delta_0$ for d-wave S.[14] However, in a thin F/S bi-layer, the effective exchange field h_0 and effective superconducting order parameter Δ may be of the same order of magnitude and thus the coexistence between superconductivity and ferromagnetism can be realized. On the assumption that the thickness of superconducting layer is smaller than the superconducting coherent length, and the thickness of the ferromagnetic layer smaller than the length of the condensate penetration into the F, a thin F/S bi-layer can be treated as a ferromagnetic superconductor (FS) film.[15−17]

In this paper, we theoretically study the universal features of the coherent quantum transport in superconductor / insulator / normal metal / insulator /ferromagnetic superconductor (S/I/N/I/FS) junctions. Solving the Bogoliubov de Gennes (BdG) equation,[18] the superconducting order parameter $\Delta(T, h_0)$ and the energy spectrum are obtained for the FS film. We extend the approach of Blonder–Tinkham–Klapwijk (BTK)[19] to the study of the Josephson current in S/I/N/I/FS. We consider simultaneously four types

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of quasiparticle injection process: electron and hole injection from the S to the normal metal (N), and corresponding hole and electron injection from the FS to the N, as shown in Fig.1. The quasiparticle interference in the N and the resonant tunnelling play an important role and exhibit a new quantum effect on the dc Josephson current in the S/I/N/I/FS junctions. It is shown that there exists the oscillation relation between the critical Josephson current and the normal metal thickness. The oscillation amplitude decreases as the thickness of the N increases or the exchange field augments.

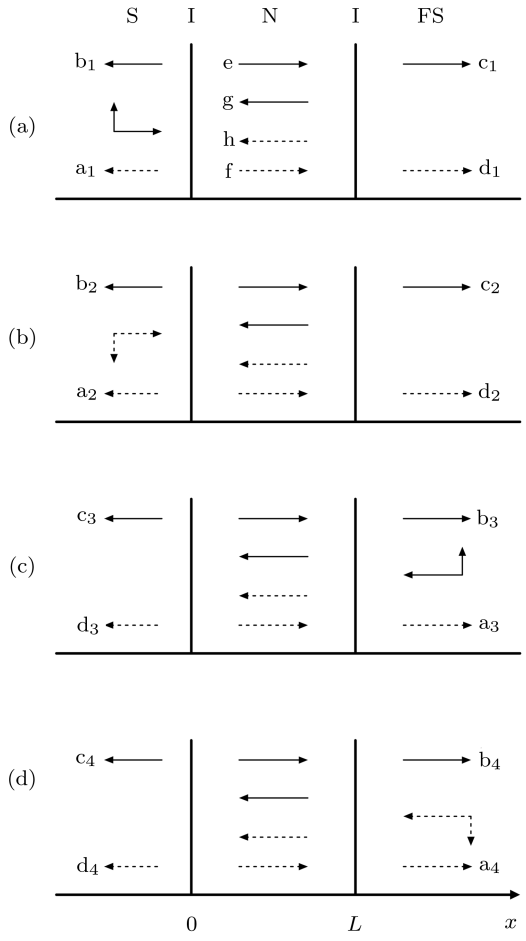


Fig.1. Schematic illustration of reflections and transmissions of quasiparticles in S/I/N/I/FS structures, in which subfigures (a) and (b), respectively, show spin-up (↑) electron and spin-down (↓) hole incident on the interface $x = 0$ from the S; subfigures (c) and (d), respectively, indicate spin-up (↑) electron and spin-down(↓) hole incident on the interface $x = L$ from the FS. Here horizontal arrows refer to quasiparticle transporting directions, solid lines represent electron in N or electronlike in S and FS, dash lines represent hole in N or holelike in S and FS.

For the model of calculation, we consider a one-dimensional S/I/N/I/FS junction with the magnetizations in the FS layer parallel to the interface. We must pay attention to the face that the quasiparticles have different potentials in the FS film. In the absence of spin-flip scattering, the four-component BdG equations are decoupled and become two sets of two-component BdG equations for the spin-up electronlike and spin-down holelike quasiparticle wave function $(u_\uparrow, v_\downarrow)$, and $(u_\downarrow, v_\uparrow)$.[20] The BdG equation is given by

$$\begin{pmatrix} H_0 - \eta_\sigma h' & \Delta(T, h_0) \\ \Delta^*(T, h_0) & -(H_0 - \eta_{\bar\sigma} h') \end{pmatrix} \begin{pmatrix} u_\sigma \\ v_{\bar\sigma} \end{pmatrix} = E \begin{pmatrix} u_\sigma \\ v_{\bar\sigma} \end{pmatrix}, \quad (1)$$

where H_0 is the single-particle Hamiltonian; the quasiparticle energy E is measured relatively to the Fermi energy E_F; $h' = h_0\Theta(x - L)$ is the effective exchange field in the FS with $\Theta(x-L)$ being the Heaviside step function; $\eta_\sigma = 1$ for $\sigma =\uparrow$; $\eta_\sigma = -1$ for $\sigma =\downarrow$; $\bar\sigma$ denotes the spin opposite to σ; $\Delta(T, h_0)$ is the effective superconducting order parameter independent of coordinates, but dependent on effective exchange field and temperature. From the BdG equations, we obtain

$$u_\sigma^2 = (1 + \sqrt{1 - \Delta^2(T, h_0)/(E + \eta_\sigma h_0)^2})/2, \quad (2)$$
$$v_{\bar\sigma}^2 = (1 - \sqrt{1 - \Delta^2(T, h_0)/(E - \eta_{\bar\sigma} h_0)^2})/2. \quad (3)$$

The wave vectors of the electronlike and holelike quasiparticles in the FS are given respectively by

$$k_\sigma^\mathrm{e} = \sqrt{(2m/\hbar^2)[E_\mathrm{F} + \sqrt{(E + \eta_\sigma h_0)^2 - \Delta^2(T, h_0)}]}, \quad (4)$$
$$k_{\bar\sigma}^\mathrm{h} = \sqrt{(2m/\hbar^2)[E_\mathrm{F} - \sqrt{(E - \eta_{\bar\sigma} h_0)^2 - \Delta^2(T, h_0)}]}. \quad (5)$$

The order parameter $\Delta(T, h)$ of the FS film is determined by the self-consistent equation $\Delta = U < \Psi_\uparrow \Psi_\downarrow >$, where U is the effective attractive potential between electrons, and $\Psi_\uparrow = \sum_k (\gamma_{k\uparrow} u_{k\uparrow} - \gamma_{k\downarrow}^+ v_{k\uparrow}^*)$ and $\Psi_\downarrow = \sum_k (\gamma_{k\downarrow} u_{k\downarrow} + \gamma_{k\uparrow}^+ v_{k\downarrow}^*)$ with $\gamma_{k\uparrow}$ and $\gamma_{k\downarrow}$ are the Bogoliubov transformative operators.[18] With expressions (2)

and (3) as well as $\gamma_{k\uparrow}$ and $\gamma_{k\downarrow}$,[18] we obtain

$$1=\frac{U}{2}\sum_{k}\left(\frac{1-f_{k\uparrow}}{\sqrt{\varepsilon_{k\uparrow}^{2}+\Delta^{2}(T,h_{0})}}-\frac{f_{k\downarrow}}{\sqrt{\varepsilon_{k\downarrow}^{2}+\Delta^{2}(T,h_{0})}}\right), \tag{6}$$

where

$$f_{k\sigma}=\frac{1}{\mathrm{e}^{[\sqrt{\varepsilon_{k\sigma}^{2}+\Delta^{2}(T,h_{0})}-\eta_{\sigma}h_{0}]\beta}+1} \tag{7}$$

with $\varepsilon_{k\sigma}=\dfrac{\hbar^{2}(k_{\sigma}^{\mathrm{e}})^{2}}{2m}-E_{\mathrm{F}}$, and $\beta=1/(k_{\mathrm{B}}T)$ the inverse temperature. From expressions (6) and (7), the self-consistent equation for $\Delta(T,h)$ is obtained as

$$\ln\left(\frac{\Delta_{0}}{\Delta(T,h_{0})}\right)=\int_{0}^{\hbar\omega_{\mathrm{D}}}\frac{\mathrm{d}\varepsilon_{k}}{\sqrt{\varepsilon_{k}^{2}+\Delta^{2}(T,h_{0})}}\left(\frac{1}{\mathrm{e}^{[\sqrt{\varepsilon_{k}^{2}+\Delta^{2}(T,h_{0})}-h_{0}]\beta}+1}+\frac{1}{\mathrm{e}^{[\sqrt{\varepsilon_{k}^{2}+\Delta^{2}(T,h_{0})}+h_{0}]\beta}+1}\right), \tag{8}$$

where $\Delta_{0}=\Delta$ (0,0) is the BCS gap at zero temperature and in the absence of exchange field; ω_{D} is the Debye frequency. By solving Eq.(8) self-consistently, the dependence of the effective order parameter $\Delta(T,h_{0})$ in the FS layer on the temperature is obtained.

In the S/I/N/I/FS junction under consideration, the two layers of insulator are located at x=0 and $x=L$ separately. The insulating barriers may be modelled by a δ-type potential $V(\delta(x)+\delta(x-L))$, where V indicates the barrier strength. The assumed spatial dependence of the effective pair potential is described as

$$\Delta=\begin{cases}\Delta_{\mathrm{s}}(T)\mathrm{e}^{\mathrm{i}\varphi_{\mathrm{L}}}, & x<0,\\ 0, & 0<x<L,\\ \Delta(T,h_{0})\mathrm{e}^{\mathrm{i}\varphi_{\mathrm{R}}}, & x>L.\end{cases} \tag{9}$$

The temperature dependence of the magnitude of the pair potential $\Delta_{\mathrm{s}}(T)$ is assumed to obey the BCS relation. We assume that $\Delta_{\mathrm{s}}(0)=\Delta_{\mathrm{s0}}$ in the S is equal to $\Delta(0,0)$ in the FS. The macroscopic phase $\varphi_{\mathrm{L}}(\varphi_{\mathrm{R}})$ of the S (FS) is measured relatively to the normal to the interface of the superconductor.

Consider an electronlike quasiparticle incident on the insulating barrier at $x=0$ from the S. With general solutions of the Eq.(1), the wavefunctions are given by

$$\Psi_{1\sigma}(x)=\mathrm{e}^{\mathrm{i}k_{\mathrm{L}}^{+}x}\begin{pmatrix}u_{\mathrm{L}}\mathrm{e}^{\mathrm{i}\phi_{\mathrm{L}}/2}\\ v_{\mathrm{L}}\mathrm{e}^{-\mathrm{i}\phi_{\mathrm{L}}/2}\end{pmatrix}+a_{\bar{\sigma}}\mathrm{e}^{\mathrm{i}k_{\mathrm{L}}^{-}x}\begin{pmatrix}v_{\mathrm{L}}\mathrm{e}^{\mathrm{i}\phi_{\mathrm{L}}/2}\\ u_{\mathrm{L}}\mathrm{e}^{-\mathrm{i}\phi_{\mathrm{L}}/2}\end{pmatrix}+b_{\sigma}\mathrm{e}^{-\mathrm{i}k_{\mathrm{L}}^{+}x}\begin{pmatrix}u_{\mathrm{L}}\mathrm{e}^{\mathrm{i}\phi_{\mathrm{L}}/2}\\ v_{\mathrm{L}}\mathrm{e}^{-\mathrm{i}\phi_{\mathrm{L}}/2}\end{pmatrix},\quad x<0, \tag{10}$$

$$\Psi_{2\sigma}(x)=e_{\sigma}\mathrm{e}^{\mathrm{i}q^{+}x}\begin{pmatrix}1\\0\end{pmatrix}+f_{\bar{\sigma}}\mathrm{e}^{-\mathrm{i}q^{-}x}\begin{pmatrix}0\\1\end{pmatrix}+g_{\sigma}\mathrm{e}^{-\mathrm{i}q^{+}x}\begin{pmatrix}1\\0\end{pmatrix}+h_{\bar{\sigma}}\mathrm{e}^{\mathrm{i}q^{-}x}\begin{pmatrix}0\\1\end{pmatrix},\quad 0<x<L, \tag{11}$$

$$\Psi_{3\sigma}(x)=c_{\sigma}\mathrm{e}^{\mathrm{i}k_{\sigma}^{\mathrm{e}}x}\begin{pmatrix}u_{\sigma}\mathrm{e}^{\mathrm{i}\phi_{\mathrm{R}}/2}\\ v_{\bar{\sigma}}\mathrm{e}^{-\mathrm{i}\phi_{\mathrm{R}}/2}\end{pmatrix}+\mathrm{d}_{\bar{\sigma}}\mathrm{e}^{-\mathrm{i}k_{\bar{\sigma}}^{h}x}\begin{pmatrix}v_{\bar{\sigma}}\mathrm{e}^{\mathrm{i}\phi_{\mathrm{R}}/2}\\ u_{\sigma}\mathrm{e}^{-\mathrm{i}\phi_{\mathrm{R}}/2}\end{pmatrix},\quad x>L, \tag{12}$$

where $a_{\sigma},b_{\sigma},c_{\sigma}$ and d_{σ} are, respectively, coefficients of the Andreev reflection, normal reflection, transmission to the FS as electronlike quasiparticles and transmission as holelike quasiparticles. Amplitudes of electrons and holes propagating in the N are given by the coefficients $e_{\sigma},f_{\sigma},g_{\sigma}$ and h_{σ}. $q^{\pm}=[k_{\mathrm{F}}^{2}\pm 2mE/\hbar^{2}]^{1/2}$ are the wave vectors of the electron and hole in the N. The wave vectors of the electronlike and holelike quasiparticles in the s-wave S are given by

$$k_{\rm L}^{\pm} = [k_{\rm F}^2 \pm 2m\sqrt{(E^2 - |\Delta_{\rm s}(T)|^2)}/\hbar^2]^{1/2}, \tag{13}$$

and

$$u_{\rm L}^2 = \frac{1}{2}\left[1 + \sqrt{1 - |\Delta_{\rm s}(T)|^2/E^2}\right], \tag{14}$$

$$v_{\rm L}^2 = \frac{1}{2}\left[1 - \sqrt{1 - |\Delta_{\rm s}(T)|^2/E^2}\right]. \tag{15}$$

The wavefunctions must satisfy the following boundary conditions:[19]

$$\Psi_{1\sigma}(x=0) = \Psi_{2\sigma}(x=0), \tag{16}$$

$$\left(\frac{{\rm d}\Psi_{2\sigma}}{{\rm d}x}\right)_{x=0} - \left(\frac{{\rm d}\Psi_{1\sigma}}{{\rm d}x}\right)_{x=0} = \frac{2mV}{\hbar^2}\Psi_{2\sigma}(x=0), \tag{17}$$

$$\Psi_{2\sigma}(x=L) = \Psi_{3\sigma}(x=L), \tag{18}$$

$$\left(\frac{{\rm d}\Psi_{3\sigma}}{{\rm d}x}\right)_{x=L} - \left(\frac{{\rm d}\Psi_{2\sigma}}{{\rm d}x}\right)_{x=L} = \frac{2mV}{\hbar^2}\Psi_{3\sigma}(x=L). \tag{19}$$

From Eqs.(10)–(12) and (16)–(19), we obtain

$$a_\sigma(\phi, E) = \frac{8[u_\sigma v_{\bar\sigma}(\cos\phi - {\rm i}(u_{\rm L}^2 - v_{\rm L}^2)\sin\phi] - u_{\rm L}v_{\rm L}[8\cos L(q^+ - q^-) + {\rm i}(8-4z^2)(u_\sigma^2 - v_{\bar\sigma}^2)\sin L(q^+ - q^-)]}{4\cos L(q^+ - q^-) - 16u_{\rm L}v_{\rm L}u_\sigma v_{\bar\sigma}\cos\phi - {\rm i}4(2+z^2)(u_\sigma^2 - v_{\rm L}^2)\sin L(q^+ - q^-) + R(u_{\rm L}^2 - v_{\rm L}^2)(u_\sigma^2 - v_{\bar\sigma}^2)}, \tag{20}$$

where $R = 4[z^2\cos L(q^+ + q^-) + \cos L(q^+ - q^-)] + 4z^3\sin L(q^+ + q^-) + z^4[\cos L(q^+ - q^-) - \cos L(q^+ + q^-)]$, $\phi = \phi_{\rm R} - \phi_{\rm L}$, and $z = 2mV/\hbar^2 k_{\rm F}$. Since h_0 is much smaller than $E_{\rm F}$, we have made the Andreev approximation in part.

With the obtained $a_\sigma(\phi, E)$, the Josephson current can be calculated using the generalized coefficient of Andreev reflection $a_\sigma(\phi, {\rm i}\omega_n)$, which is obtained by analytic continuation of E to ${\rm i}\omega_{\rm n}$, yielding [21,22]

$$I = \frac{ek_{\rm B}T\Delta_{\rm s}(T)}{2\hbar}{\rm Re}\sum_{\omega_{\rm n}}\left[\frac{a_\uparrow(\phi, {\rm i}\omega_{\rm n}) - a_\uparrow(-\phi, {\rm i}\omega_{\rm n})}{\Omega_{\rm nL}} + \frac{a_\downarrow(\phi, {\rm i}\omega_{\rm n}) - a_\downarrow(-\phi, {\rm i}\omega_{\rm n})}{\Omega_{\rm nL}}\right] \tag{21}$$

with $\Omega_{\rm nL} = \sqrt{\omega_{\rm n}^2 + \Delta_{\rm s}^2(T)}$, and the Matsubara frequency $\omega_{\rm n} = 2\pi k_{\rm B}T(n + 1/2)$.

Figures 2, 3 and 4 show the numerical results for the Josephson critical currents $I_{\rm c}$ as a function of the normal metal thickness L for the different values of exchange field h_0, barrier strength z and temperature T. Several features can be found from the figures. First, there exists an oscillation relation between the Josephson critical current $I_{\rm c}$ and corresponding normal metal thickness L. The amplitude of the oscillation decreases as L increases. Second, the exchange field and the temperature always suppress the Andreev reflection in the S/I/N/I/FS junction. The Josephson critical current is reduced as the exchange energy increases or the temperature is elevated as shown in Figs.2 and 3. Third, for $z = 0$, Josephson critical current $I_{\rm c}$ dose not exhibit the oscillation behaviour. The value of $I_{\rm c}$ is reduced with thickness L of the N increasing. The maximum values of oscillation currents for $z \neq 0$ are equal to the value of the current at the same thickness L for $z = 0$.

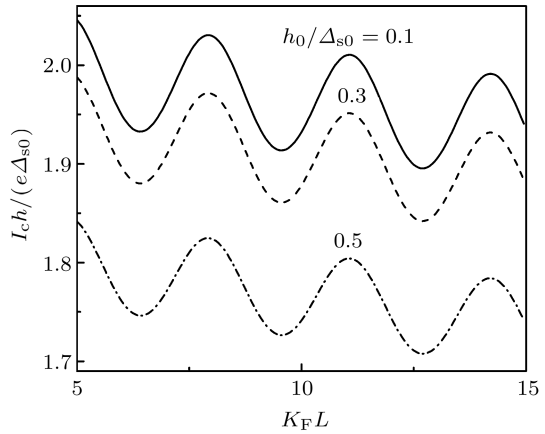


Fig.2. Critical current as a function of normal metal thickness for different exchange fields, where $\phi = \pi/2$, $k_{\rm B}T/\Delta_{\rm s0} = 0.2$ and $z = 0.2$.

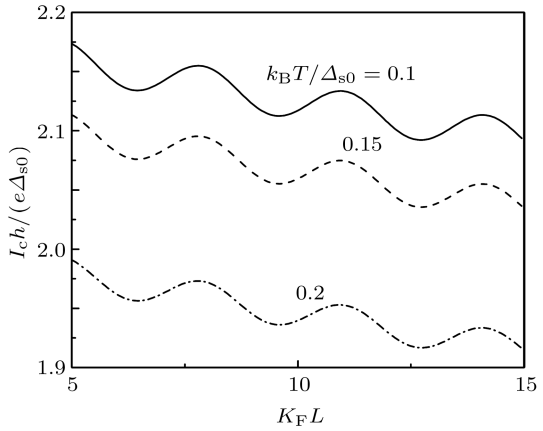


Fig.3. Critical current as a function of normal metal thickness for different temperatures, where $\phi = \pi/2$, $z = 0.1$, and $h_0/\Delta_{s0} = 0.3$.

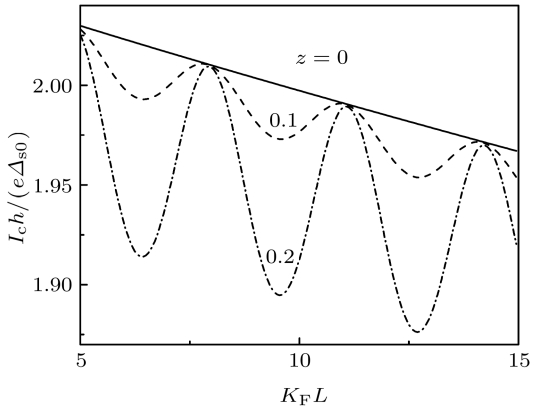


Fig.4. Critical current as a function of normal metal thickness for different barrier strengths, where $\phi = \pi/2$, $k_BT/\Delta_{s0} = 0.2$, and $h_0/\Delta_{s0} = 0.2$.

In Fig.5 the temperature dependences of the Josephson critical current I_c are plotted for different values of h_0. With increasing the temperature, I_c decreases and drops to zero at $T_c(h_0)$, which is lower than $T_c(h_0 = 0)$. Such a sudden drop in I_c in the presence of h_0 is due to the drop of $\Delta(T, h_0)$ from a finite value to zero at $T_c(h_0)$ (Ref.[16]). In Fig.6, the phase dependences of the Josephson critical current I_c are plotted for different values of L. We see that the shape of current-phase relationship is sensitive to the normal metal thickness L. Besides, the Josephson currents do not change from 0 state to π state since

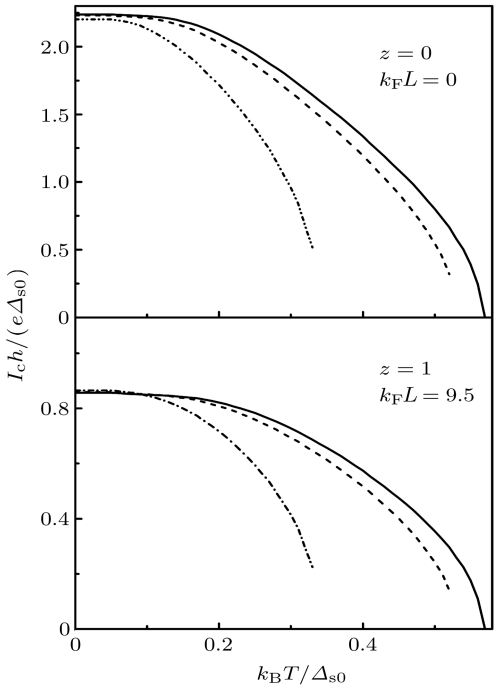


Fig.5. Critical current as a function of normal metal thickness for different temperatures, where $\phi = \pi/2$, $h_0/\Delta_{s0} = 0$ (solid line), $h_0/\Delta_{s0} = 0.3$ (dashed line), and $h_0/\Delta_{s0} = 0.6$ (dash dotted line).

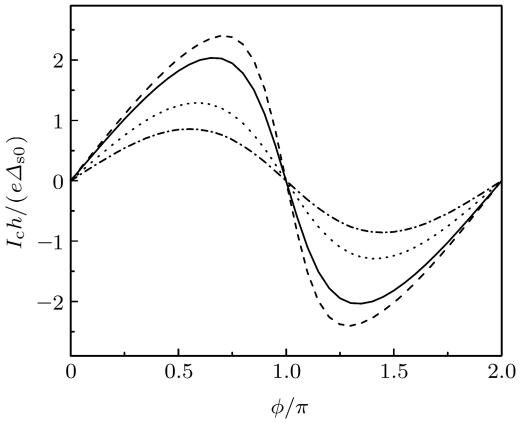


Fig.6. Dependence of current on phase difference for four values of k_FL, where $z = 1$, $h_0/\Delta_{s0} = 0.1$, $k_BT/\Delta_0 = 0.1$, $k_FL = 8$ (solid line), $k_FL = 8.5$ (dash line), $k_FL = 9$ (dot line), and $k_FL = 9.5$ (dash dot line).

h_0 is much smaller than E_F.

In summary we have studied the coherent Josephson current in an S/I/N/I/FS junction. The superconducting order parameter and the energy spectrum of the Bogoliubov excitations are derived from the BdG equation for the FS. The Josephson currents in the S/I/N/I/FS junction are calculated as a function of normal metal thickness, exchange field, temperature, and insulating barrier strength. The oscillation relation between the tunnelling current and corresponding thickness of normal layer is due to the quantum interference effect of quasiparticles in the N interlayer. The amplitude of oscillation of the Josephson current is reduced with increasing normal metal thickness L. The Josephson critical current is suppressed in the presence of exchange field and depends on temperature. It is expected that the theoretical result obtained will be confirmed in the future experiments.

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