Incoherent Pair Tunneling as a Probe of the Cuprate Pseudogap

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We argue that incoherent pair tunneling in a cuprate superconductor junction with an optimally doped superconducting and an underdoped normal lead can be used to detect the presence of pairing correlations in the pseudogap phase of the underdoped lead. We estimate that the junction characteristics most suitable for studying the pair tunneling current are close to recently manufactured cuprate tunneling devices. [S0031-9007(99)09210-8]

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The pseudogap—a depletion of the single particle spectral weight around the Fermi energy—is considered to be one of the most convincing manifestations of the unconventional nature of cuprate superconductivity [1]. The pseudogap regime sets in as the temperature is lowered below a crossover temperature T^*, and extends over a wide range of temperatures in *underdoped* samples [2]. While the pseudogap is clearly present in the spin channel [3], optical conductivity data [4] suggest that the same mechanism is responsible for the gapping of the charge degrees of freedom as well. In addition, specific heat data [5] also provide evidence that a gap opens below T^*. It has been suggested [6,7] that precursor superconducting pairing fluctuations may be responsible for these phenomena; the observations of a smooth crossover from the pseudo- to superconducting gap seen in angle-resolved photoemission [8] and scanning tunneling spectroscopy (STS) [9] lend support to this idea. There are, however, several other competing proposals that do not necessarily involve charge $2e$ pairing [10]. It is therefore of interest to find an experiment which can provide a direct test of the superconducting precursor scenario. Here we propose and analyze an experiment involving incoherent pair tunneling which provides such a test [11].

The measurement of the pair susceptibility in the normal state of a superconductor, is in principle similar [12] to other—say, magnetic—susceptibility measurements: we are interested in finding out the linear response of the system to a polarizing external field. In the present case the role of the external field is played by the rigid pair field of a second superconductor below its transition temperature, which couples to the fluctuating pair field of the normal lead. This coupling leads to an observable contribution to the tunneling current—the incoherent pair tunneling current—provided that the normal state has sizable pairing correlations.

The basic experimental configuration is illustrated in Fig. 1(a). An I-V measurement is made on a tunnel junction formed from an optimally doped cuprate superconductor A and a nonoptimally doped material B, in a temperature range $T_c^A > T > T_c^B$. The c axis is perpendicular to the A and B layers which are separated by an insulating layer. Such a structure could be obtained by varying the doping concentration of a crystal during a layer-by-layer deposition [13]. If the B lead is underdoped, as indicated in Fig. 1(b), there will be a substantial temperature region above T_c^B in which B will have a pseudogap, while if the normal lead is overdoped [denoted by B' in Fig. 1(b)], this pseudogap region will be significantly narrower. Now, for $T_c^A > T > T_c^B$ we can use the superconducting pair field of the optimally doped

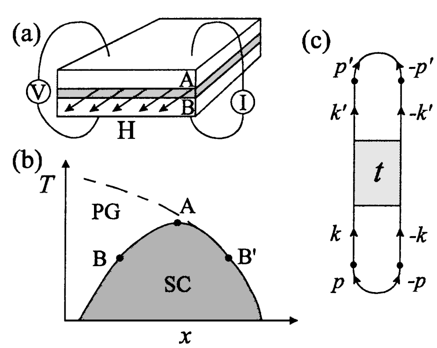

FIG. 1. (a) Proposed experimental configuration for a junction involving two cuprate leads A and B, with transition temperatures as indicated by the phase diagram of (b). (c) Diagrammatic representation of the incoherent pair tunneling current contribution. Lines ($\pm k, \pm k'$) and ($\pm p, \pm p'$) correspond to single-electron propagators of the normal pseudogapped (PG) B lead and anomalous Gor'kov propagators of the superconducting A lead, respectively. The dots ($\bullet$) represent tunneling matrix elements $V_{\mathbf{pk}}$, and the box stands for the particle-particle t matrix of B.

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superconductor to directly probe the strength of the pairing fluctuations in B, by measuring the incoherent pair tunneling contribution $I_p(V)$ to the total tunneling current $I(V)$ [12,14]. If pseudogap behavior is associated with strong precursor superconducting pairing, the contribution from the incoherent pair tunneling $I_p(V)$ should extend over a much wider temperature range than for the overdoped B' lead, even if T_c^B, for the two are equal. Furthermore, if indeed the pseudogap region is characterized by precursor pairing, the voltage structure of $I_p(V)$ provides a measure of the frequency dependence of the imaginary part of the particle-particle t matrix in the pseudogap regime.

The incoherent pair tunneling contribution to the total tunnel current $I(V)$ is shown diagrammatically in Fig. 1(c) [12]. It should be stressed that, due to the very short c-axis coherence length in the cuprates, the pair tunneling takes place between the two cuprate layers on either side of the insulating barrier. In order to analyze the experimental requirements, we consider a circular Fermi surface and assume that the pairing instability occurs in the d-wave channel with a t matrix given by $t_{\mathbf{k},\mathbf{k}',\mathbf{q}}(i\omega_m) = t_{\mathbf{q}}(i\omega_m)\cos\times(2\varphi_{\mathbf{k}})\cos(2\varphi_{\mathbf{k}'})$. Here $i\omega_m = 2m\pi T$ is the bosonic Matsubara frequency (unless noted otherwise, we take $\hbar = 1, k_B = 1$), and $\varphi_k = \arctan(k_y/k_x)$. In the absence of an external magnetic field, the incoherent pair tunneling contribution $I_p(V)$ is given by $4e$ times the imaginary part of the diagram shown in Fig. 1(c) [12],

$$I_p(V) = 4eC^2Sa^2\,\mathrm{Im}\,t_{\mathbf{q}=\mathbf{0}}(i\omega_m \to 2eV + i\delta), \qquad (1)$$

where S is the junction area, a is the lattice spacing, and the coefficient C—which determines the magnitude of the pair current—is given by the following expression:

$$C = \frac{n_i}{N^2}T\sum_{n,\mathbf{p},\mathbf{k}} F_{\mathbf{p}}(i\omega_n)G_{\mathbf{k}}(i\omega_n)G_{-\mathbf{k}}(-i\omega_n) \times \langle|V_{\mathbf{k},\mathbf{p}}|^2\rangle_{\mathrm{imp}}\cos(2\varphi_{\mathbf{k}}). \qquad (2)$$

Here we presume that the mechanism for electron transfer from A to B derives from impurity assisted hopping in the insulating layer separating A and B. We define n_i to be the number of impurity scattering sites per unit area of the insulating layer, N is the number of sites of a layer, and $\langle|V_{\mathbf{pk}}|^2\rangle_{\mathrm{imp}}$ is the impurity averaged single-electron transfer. The momenta $\mathbf{p}$ and $\mathbf{k}$ are two-dimensional vectors. In Eq. (2) we have neglected the weak voltage dependence of C which is justified in the regimes we will be studying where the t matrix dominates the voltage dependence of the pair current.

To estimate the size of the pair current we have used the Bardeen-Cooper-Schrieffer (BCS) form for the Gor'kov function to describe the superconducting A lead,

$$F_{\mathbf{p}}(i\omega_n) = \frac{\Delta_A\cos(2\varphi_{\mathbf{p}})}{\omega_n^2 + \epsilon_{\mathbf{p}}^2 + \Delta_A^2\cos^2(2\varphi_{\mathbf{p}})}. \qquad (3)$$

Here Δ_A is the maximum of the associated d-wave superconducting gap. The detailed nature of the Green's functions in B are of course very important in determining the particle-particle t matrix. However, the coefficient C is obtained by summing over both the momentum and frequency variables of the propagators. Thus C is only marginally affected by the precise form of the single particle propagators. Whether one replaces the product of the B Green's functions by their noninteracting form $G_{\mathbf{k}}(i\omega_n)G_{-\mathbf{k}}(-i\omega_n) = (\omega_n^2 + \epsilon_{\mathbf{k}}^2)^{-1}$, or whether one uses an extreme limit of pseudogap theories [7] $G_{\mathbf{k}}(i\omega_n)G_{-\mathbf{k}}(-i\omega_n) = [\omega_n^2 + \epsilon_{\mathbf{k}}^2 + \Delta_B^2\cos^2(2\varphi_{\mathbf{k}})]^{-1}$, changes the estimate of C only by factors of order unity. In these expressions $\epsilon_{\mathbf{p}}$ and $\epsilon_{\mathbf{k}}$ are the single particle energies in A and B, respectively.

We will assume that the insulating layer gives rise to a diffuse [15] electron transfer with

$$\langle|V_{\mathbf{pk}}|^2\rangle_{\mathrm{imp}} = |V_0|^2 + |V_1|^2\cos(2\varphi_{\mathbf{p}})\cos(2\varphi_{\mathbf{k}}). \qquad (4)$$

More generally one could imagine expanding the impurity averaged single-electron transfer $\langle|V_{\mathbf{pk}}|^2\rangle_{\mathrm{imp}}$ in two-dimensional crystal harmonics. In Eq. (4) we have kept only the uniform and d-wave pair transfer parts. It is the second term in Eq. (4) that will enter in our calculations. The required size of V_1 will be discussed below, together with other junction requirements.

Using Eqs. (2)–(4) one finds that

$$C = \pi^2 n_i N_A(0)N_B(0)\,|V_1|^2\Delta_A T\sum_n I_A(\omega_n)I_B(\omega_n), \qquad (5)$$

where $N_A(0)$ and $N_B(0)$ are the single particle density of states per spin, per site for layer A and B, respectively, and

$$I_{A,B}(\omega_n) = \int \frac{d\varphi_{\mathbf{p}}}{2\pi}\frac{\cos^2(2\varphi_{\mathbf{p}})}{\sqrt{\omega_n^2 + \Delta_{A,B}^2\cos^2 2\varphi_{\mathbf{p}}}} = \frac{2}{\pi\sqrt{\omega_n^2 + \Delta_{A,B}^2}} \times\left[E(k_n) + \left(\frac{\omega_n}{\Delta_{A,B}}\right)^2[E(k_n) - K(k_n)]\right]. \qquad (6)$$

In the above expression $k_n^2 = \Delta_{A,B}^2/(\omega_n^2 + \Delta_{A,B}^2)$, and K and E are the complete elliptic integrals of the first and second kind, respectively. Carrying out the Matsubara sum in Eq. (5), we find that to within numerical factors of order unity,

$$C \simeq \frac{\pi^2}{4}n_i N_A(0)N_B(0)\,|V_1|^2. \qquad (7)$$

Now at low temperatures, where A and B are both superconducting, a similar calculation shows that the Josephson critical current is given by $I_c = 2eC'\Delta_B S$, with Δ_B the low temperature maximum gap in B and the coefficient C' is closely related to C given by Eq. (2) [16].

Using this to normalize the strength of $I_p(V)$ we have

$$\frac{I_p(V)}{I_c} \approx \frac{E_J}{E_c} \,\mathrm{Im}\bar{t}(2eV)\,. \tag{8}$$

Here $E_J = \hbar I_c/2e$ is the zero temperature Josephson coupling energy between A and B, $E_c = (S/a^2)N_B(0)\Delta_B^2/2$ is the condensation energy of the B cuprate layer, and $\mathrm{Im}\bar{t}(\omega) = N_B(0)\,\mathrm{Im}t_0(\omega)$ is a dimensionless form of the t matrix for $\mathbf{q} = 0$.

It can be seen from Eq. (8) that the important quantity measured in a pair tunneling experiment is $\mathrm{Im}\bar{t}(\omega)$. The form of this function varies depending on the particular scenario adopted for describing the pseudogap. For a wide class of theories $\mathrm{Im}t_{\mathbf{q}}(\omega)$ can be expressed in terms of the pair susceptibility $\chi_{\mathbf{q}}(\omega)$ and the pairing coupling constant g [7],

$$\mathrm{Im}t_{\mathbf{q}}(\omega) = \frac{-g^2\,\mathrm{Im}\chi_{\mathbf{q}}(\omega)}{[1 + g\,\mathrm{Re}\chi_{\mathbf{q}}(\omega)]^2 + [g\,\mathrm{Im}\chi_{\mathbf{q}}(\omega)]^2}\,. \tag{9}$$

A useful form of Eq. (9) for experimental comparison is discussed in Ref. [7], although other alternatives may eventually be proposed using different precursor scenarios [6,17]. The approach of Ref. [7] provides a concrete diagrammatic prescription for computing χ. For $\mathbf{q} = 0$ and sufficiently low frequencies, $1 + g\,\mathrm{Re}\chi_0(\omega) \approx (\alpha/\gamma) \times (\omega - \omega_0)$ and $g\,\mathrm{Im}\chi_0(\omega) \approx \omega/\gamma$, where (for T close to T_c^B), $\omega_0 = (\gamma/\alpha)(T/T_c^B - 1)$. The values of α and γ depend on g. Under these conditions, Eq. (9) yields $\mathrm{Im}\bar{t}(\omega) \simeq \gamma\omega/[\alpha^2(\omega - \omega_0)^2 + \omega^2]$. In the weak-coupling limit, where the dimensionless parameter $\alpha \simeq 0$, this yields the well-known result [12,14] $\mathrm{Im}\bar{t}(\omega) = (\omega/\gamma)/[(T/T_c^B - 1)^2 + (\omega/\gamma)^2]$. In this regime the pairing fluctuations are associated with critical behavior and are essentially diffusive in nature. This case was addressed in earlier incoherent pair tunneling experiments [14] on conventional superconductors. By contrast, in the intermediate coupling regime, which corresponds to $\alpha \simeq 1$, the value of ω_0 is strongly reduced, resulting in a pronounced *resonance* in $\mathrm{Im}\bar{t}(\omega)$ at this frequency. Thus, the pair fluctuations acquire a propagating nature [7].

These two theoretical limits are illustrated in Fig. 2 which presents the self-consistently calculated t matrix [7] in weak and intermediate coupling, corresponding to B' and B, respectively. Here $T/T_c^B = 1.1$. Notice the asymmetry $\mathrm{Im}\bar{t}(\omega) \neq \mathrm{Im}\bar{t}(-\omega)$ in the second case which provides a strong signature for pair resonance effects. This asymmetry is, in turn, related to an asymmetric density of states [7], which may be associated with that observed in STS experiments [9]. This figure also reflects the predicted voltage dependence of the pair tunneling current. Within the superconducting pairing fluctuation scenarios of the type discussed in Ref. [7] a prominent peak in $\mathrm{Im}\bar{t}(\omega)$ is expected to persist in underdoped cuprates to temperatures of the order of $T \sim T^*$, considerably higher than T_c^B [18]. Alternative scenarios [6,17] can be used, presumably, to provide analogous signatures, within their respective theoretical framework. The importance of the incoherent pair tunneling experiment lies in its ability to detect the temperature and voltage dependence of such features and therefore to confirm or falsify different classes of pseudogap scenarios.

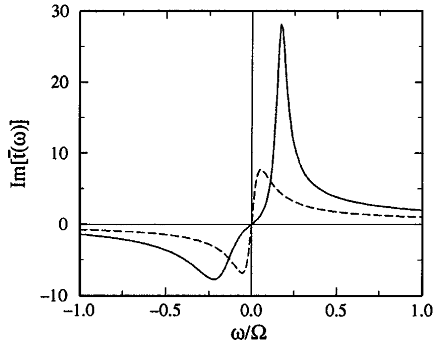

FIG. 2. Predicted voltage dependence of the pair tunneling current; following [7] the solid and dashed curves correspond to under (B)—and overdoped (B') leads at $T/T_c = 1.1$; the slightly doping dependent energy scale Ω is of the order of 100 meV. The dashed curve is similar to that of the conventional fluctuation picture [12]. The asymmetry of the solid curve is an important signature which should be noted.

Let us now estimate the size of the pair current I_p given by Eq. (8). The condensation energy density can be inferred from heat capacity measurements [5]: For an underdoped $YBa_2Cu_3O_{6+x}$ of $T_c \sim 60$ K we obtain $\epsilon_{\mathrm{cond}} \sim 2 \times 10^4$ J/m^3. Using typical values for the junction surface area $S \sim 10^{-8}$ m^2 [19] and taking the layer thickness of order $\ell_c \sim 10$ Å, we find $E_{\mathrm{c}} = \epsilon_{\mathrm{cond}} S \ell_c \sim 2 \times 10^{-13}$ J $\approx 10^6$ eV. For a typical critical current of order $I_c \sim 10$ mA, the corresponding Josephson coupling energy is $E_J = \hbar I_c/2e \sim 3 \times 10^{-18}$ J ≈ 20 eV. Consequently, $I_p \sim I_c(E_J/E_c) \sim 0.2\ \mu$A, which is of the same order as the pair currents detected in conventional superconductors [14].

Thus it is important to fabricate junctions with c axis Josephson current density in the range of 10^2 A/cm^2. Critical current densities sustained by recently fabricated trilayer junctions [13] are in this range. Further complications might be caused by thermal voltage noise [14,20] in the junction circuit due to the relatively elevated temperatures at which these measurements need to be carried out. One possibility would be to use the single-layer Bi2201 compound for both leads: this material has a phase diagram similar to that in Fig. 1(b), but with a relatively low optimal T_c. If achievable, a combined junction—with optimally doped Bi2212 and underdoped Bi2201 as electrodes—might permit extending the temperature window of operation to T^* of the Bi2201 compound. $I_p(V)$ can be suppressed [14] by turning on a magnetic field H in the plane of the junction [see Fig. 1(a)]. The suppression is given by [12] $r = 2eH\lambda_{\mathrm{ab}}^A\xi_{\mathrm{ab}}/\hbar c$, provided that $H < H_{c1}^A$. Using typical numbers for the A-side penetration depth $\lambda_{\mathrm{ab}}^A \sim 10^3$ Å, B-side coherence length near

T_c^B, $\xi_{\rm ab}^B \sim 10^2$ Å, and field $H \sim 10^2$ G, we find $r \sim 0.1$. Once the characteristic voltage feature due to $I_p(V)$ was identified near T_c^B, it could be traced to higher temperatures. The detection of $I_p(V)$ is helped by the fact that typical $Bi_2Sr_2CaCuO_{8+\delta}$ samples act as a stack of N $\sim 10^3$ intrinsic Josephson junctions [21]: the effective quasiparticle gap is NΔ_A and the subgap conductance is suppressed by 1/N.

In conclusion, we have argued that the measurement of the pair tunneling current between an optimally doped and an underdoped cuprate can be used to probe the pairing fluctuations in the pseudogap state. This experiment has, in principle, the potential to reveal whether the pseudogap state is in fact due to pairing fluctuations. Indeed, strong pairing correlations in the pseudogap state will be manifest in a large pair current, as compared to the pair current of a junction where an overdoped lead of the same T_c is used. No such strong doping dependence of the pair current is expected within pseudogap scenarios that do not invoke the onset of strong pairing correlations below T^*. To illustrate this experiment, we have chosen the particular case of a c-axis junction geometry and identified the region of the phase diagram where the experiment should be performed. We have also discussed the range of basic junction parameters suitable for observing the pair current.

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