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Pseudogap-induced asymmetric tunneling in cuprate superconductors

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ABSTRACT

Within the framework of the kinetic energy driven superconducting mechanism, the doping and temperature dependence of the asymmetric tunneling in cuprate superconductors is studied by considering the interplay between the superconducting gap and normal-state pseudogap. It is shown that the asymmetry of the tunneling spectrum in the underdoped regime weakens with increasing doping, and then the symmetric tunneling spectrum recovers in the heavily overdoped regime. The theory also shows that the asymmetric tunneling is a natural consequence due to the presence of the normal-state pseudogap.

Cuprate superconductors are complex materials that exhibit a variety of phases determined not only by temperature but also by charge carrier doping [1–3]. The pairing of electrons in the conventional superconductors [4] occurs at the superconducting (SC) transition temperature T_c, creating an energy gap in the electron excitation spectrum that serves as the SC order parameter. However, in cuprate superconductors, the normal-state pseudogap exists between T_c and the temperature T^*, with T^* is called as the normal-state pseudogap crossover temperature [1–3]. Although T_c takes a domelike shape with the underdoped and overdoped regimes on each side of the optimal doping, where T_c researches its maximum [5], T^* is much larger than T_c in the underdoped regime, then it monotonically decreases with increasing doping, and seems to merge with the T_c in the overdoped regime, eventually disappearing together with superconductivity at the end of the SC dome [1–3]. After intensive investigations over more than two decades, it has become clear that many of the unusual physical properties in cuprate superconductors can be attributed to the emergence of the normal-state pseudogap [1–3].

The complexity in cuprate superconductors is reflected in the quasiparticle excitation spectra [6–8]. The scanning tunneling microscopy/spectroscopy (STM/STS) is a powerful tool to study the quasiparticle properties in cuprate superconductors [7,8], since its remarkable energy and spatial resolution makes it particularly well suited for cuprate superconductors, which are characterized by small energy and short length scales. More accurately, the STM/STS data are proportional to the local density of quasiparticle excitations, and the accounting of their distribution can provide important insight into the nature of cuprate superconductors. In the conventional superconductors, the most complete and convincing evidence for the electron–phonon SC mechanism came from the tunneling spectrum [4,9]. During the last two decades, the tunneling study of cuprate superconductors has revealed many crucial results [7,8,10–13], where the main feature of the differential tunneling conductance spectrum is the quasiparticle excitation gap. Moreover, the presence of the excitations within the SC gap, linearly increasing with energy around $V = 0$, indicates that the SC gap has nodes, and therefore presumably d-wave symmetry. In particular, the most remarkable feature about the tunneling in cuprate superconductors is the fact that the tunneling conductivity between a metallic point and a cuprate superconductor is markedly asymmetric between positive and negative voltage biases [14].

A challenging issue for theory is to explain the asymmetric tunneling in cuprate superconductors. Recently, we [15] have discussed the interplay between the normal-state pseudogap state and superconductivity in cuprate superconductors within the framework of the kinetic energy driven SC mechanism [16], where both the charge carrier pairing state in the particle–particle channel and normal-state pseudogap state in the particle-hole channel arise from the same interaction that originates directly from the kinetic energy by exchanging spin excitations, then there is a coexistence of the SC gap and normal-state pseudogap in the whole SC dome. Furthermore, both the normal-state pseudogap and the SC gap are dominated by one energy scale, and they are the result of the strong electron correlation. Within this microscopic SC theory, some unusual properties of cuprate superconductors in the pseudogap phase have been studied [17], including the humplike anomaly of the specific-heat, the particle-hole asymmetry electronic state, and the unusual evolution of the Fermi arc length with

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doping and temperature, and the results are qualitatively consistent with the experimental results. In this paper, we study the doping and temperature dependence of the asymmetric tunneling in cuprate superconductors along with this line. by considering the interplay between the normal-state pseudogap state and superconductivity, we qualitatively reproduce some main features of the STM/STS measurements on cuprate superconductors in the whole doping range from the underdoped to heavily overdoped [7,8,10–13]. In particular, we show that the remarkably asymmetric tunneling in cuprate superconductors is a natural consequence due to the presence of the normal-state pseudogap.

Although there are hundreds of cuprate SC compounds, they all share a layered structure which contains one or more copper–oxygen planes [6]. In this case, it has been argued strongly [18] that the low-energy physics of these planes is described by the two-dimensional t–J model acting on the Hilbert space with no doubly occupied sites, where the kinetic energy includes the nearest-neighbor (NN) and next NN hopping on a square lattice with the matrix elements denoted as t and t', respectively, while the antiferromagnetic (AF) Heisenberg term with the exchange coupling constant J describes the AF coupling between localized spins. To incorporate the electron motion within the restricted Hilbert space without double electron occupancy, we [19] have developed the charge-spin separation (CSS) fermion-spin theory, where the constrained electron operators are decoupled as $C_{l\uparrow} = h^{\dagger}_{l\uparrow}S^{-}_{l}$ and $C_{l\downarrow} = h^{\dagger}_{l\downarrow}S^{+}_{l}$, with the spinful fermion operator $h_{l\sigma} = e^{-i\Phi_{l\sigma}}h_l$ that describes the charge degree of freedom of the electron together with some effects of spin configuration rearrangements due to the presence of the doped hole itself (charge carrier), while the spin operator S_l keeps track of the spin degree of freedom of the electron, then the electron single occupancy local constraint is satisfied in analytical calculations. In this CSS fermion-spin representation, the t–J model can be expressed explicitly as,

$$H = t\sum_{l\hat{\eta}}(h^{\dagger}_{l+\hat{\eta}\uparrow}h_{l\uparrow}S^{+}_{l}S^{-}_{l+\hat{\eta}} + h^{\dagger}_{l+\hat{\eta}\downarrow}h_{l\downarrow}S^{-}_{l}S^{+}_{l+\hat{\eta}}) - t'\sum_{l\hat{\tau}}(h^{\dagger}_{l+\hat{\tau}\uparrow}h_{l\uparrow}S^{+}_{l}S^{-}_{l+\hat{\tau}} + h^{\dagger}_{l+\hat{\tau}\downarrow}h_{l\downarrow}S^{-}_{l}S^{+}_{l+\hat{\tau}}) - \mu\sum_{l\sigma}h^{\dagger}_{l\sigma}h_{l\sigma} + J_{\rm eff}\sum_{l\hat{\eta}}\mathbf{S}_l\cdot\mathbf{S}_{l+\hat{\eta}}, \qquad (1)$$

where the summations $l\hat{\eta}$ and $l\hat{\tau}$ are carried over NN and next NN bonds, respectively, $\mathbf{S}_l = (S^x_l, S^y_l, S^z_l)$ are spin operators, S^-_l and S^+_l are the spin-lowering and spin-raising operators for the spin $S = 1/2$, respectively, μ is the chemical potential, $J_{\rm eff} = (1-\delta)^2 J$, and $\delta = \langle h^{\dagger}_{l\sigma}h_{l\sigma}\rangle = \langle h^{\dagger}_{l}h_{l}\rangle$ is the charge carrier doping concentration.

Superconductivity, the dissipationless flow of electrical current, is a striking manifestation of a subtle form of quantum rigidity on the macroscopic scale, where a central question is how the SC-state forms? It is all agreed that the electron Cooper pairs are crucial for the form of the SC-state because these electron Cooper pairs behave as effective bosons, and can form something analogous to a Bose condensate that flows without resistance [4,20]. This follows a fact that although electrons repel each other because of the Coulomb interaction, at low energies there can be an effective attraction that originates by the exchange of bosons. In the conventional superconductors, these exchanged bosons are *phonons* that act like a bosonic *glue* to hold the electron pairs together, then these electron Cooper pairs condense into a coherent macroscopic quantum state that is insensitive to impurities and imperfections and hence conducts electricity without resistance [4]. For cuprate superconductors, we [16] have shown in terms of Eliashberg's strong coupling theory [21] that in the doped regime without an AF long-range order the charge carriers are held together in pairs in the particle–particle channel by the effective interaction that originates directly from the kinetic energy of the t–J model (1) by the exchange of *spin excitations*, then the electron Cooper pairs originating from the charge carrier pairing state are due to the charge-spin recombination, and their condensation reveals the SC ground-state. In particular, this SC-state is controlled by both SC gap and quasiparticle coherence, which leads to that the maximal T_c occurs around the optimal doping, and then decreases in both underdoped and overdoped regimes. Furthermore, this same interaction also induces the normal-state pseudogap state in the particle-hole channel [15]. Since this normal-state pseudogap is closely related to the quasiparticle coherent weight, and therefore it suppresses the spectral weight. Following our previous discussions [15,16], the full charge carrier diagonal and off-diagonal Green's functions of the t–J model (1) in the SC-state are evaluated as,

$$g(\mathbf{k},\omega) = \frac{1}{\omega - \xi_{\mathbf{k}} - \Sigma^{(\rm h)}_1(\mathbf{k},\omega) - \bar{\Delta}^2_{\rm h}(\mathbf{k})/[\omega + \xi_{\mathbf{k}} + \Sigma^{(\rm h)}_1(\mathbf{k},-\omega)]}, \qquad (2a)$$

$$\Gamma^{\dagger}(\mathbf{k},\omega) = -\frac{\bar{\Delta}_{\rm h}(\mathbf{k})}{[\omega - \xi_{\mathbf{k}} - \Sigma^{(\rm h)}_1(\mathbf{k},\omega)][\omega + \xi_{\mathbf{k}} + \Sigma^{(\rm h)}_1(\mathbf{k},-\omega)] - \bar{\Delta}^2_{\rm h}(\mathbf{k})}, \qquad (2b)$$

where $\xi_{\mathbf{k}} = Zt\chi_1\gamma_{\mathbf{k}} - Zt'\chi_2\gamma'_{\mathbf{k}} - \mu$ is the mean-field (MF) charge carrier spectrum with the spin correlation functions $\chi_1 = \langle S^{+}_{l}S^{-}_{l+\hat{\eta}}\rangle$ and $\chi_2 = \langle S^{+}_{l}S^{-}_{l+\hat{\tau}}\rangle$, $\gamma_{\mathbf{k}} = (1/Z)\sum_{\hat{\eta}}e^{i\mathbf{k}\cdot\hat{\eta}}$, $\gamma'_{\mathbf{k}} = (1/Z)\sum_{\hat{\tau}}e^{i\mathbf{k}\cdot\hat{\tau}}$, Z is the number of the NN or next NN sites on a square lattice, the effective charge carrier pair gap $\bar{\Delta}_{\rm h}(\mathbf{k})$ is closely associated with the self-energy $\Sigma^{(\rm h)}_2(\mathbf{k},\omega)$ in the particle–particle channel as $\bar{\Delta}_{\rm h}(\mathbf{k}) = \Sigma^{(\rm h)}_2(\mathbf{k},\omega = 0)$, and can be expressed explicitly as a d-wave form $\bar{\Delta}_{\rm h}(\mathbf{k}) = \bar{\Delta}_{\rm h}\gamma^{(\rm d)}_{\mathbf{k}}$ with $\gamma^{(\rm d)}_{\mathbf{k}} = (\cos k_x - \cos k_y)/2$, while the self-energy $\Sigma^{(\rm h)}_1(\mathbf{k},\omega)$ in the particle-hole channel renormalizes the MF charge carrier spectrum, and can be rewritten approximately as $\Sigma^{(\rm h)}_1(\mathbf{k},\omega) \approx [2\bar{\Delta}_{\rm pg}(\mathbf{k})]^2/[\omega + M_{\mathbf{k}}]$, where $M_{\mathbf{k}}$ is the energy spectrum of $\Sigma^{(\rm h)}_1(\mathbf{k},\omega)$, and $\bar{\Delta}_{\rm pg}(\mathbf{k})$ is the effective normal-state pseudogap. With these above definitions, the Green's functions in Eq. (2) are obtained explicitly as [15],

$$g(\mathbf{k},\omega) = \sum_{\nu=1,2}\left(\frac{U^2_{\nu h\mathbf{k}}}{\omega - E_{\nu h\mathbf{k}}} + \frac{V^2_{\nu h\mathbf{k}}}{\omega + E_{\nu h\mathbf{k}}}\right), \qquad (3a)$$

$$\Gamma^{\dagger}(\mathbf{k},\omega) = \sum_{\nu=1,2}(-1)^{\nu}\frac{\alpha_{\nu\mathbf{k}}\bar{\Delta}_{\rm h}(\mathbf{k})}{2E_{\nu h\mathbf{k}}}\left(\frac{1}{\omega - E_{\nu h\mathbf{k}}} - \frac{1}{\omega + E_{\nu h\mathbf{k}}}\right), \qquad (3b)$$

where $\nu = 1, 2$, $\alpha_{\nu\mathbf{k}}$, $M_{\mathbf{k}}$, $\bar{\Delta}_{\rm pg}(\mathbf{k})$, $\bar{\Delta}_{\rm h}$, the coherence factors $U_{\nu h\mathbf{k}}$ and $V_{\nu h\mathbf{k}}$, and the charge carrier quasiparticle spectrum $E_{\nu h\mathbf{k}}$ have been given in Ref. [15].

In the framework of the CSS fermion-spin theory [19], the physical electron operator is given by a composite one. In this case, the d-wave charge carrier pairing state based on the exchange of the spin excitations also leads to form a d-wave electron Cooper pairing state [16] due to the charge-spin recombination [22]. This follows a fact that the electron Green's function is a convolution of the spin Green's function and charge carrier Green's function in the CSS fermion-spin representation [23]. In particular, the electron diagonal Green's function in the present case is evaluated explicitly in terms of the charge carrier diagonal Green's function (3a) and spin Green's function $D^{(0)-1}(\mathbf{k},\omega) = (\omega^2 - \omega^2_{\mathbf{k}})/B_{\mathbf{k}}$ as [17],

$$G(\mathbf{k},\omega) = \frac{1}{N}\sum_{\mathbf{p},\nu=1,2}\frac{B_{\mathbf{p}+\mathbf{k}}}{2\omega_{\mathbf{p}+\mathbf{k}}}\left[U^2_{\nu h\mathbf{p}}\left(\frac{n_{\rm F}(E_{\nu h\mathbf{p}}) + n_{\rm B}(\omega_{\mathbf{p}+\mathbf{k}})}{\omega + E_{\nu h\mathbf{p}} - \omega_{\mathbf{p}+\mathbf{k}}} + \frac{1 - n_{\rm F}(E_{\nu h\mathbf{p}}) + n_{\rm B}(\omega_{\mathbf{p}+\mathbf{k}})}{\omega + E_{\nu h\mathbf{p}} + \omega_{\mathbf{p}+\mathbf{k}}}\right) + V^2_{\nu h\mathbf{p}}\left(\frac{1 - n_{\rm F}(E_{\nu h\mathbf{p}}) + n_{\rm B}(\omega_{\mathbf{p}+\mathbf{k}})}{\omega - E_{\nu h\mathbf{p}} - \omega_{\mathbf{p}+\mathbf{k}}} + \frac{n_{\rm F}(E_{\nu h\mathbf{p}}) + n_{\rm B}(\omega_{\mathbf{p}+\mathbf{k}})}{\omega - E_{\nu h\mathbf{p}} + \omega_{\mathbf{p}+\mathbf{k}}}\right)\right], \qquad (4)$$

where $n_{\rm B}(\omega)$ and $n_{\rm F}(\omega)$ are the boson and fermion distribution functions, respectively, while the spin excitation spectrum $\omega_{\mathbf{p}}$ and the function $B_{\mathbf{p}}$ have been given in Ref. [23].