The charge asymmetry in superconductivity of hole- and electron-doped cuprates

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Within the $t$-$t'$-$J$ model, the charge asymmetry in superconductivity of hole- and electron-doped cuprates is studied based on the kinetic energy driven superconducting mechanism. It is shown that superconductivity appears over a narrow range of doping in electron-doped cuprates, and the superconducting transition temperature displays the same kind of the doping controlled behavior that is observed in the hole-doped case. However, the maximum achievable superconducting transition temperature in the optimal doping in electron-doped cuprates is much lower than that of the hole-doped case due to the electron-hole asymmetry.

The parent compounds of cuprate superconductors are believed to belong to a class of materials known as Mott insulators with the antiferromagnetic (AF) long-range order (AFLRO), then superconductivity occurs by the electron or hole doping. Both hole-doped and electron-doped cuprate superconductors have the layered structure of the square lattice of the CuO$_2$ plane separated by insulating layers. It has been found from experiments that only an approximate symmetry in the phase diagram exists about the zero doping line between electron- and hole-doped cuprates. For hole-doped cuprates, AFLRO disappears rapidly with doping, and is replaced by the charge-spin separation (CSS) mechanism and magnetic excitations, and show that both hole- and electron correlation is common for both hole-doped and electron-doped cuprates.

The charge asymmetry in superconductivity of hole- and electron-doped cuprates appears over a narrow range of doping, and then decreases in both underdoped and overdoped regimes. In this paper, we study the charge asymmetry in superconductivity of hole- and electron-doped cuprates along with this line. We show that superconductivity appears over a narrow range of doping in electron-doped cuprates, and the maximum achievable SC transition temperature in the optimal doping is lower than that of the hole-doped case due to the electron-hole asymmetry.

In both hole- and electron-doped cuprates, the characteristic feature is the presence of the two-dimensional CuO$_2$ plane as mentioned above, and it seems evident that the unusual behaviors are dominated by this plane. Although the $t$-$J$ model captures the essential physics of the doped CuO$_2$ plane, the electron-hole asymmetry may be accounted for by including further neighbor hopping. Therefore we start from the $t$-$t'$-$J$ model,

$$H = -t \sum_{\nu \tau} C_{i \sigma}^\dagger C_{i+\nu \tau} + t' \sum_{i \tau} C_{i \sigma}^\dagger C_{i+\tau \sigma} + \mu \sum_{i \sigma} C_{i \sigma}^\dagger C_{i \sigma} + J \sum_{\nu \eta} S_i \cdot S_{i+\nu \eta},$$

with $\eta = \pm \hat{x}, \pm \hat{y}$, $\hat{\tau} = \pm \hat{x} \pm \hat{y}$, $C_{i \sigma}^\dagger \sigma C_{i \sigma}$ is the electron creation (annihilation) operator, $S_i = C_i^\dagger \sigma \sigma C_i / 2$ is spin operator with $\sigma = (\sigma_x, \sigma_y, \sigma_z)$ as Pauli matrices, and $\mu$ is the chemical potential. For the electron doping, we can perform a particle-hole transformation $C_{i \sigma} \rightarrow C_{i-\sigma}^\dagger$, so that the difference between hole and electron doping is expressed as the sign difference of the hopping parameters, i.e., $t > 0$ and $t' > 0$ for hole doping and $t < 0$ and $t' < 0$ for electron doping, then the $t$-$t'$-$J$ model (1) in both hole- and electron-doped cases is always subject to an important on-site local constraint to avoid the double
occupancy, i.e., $\sum_\sigma C_{i\sigma}^\dagger C_{i\sigma} \leq 1$. This single occupancy local constraint can be treated properly within the CSS fermion-spin theory\textsuperscript{19}, where the constrained electron operators are decoupled as, $C_{i\uparrow} = h_{i\uparrow}^\dagger S_i^- + C_{i\downarrow} = h_{i\downarrow}^\dagger S_i^+$, with the spinful fermion operator $h_{i\sigma} = e^{-i\Phi_{i\sigma}} h_i$ describes the charge degree of freedom together with some effects of the spin configuration rearrangements due to the presence of the doped charge carrier itself (dressed charge carrier), while the spin operator $S_i$ describes the spin degree of freedom (dressed spin), then the electron local constraint for the single occupancy is satisfied in analytical calculations, and low-energy behavior of the $t$-$t'$-$J$ model (1) in this CSS fermion-spin representation can be expressed as\textsuperscript{17-19},

\[
H = -t \sum_{i\bar{i}} \langle h_{i\uparrow}^\dagger h_{i\uparrow}^\dagger S_{i\bar{i}}^- + h_{i\downarrow}^\dagger S_{i\bar{i}}^- h_{i\bar{i}\uparrow}^\dagger S_{i\bar{i}}^+ + t' \sum_{i\bar{i}} \langle h_{i\uparrow}^\dagger S_{i\bar{i}}^+ h_{i\uparrow}^\dagger S_{i\bar{i}}^- + h_{i\downarrow}^\dagger S_{i\bar{i}}^+ h_{i\bar{i}\uparrow}^\dagger S_{i\bar{i}}^- \rangle
- \mu \sum_{i\bar{i}} h_{i\sigma} h_{i\bar{i}\sigma} + J_{\text{eff}} \sum_{i\bar{i}} S_i \cdot S_{i\bar{i}},
\]

with $J_{\text{eff}} = (1 - x)^2 J$, and $\delta = \langle h_{i\uparrow}^\dagger h_{i\uparrow}^\dagger \rangle = \langle h_{i\downarrow}^\dagger h_{i\downarrow}^\dagger \rangle$ is the doping concentration. As a consequence, the kinetic energy terms in the $t$-$t'$-$J$ model have been expressed as the interactions between dressed charge carriers and spins, which reflects that even the kinetic energy terms in the $t$-$t'$-$J$ Hamiltonian have strong Coulombic contributions due to the restriction of no doubly occupancy of a given site. These interactions between dressed charge carriers and spins are quite strong, and we\textsuperscript{17,18} have shown that in the case without AFLRO, these interactions can induce the dressed charge carrier pairing state (then the electron Cooper pairing state) by exchanging dressed spin excitations in the higher power of the doping concentration $\delta$. Since the SC state in both hole- and electron-doped cuprates is characterized by the electron Cooper pairs, forming SC quasiparticles\textsuperscript{9,12}, and in the real space the gap function and pairing force have range of one lattice spacing\textsuperscript{23}, therefore the order parameter for the electron Cooper pair can be expressed as, \[
\Delta = (C_{i\uparrow}^\dagger C_{i\uparrow}^\dagger - C_{i\downarrow}^\dagger C_{i\downarrow}^\dagger) = \langle h_{i\uparrow}^\dagger h_{i\uparrow}^\dagger S_i^+ h_{i\downarrow}^\dagger S_i^- - h_{i\downarrow}^\dagger h_{i\downarrow}^\dagger S_i^+ h_{i\uparrow}^\dagger S_i^- \rangle = -\langle S_i^+ S_i^- \rangle = \Delta_h, \]
with the dressed charge carrier pairing order parameter $\Delta_h = \langle h_{i\uparrow}^\dagger h_{i\uparrow}^\dagger - h_{i\downarrow}^\dagger h_{i\downarrow}^\dagger \rangle$, which shows that the SC order parameter is closely related to the dressed charge carrier pairing amplitude, and is proportional to the number of charge carriers, and not to the number of electrons. Following the Eliashberg’s strong coupling theory\textsuperscript{24}, we obtain the self-consistent equations that satisfied by the full dressed charge carrier diagonal and off-diagonal Green’s functions\textsuperscript{17},

\[
g(k) = g^{(0)}(k) + g^{(0)}(k) \Sigma^{(h)}(k) g(k) - \Sigma^{(h)}(-k) \Sigma^{(h)}(k),
\]

\[
\Sigma^{(h)}(k) = g^{(0)}(-k) \Sigma^{(h)}(-k) \Sigma^{(h)}(-k)
\]

respectively, where the four-vector notation $k = (k, \mathbf{p}, \omega_n)$, the dressed charge carrier mean-field (MF) diagonal Green’s function\textsuperscript{17} $g^{(0)}(k) = i\omega_n - \xi_k$, the MF dressed charge carrier excitation spectrum $\xi_k = Z \chi_{\uparrow \uparrow} - Z t' \chi_{\uparrow \downarrow} - \mu$, with \(\chi = (1/Z) \sum_\eta e^{i\Phi_\eta}, \gamma_k' = (1/Z) \sum_\tau e^{i\tau k}, Z \) is the number of the nearest neighbor or second-nearest neighbor sites, the dressed spin correlation functions $\chi_1 = (S_{n+1}^+ S_{n+1}^-)$ and $\chi_2 = (S_{n+1}^+ S_{n+1}^+)$, and the dressed charge carrier self-energy functions\textsuperscript{17,18},

\[
\Sigma^{(h)}(k) = \frac{1}{N^2} \sum_{p\cdot p'} (Z t' g_{p+p'+k} - Z t' g_{p+p'+k})^2 
\times \frac{1}{\beta} \sum_{p\cdot p'} g(p + k) \frac{1}{\beta} \sum_{p\cdot p'} D^{(0)}(p') D^{(0)}(p' + p),
\]

\[
\Sigma^{(h)}(k) = \frac{1}{N^2} \sum_{p\cdot p'} (Z t' g_{p+p'+k} - Z t' g_{p+p'+k})^2 
\times \frac{1}{\beta} \sum_{p\cdot p'} g(-p - k) \frac{1}{\beta} \sum_{p\cdot p'} D^{(0)}(p') D^{(0)}(p' + p),
\]

where $p = (p, i\omega_n)$, $p' = (p', i\omega_n)$, and the MF dressed spin Green’s function\textsuperscript{17},

\[
D^{(0)}(p) = \frac{B_p}{(i\omega_n)^2 - \omega_p^2},
\]

with $B_p = \lambda_1 [2 \chi_1 (\epsilon_p - 1) - \lambda_2 (2 \chi_2 \gamma_k' - \chi_2), \lambda_1 = 2 Z J_{\text{eff}}, \lambda_2 = 4 Z \phi_0 t', \epsilon = 1 + 2t_0 / J_{\text{eff}}$, and the MF dressed spin excitation spectrum $\omega_p^2 = A_1 (\gamma_k')^2 + A_2 (\gamma_k')^2 + A_3 \gamma_k' \gamma_k' + A_4 \gamma_k' + A_5 \gamma_k' + A_6$, with $A_1 = \alpha \lambda_1^2 (\chi_1^2 + \chi_2^2)/2, A_2 = \alpha \lambda_2^2 \chi_2^2, A_3 = -\alpha \lambda_1 \lambda_2 (\chi_1^2 + \chi_2^2) \chi_1^2 / 2, A_4 = -\alpha \lambda_1 \chi_2 (\chi_1^2 + \chi_2^2) / 2, A_5 = -\chi_1^2 (\chi_1^2 + \chi_2^2) / 2, A_6 = \lambda_1^2 (\chi_1^2 + \chi_2^2) / 2, \gamma_k' / 2, A_7 = \alpha \lambda_1 \lambda_2 (\chi_1^2 + \chi_2^2) / 2, A_8 = \lambda_1^2 (\chi_1^2 + \chi_2^2) / 2, A_9 = \alpha \lambda_1 \lambda_2 \chi_2, A_{10} = \lambda_2^2 (\chi_1^2 + \chi_2^2) / 2, A_{11} = -\alpha \lambda_1 \lambda_2 (\chi_1^2 + \chi_2^2) / 2, A_{12} = \lambda_1^2 (\chi_1^2 + \chi_2^2) / 2, A_{13} = -\alpha \lambda_1 \lambda_2 \chi_2, A_{14} = \lambda_2^2 (\chi_1^2 + \chi_2^2) / 2$. In order to satisfy the sum rule of the correlation function \(\sum_{n\cdot n'} (S_{n+n}^+ S_{n+n}^-) = 1/2 \) in the case without AFLRO, the important decoupling parameter $\alpha$ has been introduced in the MF calculation\textsuperscript{17}, which can be regarded as the vertex correction.

In Eq. (4), the self-energy function $\Sigma^{(h)}(k)$ is called as the effective dressed charge carrier gap function since it contains both pairing force and dressed charge carrier gap function, while the self-energy function $\Sigma^{(h)}(k)$ renormalizes the MF dressed charge carrier spectrum, and there-
fore it describes the single particle (quasiparticle) coherence. In particular, \(\Sigma^{(h)}(k)\) is an even function of \(i\omega_n\), while \(\Sigma^{(h)}_1(k)\) is not. For the convenience of discussions, we separate \(\Sigma^{(h)}_1(k)\) into its symmetric and antisymmetric parts as, \(\Sigma^{(h)}_1(k) = \Sigma^{(h)}_{1s}(k) + i\omega_n\Sigma^{(h)}_{1a}(k)\), then \(\Sigma^{(h)}_1(k)\) and \(\Sigma^{(h)}_{1a}(k)\) are both even functions of \(i\omega_n\). According to the Eliashberg’s strong coupling theory, we can define the charge carrier single particle (quasiparticle) coherent weight \(Z^{(1)}_F(k) = 1 - \Sigma^{(h)}_{1a}(k)\). On the other hand, the retarded function \(\text{Re}\Sigma^{(h)}(k)\) may be a constant, independent of \((k, \omega)\). It just renormalizes the chemical potential, and therefore can be neglected. Furthermore, we only study the static limit of the effective dressed charge carrier gap function and single particle coherent weight, i.e., \(\Sigma^{(h)}_2(k) = \Delta(k)\). \(Z^{(1)}_F(k)\) is a function of \(k\), the wave vector dependence is unimportant, since everything happens at the electron Fermi surface. As in the previous discussions within the t-J model, the special wave vector can be estimated qualitatively from the electron momentum distribution as \(k_0 = k_0 - k_F\) with \(k_A = [\pi, \pi]\) and \(k_F \approx (1 - x)\pi/2, (1 - x)\pi/2\), which guarantees \(Z_F = Z_F(k_0)\) near the electron Fermi surface. In this case, the dressed charge carrier diagonal and off-diagonal Green’s functions in Eq. (3) can be rewritten explicitly as,

\[
g(k) = \frac{1}{2}\left(1 + \frac{\xi_k}{E_k}\right)\frac{Z_F}{i\omega_n - E_k} + \frac{1}{2}\left(1 - \frac{\xi_k}{E_k}\right)\frac{Z_F}{i\omega_n + E_k},
\]

\[
\gamma(k) = -\frac{1}{2} \frac{\Delta_{hZ}(k)}{E_k}\left(\frac{Z_F}{i\omega_n - E_k} - \frac{Z_F}{i\omega_n + E_k}\right),
\]

with \(\xi_k = Z_F\xi_k\), \(\Delta_{hZ}(k) = Z_F\Delta_{h}(k)\), and the dressed charge carrier quasiparticle spectrum \(E_k = \sqrt{E_k^2 + (\Delta_{hZ}(k))^2}\).

Although superconductivity with both d-wave and s-wave symmetries appear within the t-J model, the SC state has a dominant d-wave symmetry in the optimal doping. Moreover, we have discussed the effect of the additional second neighbor hoping on superconductivity, and found that the d-wave SC pairing correlation is enhanced, while the s-wave SC pairing correlation is heavily suppressed. In this paper, we are interested in the charge asymmetry in superconductivity of hole- and electron-doped cuprates. To make the discussions simpler, we only consider the d-wave case, i.e., \(\Delta_{hZ}(k) = \Delta_{hZ}(k)^{d}\), with \(\Delta_{hZ}(k) = (\cos k_x - \cos k_y)/2\). In this case, the dressed charge carrier effective gap parameter and single particle coherent weight in Eq. (4) satisfy the following equations,

\[
1 = \frac{1}{N^2} \sum_{k, q, p} (Z\gamma_{k+q} - Z\gamma_{k+q}^\dagger \gamma_{k-p+q} \gamma_{k-p+q}^\dagger \frac{Z_F}{E_k} \frac{B_q B_p}{\omega_q \omega_p}),
\]

\[
Z_F = 1 + \frac{1}{N^2} \sum_{q, p} (Z\gamma_{k+q+p} - Z\gamma_{k+q+p}^\dagger \gamma_{k-p+q+p} \gamma_{k-p+q+p}^\dagger \frac{Z_F}{E_k} \frac{B_q B_p}{4\omega_q \omega_p}),
\]

\[
\times \left(\frac{F_1^{(1)}(k, q, p)}{(\omega_p - \omega_q - E_k)^2} + \frac{F_1^{(2)}(k, q, p)}{(\omega_p + \omega_q - E_k)^2}\right),
\]

\[
\times \left(\frac{F_2^{(1)}(k, q, p)}{(\omega_p - \omega_q - E_{k-q-p})^2} + \frac{F_2^{(2)}(k, q, p)}{(\omega_p + \omega_q + E_{k-q-p})^2}\right),
\]

\[
\times \left(\frac{F_2^{(3)}(k, q, p)}{(\omega_p + \omega_q - E_{k-q-p})^2} + \frac{F_2^{(4)}(k, q, p)}{(\omega_p + \omega_q + E_{k-q-p})^2}\right),
\]

\[
\phi_1 = \frac{1}{2N} \sum_k \gamma_k \left(1 - \frac{\xi_k}{E_k} \coth \frac{1}{2} \frac{\beta E_k}{}\right),
\]

\[
\phi_2 = \frac{1}{2N} \sum_k \gamma_k' \left(1 - \frac{\xi_k}{E_k} \coth \frac{1}{2} \frac{\beta E_k}{}\right),
\]

\[
\delta = \frac{1}{2N} \sum_k \left(1 - \frac{\xi_k}{E_k} \coth \frac{1}{2} \frac{\beta E_k}{}\right),
\]

\[
\chi_1 = \frac{1}{N} \sum_k \frac{B_k}{\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
\chi_2 = \frac{1}{N} \sum_k \frac{B_k}{\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
C_1 = \frac{1}{N} \sum_k \frac{B_k}{\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
C_2 = \frac{1}{N} \sum_k \frac{B_k}{\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
C_3 = \frac{1}{N} \sum_k \frac{B_k'}{\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
\chi_1' = \frac{1}{N} \sum_k \frac{B_k}{2\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]

\[
\chi_2' = \frac{1}{N} \sum_k \frac{B_k}{2\omega_k} \coth \frac{1}{2} \frac{\beta \omega_k}{},
\]
\[ C_i^j = \frac{1}{N} \sum_k \gamma_k^2 B_z(k) \coth \left( \frac{1}{2} \beta \omega_z(k) \right), \]
\[ C_3^3 = \frac{1}{N} \sum_k \gamma_k \gamma_k^* B_z(k) \coth \left( \frac{1}{2} \beta \omega_z(k) \right), \]
then all the above order parameters, decoupling parameter \( \alpha \), and chemical potential \( \mu \) are determined by the self-consistent calculation.\textsuperscript{17, 18}

It has been shown\textsuperscript{17} that the dressed charge carrier pairing state originating from the kinetic energy terms by exchanging dressed spin excitations can lead to form the electron Cooper pairing state, where the SC gap function is obtained from the electron off-diagonal Green’s function \( \Gamma^i(i, t, t - t') = \langle \langle C_i^1(t) C_j^1(t') \rangle \rangle \), which is a convolution of the dressed spin Green’s function and dressed charge carrier off-diagonal Green’s function, and reflects the charge-spin recombination.\textsuperscript{20} In the present case, this electron off-diagonal Green’s function can be obtained in terms of the MF dressed spin Green’s function (5) and dressed charge carrier off-diagonal Green’s function (6b) as,
\[
\Gamma^i(k) = \frac{1}{N} \sum_p Z_F \Delta_{hZ}(p - k) B_p \frac{E_p - k}{2 \omega_p} \times \left( \frac{\omega_p + E_{p-k}}{\omega_n} \right) \left( \frac{n_B(\omega_p) + n_F(-E_{p-k})}{(\omega_n)^2 - (\omega_p + E_{p-k})^2} \right)
\]
\[
- \left( \frac{\omega_p - E_{p-k}}{\omega_n} \right) n_B(\omega_p) + n_F(E_{p-k}) \left( \frac{(\omega_n)^2 - (\omega_p - E_{p-k})^2}{(\omega_n)^2 - (\omega_p - E_{p-k})^2} \right).\]

With the help of this electron off-diagonal Green’s function, the SC gap function is obtained as \( \Delta(k) = -(1/\beta) \sum_{i \omega_n} \Gamma^i(k, i \omega_n) \), and can be evaluated explicitly,
\[
\Delta(k) = \frac{1}{N} \sum_p Z_F \Delta_{hZ}(p - k) \frac{E_{p-k}}{2 \omega_p} \times \tan \left( \frac{1}{2} \beta E_{p-k} \right) \frac{B_p}{2 \omega_p} \coth \left( \frac{1}{2} \beta \omega_p \right),
\]
which shows that the SC transition temperature \( T_c \) occurring in the case of the SC gap parameter \( \Delta = 0 \) is identical to the dressed charge carrier pair transition temperature occurring in the case of the effective dressed charge carrier pairing gap parameter \( \Delta_{hZ} = 0 \). Since the absolute values of \( t \) and \( t' \) are almost same for both hole- and electron-doped cuprates\textsuperscript{21}, and therefore in this paper, the commonly used parameters are chosen as \( t/J = 2.5 \) and \( t'/J = 0.3 \) for the hole doping, and \( t/J = -2.5 \) and \( t'/J = 0.3 \) for the electron doping. In Fig. 1, we plot the SC transition temperature \( T_c \) as a function of the doping concentration \( \delta \) for (a) the electron doping and (b) the hole doping in comparison with the corresponding experimental results of Pr\(_{2-x}\)Ce\(_x\)CuO\(_{4-y}\)\textsuperscript{8} and La\(_{2-x}\)Sr\(_x\)CuO\(_4\)\textsuperscript{8} (inset). Our results indicate that for the hole-doped case, superconductivity appears over a wide range of doping, where the maximal SC transition temperature \( T_c \) occurs around the optimal doping concentration \( \delta_{opt} \approx 0.15 \), and then decreases in both underdoped and overdoped regimes. In analogy to the phase diagram of the hole-doped case, superconductivity appears over a narrow range of doping in the electron-doped side, where the SC transition temperature \( T_c \) increases sharply with increasing doping in the underdoped regime, and reaches a maximum in the optimal doping \( \delta_{opt} \approx 0.14 \), then decreases sharply with increasing doping in the overdoped regime. However, the maximum achievable SC transition temperature in the optimal doping in electron-doped cuprates is much lower than that of the hole-doped case due to the electron-hole asymmetry. Using an reasonably estimative value of \( J \sim 800K \) to \( 1200K \) in doped cuprates, the SC transition temperature in the optimal doping is \( T_c \approx 0.22J \approx 176K \sim 264K \) for the hole-doped case, and \( T_c \approx 0.136J \approx 108K \sim 163K \) for the electron-doped case, in qualitative agreement with the corresponding experimental data.\textsuperscript{3, 4, 8}

The essential physics of the doping dependent SC transition temperature in the electron-doped case is almost the same as in the hole-doped side, and detailed explanations have been given in Ref. [18]. In the framework of the kinetic energy driven superconductivity,\textsuperscript{17} the self-energy function \( \Sigma_{\nu Z}(k) \) describes the single particle (quasiparticle) coherence, and therefore \( Z_F \) is closely related to the quasiparticle density. Since the SC state is established through an emerging quasiparticle,\textsuperscript{27} then the SC state is controlled by both gap function and quasiparticle coherence, which is reflected explicitly in the self-consistent equations (7a) and (7b). It has been shown that the doping dependent behavior of the single particle coherent weight resembles that of the superfluid density in doped cuprates,\textsuperscript{18} i.e., \( Z_F \) grows linearly with the doping concentration in the underdoped and optimally doped regimes, and then decreases with increasing doping in the overdoped regime, which leads to that the SC transition temperature reaches a maximum in the optimal doping.

**FIG. 1.** The superconducting transition temperature as a function of the doping concentration with (a) \( t/J = -2.5 \) and \( t'/J = 0.3 \) for the electron doping and (b) \( t/J = 2.5 \) and \( t'/J = 0.3 \) for the hole doping. Inset: the corresponding experimental results of Pr\(_{2-x}\)Ce\(_x\)CuO\(_{4-y}\) taken from Ref. [8] and La\(_{2-x}\)Sr\(_x\)CuO\(_4\) from Ref. [3].
and then decreases in both underdoped and overdoped regimes. On the other hand, it has been shown\textsuperscript{21} that AFLRO can be stabilized by the $t'$ term for the electron-doped case, which may lead to the charge carrier's localization over a broader range of doping, this is also why superconductivity appears over a narrow range of doping in electron-doped cuprates.

In the CSS fermion-spin theory\textsuperscript{19}, the physical electron isdecoupled as the dressed charge carrier $h_{i\sigma} = e^{-i\Phi_{i\sigma}} h_i$ and spin $S_i$. Since the phase factor $\Phi_{i\sigma}$ in the dressed charge carrier is separated from the bare spinon operator, and then it describes a spin cloud\textsuperscript{19}. Therefore the dressed charge carrier $h_{i\sigma}$ is a spinless fermion $h_i$ (bare charge carrier) incorporated the spin cloud $e^{-i\Phi_{i\sigma}}$ (magnetic flux), thus is a magnetic dressing. In other words, the dressed charge carrier carries some spin messages, i.e., it shares its nontrivial spin environment. It has been shown\textsuperscript{19} that these dressed charge carrier and spin are gauge invariant under a local U(1) gauge transformation, and in this sense, they are real and can be interpreted as the physical excitations. In doped cuprates, the normal-state above the SC transition temperature exhibits a number of anomalous properties which is due to CSS\textsuperscript{20}; while the SC state is characterized by the charge-spin recombination\textsuperscript{26}. Based on the CSS fermion-spin theory, we\textsuperscript{28} have discussed the charge dynamics of the underdoped cuprates in the normal-state, and show that under temperature $T^*$, the magnetic fluctuation is strong enough to lead to a pseudogap. This pseudogap would reduce the charge carrier scattering and thus is responsible for the temperature linear to the nonlinear range in the in-plane resistivity and the crossovers to theinsulating-like range in the c-axis resistivity. Furthermore, the temperature $T^*$ is doping dependent, and grows monotonously as the doping concentration decreases, and disappear in higher doping\textsuperscript{28}. It has been shown\textsuperscript{28} that this pseudogap (then the temperature $T^*$) is obtained from the charge carrier Green’s function in the normal state by considering the second-order correction due to the spin pair bubble. In the kinetic energy driven SC mechanism\textsuperscript{17,18}, the charge carrier pairing state (then the electron SC-state and SC transition temperature $T_c$) occurs directly through the kinetic energy by exchanging spin excitations, and is controlled by both charge carrier gap function and single particle coherent weight. This single particle coherent weight $Z_F(T_c)$ resembles that of the superfluid density. In other words, $T^*$ is closely related to the spin fluctuation, while $T_c$ is self-consistently governed by the single particle coherence and dressed charge carrier pair gap function, this is why there are some differences between $T^*$ and $T_c$.

In summary, within the framework of the kinetic energy driven the SC mechanism\textsuperscript{17}, we have discussed the charge asymmetry in superconductivity of hole- and electron-doped cuprates based on the $t$-$t'$-$J$ model. Our results show that for the hole-doped case, superconductivity appears over a wide range of doping, where the maximal SC transition temperature occurs around the optimal doping concentration, and then decreases in both underdoped and overdoped regimes. In analogy to the phase diagram of hole-doped case, superconductivity appears over a narrow range of doping in the electron-doped side, where the SC transition temperature increases sharply with increasing doping in the underdoped regime, and reaches a maximum in the optimal doping, then decreases sharply with increasing doping in the overdoped regime. However, the maximum achievable SC transition temperature in the optimal doping in the electron-doped case is much lower than that of the hole-doped side due to the electron-hole asymmetry. Our these results are in qualitative agreement with the experimental observations.

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