Asymmetry of the electron spectrum in hole-doped and electron-doped cuprates

Huaiming Guo and Shiping Feng

Department of Physics, Beijing Normal University, Beijing 100875, China

Within the $t'$-$J$ model, the asymmetry of the electron spectrum and quasiparticle dispersion in hole-doped and electron-doped cuprates is discussed. It is shown that the quasiparticle dispersions of both hole-doped and electron-doped cuprates exhibit the flat band around the $(\pi, 0)$ point below the Fermi energy. The lowest energy states are located at the $(\pi/2, \pi/2)$ point for the hole doping, while they appear at the $(\pi, 0)$ point in the electron-doped case due to the electron-hole asymmetry. Our results also show that the unusual behavior of the electron spectrum and quasiparticle dispersion is intrinsically related to the strong coupling between the electron quasiparticles and collective magnetic excitations.

The parent compounds of cuprate superconductors are believed to belong to a class of materials known as Mott insulators with an antiferromagnetic (AF) long-range order (AFLRO), then superconductivity emerges when charge carriers, holes or electrons, are doped into these Mott insulators. Although both hole-doped and electron-doped cuprates have the layered structure of the square lattice of the CuO$_2$ plane separated by insulating layers, the significant difference of the electronic states between hole-doped and electron-doped cuprates is observed, which reflects the electron-hole asymmetry. For the hole-doped cuprates, AFLRO is reduced dramatically with doping, and vanished around the doping $\delta \sim 0.05$. But a series of inelastic neutron scattering measurements show that the incommensurate short AF correlation persists in the underdoped, optimally doped, and overdoped regimes, then in low temperatures, the systems become superconducting (SC) over a wide range of the hole doping concentration $\delta$, around the optimal doping $\delta \sim 0.15$. However, AFLRO survives until superconductivity appears over a narrow range of $\delta$, around the optimal doping $\delta \sim 0.15$ in the electron-doped cuprates. In particular, the optimal transition temperature in the electron-doped cuprates is much lower than that in the hole-doped case, and the commensurate spin response in the SC-state is observed. These experimental observations show that the unconventional physical properties of both hole-doped and electron-doped cuprates mainly depend on the extent of the doping concentration. Since many of the unconventional physical properties, including the relatively high SC transition temperature, have often been attributed to particular characteristics of low energy excitations determined by the electronic structure, then a central issue to clarify the nature of the unconventional physical properties is how the electronic structure evolves with the doping concentration.

From the angle-resolved photoemission spectroscopy (ARPES) measurements, it has been shown that the electron spectral function $A(k, \omega)$ in doped cuprates is strongly momentum and doping dependent. For the hole doping, the charge carriers doped into the parent Mott insulators first enter into the $k = [\pi/2, \pi/2]$ (in units of inverse lattice constant) point in the Brillouin zone, while the charge carriers are accommodated at the $k = [\pi, 0]$ point in the electron-doped case. Moreover, $A(k, \omega)$ has a flat band form as a function of energy $\omega$ for $k$ in the vicinity of the $[\pi, 0]$ point, which leads to the unusual quasiparticle dispersion around the $[\pi, 0]$ point with anomalously small changes of electron energy as a function of momentum. In particular, this flat band is just below the Fermi energy for the hole-doped cuprates, while it is located well below the Fermi energy for the electron-doped case. The flat band reflects the underlying electronic structure near a band saddle point, and is manifestation of a strong coupling between the electron quasiparticles and collective magnetic excitations. Since the normal-state pseudogap starts growing first in the single-particle excitations around the $[\pi, 0]$ point, and then exists in a wide range of the doping concentration, therefore the broad feature in the electron spectrum around the $[\pi, 0]$ point has a particular importance in the mechanism of the normal-state pseudogap formation, and is responsible for the unconventional normal-state properties. Recently, a new low photon energy regime of ARPES has been accessed with lasers and used to study cuprate superconductors, where the clearest evidence for the existence of the electron quasiparticles in the normal-state has been observed. Therefore these ARPES experiments have produced some interesting data that introduce important constraints on the model and theory of cuprate superconductors.

The doping evolution of the normal-state electron spectrum in the hole-doped cuprates has been extensively studied within some strongly correlated models. The most striking aspect is that the unusual quasiparticle dispersion can not be explained by either of the band theory scenarios. The numerical calculation of the electron spectrum based on the large U Hubbard model for the hole-doped case shows the quasiparticle dispersion similar to those observed in experiments. In this calculation, the flat band around the $[\pi, 0]$ point arises from the large Coulomb interaction U. The unusual quasiparticle dispersion in the hole-doped case has been studied numerically within the $t$-$J$ type model, and the result of the quasiparticle dispersion along the $[0, 0]$ to...
[\pi, 0] \text{ direction is quantitative agreement with the experiments. This study also shows unambiguously that the energy scale of the quasiparticle band is controlled by the magnetic interaction } J. \text{ In particular, the electron-hole asymmetry in hole-doped and electron-doped cuprates has been discussed based on numerically exact diagonalization methods}^{21}, \text{ it is shown that the electron-hole asymmetry comes from the coupling of the charge carriers with the spin background. Moreover, many authors}^{22,23} \text{ suggest that the unusual electron spectrum in doped cuprates is a natural consequence of the charge-spin separation (CSS). Although the exact origin of the striking behavior of the electron spectrum still is controversial, a strongly correlated many-body like approach may be appropriate to describe the electronic structure of doped cuprates. Recently, we}^{24} \text{ have developed a CSS fermion-spin theory for description of the unconventional physical properties in doped cuprates, where the electron operator is decoupled as a gauge invariant dressed charge carrier and spin. Within this framework, we have shown that the charge transport is mainly governed by the scattering from the dressed charge carriers due to the spin fluctuation, and the scattering from the spins due to the dressed charge carrier fluctuation dominates the spin response}^{24}. \text{ In particular, the charge-spin recombination of the dressed charge carrier and spin automatically gives the electron quasiparticle character}^{25}. \text{ In this paper, we study the asymmetry of the electron spectrum in hole-doped and electron-doped cuprates along with this line. Our results show that the quasiparticle dispersions of both hole-doped and electron-doped cuprates exhibit the flat band around the } [\pi, 0] \text{ point below the Fermi energy. The lowest energy states are located at the } [\pi/2, \pi/2] \text{ point for the hole doping, while they appear at the } [\pi, 0] \text{ point in the electron-doped case due to the electron-hole asymmetry. Our results also show that the striking behavior of the electron spectrum in doped cuprates is intrinsically related to the strong coupling between the electron quasiparticles and collective magnetic excitations.}

\text{In both hole-doped and electron-doped cuprates, the characteristic feature is the presence of the two-dimensional CuO}_2 \text{ plane}^{2,3} \text{ as mentioned above, and it seems evident that the unusual behaviors are dominated by this CuO}_2 \text{ plane. Although the } t-J \text{ model captures the essential physics of the doped CuO}_2 \text{ plane}^{26}, \text{ the electron-hole asymmetry may be properly accounted by generalizing the } t-J \text{ model to include the second-nearest neighbors hopping terms } t''^{27}. \text{ In this case, we start from the } t-t'J \text{ model on a square lattice,}

\begin{align}
H = -t \sum_{i, j} C^\dagger_{i \alpha} C^\dagger_{i + \hat{x}, \alpha} + t' \sum_{i, j} C^\dagger_{i \alpha} C^\dagger_{i + \hat{y}, \alpha} \\
+ \mu_0 \sum_{i, \sigma} C^\dagger_{i \sigma} C_{i \sigma} + J \sum_{i} \mathbf{S}_i \cdot \mathbf{S}_{i + \hat{\eta}},
\end{align}

with } \hat{\eta} = \pm \hat{x}, \pm \hat{y}, \hat{r} = \pm \hat{x} \pm \hat{y}, \text{ and } C^\dagger_{i \sigma} (C_{i \sigma}) \text{ is the electron creation (annihilation) operator, } \mathbf{S}_i = C^\dagger_{i} \mathbf{\tau} C_{i}/2 \text{ is spin operator with } \mathbf{\tau} = (\sigma_x, \sigma_y, \sigma_z) \text{ as Pauli matrices, and } \mu_0 \text{ is the chemical potential. For the electron-doped case, we can perform a particle-hole transformation } C_{i \sigma} \rightarrow C^\dagger_{i \sigma}, \text{ so that the difference between hole-doped and electron-doped cases is expressed as the sign difference of the hopping parameters, i.e., } t > 0 \text{ and } t' > 0 \text{ for the hole doping and } t < 0 \text{ and } t' < 0 \text{ for the electron doping}^{28}. \text{ In this case, the } t-t'J \text{ model (1) in both hole-doped and electron-doped cases is always subject to an important on-site local constraint to avoid the double occupancy, i.e., } \sum_{\sigma} C^\dagger_{i \sigma} C_{i \sigma} \leq 1. \text{ Therefore the strong electron correlation in the } t-t'J \text{ model (1) manifests itself by this single occupancy local constraint}^{26}. \text{ It has been shown that this local constraint can be treated properly in analytical calculations within the CSS fermion-spin theory}^{24}, \text{ with}

\begin{align}
C_{i \uparrow} = h^\dagger_{i \uparrow} S^\dagger_{i \uparrow} \text{ and } C_{i \downarrow} = h^\dagger_{i \downarrow} S^\dagger_{i \downarrow}, \text{ where the spinful fermion operator } h_{i \sigma} = e^{-i\Phi_{i \sigma}} h_i \text{ 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 \text{ describes the spin degree of freedom, then the local constraint for the single occupancy, } \sum_{\sigma} C^\dagger_{i \sigma} C_{i \sigma} = S^\dagger_{i \uparrow} h^\dagger_{i \uparrow} S^\dagger_{i \uparrow} + S^\dagger_{i \downarrow} h^\dagger_{i \downarrow} S^\dagger_{i \downarrow} = h^\dagger_{i \uparrow} (S^\dagger_{i \uparrow} S^\dagger_{i \downarrow} + S^\dagger_{i \downarrow} S^\dagger_{i \uparrow}) = 1, \text{ is satisfied in analytical calculations. These dressed charge carrier and spin are gauge invariant}^{24}, \text{ and in this sense, they are real and can be interpreted as the physical excitations}^{29}. \text{ Although in common sense } h_{i \sigma} \text{ is not a real spinful fermion, it behaves like a spinful fermion. In this CSS fermion-spin representation, the low-energy behavior of the } t-t'J \text{ model (1) can be expressed as}^{24},

\begin{align}
H &= -t \sum_{i \hat{\eta}} (h_{i \uparrow} S^\dagger_{i \uparrow} h^\dagger_{i \uparrow} S^\dagger_{i \uparrow} + h_{i \downarrow} S^\dagger_{i \downarrow} h^\dagger_{i \downarrow} S^\dagger_{i \downarrow}) \\
+ t' \sum_{i \hat{\tau}} (h_{i \uparrow} S^\dagger_{i \uparrow} h^\dagger_{i \uparrow} S^\dagger_{i \uparrow} + h_{i \downarrow} S^\dagger_{i \downarrow} h^\dagger_{i \downarrow} S^\dagger_{i \downarrow}) \\
- \mu_0 \sum_{i} h^\dagger_{i \sigma} h_{i \sigma} + J_{\text{eff}} \sum_{i \hat{\eta}} \mathbf{S}_i \cdot \mathbf{S}_{i + \hat{\eta}},
\end{align}

with } J_{\text{eff}} = (1 - \delta)^2 J \text{ and } \delta = \langle h^\dagger_{i \uparrow} h_{i \uparrow} \rangle = \langle h^\dagger_{i \downarrow} h_{i \downarrow} \rangle \text{ is the doping concentration. As a consequence, the magnetic energy } (J) \text{ term in the } t-t'J \text{ model is only to form an adequate spin configuration}^{29}, \text{ while the kinetic energy part has been expressed as the interaction between the dressed charge carriers and spins, which reflects that even the kinetic energy part in the } t-t'J \text{ Hamiltonian has the strong Coulombic contribution due to the restriction of no doubly occupancy of a given site, and therefore dominates the essential physics in doped cuprates.}

\text{For the discussions of the electron spectrum, we need to calculate the electron Green’s function } G(i - j, t - t') = \langle \langle C_{i \sigma}(t); C^\dagger_{j \sigma}(t') \rangle \rangle, \text{ which is a convolution}^{29} \text{ of the spin Green’s function } D(i - j, t - t') = \langle \langle S^\dagger_{i \sigma}(t); S_{j \sigma}(t') \rangle \rangle \text{ and dressed charge carrier Green’s function } g(i - j, t - t') = \langle \langle h_{i \sigma}(t); h^\dagger_{j \sigma}(t') \rangle \rangle, \text{ and can be formally expressed in terms of the spectral representation as,}
for hole-doped and electron-doped cases, respectively, where the dressed charge carrier spectral function
\[ A_k(q, \omega) = -2 \text{Im} q(q, \omega), \]
the spin spectral function
\[ A_s(k, \omega) = -2 \text{Im} D(k, \omega), \]
and \( n_F(\omega) \) and \( n_F(\omega) \) are the boson and fermion distribution functions, respectively. This convolution reflects the charge-spin recombination.

In the calculation of the electron Green’s function (3b) for the electron-case, the particle-hole transformation \( C_{i\sigma} \rightarrow C_{i\sigma}^\dagger \) as mentioned above has been considered. Since the quantum spin operators obey the Pauli spin algebra, i.e., the spin one-half raising and lowering operators \( S^+ \) and \( S^- \) behave as fermions on the same site and as bosons on different sites, then this problem can be discussed in terms of the equation of motion method. It has been shown that in the mean-field (MF) level, the spin system is an anisotropic way from the half-filling. In this case, we need to define another spin Green’s function \( D_z(i - j, t - t') = \langle (S^i(t); S^j(t')) \rangle \), and then both spin Green’s functions \( D_p(p, \omega) \) and \( D_z(p, \omega) \) describe the spin propagations. In the doped regime without AFLRO, i.e., \( S^z_i \) = 0, a MF theory of the \( t-J \) model has been developed. Following their discussions, the MF dressed charge carrier and spin Green’s functions of the \( t-t' \)-J model have been obtained as:

\[
g^{(0)}(k, \omega) = \frac{1}{\omega - \xi_k}, \quad D^0(p, \omega) = \frac{B_p}{2\omega_p} \left( \frac{1}{\omega - \omega_p} - \frac{1}{\omega + \omega_p} \right), \quad D^z(0, \omega) = \frac{B_p}{2\omega_p} \left( \frac{1}{\omega - \omega_p} - \frac{1}{\omega + \omega_p} \right),
\]

where \( B_p = 2\lambda_1(\gamma_p - A_2 - A_2\beta(2\gamma_p' - \chi_2)), \)
\( B_{2p} = \epsilon_\chi(\gamma_p - A_2 - A_2\beta(2\gamma_p' - \chi_2)), \)
\( \lambda_1 = 2ZJ_{J_{\text{eff}}}, \quad \lambda_2 = 4Zg_\epsilon, \)
\( A_1 = \gamma_1 + \chi_2/2, \quad A_2 = \gamma_2 + \gamma_2/2, \quad \epsilon = 1 + 2t\phi_1/J_{\text{eff}} \),
the spin correlation functions \( \chi_1 = \langle S^i_z S^j_{-z} \rangle, \chi_2 = \langle S^i_z S^j_{-z} \rangle, \chi_3 = \langle S^i_z S^j_{-z} \rangle \),
the dressed charge carrier’s particle-hole parameters \( \phi_1 = \langle h^i(h_{i-\sigma}) \rangle \) and \( \phi_2 = \langle h^i(h_{i-\sigma}) \rangle \),
the number of the nearest or second-nearest neighbor sites, while the MF dressed charge carrier and spin excitation spectra are given by:

\[
\xi_k = \epsilon_k - \mu_0.
\]
\( g(\mathbf{k}, i\omega_n) = g^{(0)}(\mathbf{k}, i\omega_n) + g^{(0)}(\mathbf{k}, i\omega_n) \Sigma^{(h)}(\mathbf{k}, i\omega_n) g(\mathbf{k}, i\omega_n), \)  

(6)

with the dressed charge carrier self-energy is evaluated from the spin pair bubble as,

\[
\Sigma^{(h)}(\mathbf{k}, i\omega_n) = \frac{1}{N \beta} \sum_{\mathbf{p}, \mathbf{p}'} \Lambda(\mathbf{k}, \mathbf{p}, \mathbf{p}') \frac{1}{\beta} \sum_{\mathbf{p}'_m} g(\mathbf{p} + \mathbf{k}, i\mathbf{p}'_m + i\omega_n) 
\times \frac{1}{\beta} \sum_{\mathbf{p}'_m} D^{(0)}(\mathbf{p}', i\mathbf{p}'_m) 
\times D^{(0)}(\mathbf{p}' + \mathbf{k}, i\mathbf{p}'_m + i\omega_n), 
\]

(7)

where \( \Lambda(\mathbf{k}, \mathbf{p}, \mathbf{p}') = [Zt'_{\mathbf{p}+\mathbf{p}'+\mathbf{k}} - Z t'_{\mathbf{p}+\mathbf{p}'+\mathbf{k}}]^2 \). This self-energy function \( \Sigma^{(h)}(\mathbf{k}, \omega) \) renormalizes the MF dressed charge carrier spectrum, and therefore it describes the quasiparticle coherence. In particular, \( \Sigma^{(h)}(\mathbf{k}, \omega) \) is not even function. For the convenience of the discussions, \( \Sigma^{(h)}(\mathbf{k}, \omega) \) can be broken up into its symmetric and antisymmetric parts as, \( \Sigma^{(h)}(\mathbf{k}, \omega) = \Sigma^{(s)}(\mathbf{k}, \omega) + \omega \Sigma^{(a)}(\mathbf{k}, \omega) \), therefore both \( \Sigma^{(s)}(\mathbf{k}, \omega) \) and \( \Sigma^{(a)}(\mathbf{k}, \omega) \) are even functions of \( \omega \). Now we define the quasiparticle coherent weight as \( Z_F^{-1}(\mathbf{k}, \omega) = 1 - \Sigma^{(s)}(\mathbf{k}, \omega) \), then the full dressed charge carrier Green’s function in Eq. (6) can be written as,

\( g(\mathbf{k}, \omega) = \frac{Z_F(\mathbf{k}, \omega)}{\omega - Z_F(\mathbf{k}, \omega)[\xi_\mathbf{k} + \Sigma^{(s)}(\mathbf{k}, \omega)]}. \)

(8)

Since we only discuss the low-energy behavior of doped cuprates, then the quasiparticle coherent weight can be discussed in the static limit, i.e., \( Z_F^{-1}(\mathbf{k}) = 1 - \Sigma^{(s)}(\mathbf{k}, \omega) \mid_{\omega=0} \) and \( \Sigma^{(s)}(\mathbf{k}, \omega) = \Sigma^{(s)}(\mathbf{k}, \omega) \mid_{\omega=0} \). Although \( Z_F(\mathbf{k}) \) and \( \Sigma^{(s)}(\mathbf{k}) \) still are a function of \( \mathbf{k} \), the wave vector dependence is unimportant. It has been shown from ARPES experiments that in the normal-state, the lowest energy states are located at the \( [\pi/2, \pi/2] \) point for the hole-doped cuprates, and the \( [\pi, 0] \) point in the electron-doped case, which indicates that the majority contribution for the electron spectrum comes from the \( [\pi/2, \pi/2] \) point for the hole doping, and the \( [\pi, 0] \) point for the electron doping. In this case, the wave vector \( \mathbf{k} \) in \( Z_F(\mathbf{k}) \) and \( \Sigma^{(s)}(\mathbf{k}) \) can be chosen as \( Z_F^{-1} = 1 - \Sigma^{(s)}(\mathbf{k}) \mid_{\mathbf{k}=\mathbf{k}_h=[\pi/2, \pi/2]} \) and \( \Sigma^{(s)} = \Sigma^{(s)}(\mathbf{k}) \mid_{\mathbf{k}=\mathbf{k}_h=[\pi/2, \pi/2]} \) for the hole doping, and \( Z_F^{-1} = 1 - \Sigma^{(s)}(\mathbf{k}) \mid_{\mathbf{k}=\mathbf{k}_e=[\pi, 0]} \) and \( \Sigma^{(s)} = \Sigma^{(s)}(\mathbf{k}) \mid_{\mathbf{k}=\mathbf{k}_e=[\pi, 0]} \) for the electron doping, then the dressed charge carrier Green’s function in Eq. (8) can be expressed explicitly as,

\( g(\mathbf{k}, \omega) = \frac{Z_F(\mathbf{k}, \omega)}{\omega - \xi_\mathbf{k}}. \)

(9)

where the renormalized dressed charge carrier quasiparticle spectrum \( \tilde{\xi}_\mathbf{k} = \xi_\mathbf{k} - \mu \), with \( \xi_\mathbf{k} = Z_F\xi_\mathbf{k} \) and renormalized chemical potential \( \mu = Z_F(\mu_0 - \Sigma^{(s)}(\mathbf{k})) \). As we will see later, this \( Z_F \) reduces the dressed holon (then electron quasiparticle) bandwidth, and then the energy scale of the electron quasiparticle band is controlled by the magnetic interaction \( J \), while \( \Sigma^{(s)}(\mathbf{k}) \) renormalizes the chemical potential, and therefore plays an important role in qualitatively determining the positions of peaks from the doping dependence of the electron spectrum. In this case, the quasiparticle coherent weight \( Z_F \) satisfies the following equation,

\[
Z_F^{-1} = 1 + \frac{1}{N^2} \sum_{\mathbf{p}, \mathbf{p}'} \Lambda(\mathbf{k}, \mathbf{p}, \mathbf{p}') Z_F B_{\mathbf{p}} B_{\mathbf{p}'} \frac{1}{4\omega_{\mathbf{p}} \omega_{\mathbf{p}'+\mathbf{k}}} \times \left( F_1(\mathbf{k}, \mathbf{p}, \mathbf{p}') \frac{(\omega_{\mathbf{p}+\mathbf{p}'} - \omega_{\mathbf{p}' - \xi_\mathbf{p} + \mathbf{k}})^2}{(\omega_{\mathbf{p}'} - \omega_{\mathbf{p}'} + \xi_\mathbf{p} + \mathbf{k})^2} + \frac{F_2(\mathbf{k}, \mathbf{p}, \mathbf{p}')}{(\omega_{\mathbf{p}'} + \omega_{\mathbf{p}'} - \xi_\mathbf{p} + \mathbf{k})^2} + \frac{F_3(\mathbf{k}, \mathbf{p}, \mathbf{p}')}{(\omega_{\mathbf{p}'} + \omega_{\mathbf{p}'} + \xi_\mathbf{p} + \mathbf{k})^2} \right),
\]

(10)
\[ C_i^r = \frac{1}{N} \sum_k \gamma_k \frac{B_{k}^{2}}{2\omega_{k}} \coth \left( \frac{1}{2} \beta \omega_{k} \right), \quad (11) \]

\[ C_j^r = \frac{1}{N} \sum_k \gamma_k \frac{B_{k}^{2}}{2\omega_{k}} \coth \left( \frac{1}{2} \beta \omega_{k} \right), \quad (11m) \]

then all order parameters, decoupling parameter \( \alpha \), and chemical potential \( \mu \) are determined by the self-consistent calculation. In this sense, our above calculations are exact without using adjustable parameters, in other words, they are controllable.

With the help of the dressed charge carrier Green's function \( g(\mathbf{k}, \omega) \) in Eq. (9) and MF spin Green's function \( D^{(0)}(\mathbf{p}, \omega) \) in Eq. (4b), the electron Green's function in Eq. (3) can be evaluated explicitly as,

\[ G_{p\text{-type}}(\mathbf{k}, \omega) = \frac{1}{N} \sum_{\mathbf{p}} Z_F \frac{B_{\mathbf{p} + \mathbf{k}}}{2\omega_{\mathbf{p} + \mathbf{k}}} \left( \frac{L_1(\mathbf{k}, \mathbf{p})}{\omega + \xi_{\mathbf{p}} - \omega_{\mathbf{p} + \mathbf{k}}} + \frac{L_2(\mathbf{k}, \mathbf{p})}{\omega + \xi_{\mathbf{p}} + \omega_{\mathbf{p} + \mathbf{k}}} \right), \quad (12a) \]

\[ G_{n\text{-type}}(\mathbf{k}, \omega) = \frac{1}{N} \sum_{\mathbf{p}} Z_F \frac{B_{\mathbf{p} + \mathbf{k}}}{2\omega_{\mathbf{p} + \mathbf{k}}} \left( \frac{L_1(\mathbf{k}, \mathbf{p})}{\omega - \xi_{\mathbf{p}} + \omega_{\mathbf{p} + \mathbf{k}}} + \frac{L_2(\mathbf{k}, \mathbf{p})}{\omega - \xi_{\mathbf{p}} - \omega_{\mathbf{p} + \mathbf{k}}} \right), \quad (12b) \]

for hole-doped and electron-doped cases, respectively, where \( L_1(\mathbf{k}, \mathbf{p}) = n_F(\xi_{\mathbf{p}}) + n_B(\omega_{\mathbf{p} + \mathbf{k}}) \) and \( L_2(\mathbf{k}, \mathbf{p}) = 1 - n_F(\xi_{\mathbf{p}}) + n_B(\omega_{\mathbf{p} + \mathbf{k}}) \), then the electron spectral function \( A(\mathbf{k}, \omega) = -2\text{Im}G(\mathbf{k}, \omega) \) is obtained from the above corresponding electron Green's function as,

\[ A_{p\text{-type}}(\mathbf{k}, \omega) = 2\pi \frac{1}{N} \sum_{\mathbf{p}} Z_F \frac{B_{\mathbf{p} + \mathbf{k}}}{2\omega_{\mathbf{p} + \mathbf{k}}} \]
\[ \times \left[ L_1(\mathbf{k}, \mathbf{p}) \delta(\omega + \xi_{\mathbf{p}} - \omega_{\mathbf{p} + \mathbf{k}}) \right. \]
\[ + \left. L_2(\mathbf{k}, \mathbf{p}) \delta(\omega + \xi_{\mathbf{p}} + \omega_{\mathbf{p} + \mathbf{k}}) \right], \quad (13a) \]

\[ A_{n\text{-type}}(\mathbf{k}, \omega) = 2\pi \frac{1}{N} \sum_{\mathbf{p}} Z_F \frac{B_{\mathbf{p} + \mathbf{k}}}{2\omega_{\mathbf{p} + \mathbf{k}}} \]
\[ \times \left[ L_2(\mathbf{k}, \mathbf{p}) \delta(\omega - \xi_{\mathbf{p}} + \omega_{\mathbf{p} + \mathbf{k}}) \right. \]
\[ + \left. L_1(\mathbf{k}, \mathbf{p}) \delta(\omega - \xi_{\mathbf{p}} - \omega_{\mathbf{p} + \mathbf{k}}) \right]. \quad (13b) \]

We are now ready to discuss the electron spectrum and quasiparticle dispersion in doped cuprates. Since the absolute values of \( t \) and \( t' \) are almost same for both hole-doped and electron-doped cuprates\(^{27}\), therefore the commonly used parameters in this paper are chosen as \( t/J = 2.5 \) and \( t'/J = 0.375 \) for the hole doping, and \( t/J = -2.5 \) and \( t'/J = -0.375 \) for the electron doping.

We have performed the calculation for the electron spectral function in Eq. (13), and the results at the \([\pi, 0]\) point for (a) the hole doping and (b) electron doping, and at the \([\pi/2, \pi/2]\) point for (c) the hole doping and (d) electron doping with temperature \( T = 0.1J \) in the doping concentration \( \delta = 0.09 \) (solid line), \( \delta = 0.12 \) (dashed line), and \( \delta = 0.15 \) (dotted line) are plotted in Fig. 1. For the hole doping, although both positions of the quasiparticle peaks at the \([\pi, 0]\) and \([\pi/2, \pi/2]\) points are below the Fermi energy, the position of the quasiparticle peak at the \([\pi/2, \pi/2]\) point is more close to the Fermi energy, which indicates that the lowest energy states are located at the \([\pi/2, \pi/2]\) point. In other words, the low energy spectral weight with the majority contribution to the low-energy properties of the hole-doped cuprates comes from the \([\pi/2, \pi/2]\) point. However, only the position of the quasiparticle peak at the \([\pi, 0]\) point for the electron doping is below the Fermi energy, and in contrast to the hole-doped case, the position of the quasiparticle peak at the \([\pi/2, \pi/2]\) point is above the Fermi energy, which mean that the lowest energy states appear at the \([\pi, 0]\) point, i.e., only states around the \([\pi, 0]\) point have the majority contribution to the low-energy properties of the electron-doped cuprates. These behaviors reflect the electron-hole asymmetry in hole-doped and electron-doped cuprates. Moreover, the electron spectrum is doping dependence. The quasiparticle peaks at the \([\pi, 0]\) and \([\pi/2, \pi/2]\) points for the hole doping and at the \([\pi, 0]\) point for the electron doping become sharper, while the spectral weight of these peaks increases in intensity with increasing doping. Furthermore, we have also discussed the temperature dependence of the electron spectrum, and the results show that the spectral weight is suppressed with increasing temperatures. Our these results are qualitatively consistent with the ARPES experimental data\(^{4,11-15,17}\).

For a better understanding of the anomalous form

![Fig. 1. The electron spectral function $A(k, \omega)$ at $\pi, 0$ point for (a) the hole doping and (b) electron doping, and at $\pi/2, \pi/2$ point for (c) the hole doping and (d) electron doping with $T = 0.1J$ in the doping concentration $\delta = 0.09$ (solid line), $\delta = 0.12$ (dashed line), and $\delta = 0.15$ (dotted line), where the commonly used parameters are chosen as $t/J = 2.5$ and $t'/J = 0.375$ for the hole doping, and $t/J = -2.5$ and $t'/J = -0.375$ for the electron doping.](image-url)
of the electron spectrum $A(k,\omega)$ as a function of energy $\omega$ for $k$ in the vicinity of the $[\pi,0]$ point, we have made a series of calculations for $A(k,\omega)$, and the results for (a) the hole doping and (b) electron doping with $T = 0.1J$ in $\delta = 0.15$ at $[0.9\pi,0]$ (solid line), $[0.95\pi,0]$ (long dashed line), $[\pi,0]$ (short dashed line), $[\pi,0.05\pi]$ (dash-dotted line), and $[\pi,0.1\pi]$ (dotted line) points are plotted in Fig. 2. Obviously, the positions of these peaks of the electron spectral function $A(k,\omega)$ around the $[\pi,0]$ point are almost not changeable, which leads to the unusual quasiparticle dispersion around the $[\pi,0]$ point. Furthermore, the lowest energy peaks are well defined at all momenta. To show the broad feature in the electron spectrum around the $[\pi,0]$ point clearly, we plot the positions of the lowest energy quasiparticle peaks in $A(k,\omega)$ as a function of momentum along the high symmetry directions for (a) the hole doping and (b) electron doping with $T = 0.1J$ at $\delta = 0.15$ in Fig. 3. For comparison, the corresponding results of the bare electron dispersion of the $t$-$t'$ model (dotted line), and experimental results (inset) of the electron dispersion from the hole-doped cuprate$^{12}$ Bi$_2$Sr$_2$CaCu$_2$O$_8+\delta$ and electron-doped cuprate$^{11}$ Nd$_{2-x}$Ce$_x$CuO$_{4+\delta}$ are also shown in Fig. 3. Our these results show that in accordance with the anomalous property of the electron spectrum in Fig. 2, the electron quasiparticles around the $[\pi,0]$ point disperse very weakly with momentum, and then the unusual flat band appears, while the Fermi energy is only slightly above this flat band, in qualitative agreement with these obtained from ARPES experimental measurements on doped cuprates$^{4,11-15}$.

Since the full electron Green’s function (then the electron spectral function) is obtained beyond the MF approximation by considering the fluctuation due to the

![Figure 2](image1.png)

**FIG. 2.** The electron spectral function $A(k,\omega)$ for (a) the hole doping and (b) electron doping with $T = 0.1J$ in $\delta = 0.15$ at $[0.9\pi,0]$ (solid line), $[0.95\pi,0]$ (long dashed line), $[\pi,0]$ (short dashed line), $[\pi,0.05\pi]$ (dash-dotted line), and $[\pi,0.1\pi]$ (dotted line) points, where the commonly used parameters are chosen as $t/J = 2.5$ and $t'/J = 0.375$ for the hole doping, and $t/J = -2.5$ and $t'/J = -0.375$ for the electron doping.

![Figure 3](image2.png)

**FIG. 3.** The position of the lowest energy quasiparticle peaks in $A(k,\omega)$ as a function of momentum for (a) the hole doping with $t/J = 2.5$ and $t'/J = 0.375$, and (b) electron doping with $t/J = -2.5$ and $t'/J = -0.375$, in $T = 0.1J$ at $\delta = 0.15$. The dotted line are corresponding results of the bare electron dispersion of the $t$-$t'$ model. Inset: the corresponding experimental results of the hole-doped cuprate Bi$_2$Sr$_2$CaCu$_2$O$_8+\delta$ and electron-doped cuprate Nd$_{2-x}$Ce$_x$CuO$_{4+\delta}$ taken from Refs. [12] and [11], respectively.
spin pair bubble, therefore the nature of the electron spectrum is closely related to the strong coupling between the dressed charge carriers (then electron quasiparticles) and collective magnetic excitations. This can be understood from a comparison between the bare electron dispersion of the $t$-$t'$ model and renormalized electron quasiparticle dispersion of the $t$-$t'$-$J$ model in Fig. 3. Our results show that the single-particle hopping in the $t$-$t'$-$J$ model is strongly renormalized by the magnetic interaction. As a consequence, the quasiparticle bandwidth is reduced to the order of (a few) $J$, and therefore the energy scale of the quasiparticle band is controlled by the magnetic interaction. This renormalization due to the strong interaction is then responsible for the unusual electron quasiparticle spectrum and production of the flat band. On the other hand, although $t'$ does not change spin configuration, the interplay of $t'$ with $t$ and $J$ causes a further weakening of the AF spin correlation for the hole doping, and enhancing the AF spin correlation for the electron doping, which shows that the AF spin correlations in the electron doping is stronger than these in the hole-doped side, and leads to the asymmetry of the electron spectrum in hole-doped and electron-doped cuprates. This is why $t'$ term plays an important role in explaining the difference between electron and hole doping. Moreover, our present results also show that the electron quasiparticle excitations originating from the dressed holons and spins are due to the charge-spin recombination, the composite nature of the electron quasiparticle excitations, and then the unconventional normal-state properties in doped cuprates are attributed to the presence of the dressed charge carriers, spin, and electron quasiparticle excitations.

Finally, we have noted that an obvious weakness of the present results is that the flat band for the electron doping is not well below the Fermi energy. In the above calculation, the spin part has been limited to the MF level, and therefore the spin fluctuations beyond the MF level is not considered. Since the AF spin correlations in the electron doping is stronger than these in the hole doping as mentioned above, it is then possible that the weakness perhaps due to neglecting the spin fluctuations in the present case may be cured by considering them, and these and other related issues are under investigation now.

In summary, we have studied the asymmetry of the electron spectrum and quasiparticle dispersion in hole-doped and electron-doped cuprates based on the $t$-$t'$-$J$ model. Our results show that the quasiparticle dispersions of both hole-doped and electron-doped cuprates exhibit the flat band around the $[\pi, 0]$ point below the Fermi energy. The lowest energy states are located at the $[\pi/2, \pi/2]$ point for the hole doping, while they appear at the $[\pi, 0]$ point for the electron doping due to the electron-hole asymmetry. Our results also show that the unusual behavior of the electron spectrum and quasiparticle dispersion is intriguingly related to the strong coupling between the electron quasiparticles and collective magnetic excitations. Within the CSS fermion-spin theory, we have developed a kinetic energy driven SC mechanism, where the dressed charge carriers interact directly through the kinetic energy by exchanging the spin excitations, leading to a net attractive force between the dressed charge carriers, then the electron Cooper pairs originating from the dressed charge carrier pairing state are due to the charge-spin recombination, and their condensation reveals the SC ground-state. Based on this SC mechanism, we have discussed the doping and temperature dependence of the electron spectrum of hole-doped and electron-doped cuprates in the SC-state, and related theoretical results will be presented elsewhere.

ACKNOWLEDGMENTS

The authors would like to thank Dr. T.X. Ma and Dr. Y. Lan for the helpful discussions. This work was supported by the National Natural Science Foundation of China under Grant Nos. 10125415 and 90403005, and the 973 project from the Ministry of Science and Technology of China under Grant No. 2006CB601002.

1 See, e.g., M.A. Kastner, R.J. Birgeneau, G. Shiran, and Y. Endoh, Rev. Mod. Phys. 70, 897 (1998), and references therein.
2 J.G. Bednorz and K.A. Müller, Z. Phys. B 64, 189 (1986).
3 Y. Tokura, H. Takagi, and S. Uchida, Nature 337, 345 (1989).
4 See, e.g., A. Damascelli, Z. Hussain, and Z.-X. Shen, Rev. Mod. Phys. 75, 475 (2003); Z.X. Shen, and D.S. Dessau, Phys. Rep. 70, 253 (1995), and references therein.
5 C.C. Homes, B.P. Clayman, J.L. Peng, and R.L. Greens, Phys. Rev. B 56, 5525 (1997).
6 K. Yamada, C.H. Lee, K. Kurahashi, J. Wada, S. Waki moto, S. Ueki, H. Kimura, Y. Endoh, S. Hosoya, and G. Shirane, Phys. Rev. B 57, 6165 (1998).
7 See, e.g., J.L. Tallon, J.W. Loram, J.R. Cooper, C. Panagopoulou, and C. Bernhard, Phys. Rev. B 68, 180501 (2003).
8 H. Takagi, S. Uchida, and Y. Tokura, Phys. Rev. Lett. 62, 1197 (1989).
9 J.L. Peng, E. Maiser, T. Venkatesan, R.L. Greene, and G. Czjzek, Phys. Rev. B 55, R6145 (1997).
10 K. Yamada, K. Kurahashi, T. Uefuji, S. Park, S.H. Lee, and Y. Endoh, Phys. Rev. Lett. 90, 137004 (2002).
11 Z.X. Shen, W.E. Spicer, D.M. King, D.S. Dessau, and B.O. Wells, Science 267, 343 (1995); C. Kim, P.J. White, Z.X. Shen, T. Tohyama, Y. Shibata, S. Maekawa, B.O. Wells, Y.J. Kim, R.J. Birgeneau, and M.A. Kastner, Phys. Rev. Lett. 80, 4245 (1998).
12 D.S. Dessau, Z.X. Shen, D.M. King, D.S. Marshall, L.W. Lombardo, P.H. Dickinson, A.G. Loeser, J. DiCarlo, C.H.
Park, A. Kapitulnik, and W.E. Spicer, Phys. Rev. Lett. 71, 2781 (1993).

13 B.O. Wells, Z.X. Shen, A. Matsuura, D.M. King, M.A. Kastner, M. Greven, and R.J. Birgeneau, Phys. Rev. Lett. 74, 964 (1995); D.S. Marshall, D.S. Dessau, A.G. Loeser, C.H. Park, A. Matsuura, J.N. Eckstein, I. Bozovic, P. Fournier, A. Kapitulnik, W.E. Spicer, and Z.X. Shen, Phys. Rev. Lett. 76, 4841 (1996).

14 N.P. Armitage, F. Ronning, D.H. Lu, C. Kim, A. Damascelli, K.M. Shen, D.L. Feng, H. Eisaki, Z.X. Shen, P.K. Mang, N. Kaneko, M. Greven, Y. Onose, Y. Taguchi, and Y. Tokura, Phys. Rev. Lett. 88, 227001 (2002); N.P. Armitage, D.H. Lu, C. Kim, A. Damascelli, K.M. Shen, F. Ronning, D.L. Feng, P. Bogdanov, Z.X. Shen, Y. Onose, Y. Taguchi, P.K. Mang, N. Kaneko, and M. Greven, Phys. Rev. Lett. 87, 147003 (2001).

15 D.M. Kim, Z.X. Shen, D.S. Dessau, B.O. Wells, W.E. Spicer, A.J. Arko, D.S. Marshall, J. DiCarlo, A.G. Loeser, C.H. Park, E.R. Ratner, J.L. Peng, Z.Y. Li, and R.L. Greene, Phys. Rev. Lett. 70, 3159 (1993); R.O. Anderson, R. Claessen, J.W. Allen, C.G. Olson, C. Janowitz, L.Z. Liu, J.H. Park, M.B. Maple, Y. Dalichaouch, M.C. de Andrade, R.F. Jardim, E.A. Early, S.J. Oh, and W.P. Ellis, Phys. Rev. Lett. 70, 3163 (1993).

16 Z.X. Shen and J.R. Schrieffer, Phys. Rev. Lett. 78, 1771 (1997).

17 J.D. Koralek, J.F. Douglas, N.C. Plumb, Z. Sun, A. Fedorov, M. Murmane, H. Kapteyn, S. Cundiff, Y. Aiura, K. Oka, H. Eisaki, and D.S. Dessau, cond-mat/0508404.

18 See, e.g., E. Dagotto, Rev. Mod. Phys. 66, 763 (1994), and references therein.

19 N. Bulut, D.J. Scalapino, and S.R. White, Phys. Rev. B 50, 7215 (1994).

20 E. Dagotto and A. Nazarenko, Phys. Rev. Lett. 73, 728 (1994); T. Xiang and J.M. Wheatley, Phys. Rev. B 54, R12653 (1996).

21 R.J. Gooding, K.J.E. Vos, and P.W. Leung, Phys. Rev. B 50, 12866 (1994).

22 P.W. Anderson, J. Phys. Chem. Solids 54, 1073 (1993).

23 R.B Laughlin, Phys. Rev. Lett. 79, 1726 (1997); J. Low. Temp. Phys. 99, 443 (1995).

24 Shiping Feng, Jihong Qin, and Tianxing Ma, J. Phys. Condens. Matter 16, 343 (2004); Shiping Feng, Tianxing Ma, and Jihong Qin, Mod. Phys. Lett. B17, 361 (2003); Shiping Feng, Z.B. Su, and L. Yu, Phys. Rev. B 49, 2368 (1994).

25 Shiping Feng and Yun Song, Phys. Rev. B55, 642 (1997); Feng Yuan and Shiping Feng, Phys. Lett. A 271, 429 (2000).

26 P.W. Anderson, Science 235, 1196 (1987).

27 M.S. Hybertson, E. Stechel, M. Schuter, and D. Jennison, Phys. Rev. B 41, 11068 (1990).

28 Tianxing Ma, Huaiming Guo, and Shiping Feng, Phys. Lett. A 337, 61 (2005); Bin Liu, Ying Liang, Shiping Feng and Wei Ye Chen, Phys. Rev. B 69, 224506 (2004).

29 P.W. Anderson, Phys. Rev. Lett. 67, 2092 (1991); Science 288, 480 (2000); cond-mat/0108522.

30 Shiping Feng and Zhongbing Huang, Phys. Lett. A 232, 293 (1997); Feng Yuan, Jihong Qin, Shiping Feng, and W.Y. Chen, Phys. Rev. B 67, 134505 (2003).

31 D.M. Newns et al., Comments Condens. Matter Phys. 15, 273 (1992), and references therein.

32 See, e.g., G.D. Mahan, Many Particle Physics, (Plenum Press, New York, 1981), Chapter 9.

33 Shiping Feng, Phys. Rev. B68, 184501 (2003); Shiping Feng and Tianxing Ma, Phys. Lett. A 350, 138 (2006); Shiping Feng, Tianxing Ma, and Huaiming Guo, Physica C 436, 14 (2006); Shiping Feng and Tianxing Ma, in New Frontiers in Superconductivity Research, edited by B.P. Martins (Nova Science Publishers, New York, 2006), in press.