Spin-gap effect on resistivity in the t-J model

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We calculate the spin-gap effect on dc resistivity in the t-J model of high-$T_c$ cuprates by using the Ginzburg-Landau theory coupled with a gauge field as its effective field theory to get $\rho(T) \propto T(1 - c(T^* - T)^d)$, where $T^*$ is the spin-gap onset temperature. By taking the compactness of massive gauge field into account, the exponent $d$ deviates from its mean-field value 1/2 and becomes a nonuniversal $T$-dependent quantity, which improves the correspondence with the experiments.

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The idea of charge-spin separation (CSS) by Anderson which accounts for the anomalous behavior of various normal-state properties of high-$T_c$ cuprates allows us to treat holons and spinons introduced in the slave-boson (SB) mean-field theory (MFT) of the t-J model as quasi-free particles. Fluctuations around MFT are described by gauge fields coupled to holons and spinons, the effects of which may be calculated in perturbation theory. Actually, by using a massless gauge field, Nagaosa and Lee obtained the dc resistivity $\rho(T) \propto T$, which agrees with the experiment for $T > T^*$ where $T^*$ is the onset temperature of spin gap.

For $T < T^*$, the experimentally observed $\rho(T)$ reduces from this T-linear behavior. Because the gauge field is expected to acquire a mass $m_A$ in the spin-gap state, this reduction could be understood as a mass effect; the fluctuations of the gauge field become weaker and the scatterings between holons (the carriers of charge) and gauge bosons are reduced. Actually, in the previous paper, we obtained the following result:

$$\rho(T) \simeq \frac{3\pi^2m}{3T} \left[1 + X(T) - \sqrt{1 + X(T)^2} - 1\right],$$

$$X(T) = \frac{m_A^2(T)}{8\pi^2n_B(T)}, \quad \frac{1}{\bar{m}} = \frac{1}{m_F} + \frac{f_B(-\mu_B)}{2m_B},$$

where $m_F$ is the spinon (holon) mass, $f_B(-\mu_B) = \exp(-\beta\mu_B) - 1)^{-1}$, with the holon chemical potential $\mu_B$, $n_B$ is the holon density, and $\bar{m}_B = n_B/f_B(-\mu_B)$. The factor in the square brackets represents the reduction rate due to $m_A^2(T)$. By assuming the behavior $m_A(T) \propto (T^* - T)^d (d > 0)$ near $T^*$, and ignoring the weak $T$-dependence in $\bar{m}(T)$ and $\bar{n}_B(T)$, we get

$$\rho(T) \propto T[1 - c(T^* - T)^d].$$

The MF value $d = 1/2$ is excluded since the experiment shows smooth deviations from the $T$-linear form, which requires $d > 1$. In this letter, we calculate $m_A$ by setting up the effective field theory and taking the compactness of gauge field into account, finding that $d$ is not a universal constant but has the $T$-dependence.

The effective theory is the Ginzburg-Landau (GL) theory $L(\lambda, A_i)$ coupled with a gauge field $A_i$, where the complex scalar $\lambda$ represents the $d$-wave spinon pairing. Halperin, Lubensky, and Ma considered a similar system in 3D, where $A_i$ corresponds to the electromagnetic field and $\lambda$ to the Cooper-pair field. They calculated the effect of $A_i$ on $\lambda$, which converted the second-order phase transition to first-order. More recently, Ubbens and Lee calculated the one-loop effect of $A_i$ in the SB MFT of the 2D t-J model, and concluded again that the pairing transition at $T^*$ becomes first order. However, their $T^*$ appears below the superconducting transition temperature $T_c$, so they concluded that the spin-gap phase is completely destroyed by gauge-field fluctuations. In the present study, we take the compactness of $A_i$ into account, which originates from the t-J model defined on the lattice itself and gives rise to interactions like $\lambda^2\cos(2\alpha a A_i)$, where $\alpha$ is the lattice constant. Even in the CSS state, it generates nontrivial vertices that are missing in the usual treatments which use $\lambda^2 A_i^2$. We find that the periodic interaction stabilizes the system so as to have a phase transition above $T_c$.

For the 2D GL theory coupled with a gauge field, Nagaosa and Lee argued that vortices put a phase transition into a crossover. On the other hand, the 3D system has a genuine phase transition. This is supported by Monte Carlo simulations and other studies. In the low-temperature phase (spin-gap phase, Higgs phase), vortex loops do not proliferate, while in the high-temperature phase, they do. The former phase is well described by the usual order parameter $\lambda$, while the latter is described by the disorder parameter that measures the vortex-loop density. We assume a small but finite three-dimensionality with anisotropy $\alpha > 0$ ($\alpha = 1$ for 3D and $\alpha = 0$ for 2D), so there takes place a genuine...
phase transition at $T^*$ in the present model.

For a sufficiently small $\alpha$, the 3D critical behavior is observed only in the small interval in $T$ near $T^*$ which vanishes as $\alpha \to 0$. Beyond this, calculations in the pure 2D system ($\alpha = 0$) should give a reliable result according to the general theory of critical phenomena. There is a good example of this; the antiferromagnetic transition of cuprates at $T = T_{AF}$ [13], for which the 3D behavior takes place for $|T - T_{AF}| \lesssim O(1/\ln |\alpha|^2)$. Beyond this interval, the 2D results fit the experimental data well.

Let us start with the SB t-J Hamiltonian given by

$$H = -t \sum_{\sigma} \sum_{x,i} \left( \hat{b}^\dagger_{x+i}\sigma \hat{f}_{x+i}\sigma b_x + H.c. \right) - \frac{2}{3} \sum_{x,i} \left( f^\dagger_{x+i}\sigma f^\dagger_{x+i}\sigma - f^\dagger_{x+i}\sigma f^\dagger_{x+i}\sigma \right)^2 + H_\mu, \quad H_\mu = - \sum_x \left( \mu_B \hat{b}^\dagger b_x + \mu_F \sum_\sigma f^\dagger_\sigma f_\sigma \right).$$

where $f_x$ is the fermionic spinon operator with spin $\sigma$ ($\uparrow, \downarrow$) at site $x$ of a 2D lattice, and $b_x$ is the bosonic holon operator. The direction index $i = 1, 2$ is also used as unit vectors. $\mu_{B,F}$ are the chemical potentials to enforce $\langle \hat{b}_i \hat{b}_i \rangle = 0$, and $\sigma f^\dagger_{x+i}\sigma = 1 - \delta$ where $\delta$ is the doping parameter [13]. We introduce the complex auxiliary fields $\chi_{xi}$ and $\lambda_{xi}$ on the link $(x, x + i)$ to decouple both $t$ and $J$ terms as

$$H_{MF} = \sum_{x,i} \left( \frac{3J}{8} |\chi_{xi}|^2 + \frac{2}{3J} |\lambda_{xi}|^2 ight) - \left\{ \lambda_{xi} \left( \frac{3J}{8} \sum_\sigma f^\dagger_{x+i}\sigma f_{x}\sigma + tb^\dagger_{x+i}b_x \right) + H.c. \right\}, \quad (4)$$

$\chi_{xi}$ describes hoppings of holons and spinons, while $\lambda_{xi}$ describes the resonating-valence-bond (singlet spin-pair) amplitude. In path integral formalism, the partition function $Z(\beta)$ [$\beta = (k_B T)^{-1}$] is given by

$$Z = \int [db][df][d\lambda] \exp(-S), \quad S = \int_0^\beta d\tau \left[ \sum_x \left( \hat{b}_x \hat{b}_x + \sum_\sigma \bar{f}_\sigma f_\sigma \right) + H_{MF} \right], \quad (5)$$

where $\tau$ is the imaginary time and $\hat{b}_x = \partial \hat{b}_x(\tau)/\partial \tau$, etc. Let us consider the low-energy effective theory at temperature $T \ll T_{CSS}$, where $T_{CSS}$ is the critical temperature below which the CSS takes place as a deconfinement phenomenon [13]. It is reasonable to translate the lattice variables to the continuum fields like $f_{x\sigma} \to a_{x\sigma}(x)$, etc. using the lattice constant. The Hamiltonian of the continuum field theory is given by

$$H = N_{site} \frac{3}{4} f \lambda^2 + m_F \chi \int d^d x \left[ |\lambda_a(x)|^2 + |\lambda_d(x)|^2 \right] + \int d^d x \left[ \frac{1}{2m_B} \sum_i \left| D_i b(x) \right|^2 - \mu_B \left| b(x) \right|^2 \right] \quad (6)$$

$$+ \int d^d x \left[ \frac{1}{2m_F} \sum_i \left| D_i f_\sigma(x) \right|^2 - \mu_F \left| f_\sigma(x) \right|^2 \right] + \int d^d k \frac{d^d q}{(2\pi)^4} \left[ \Delta_{SG}(k, q) f^\dagger_\sigma(k) f^\dagger_\sigma(-k + \frac{q}{2}) + H.c. \right],$$

where $\mu_{B,F}$ are the chemical potentials for the continuum theory, $f_\sigma(k)$ is the Fourier transform of $f_\sigma(x)$, and

$$\lambda_{a,d} = \frac{1}{2}(\lambda_1 \pm \lambda_2), \quad \frac{1}{2m_B} = t \chi a^2, \quad \frac{1}{2m_F} = \frac{3}{8} J \chi a^2, \quad k_F^2 = \frac{2\pi}{a^2}(1 - \delta), \quad \Delta_{SG}(k, q) = 2(1 - \delta) \left( \frac{k_F^2 - k_{\lambda}^2}{k_F^2} \right) \lambda_d(q) - 2 \delta \lambda_{a}(q). \quad (7)$$

To obtain $H$, we modified the dispersions of holons and spinons from cosine form to the quadratic one. $D_i \equiv \partial_i - iA_i$ is the covariant derivative with the gauge field $A_i(x)$. Here we introduced $A_i(x)$ by the correspondence $\chi_{xi} \equiv \chi \exp[i\pi A_i(x)]$, where $\chi$ is the radial part of $\chi_{xi}$. We ignore the fluctuations of $|\chi_{xi}|$ since $T \ll T_{c}$, where $T_c$ is the onset temperature of $\chi$. In the SB MFT, $\chi$ is estimated at small $\delta$'s as

$$\chi \simeq \left( \sum_{\sigma} f^\dagger_{x+i+\sigma} f_{x+\sigma} + \frac{8t}{3J} b^\dagger_{x+i} b_x \right)_{MF} \simeq \frac{4}{\pi^2} \sin^2 \left( \frac{\pi}{2} \sqrt{1 - \delta} \right) + \frac{8|t|}{3J} \delta, \quad (8)$$

if the spinon pairing $\chi_{xi}$ is neglected.

To obtain the effective action of $\lambda_1$ and $A_i$, $b$ and $f_\sigma$ are integrated by the standard bilinear integrations. This procedure generates dissipative terms of $\lambda_1$ and $A_i$. The most singular contributions to $Z$ from the integrations over $\lambda_1$ and $A_i$ come from their static $(\tau$-independent) modes, so we keep only the static modes in the effective Lagrangian density, which is given up to the fourth-order in fields and derivatives by

$$L_{eff} = a_s |\lambda_s|^2 + a_d |\lambda_d|^2 + 4b \delta^4 |\lambda_s|^4 + \frac{3}{2} b(1 - \delta)^4 |\lambda_d|^4$$

$$+ 2b \delta^2 (1 - \delta)^2 \left( 4|\lambda_s|^2 |\lambda_d|^2 + 2\lambda_s^2 \lambda_d^2 + 2\lambda_s \lambda_d^2 \right)$$

$$+ c \sum_i \left( 2\delta^2 |D_i \lambda_s|^2 + (1 - \delta)^2 |D_i \lambda_d|^2 \right)$$

$$+ c \delta (1 - \delta) \left( |D_1 \lambda_s D_1 \lambda_d - D_2 \lambda_s D_2 \lambda_d + H.c. \right)$$

$$+ \frac{1}{12 \pi \tilde{m}} \sum_{ij} \frac{1}{4} F_ijkl F_{ij}, \quad (9)$$

where $\tilde{m}$ is defined in Eq.1, and
where we introduced the Proca (massive vector) field
\[ \delta L \]

From the above \( T \) is the Euler number. \( \omega \delta \) is the cutoff of the spinon energy \( \xi \equiv k^2/(2m_F) - \mu_F \) in the one-loop integrals representing spinon pairings, and is estimated as \( \omega \delta \sim \Omega(\mu_F) \). From the potential energy of \( \lambda_i(x) \), the system favors the d-wave state at small \( \delta \)'s, and the s-wave state at large \( \delta \)'s. Let us focus on small \( \delta \)'s by parameterizing \( \lambda_i(x) = \lambda \exp[i\theta(x)] \), \( \lambda_2(x) = -\lambda \exp[i\theta(x)] \). Here we introduced \( \lambda \), the spin-gap amplitude, for the radial parts of \( \lambda_i(x) \), ignoring their fluctuations. Then we have the effective Lagrangian density,
\[ L_{\text{eff}} = L_\lambda + L_A, \]
\[ L_\lambda = ad \lambda^2 + \frac{3}{2} \beta (1 - \delta)^4 \lambda^4, \]
\[ L_A = c (1 - \delta)^2 \lambda^2 (\partial_i \theta - 2A_i) \]
\[ + \frac{1}{12 \pi m} \sum_i \frac{1}{4} F_{ij} F_{ij}. \]
From the above \( L_\lambda \), the MF result is obtained as
\[ k_B T_\lambda = \frac{2 e^7}{\pi} \omega \lambda \exp \left[ -\frac{\pi \lambda}{(1 - \delta)^2} \right], \]
\[ \frac{7 \zeta(3)}{8 \pi^2} (1 - \delta)^2 \left( \frac{\lambda(T)}{k_B T} \right)^2 \approx \frac{1}{3} \left( 1 - \frac{T}{T_\lambda} \right), \]
where \( T_\lambda \) is the critical temperature below which \( \lambda \) develops, and should be identified as \( T^* \). The second result is reliable at \( T \) near \( T_\lambda \).

Let us consider the effect of \( L_A \) on \( L_\lambda \) by integrating over \( A_i \). We have treated \( A_i \) as a noncompact gauge field although it was originally compact. This procedure is appropriate for the kinetic term of \( A_i \) because we consider the region \( T \ll T_{\text{CSS}} \). We will respect the compactness of \( A_i \) and the angle-nature of \( \theta \) by considering the following new Lagrangian \( L_B \) with the periodic mass term;
\[ L_B = \frac{1}{12 \pi m} \sum_i \frac{1}{4} F_{ij} F_{ij} \]
\[ + c (1 - \delta)^2 \lambda^2 \frac{1}{a^2} \left| 4 - 2 \sum_i \cos \left( 2 B_i \right) \right|, \]
where we introduced the Proca (massive vector) field \( B_i \equiv \partial_i \theta/2, \quad (F_{ij} = \partial_i B_j - \partial_j B_i) \). Let us take the unitary gauge. Then the integrals reduce as \( [d\theta][dA_i] \equiv [d\theta][dB_i] \rightarrow [dB_i] \).

Let us estimate the gauge-field mass by the variational method. We choose the variational Lagrangian \( L_B \) for \( L_B \)
\[ L_B' = \frac{1}{12 \pi m} \left( \sum_i \frac{1}{4} F_{ij} F_{ij} + \sum_i \frac{m_A^2}{2} B_i B_i \right), \]
where \( m_A \) is a variational parameter. The variational free energy density \( F_B = F_B' + (L_B - L_B') \) is given by
\[ F_B(m_A) = F_B(0) + \frac{k_B T}{8 \pi} m_A^2 - \frac{4 c (1 - \delta)^2 \lambda^2}{a^2} \left( \frac{m_A^2}{q_c^2} \right) \frac{2}{\pi T}, \]
\[ k_B \Theta(T) \equiv \frac{1}{3 a^2 m} = \lambda \left( \frac{f_B}{4} + \frac{1}{3} f_B (-\mu_B) \right), \]
where \( q_c \), the momentum cutoff of \( B_i \), is \( O(a^{-1}) \). We have omitted the higher-order terms of \( O(m_A^4/q_c^4) \). Note that we took the definition of the propagator at the same point and the trace of a functional operator \( \Theta \) as
\[ \langle B_i(x) B_j(x) \rangle \equiv \lim_{y \to x} \langle B_i(x) B_j(y) \rangle, \]
\[ \Theta (T) = \int d^2 x \lim_{y \to x} \langle \Theta(y) \rangle, \]
By minimizing \( F_B(m_A) \), we get
\[ m_A^2(T) = q_c^2 \left( \frac{96 \pi m c (1 - \delta)^2 \lambda^2}{q_c^2} \right)^{2/(3)}, \]
\[ d(T) = \frac{\Theta(T)}{2 \Theta(T) - T}. \]

We note that the fluctuations of \( A_i \) do not affect the MF result \( T_{\text{CSS}} \) as long as \( d(T) > 1 \) since then the order of the corrections becomes higher than \( \lambda^4 \). Thus the gauge-field mass \( m_A(T) \) starts to develop continuously at \( T_\lambda \) as 
\[ m_A(T) \propto (T^* - T)^{d(T)} \]. That is, the exponent \( d \) is neither 1/2 nor a constant, and drastically changes especially when \( T_\lambda < T \), where \( T_\lambda \) is a root of the equation \( T_A = \Theta(T_A) \), at which \( d(T) \) diverges. This is in strong contrast with the noncompact case \([10]\).

If we write \( q_c^2 = \epsilon a^2 \), we have 
\[ m_A^2(T) = q_c^2 (1 - T/T_{\text{CSS}})^{2/(3)}, \]
with \( z = [16(1 - \delta) \lambda] m/(3 m_F) \). A straightforward estimation, \( \pi a^2 = (2 \pi/a)^2 \), by keeping the area of momentum space, gives \( \epsilon = 4 \pi \). However, this gives rise to a nonrealistic curve of \( \rho(T) \) that deviates from the \( T \)-linear behavior and decreases too rapidly, due to the large factor \( z \approx 4 \). Thus we regard \( q_c \) to be a parameter of the effective theory, and choose \( \epsilon \) so as to obtain a reasonable \( \rho(T) \). For example, we require \( z = 1 \), which implies \( \epsilon = 16 \pi (1 - \delta) m/m_F \).

Finally, we need to consider the renormalization effect of the hopping parameter \( t \). We assume that the 3D system exhibits Bose condensation at the temperature scale of \( T_{\text{CSS}} \sim 2 \pi n_B m_B = 4 \pi t \chi \delta \), and regard \( T_B \) to be the observed \( T_c \) in the lightly-doped region. Since \( t \sim 0.3 \) eV gives rise to \( T_B \sim 3000 \) K at \( \delta \sim 0.15 \), one needs to use an effective \( t^* \sim 0.01 \) eV in place of \( t \) so as to obtain a realistic \( T_c \sim 100 \) K \([10]\).

We show in Fig. 3 the phase diagram with the spin-gap on-set temperature \( T_{\lambda, A} \) at which the mass exponent \( d \) diverges, and \( T_B \). In Fig. 3 we plot \( \rho(T) \). As
explained, the curves reproduce the experimental data much better than those with $d = 1/2$ of the MF result, showing smooth departures from the $T$-linear curves, i.e., $d(T_\lambda) > 1$ for the region of interest, $0.05 \lesssim \delta \lesssim 0.15$.

The present results of $\rho(T)$ support that our treatment of gauge-field fluctuations by the variational treatment of compactness is suitable to describe the spin-gap state in the t-J model, although more investigation is certainly necessary.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig1.png}
\caption{Mean-field phase diagram of the t-J model. $T_B$ is the Bose condensation temperature. $T_A$ is the spin-gap onset temperature. $T_A$ is the root of $T = \Theta(T)$ at which $d(T)$ diverges. We chose $t^* = 0.01$ eV, $J = 0.15$ eV, and $\omega_\lambda = \pi J \chi/(2e^2)$.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig2.png}
\caption{Resistivity $\rho(T)$ in $h/e^2$ for several $\delta$'s with the parameters chosen in Fig.1. The dotted lines represent the case of $X(T) = 0$ in (1). The exponent $d(T_\lambda)$ decreases as $16.4, 4.8, 2.8, 2.0$, as $\delta$ increases.}
\end{figure}

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