Thermodynamics and Phase Diagram of High Temperature Superconductors

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Thermodynamic quantities are derived for superconducting and pseudogap regimes by taking into account both amplitude and phase fluctuations of the pairing field. In the normal (pseudogap) state of the underdoped cuprates, two domains have to be distinguished: near the superconducting region, phase correlations are important up to temperature \( T_\phi \). Above \( T_\phi \), the pseudogap region is determined only by amplitudes, and phases are uncorrelated. Our calculations show excellent quantitative agreement with specific heat and magnetic susceptibility experiments on cuprates. We find that the mean field temperature \( T_0 \) has a similar doping dependence as the pseudogap temperature \( T^* \), whereas the pseudogap energy scale is given by the average amplitude above \( T_c \).

One of the most intriguing problems in high temperature superconductivity is the presence of a region above the critical temperature \( T_c \) and below a temperature \( T^* \) where observable quantities deviate from Fermi liquid behavior. This region is called the pseudogap region because it contains effects similar to superconductivity such as a partial suppression of electronic density of states.

The origin of such a pseudogap above \( T_c \) is unclear. There are four major approaches concerning its theoretical understanding: the first is based on the formation of incoherent Cooper pairs above \( T_c \). Phase order [3] or Bose condensation [4] would then establish superconductivity at \( T_c \). The second assumes that the pseudogap is induced by antiferromagnetic fluctuations [5]. The third approach is based on spin-charge separation where spins bind together to form spin singlets and the energy needed to split them apart leads to the formation of a "spin gap" [6]. The fourth assumes the existence of a quantum critical point [7] but the latter has never been observed. However, these approaches seem to be unable to describe specific heat and magnetic susceptibility.

The main aim of this Letter is to show that various experimental observations can indeed be interpreted in terms of fluctuations of the pairing field \( \psi = |\psi| e^{i\phi} \) and that two temperature regions have to be distinguished (see Fig. 5): for a relatively small temperature interval \( T_c < T < T_\phi \) the phase of \( \psi \) is still uncorrelated in space over some correlation length \( \xi \) (the Kosterlitz-Thouless correlation length in 2D). Thus, in this regime, observables are governed by correlated phase fluctuations described by the XY model. For \( T_\phi < T < T^* \), phases of \( \psi \) are essentially uncorrelated (\( \xi \) is on the order of the lattice constant), but \( |\psi| \) is still fluctuating and nonzero, signaling local pair fluctuations. This explains the wide hump seen in specific heat experiments [1], the depression of the spin susceptibility [8], and the persistence of the pseudogap for \( T < T^* \).

Our method has a major difference with respect to the Emery and Kivelson phase fluctuations scenario [3] of the pseudogap regime: we show that phase fluctuations influence the pseudogap only up to a temperature \( T_\phi \), which is much smaller than \( T^* \). Above \( T_\phi \), observables are thus determined only by the amplitude of the pairing field.

The picture of two different regimes above \( T_c \) is also supported by other experiments: (a) Demsar et al. [9], interpreting the real-time measurements of the quasiparticle relaxation dynamics, find a temperature interval of only a few K in which pair fluctuations associated with their collective phase are important, whereas the pseudogap persists to much higher temperatures. (b) Hall effect measurements [10] in underdoped GdBa$_2$Cu$_3$O$_{y-x}$ show a characteristic temperature \( T' \) between \( T_c \) and \( T^* \), at which the temperature dependence of \( \cot(\Theta_{\mu}) \) deviates from \( T^2 \) and the Hall coefficient has a peak. A possible explanation consists in considering vortex excitations as scattering centers modifying the Hall angle \( \Theta_{\mu} \) and Hall coefficient, which would again suggest identifying \( T' \) with our \( T_\phi \), below which correlated vortices exist.

We base our calculations on a \( d \)-wave attractive Hubbard model

\[
H = -t \sum_{\langle i,j \rangle \sigma} c^\dagger_{i\sigma} c^{}_{j\sigma} - U \sum_i Q^\dagger_{d}\,(i)Q_d\,(i),
\]

with a hopping \( t \) between nearest neighbor sites \( i \) and \( j \) on a square lattice. The interaction favors the formation of onsite \( d \)-wave pairs since \( Q^\dagger_{d}\,(i) = \sum_j D_{ij} Q^\dagger_{d}\,(j) \), where \( D_{ij} \) is \( 1, (-1) \) for \( i \) being the nearest neighbor site of \( j \) in the horizontal (vertical) direction. \( Q^\dagger_{d}\,(i) = (c^\dagger_{i\uparrow} c^\dagger_{j\downarrow} - c^\dagger_{i\downarrow} c^\dagger_{j\uparrow}) / \sqrt{2} \) is the singlet pair operator of neighboring sites. Decoupling the interaction with the help of a Stratonovich-Hubbard transformation, the partition function \( Z = \text{Tr} e^{-\beta H} \) is then

\[
Z = Z_n \int D^2 \psi <\mathcal{T} e^{-\int_0^\beta d\tau \sum_i (1/2)|\phi|^2 + \phi Q_d\,(i) + \text{H.c.})>_H = \int D^2 \psi <\mathcal{T} e^{-\int_0^\beta d\tau \sum_i (1/2)|\phi|^2 + \phi Q_d\,(i) + \text{H.c.})>_H.
\]
where \( \psi = \psi(i, \tau) = |\psi(i, \tau)|e^{i\phi(i, \tau)} \), and \( H_n \) is the non-interacting part. The trace over the fermionic operators can be evaluated, yielding

\[
Z = \int D^2\psi e^{-\int_0^\beta d\tau \sum_i [\frac{1}{2}\psi(\tau)^4 + \text{Tr} \ln G]}.
\]  

(2)

Here \( G \) is a Nambu matrix of one-electron Green functions for fermions interacting with a given space and time dependent pairing field \( \psi(i, \tau) \).

Expanding (2) in powers of \( \nabla \psi \), \( Z \) can be written as a functional integral involving an action \( S[\psi] \) for a field \( \psi \) that changes slowly in space and that can be taken time independent:

\[
S[\psi] = \int d^d r [S_0(|\psi|) + S_1(\nabla \psi)],
\]

(3)

where \( S_0 \) is a local function of \( |\psi(r)| \), and \( S_1 = c|\nabla \psi|^2/2 \) where \( c \) is a constant. \( d \) is the dimension.

Now we compute observables such as energy, specific heat, and spin susceptibility. Our main strategy is to treat separately amplitude and phase fluctuations. In this spirit, two different approaches are possible: (i) The amplitude is fixed and determined by a suitable variational equation. (ii) The energy is expanded around the average amplitude.

Variational method.—We can neglect amplitude correlations since simulations show that they are weak between different sites \( i, j \): \( \langle |\psi_i|/|\psi_j| \rangle - \langle |\psi_i|^2 \rangle = 0 \) (i.e., the amplitude is always positive and shows no critical behavior). Rewriting the free energy \( F \) in terms of a constant amplitude \( |\tilde{\psi}| \) yields

\[
F = -\frac{1}{\beta} \log \int D\phi e^{-\beta[V S_0(|\tilde{\psi}|)]/V/\beta + \int d^d r S_1},
\]

where the Jacobian of the polar transformation is put into the exponential, and \( V \) is the volume. Equating the derivative of \( F \) with respect to \( |\tilde{\psi}| \) to zero leads to the self-consistent equation:

\[
\frac{\partial S_0(|\tilde{\psi}|)}{\partial |\tilde{\psi}|} - \frac{1}{\beta |\tilde{\psi}|} + c|\tilde{\psi}|(\nabla e^{i\phi})^2_{XY} = 0.
\]

(4)

The first term, called the amplitude contribution, leads to the BCS gap equation \([11]\) if other contributions are neglected. The second comes from the Jacobian and implies that the amplitude is never zero. The third term is the expectation value of the energy in the \( XY \) model depending on the constant coupling \( c|\tilde{\psi}|^2/2 \). This contribution characterizes the influence of the phase fluctuations. Solutions of Eq. (4) are reliable for all temperatures except for \( T \ll T_c \). They are accurate at \( T_c \) only in the underdoped regime.

Average value method.—Now we keep the local coupling in \( S_1 \) between amplitude and phase by introducing the statistical Ginzburg-Landau (GL) model (see Ref. [12]). Averages such as \( \langle |\tilde{\psi}| \rangle \) and \( \langle S_1 \rangle \) are computed within the GL model with all fluctuations and correla-

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by requiring that the specific heat jump recovers its BCS value in the overdoped regime. We found that this ratio must be $|\psi_0|/T_0 = 3.06$ in the $s$-wave case, and $|\psi_0|/T_0 = 3.72$ for $d$-wave symmetry.

Both approaches are valid below and above the critical temperature $T_c$, which is the temperature where the phase stiffness becomes zero.

The specific heat $C$ is the sum of $C_0$ and $C_1$, amplitude and gradient contributions, respectively. Defining $Y$ for $YBa_2Cu_3O_6^{13}$, For the fit of Fig. 2, using the reasonable value $\gamma = C/(\gamma_n T)$ where $\gamma_n$ is the Sommerfeld constant, the reduced specific heat is

$$\gamma = \gamma_0(\langle|\psi|\rangle) + \gamma_1,$$

(8)

where $\gamma_1$ is divided by $T_c$ instead of $T$ since $S_1$ is classical and does not satisfy the third law of thermodynamics.

The amplitude contribution $\gamma_0$ is 1 at high temperature. The gradient contribution is normalized as:

$$C_1 = k_B C_1^{(s)} \frac{\gamma_0}{\gamma_n N k_B},$$

(9)

where $V = N \xi_0^3$, and $C_1^{(s)}/(N k_B)$ is the specific heat per number of lattice sites coming from the simulations. Experiments give $\gamma_n = 26$ mJ K$^{-1}$ mol$^{-1} = 252$ J K$^{-1}$ m$^{-3}$. For the fit of Fig. 2, using the reasonable value $\xi_0 = 16$ Å, we get the dimensionless constant $\alpha = k_B/(\xi_0 \gamma_n) = 13.5$. In Fig. 2 the experimental specific heat of $YBa_2Cu_3O_6$, [1] is fitted using the variational method reproducing the double peak structure: a sharp peak at $T_c$ coming mainly from correlated phase fluctuations and a wide hump rounded by amplitude fluctuations (see Fig. 1 for typical distributions). The crossover temperature $T_1$ corresponds to the temperature where $\gamma_1$ is less than approximately 2% of the normal specific heat. In the amplitude equation (4), a two-dimensional density of states $D(\varepsilon) = 1/W$ is used with $W = 5000$ K, $\mu = 0.25$, and $U = 959$ K. These values give $T_0 = 200$ K and $\psi_0 = 2.14 T_0$ in agreement with experiments [15]. The other parameters are $V_0 = 72$ K and $\eta = 5$.

The average value method is used to reproduce specific heat measured for different dopings in Fig. 3. For underdoped systems $x < 0.80$ we use simulations in $d = 2$. For the more overdoped, $x \geq 0.80$, simulations are done in $d = 3$. Parameters $V_0$ and $T_0$ extracted from the best fits are shown in the phase diagram of Fig. 5. $\alpha$ ranges from 7 to 15.

The magnetic susceptibility $\chi$ has two contributions: the paramagnetic spin susceptibility $\chi_p$, and the orbital diamagnetic susceptibility $\chi_d$. $\chi_p$ has been measured by Takigawa et al. [8] on powder $YBa_2Cu_3O_{6.63}$ using Cu and O NMR experiments. The direct contribution coming from phases is negligible because NMR probes essentially the presence of pairs which is related to the amplitude. Therefore, $\chi_p$ is given by the amplitude contribution:

$$\chi_p = \frac{\chi_0}{2T} \int_0^{2\pi} \frac{d\theta}{2\pi} \int_0^{\infty} d\epsilon \cosh^{-1} \left( \frac{\langle|\psi|\rangle^2_{GL} \cos^2(2\theta)}{2T} \right).$$

(10)

$\chi_0$ is the Pauli spin susceptibility. In Fig. 4, we compare the result of Eq. (10) and Takigawa's measurements on powder $YBa_2Cu_3O_{6.63}$. The best fit yields $T_0 = 159.2$ K, $V_0 = 59.3$ K. The competition between amplitude and thermal energy enters into the spin susceptibility (10) by the ratio $\langle|\psi|\rangle/T$. The temperature $T_0$ where phases start to influence the thermodynamics is the temperature where $\langle|\psi|\rangle$ deviates from the average amplitude computed for random phases. This temperature is found above $T_c$ at $T_0 = 90$ K. The orbital diamagnetic susceptibility $\chi_d$ in $YBa_2Cu_3O_{6.60}$ shows fluctuation effects up to about 15 K above $T_c$. For $T > T_0$, phase fluctuations are so strong that $\chi_d$ vanishes [16].

The Green function can be calculated in coherent potential approximation (CPA) for electrons that are scattered from spatially uncorrelated “pair impurities.” We extended the $s$-wave CPA [17] by using a $d$-wave amplitude and uncorrelated phases. We have computed the
density of states at half-filling $\mu = 0.5W$ and maximum amplitude $|\phi| = 0.8W$. The result is a pseudogap of width $\Delta = 0.2W$, although phases are completely uncorrelated.

**Phase diagram.**—Values of $T_0$, $V_0$, and $T_\phi$ extracted from specific heat fits (Fig. 3) are reported in Fig. 5. $T_1$ is the temperature where $C_i/\gamma_a = 2\%$. $T_\phi$ is computed in the same way as shown in Fig. 4 using average amplitudes from specific heat fits. Here it is the temperature where the amplitude differs by 2\% from the amplitude in a random phase field. The energy scale $E_\phi := \langle |\phi| \rangle_{T=200K}$ of the pseudogap is defined as the amplitude at $T = 200 K$. $E_\phi$ shows the same doping dependence as the one found by Tallon and Loram [18]. However, it is not related to any hidden critical point. Phase correlations above $T_c$ grow rapidly in the underdoped regime following the $T_\phi$ line and have a similar doping dependence as Nernst effect results [19]. The gradient specific heat from $\delta_1$ disappears more rapidly such as in the Hall effect [10].

**Discussion.**—Amplitude and phase fluctuations are the key for understanding the pseudogap regime of underdoped high temperature superconductors. Phase correlations disappear completely near a temperature $T_\phi$ above $T_c$, and therefore, for $T > T_\phi$, the pseudogap region is dominated by amplitude fluctuations. The mean field temperature $T_0$ has a similar doping dependence as $T^*$, signaling that the pseudogap region is due to independent fluctuating pairs. Comparisons with measured specific heat on underdoped YBCO reproduce the double peak structure: a sharp peak at $T_c$ occurring mainly from phase correlations and a separate wide hump below $T^*$ rounded by fluctuations. The spin susceptibility, related to the amplitude, recovers its normal behavior near $T^*$, whereas the orbital magnetic susceptibility, related to phases, disappears near $T_\phi$. These considerations are independent from the underlying pairing mechanism, and any microscopic theory inducing pairing should lead to similar conclusions.

All of these findings provide additional evidence for the fact that superconductivity and pseudogap have the same origin. The former is primarily related to phases of the pairing field, which are ordered below the transition temperature and whose correlations survive over a limited temperature region above $T_c$. The pseudogap regime of underdoped materials then extends to much higher temperatures due to the persisting amplitude fluctuations of the pairing field.

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