PAIRING FLUCTUATIONS IN HIGH TEMPERATURE SUPERCONDUCTORS

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Short coherence length superconductors show anomalous behaviour in various thermodynamic and transport properties above their critical temperature T_c. We interpret these observations in the framework of the attractive Hubbard model by considering the interaction between unpaired electrons and virtual pairs. Between T_c and T_ϕ, correlations among the phases of the pairs lead to an XY-type contribution to the specific heat C_V, to strong diamagnetic fluctuations, and to a well-developed pseudogap. Between T_ϕ and T^*, the pairing amplitude is visible in the spin susceptibility, in a broad hump in C_V, and in a pseudogap which fills up. We also study the internal structure of these preformed pairs.

Introduction

High temperature cuprate superconductors show various anomalous features above their transition temperature T_c which deviate from Fermi liquid behaviour. One of the most intriguing observations, made by photoemission and tunnelling, is the reduction in the electronic density of states — the pseudogap — observed between T_c and a temperature T^* which increases when doping is reduced [1]. Thermodynamic equilibrium properties (specific heat [2], spin susceptibility [3], NMR relaxation rate [3], etc.), as well as transport coefficients (resistivity [1], Hall constant [4]) also markedly differ from the behaviour seen in normal metals. The origin of these phenomena has been attributed to different mechanisms. One is the formation of incoherent Cooper pairs (the so-called "preformed" pairs) [5]. The superconducting transition is then due to the onset of phase coherence, respectively to Bose condensation of these pairs. Another approach is based on the proximity of an antiferromagnetic region in the phase diagram [6] and assumes that antiferromagnetic fluctuations, being quite generally at the basis of superconductivity, are responsible for the anomalies. Other instabilities, giving rise to structures like stripes or d-density-waves could also manifest themselves in a pseudogap. Spin-charge separation [7], well-known from interacting electrons in one dimension, could equally provide a scenario leading to the anomalous observations. Finally, the existence of a quantum critical point [8] at a very specific doping value might lead to such behaviour.

The main aim of our calculations is to show that the scenario of strong fluctuations of the superconducting pairing field Δ leading to such "preformed" pairs can indeed provide a coherent description of various anomalous phenomena observed for $T_c < T < T^*$. Our picture is, however, more refined. By analyzing the influence of both, phase and amplitude of Δ, on various observables, we infer that two temperature regions have to be distinguished. For a relatively small temperature interval $T_c < T < T_\phi$, the phase of Δ is still correlated in space over some correlation length ξ (given by the Berezinski-Kosterlitz-Thouless correlation length in $2D$), whereas the amplitude $|\Delta|$ is almost constant. Thus, in this regime, physics is governed by correlated phase fluctuations described by the XY-model. These phase correlations produce the critical part of the specific heat and an enhanced orbital susceptibility observed in some underdoped compounds. For $T_\phi < T < T^*$, the phases of Δ are essentially uncorrelated (ξ is of the order of the lattice constant), but $|\Delta|$ is still non-zero, signalling independent fluctuating local pairs. This explains the wide hump seen in specific heat experiments [2] and the depression of the spin susceptibility for $T < T^*$ [3].

The pseudogap will extend over the whole range $T_c < T < T^*$, being, respectively, phase or amplitude dominated. Thus it will only be suppressed by magnetic fields much higher than the one which destroys phase coherence, and thus superconductivity [9].

The existence of a temperature intermediate between T_c and T^* has also been found in other theoretical work: Devillard and Ranninger [10], using a "Boson-Fermion" description of pairing in cuprates, find that uncorrelated pairing of electrons leads to the opening of a pseudogap at T^*. These pairs acquire well behaved itinerant features at T^*_B, leading to partial Meissner screening, and thus to diamagnetic susceptibility, and Drude-type behaviour of the optical conductivity. As a function of lattice anisotropy (and thus of doping), T_B^* has the same tendency as T_c, whereas the higher temperature T^* has the opposite trend. Although, in [10], T_B^* is related to the temperature where the pair life time becomes long, it might be identified with our T_ϕ. Timm et al. [11] find a "strong" pseudogap with local Cooper pairs subject to strong phase fluctuations existing for $T_c < T < T_c^*$, whereas a weaker supression of $D(E)$ survives up to $T^* > T_c^*$. The intermediate temperature T_c^* can be identified with our T_ϕ, however, contrary to the present interpretation, they conclude that pairs have essentially disappeared above T_c^* and they attribute the pseudogap in this regime to spin fluctuations or the onset of stripe inhomogeneities.

1. Theoretical Framework

We base our calculations on an extended attractive Hubbard model

$$H = \sum_{<l,l'>\sigma} t_{ll'} c_{l\sigma}^\dagger c_{l'\sigma} - W \sum_{<l,l'>} Q_{ll'}^\dagger Q_{ll'} \quad (1)$$

with hopping element $t_{ll'}$ between lattice sites l and l' and on a square lattice with lattice constant a, and an interaction of strength W favoring the formation of singlet pairs on neighboring sites

$$Q_{ll'}^\dagger = \frac{1}{\sqrt{2}}[c_{l\uparrow}^\dagger c_{l'\downarrow}^\dagger - c_{l\downarrow}^\dagger c_{l'\uparrow}^\dagger]. \quad (2)$$

The second term in (2) is identical to the interaction term (antiferromagnetic exchange and density-density coupling) in the $t-J$-model. We select the d-wave part of Q, given in the reciprocal space by

$$Q_d^\dagger(\mathbf{q}) = \frac{1}{2}\sum_{\mathbf{k}}(\cos k_x a - \cos k_y a)[c_{\mathbf{k}\uparrow}^\dagger c_{\mathbf{q-k}\downarrow}^\dagger - c_{\mathbf{k}\downarrow}^\dagger c_{\mathbf{q-k}\uparrow}^\dagger]. \quad (3)$$

Decoupling the interaction with the help of the Stratonovich—Hubbard transformation, the partition function can be written as

$$Z = \mathrm{Tr} e^{-\beta H} = \int D^2\Delta \mathrm{Tr} e^{-\beta H_0} T e^{-\int_0^\beta d\tau H_1(\tau)},$$

$$H_1 = \sum_{\mathbf{q}} \left[\frac{1}{W}|\Delta(\mathbf{q},\tau)|^2 + (\Delta(\mathbf{q},\tau)Q_d^\dagger(\mathbf{q}) + \mathrm{h.c})\right]. \quad (4)$$

The trace over the fermionic operators can be evaluated yielding

$$Z = \int D^2\Delta e^{-\int_0^\beta d\tau L[\Delta]}, \quad (5)$$

with an action

$$L = \sum_{\mathbf{q}} \frac{1}{W}|\Delta(\mathbf{q},\tau)|^2 - \mathrm{Tr}\ln \mathcal{G}^{-1}. \quad (6)$$

Here $\mathcal{G}$ is the Nambu matrix of one-electron Green functions for fermions interacting with a given, space and time dependent pairing field $\Delta(l,\tau) = |\Delta(l,\tau)|e^{i\theta(l,\tau)}$:

$$\sum_{l_1} \begin{pmatrix} \left(\frac{\partial}{\partial\tau} + \mu\right)\delta_{ll_1} - t_{ll_1} & \Delta(l,\tau)D_{l,l_1} \\ \Delta(l,\tau)^* D_{l,l_1} & \left(\frac{\partial}{\partial\tau} + \mu\right)\delta_{ll_1} + t_{ll_1} \end{pmatrix} \mathcal{G}(l_1,l';\tau)$$

$$= \hat{I}\delta_{ll'}\delta(\tau) \quad (7)$$

with

$$D_{l,l_1} = \begin{cases} 1 & \text{for } l,l_1 \text{ n.n. in } x\text{-direction;} \\ -1 & \text{for } l,l_1 \text{ n.n. in } y\text{-direction;} \\ 0 & \text{otherwise;} \end{cases} \quad (8)$$

where "n.n." stands for "nearest neighbor". For simplicity, we have negleted variations of Δ inside a pair. The symbol "Tr" in (7) implies summing over lattice points, respectively wave vectors, integrating over Matsubara time τ and taking the trace over the Nambu matrix. Eckl et al. [12] have calculated $\mathcal{G}$ numerically by fixing $|\Delta|$ and by taking (time independent) phase configurations $\{\theta(l)\}$ coming from Monte Carlo simulations of the $2D$ XY-model. For an analytical approach to the free energy, we first expand $\mathcal{G}$ in powers of $e^{i\theta(l,\tau)} - 1$, the deviation from a homogeneous phase configuration. Alternatively, one can go to the continuum limit and expand in powers of the gradient of Δ [13]. Taking into account the lowest order gradient terms yields $L = L_{\mathrm{BCS}} + L_{XY}$, where L_{BCS} has the form (6) with $\mathcal{G}$ being the BCS solution of (7) for constant $|\Delta|$ and $\theta(l,\tau) = 0$ and

$$L_{XY} = \sum_{<ll'>} V_0[1 - \cos(\theta(l) - \theta(l'))]. \quad (9)$$

The first term of the action yields the free energy of a BCS-system with a constant "pseudogap" $|\Delta|$, whereas the second term represents the free energy of an XY-model. Its properties will be evaluated, either by Monte Carlo simulations (see subsection 2.) or by going over to a vortex representation (used in subsection 2.). We have neglected the time dependence of θ, which would introduce quantum fluctuation [14] and retardation effects.

The one-electron Green function $G(\mathbf{k}, z)$ of the system has to be calculated by averaging $\mathcal{G}$ over the configurations of $\Delta(l,\tau)$ using the probability distribution given by the action in the exponential in (5). The effect of scattering of the quasiparticles from phase fluctuations will then show up as a self-energy determining $\mathcal{G}$. To the lowest order, it is given by [15]

$$\sigma(\mathbf{k}, z) = \frac{1}{\beta N}\sum_{\mathbf{q},z'} G(\mathbf{k}-\mathbf{q}, z-z')\Phi(\mathbf{q}, z'), \tag{10}$$

with

$$\Phi(\mathbf{q}, z) = \int d\tau \sum_{l} e^{i\mathbf{q}\mathbf{R_l}} e^{iz\tau}\langle\Delta^*(l,\tau)\Delta(0,0)\rangle \tag{11}$$

being the phase correlator concentrated at $z = 0$ in the classical XY-system. Such a form of σ, resulting from the expansion leading to (9) and (10), is best suited for the regime of weak phase fluctuations ($T_c < T < T_\phi$). In the region $T_\phi < T < T^*$, a different approach is more appropriate in order to cope with the strong phase fluctuations when calculating $G(\mathbf{k}, z)$. Equation (7) is then interpreted to describe electrons scattered from local "pair impurities" at each lattice site l, their phases being independent and homogeneously distributed between 0 and 2π. For such a system, the "Coherent Potential Approximation" (CPA) has proven to be very efficient to obtain $G(\mathbf{k}, z)$. The system with fluctuating phases is replaced by an effective medium characterized by a Green function with a given (momentum independent) self-energy $\sigma_{eff}(z)$. This quantity is determined by replacing one site of the effective medium lattice by the "true" $\Delta(l,\tau) = |\Delta|e^{i\theta(l,\tau)}$. This "one impurity problem" can be solved and the corresponding $\mathcal{G}$ is averaged over the (constant) probability distribution for the phase $\theta(l,\tau)$. Requiring this averaged G to be equal to the Green function of the effective medium allows one to determine $\sigma_{\text{eff}}(z)$. CPA has been applied to s-wave pairing by Gyorffy et al [16]. It can easily be generalized for the d-wave case at hand.

In order to obtain some information about the internal structure of the preformed pairs, we calculate the propagator X of a singlet pair with the center of mass coordinate $\mathbf{R} = (\mathbf{R}_l + \mathbf{R}_{l'})/2$, the partners of which are at a relative distance $\boldsymbol{\rho} = \mathbf{R}_l - \mathbf{R}_{l'}$, using the operators Q in (2):

$$X(\mathbf{R}, \boldsymbol{\rho}, t; \mathbf{R}', \boldsymbol{\rho}', t') = -i\langle T[Q(\mathbf{R}, \boldsymbol{\rho}, t)Q^\dagger(\mathbf{R}', \boldsymbol{\rho}', t')]\rangle. \tag{12}$$

X corresponds to a special choice of arguments in the general two-electron Green function G_2. After the Stratonovich—Hubbard transformation, leading to the form (5) of the partition function, G_2 can — like $G(\mathbf{k}, z)$ — be expressed as an average over a probability distribution for the field Δ. Alternatively one may start from the usual Bethe—Salpeter equation [17], representing the sum over repeated two-particle scattering events due to a distance dependent interaction (restricted to nearest neighbors in (1)):

$$X(\mathbf{P}, \boldsymbol{\rho}, \boldsymbol{\rho}'; z) = B(\mathbf{P}, \boldsymbol{\rho}-\boldsymbol{\rho}', z) +$$

$$+ i\int d^2r B(\mathbf{P}, \boldsymbol{\rho}-\mathbf{r}, z)V(\mathbf{r})X(\mathbf{P}, \mathbf{r}, \boldsymbol{\rho}', z), \tag{13}$$

where $\mathbf{P}$ is the center of mass momentum and B is the "two-particle bubble",

$$B(\mathbf{P}, \boldsymbol{\rho}, z) = \int d^2k e^{-i\mathbf{k}\boldsymbol{\rho}} \times$$

$$\times \frac{1}{\beta}\sum_{z'} G(\mathbf{k}+\frac{\mathbf{P}}{2}, z')G(\mathbf{k}-\frac{\mathbf{P}}{2}, z-z'). \tag{14}$$

Using a separable potential

$$V(\mathbf{k}-\mathbf{k}') = W\varphi(\mathbf{k})\varphi(\mathbf{k}') \tag{15}$$

of d-wave symmetry

$$\varphi(\mathbf{k}) = \cos k_x a - \cos k_y a\,, \tag{16}$$

one can easily solve the integral equation :

$$X(\mathbf{P}, \boldsymbol{\rho}, \boldsymbol{\rho}', z) = iW \int d\mathbf{k}\, \varphi^*(\mathbf{k})e^{-i\mathbf{k}\boldsymbol{\rho}}B(\mathbf{P}, \mathbf{k}, z) \times$$

$$\times \int d\mathbf{k}'\varphi(\mathbf{k}')\int d\boldsymbol{\rho_1} e^{i\mathbf{k}'\boldsymbol{\rho_1}}X(\mathbf{P}, \boldsymbol{\rho_1}, \boldsymbol{\rho}', z). \tag{17}$$

Here we have subtracted the "background" contribution coming from independent particles. The function

$$g(\mathbf{P}, \boldsymbol{\rho}) = \int d\omega N_{\text{B}}(\omega)\text{Im}X(\mathbf{P}, \boldsymbol{\rho}, \boldsymbol{\rho}, \omega - i\epsilon) \tag{18}$$

with the Bose distribution $N_{\text{B}}(\omega)$ then represents the internal structure of a pair with the center of mass momentum $\mathbf{P}$. Our Bethe—Salpeter equation does not

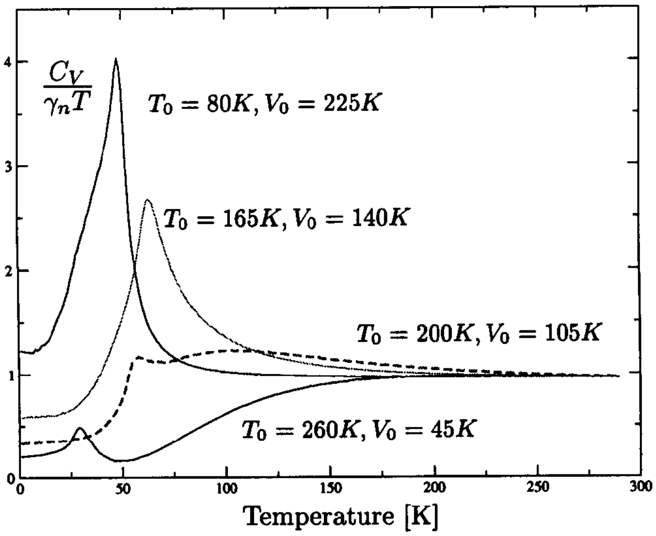

Fig. 1. Specific heat for various phase couplings V_0 and mean-field temperatures T_0

involve any "anomalous" Green functions, as they would arise below the BCS transition temperature. We are thus indeed describing "preformed" Cooper pairs. Their finite life time can be estimated from the imaginary part of the corresponding T-matrix.

2. Results

Our theoretical procedure has led to explicit expressions for the grand canonical free energy and the one-electron Green function, in particular its self-energy, and the pair correlator. For specific calculations, we use the tight binding spectrum in (1) with a nearest neighbor hopping t. All energies are given in units of the non-interacting band width $B = 8t$ in two dimensions, $E = 0$ lying in the middle of the band.

We apply the results of section 1. to the following quantities.

2.1. Specific Heat (C_V)

C_V is obtained from the free energy resulting from the functional integral (5), using the decomposition $L = L_{\mathrm{BCS}}+L_{XY}$. For this calculation, we have replaced L_{BCS} by Landau—Ginzburg form $L_{\mathrm{LG}}(\Delta) = \sigma(t|\Delta|^2+|\Delta|^4/2)$, where the factor $t = T/T_0 - 1$ introduces the mean field temperature and σ (=8 in our calculations) controls the strenght of the amplitude fluctuations. Fig. 1 shows the result for different values of T_0 and V_0. The specific heat peak located at the superconducting transition temperature comes from the XY part (9) corresponding to the fact that — in our approach — the phase transition belongs to the XY universality class. This part has been obtained by Monte Carlo simulation in $2D$, which also yields the Berezinski—Kosterlitz—Thouless (BKT) jump in the phase stiffness. The second feature, a broader hump, is the "amplitude contribution" coming from L_{BCS}. It occurs when T is of the order of the mean value of $|\Delta|$, i.e. when the thermal energy is comparable to the width of the pseudogap. The XY-coupling constant in (9) and the BCS (mean-field) temperature in (10) have been chosen in order to obtain realistic values of T_c and T^*, the temperature at which C_V/T becomes constant. There is good agreement with the measured data for underdoped YBCO [2].

C_V is a first example of a thermodynamic quantity showing the intermediate temperature T_ϕ, which can be identify as the minimum between the peak and the hump, that is between the region where phase correlations contribute to the specific heat and the region where competition between the pseudogap, given by the amplitude of the pairing field, and thermal energy starts to manifest itself. These features can be easily seen in the lowest curves in Fig. 1, for which the parameters ($T_0 = 200$ K, $V_0 = 105$ K and $T_0 = 260$ K, $V_0 = 45$ K, respectively) are adjusted to the underdoped case.

2.2. Diamagnetic Susceptibility

The orbital response χ_{dia} of the electrons to an applied magnetic field (perpendicular to the lattice planes) can be expressed by the auto-correlation function of the electronic current density $\mathbf{j}$. Since, in the functional integral representation (5), the latter involves the phase gradient, χ_{dia} is related to the vortex correlation function of the XY-model [18]. Using known forms of the latter for T above the BKT-temperature and an activated form for the number of thermal vortex excitations, the observed T-dependence of the (strongly enhanced) zero-field χ_{dia} of underdoped YBCO can be well reproduced [19], see Fig. 2. χ_{dia} shows again very clearly the role of T_ϕ. Indeed, phase fluctuations are so strong above T_ϕ that the orbital susceptibility is negligible. The observed field dependence of the magnetization is also much better reproduced than by simple Landau—Ginzburg calculations [18, 19]. It shows a relatively sharp cross-over between a large initial value (yielding the enhanced χ_{dia}) to a smaller slope at higher fields.

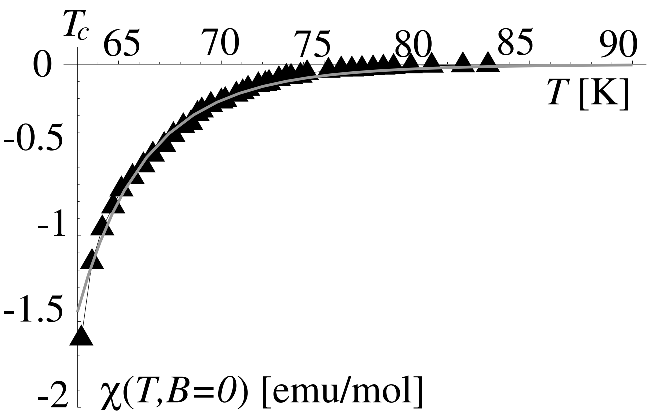

Fig. 2. Zero-field susceptibility $\chi(T, B = 0)$: triangles (▲) are the experimental data from [19] and the full line is the best fit of the theoretical expression assuming the activated behaviour

2.3. Spin Susceptibility

A Zeeman coupling of the electronic spin to an applied magnetic field is introduced into (1). In the partition function (5), its main influence is on the BCS-part of L.

The ratio between the corresponding susceptibility χ_{sp} and its free electron value is then given by

$$\frac{\chi_{sp}}{\chi_{sp}^0} = \frac{1}{2T}\int_0^{2\pi}\frac{d\theta}{2\pi}\times$$

$$\times\int_0^{\infty} d\epsilon D_0(\epsilon)\cosh^{-2}\frac{\sqrt{\epsilon^2+\Delta^2\cos^2(2\theta)}}{2T} \qquad (19)$$

with the free electron density of states D_0. It has BCS-form with a T-dependent (pseudo)gap $|\Delta|$, and the observed reduction with decreasing temperature, shown in Fig. 3, expresses the fact that preformed singlet pairs do not contribute to χ_{sp}. The curves are in good agreement with results from NMR measurements on underdoped YBCO [3]. Approximating the width of the pseudogap by a T-independent value would yield less good agreement with the experimental data.

2.4. Electronic Density of States $D(E)$

Electronic spectral functions $A(\mathbf{k},\omega)$ have been calculated by many authors (see, e.g., [14]) by using expression (11) for the self-energy. A pronounced pseudogap shows up, the width of which does not depend very much on T as in experiment, since it is essentially given by $|\Delta|$. With decreasing the phase correlation length, the pseudogap begins to fill up.

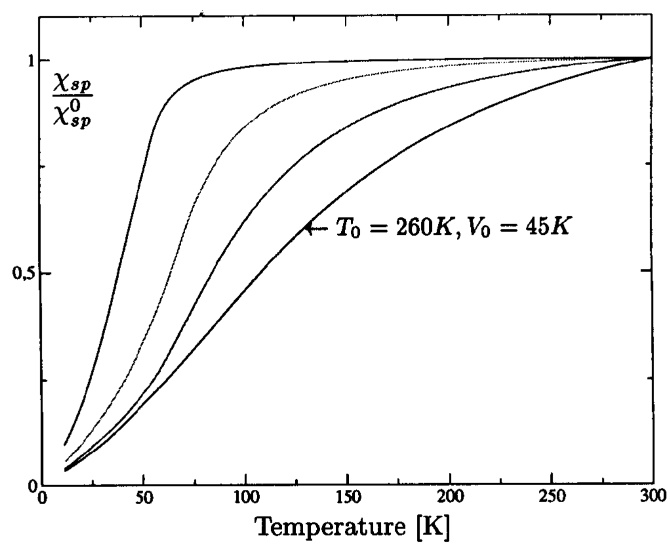

Fig. 3. Spin susceptibility for the same couplings V_0 and mean-field temperatures T_0 as for specific heat in Fig.1

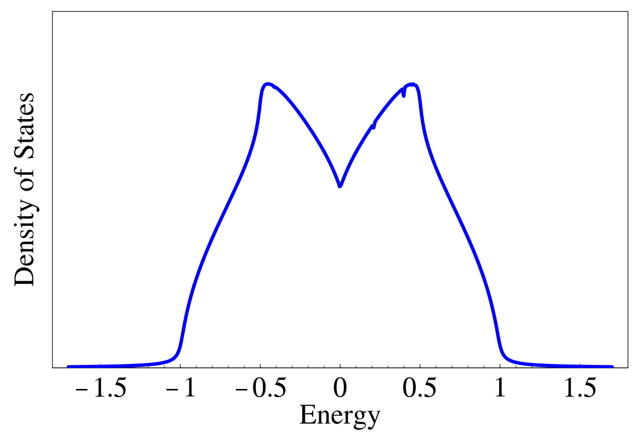

Fig. 4. Electronic density of states calculated in CPA

Here we show $D(E)$, in Fig. 4, as it is obtained from a CPA calculation. The $\mathbf{k}$-dependence of A has the d-wave symmetry imposed by the corresponding form factor in (3). CPA is a temperature independent method. The filling up of the pseudogap is due to a diminishing value of $|\Delta|$ for increasing T and to an increasing finite line width of A coming from interactions with impurities or phonons. The CPA result shows that the pseudogap will persist above T_ϕ, where the coherence of the preformed pairs has totally disappeared. A similar behaviour is also seen in the numerical results by Eckl et al. [12]: the pseudogap goes from T_c right up to T^*, with no qualitative change when T_ϕ is crossed, where the

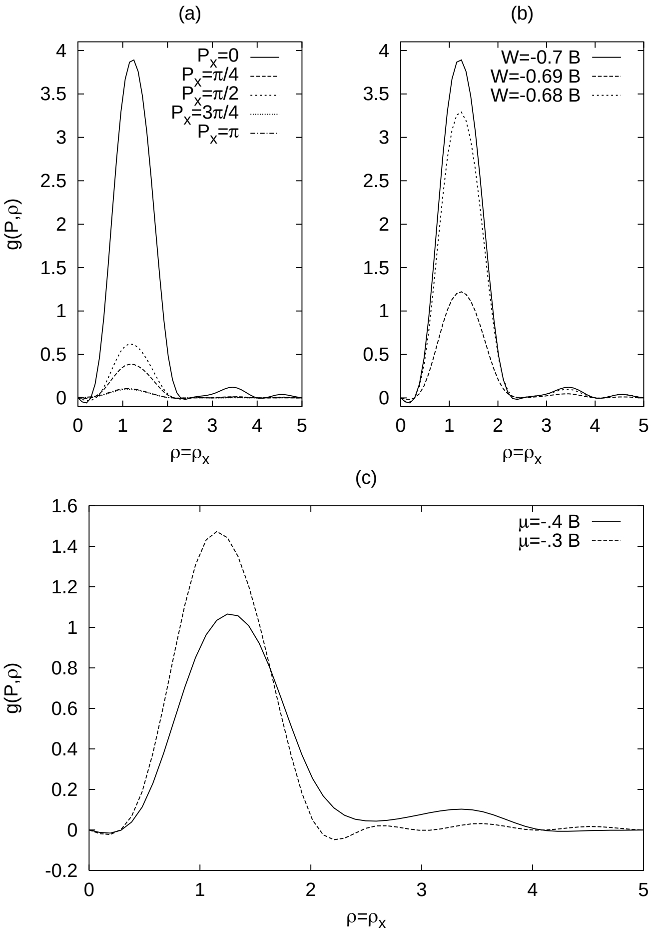

Fig. 5. "Pair structure function" $g(\mathbf{P},\rho)$ for $\mu=-0.4B$, $W=0.7B$ and $T=1.5T_c$: (*a*) $P_y=0$ varying P_x; (*b*) $P_x=0$ varying P_y; (c) for $\mu=-0.4$ and $\mu=-0.3$ and $\mathbf{P}=(0,0)$

BKT correlation length has been reduced to essentially one lattice constant. Calculating the electronic density of states of a quasi-one-dimensional Peierls system, Monnien [20] has also identified a transition from a Gaussian regime with amplitude fluctuations to a non-Gaussian regime where phase fluctuations dominate. The pseudogap evolves smoothly across this transition.

2.5. Internal Structure of Preformed Pairs

The "pair structure function" $g(\mathbf{P},\rho)$ — given by equation (18) — is shown in Fig. 5 for various momenta $\mathbf{P}$ and directions of the internal distance $\boldsymbol{\rho}$ between the partners of the preformed pairs. In Fig. 5,*a* and 5,*b*, the one-electron Green's function is given the form of a pseudogap with a width corresponding to Quantum Monte Carlo data. The curves show that pairs have typically the size of a lattice constant, which is a consequence of our nearest neighbor type attraction. Moving pairs show a slight contraction and they are less numerous than pairs at rest (Fig. 5,*a*). Increasing the interaction strength (Fig. 5,*b*) and moving the chemical potential more inside the band (Fig. 5,*c*) enhances the number of pairs. Negative parts of g indicate a "depletion" with respect to the background of noninteracting particles that has been subtracted. The form of g corresponds to Friedel oscillations, the wavelength of which depends on electronic density (this can be seen by changing the value of μ). As expected, the d-wave pairs are elongated in the direction of the coordinate axis, whereas their width shrinks to zero for $\boldsymbol{\rho}$ along the diagonal ($\rho_x=\rho_y$).

Conclusions

We have analyzed the influence of phase and amplitude fluctuations of the pairing field on various observable quantities of underdoped high temperature superconductors above their transition temperature T_c. Two temperature intervals have been identified: for $T_c < T < T_\phi$ phases are correlated across a T-dependent correlation length, whereeass phases are essentially uncorrelated in space for $T_\phi < T < T^*$, but the amplitude of the pairing field is still finite. The signature of both, phase and amplitude $|\Delta|$, can be seen in the specific heat and in the evolution of the electronic pseudogap. Other quantities are only sensitive to one of them: spin susceptibility is due to $|\Delta|$, whereas the diamagnetic response — due to orbital motion - is primarily influenced by phase correlations.

This picture of two different regimes above T_c is also supported by other experiments:
a) Demsar et al. [21], interpreting real-time measurements of the quasiparticle relaxation dynamics with femtosecond optical spectroscopy, observe a T-dependent collective gap and a T-independent pseudogap. Fluctuations of the former extend over a temperature interval of a few K into the normal region. Here, pair fluctuations associated with their collective phases are still important. The pseudogap, however, persists to much higher temperatures. They associate the latter regime with a fluctuating presence of pairs in regions of low density, where no collective effects are present.
b) Misochko et al. [22] investigate the optical reflectivity in the same femtosecond domain in near optimally doped YBCO. Their observations suggest the existence of two regimes with different dynamical and relaxational properties. At higher T, local but

uncorrelated pairs are formed, giving rise to localized excitations in the lattice. When T is decreased, these excitations disappear showing that the pairs become itinerant. Since the superconducting current is proportional to the phase gradient, itinerance becomes indeed possible when phase coherence is built up.

c) Hall effect measurements [4] in underdoped $GdBa_2O_{7-\delta}$ show a characteristic temperature T' between T_c and T^*, at which the T-dependence of $\mathrm{ctg}(\Theta_H)$ deviates from T^2 and the Hall coefficient has a peak. A possible explanation put forward by the authors consists in considering vortex excitations as scattering centers modifying thus the Hall angle Θ_H and Hall coefficient, which could again suggest identifying T' with our T_ϕ below which correlated vortices exist.

The evolution of the superconducting gap structure into the pseudogap near T_c seen in [21], as well as in photoemission [24] and tunnelling experiments [25], is another interesting theoretical challenge. In our approach, this can be done by letting the phases below T_c fluctuate around a fixed mean value, so that the average $\langle\Delta\rangle$ of the complex field Δ is non-zero. Preliminary calculations [23] show that it is then possible to reproduce the cross-over from a rather broad pseudogap structure in the electronic density of states above T_c to the sharper gap peaks below T_c.

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НАДПРОВІДНІ ФЛУКТУАЦІЇ У ВИСОКОТЕМПЕРАТУРНИХ НАДПРОВІДНИКАХ

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Надпровідники з малою кореляційною довжиною демонструють аномальну поведінку різних термодинамічних і транспортних властивостей при температурах, вищих за їхню критичну температуру T_c. Ми інтерпретуємо ці спостереження в рамках моделі Хаббарда з притяганням, розглядаючи взаємодію неспарених електронів з віртуальними парами. Між T_c і T_ϕ кореляції фаз таких локальних пар дають внесок типу XY-моделі у теплоємність C_V та приводять до сильних діамагнітних флуктуацій і добре розвиненої псевдощілини. Між T_ϕ та T^* амплітуда спарювання проявляється у спіновій сприйнятливості, у широкому максимумі на залежності $C_V(T)$ та псевдощілині, що заповнюється. Вивчається також структура локальних пар.

СВЕРХПРОВОДЯЩИЕ ФЛУКТУАЦИИ
В ВЫСОКОТЕМПЕРАТУРНЫХ СВЕРХПРОВОДНИКАХ

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Сверхпроводники с малой корреляционной длиной демонстрируют аномальное поведение различных термодинамических и транспортных свойств при температурах, выше их критической температуры T_c. Мы интерпретируем эти наблюдения в рамках модели Хаббарда с притяжением, рассматривая взаимодействие неспаренных электронов с виртуальными парами. Между T_c и T_ϕ корреляции фаз таких локальных пар приводят к вкладу типа XY-модели в теплоемкость C_V, к сильным диамагнитным флуктуациям, хорошо развитой псевдощели. Между T_ϕ и T^* амплитуда спаривания проявляется в спиновой восприимчивости, в широком максимуме на зависимости $C_V(T)$ и в заполняющейся псевдощели. Изучается также структура локальных пар.