The Stripe-Phase Quantum-Critical-Point Scenario
for High-$T_c$ Superconductors

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Abstract

A summary is given of the main outcomes of the quantum-critical-point scenario for high-$T_c$ superconductors, developed in the last few years by the Rome group. Phase separation, which commonly occurs in strongly correlated electronic systems, turns into a stripe instability when Coulomb interaction is taken into account. The stripe phase continuously connects the high doping regime, dominated by charge degrees of freedom to the low doping regime where spin degrees of freedom are most relevant. Dynamical stripe fluctuations enslave antiferromagnetic fluctuations at high doping. Critical fluctuations near the stripe instability mediate a singular interaction between quasiparticles, which is responsible for the non-Fermi liquid behavior in the metallic phase and for the Cooper pairing with $d$-wave symmetry in the superconducting phase.

I. THE FRAMEWORK

Since the discovery of superconducting (SC) copper oxides a formidable effort has been produced to provide a unified theory for the rich phase diagram of these materials.
The antiferromagnetic (AFM) phase at zero and very low doping is usually described as resulting from the strongly correlated nature of the copper-oxygen planes, within the Hubbard model or the related $t-J$ model.

As far as the SC phase is concerned the main points under investigation are the nature of the (strong) pairing mechanism, the unusual ($d$-wave) symmetry of the order parameter and the strong dependence of the critical temperature $T_c$ on the doping $x$.

The properties of the normal state are to some extent even more challenging, the standard Fermi liquid (FL) theory appearing to be violated. The copper oxides are characterized by a low dimensionality, revealed by the strong anisotropy of the transport properties. In the metallic phase above $T_c$ at optimum doping a non-FL behavior sets in, with a linear in-plane resistivity over a wide range of temperatures [2], indicating the absence of any energy scale, besides the temperature itself. In the underdoped region two new temperature scales appear above $T_c$. The higher, $T_0(x)$, marks the onset of a new regime characterized by a reduction of the quasiparticle density of states, and is mainly revealed by the presence of broad maxima in the spin susceptibility [3], and by a downward deviation of the in-plane resistivity as a function of the temperature [4]. At a lower temperature $T^*(x)$ a (local) gap in the spin and charge channels appears in ARPES [3–7], NMR [8], neutron scattering [9–12] and specific heat measurements [13].

Anderson proposed [14] to extend the $d=1$ Luttinger Liquid behavior to $d=2$ and explain the anomalies in the metallic phase.

However, no sign of such a new quantum metallic state was found within a renormalization-group approach in $d=2$ [13]. Rather a dimensional crossover drives the system to a FL state as soon as $d > 1$ in the presence of short-range forces [16]. When long-range forces are taken into account a non-FL behavior may arise in the presence of a sufficiently singular interaction $\Gamma_{eff}(q) \sim 1/q^\alpha$ with $\alpha \geq 2(d-1)$ [17].

The onset of an instability is a mechanism which provides a suitable singular scattering. Indeed critical fluctuations mediate an effective interaction between quasiparticles.
\[ \Gamma_{eff}(q, \omega) \simeq -V/[(q - Q)^2 + \kappa^2 - i\gamma\omega] \] where \( V \) is the strength of the static effective potential at criticality, \( Q \) is the critical wavevector, \( \kappa^2 \sim \xi^{-2} \) is a mass term which is related to the inverse of the correlation length and provides a measure of the distance from criticality. The characteristic time scale of the critical fluctuations is \( \gamma \). We point out that the static part of this effective interaction has the form of the Ornstein-Zernike critical correlator.

Proposals about the nature of the relevant instability include (i) an AFM Quantum Critical Point (QCP) \([18,19]\), (ii) a charge-transfer instability \([20]\), (iii) an “as-yet unidentified” QCP regulating a first-order phase-transition between the AFM state and the SC state \([21]\), (iv) an incommensurate charge-density-wave (ICDW) QCP \([22,23]\).

The theory of the AFM QCP \([19]\) is based on the hypothesis that the presence of an AFM phase at low doping is the relevant feature common to all cuprates and on the observation that strong AFM fluctuations survive at larger doping \([8-12]\). However at doping as high as the optimum doping it is likely that charge degrees of freedom play a major role whereas spin degrees of freedom follow the charge dynamics, and are enslaved \([24]\) by the charge instability controlled by the ICDW QCP \([22,23]\). The AFM fluctuations are thus extended to a region far away from the AFM QCP, due to the natural tendency of hole-poor domains towards antiferromagnetism. The strong interplay between charge and spin degrees of freedom gives rise to the “stripe phase” which continuously connects the onset of the charge instability (ICDW QCP) at high-doping with the low-doping regime characterized by the tendency of the AFM background to expel mobile holes. Because of this we shall more properly refer to the ICDW QCP as the stripe QCP. Therefore we point out that the presence of a stripe QCP is not alternative to the presence of the AFM QCP, which is found at lower doping. The two points control the behavior of the system in different regions of doping. On the other hand the existence of a QCP at optimum doping, where no other energy scale besides the temperature is present in transport measurements, is the natural explanation for the peculiar nature of this doping regime in the phase diagram of all SC copper oxides.

There is an increasing amount of theoretical and experimental evidence in favour of the presence of a QCP near optimum doping. Indeed, the instability with respect to phase
separation (PS) into hole-rich and hole-poor regions is a generic feature of models for strongly correlated electrons with short-range interactions \[25\], which is turned into a frustrated PS \[26\] or in an ICDW instability \[22\] when long-range Coulomb forces are taken into account to guarantee large-scale neutrality. Close to PS (or ICDW) there is always a region in parameter space where Cooper pair formation is present, pointing towards a connection between PS and superconductivity.

The most compelling evidence for a QCP near optimum doping is provided by the resistivity measurements. An insulator-to-metal transition is found when the SC phase is suppressed by means of a pulsed magnetic field \[27\]. When extrapolated to zero temperature such a transition takes place near optimum doping, and anyway at too high a doping to be associated to the spin-glass region \[28\] characterized by the local moment formation as seen in the muon experiments \[29\]. The spin-glass region should be instead a signature of the coexistence of superconductivity with antiferromagnetism proper of the $SO(5)$ theory \[30\]. Moreover a clear indication that this insulator-to-metal transition \[27\] is driven by some spatial charge ordering is provided by its occurrence at a much higher temperature in samples near the filling 1/8. Commensurability effects near this “magic” filling have repeatedly been reported in related compounds \[31\].

Hints for a critical behavior of the charge susceptibility come from the study of the chemical potential shift in PES and BIS experiments \[32\]. A dramatic flattening of the $\mu$ vs $x$ curve starting at $x \approx 0.15$ could be the signature of a divergent compressibility. Finally, stripes of either statical or dynamical nature are seen in neutron scattering experiments \[24\], EXAFS \[33\] and X-ray diffraction \[34\].

It must be pointed out that the characteristics of the stripe phase produced by the ICDW instability are system and model dependent. The direction of the critical wavevector $Q_c$ is diagonal in YBCO \[10\], and in nickelates \[35\] where one-hole filled domain walls are present, and vertical in Nd doped LSCO \[31\], with half-filled domain walls. It has been shown that a strong on-site Hubbard repulsion and a long-range potential stabilize vertical half-filled stripes \[36\].
The stripe-QCP scenario provides therefore a scheme to interpolate between the repulsion which gives rise to the AFM state at low doping and the attraction giving rise to SC, through the (local) PS or ICDW.

II. THE NORMAL STATE

To investigate the effect of the stripe QCP on the normal-state properties of the system we are led to consider an effective interaction between quasiparticles

$$\Gamma_{\text{eff}}(q, \omega) = -\sum_{i=c,s} \frac{V_i}{[k_i^2 + \eta_q - Q_i - i\gamma_i \omega]},$$

(1)

where $\eta_q = 2 - (\cos q_x + \cos q_y)$, which is mediated by both charge ($c$) and (enslaved) spin ($s$) fluctuations. The $q$ dependence in (1), was taken in the cos-like form to reproduce the $(q - Q_i)^2$ behavior close to the critical wavevectors $Q_i$ and to maintain the lattice periodicity near the zone boundary.

We point out that the above form (1) for the interaction mediated by charge fluctuations was found within a slave-boson approach to the Hubbard-Holstein model with long-range Coulomb interaction, close to the ICDW instability [37]. The same form mediated by spin fluctuations corresponds to the dynamic susceptibility proposed by Millis, Monien and Pines [38] to fit NMR and neutron scattering experiments, in the limit of strong damping.

We take a free-electron spectrum of the form

$$\xi_k = -2t(\cos k_x + \cos k_y) - 4t' \cos k_x \cos k_y - \mu,$$

(2)

where nearest-neighbor ($t$) and next-to-nearest-neighbor ($t'$) hopping terms are considered, to reproduce the main features of the band dispersion and the Fermi surface (FS) observed in SC copper oxides. The chemical potential is treated self-consistently within a perturbative approach, to fix the number of particles.

The first-order in perturbation theory yields an electron self-energy $\Sigma(k, \varepsilon) = \sum_{i=c,s} \Sigma_i(k, \varepsilon)$, where

$$\text{Im} \Sigma_i(k, \varepsilon) = \frac{V_i}{\gamma_i} \int_{-\pi}^{+\pi} \int \frac{d\xi_{k'}}{4\pi^2} \frac{d\xi_{k'}}{4\pi^2} \left[ f(\xi_{k'}) + b(\xi_{k'} - \varepsilon) \right] \times$$
\[
\frac{\varepsilon - \xi_{k'}}{[\varepsilon - \xi_{k'}]^2 + \Omega_i^2(k - k')},
\]

(3)

\[
\Omega_i(q) = \gamma_i^{-1}[\kappa_i^2 + \eta_{q - Q_i}],
\]

\(f(\varepsilon)\) is the Fermi function and \(b(\varepsilon)\) is the Bose function. The real part of the self-energy \(\text{Re}\Sigma(k, \varepsilon)\) is obtained by a Kramers-Kröning transformation of (3). To keep the inversion symmetry \(k \to -k\) we symmetrize the self-energies \(\Sigma_{c,s}\) with respect to \(\pm Q_{c,s}\). We assume that \(Q_s = (\pi, \pi)\), neglecting the possibility for a discommensuration of the spin fluctuations in a (dynamical) stripe phase [24]. This would introduce minor changes in our results. The relevant direction of the critical wavevector \(Q_c\) is still debated and can be material dependent [39]. We analyze here the case \(Q_c \sim (1, -1)\) which has been suggested by the analysis of ARPES experiments on Bi2212 [40].

In the absence of superconductivity the stripe QCP, located at \(\kappa_{c}^2 = 0\), is the end point of two lines which divide the \(T\) vs \(x\) plane in three regions, the ordered-phase region at lower doping and low temperature, the quantum disordered region at higher doping and low temperature and the quantum critical region around the critical doping \(x_c \simeq x_{\text{optimal}}\) where the only energy scale is \(\kappa_{c}^2 \sim T\) and the maximum violation of the FL behavior in the metallic phase is found. In this region the system is characterized by an anomalously large quasiparticle inverse scattering time \(1/\tau_k \sim \sqrt{T}\) at those point of the FS connected by the critical wavevectors (hot spots) and displays a linear-in-\(T\) resistivity in \(d = 2\) [41], which is turned to a \(T^{3/2}\) which could be a signature of a \(d = 3\) quantum critical behavior for less anisotropic systems, with the change in the temperature dependence occurring over the whole temperature range by increasing doping.

To study the effect of the singular interaction (1) on the single-particle properties in the metallic phase we calculate the spectral density \(A(k, \varepsilon) = \frac{\pi^{-1}|\text{Im}\Sigma(k, \varepsilon)|}{\{(\varepsilon - \xi_k - \text{Re}\Sigma(k, \varepsilon))^2 + [\text{Im}\Sigma(k, \varepsilon)]^2\}}\). To allow for a comparison with ARPES experiments we analyze the convoluted spectral density

\[
\tilde{A}(k, \varepsilon) = \int_{-\infty}^{+\infty} d\varepsilon' A(k, \varepsilon') f(\varepsilon') E(\varepsilon' - \varepsilon)
\]

(4)

which takes care of the absence of occupied states above the Fermi energy, through the Fermi function \(f(\varepsilon)\), and of the experimental energy resolution \(\Delta\), through a resolution function.
We take $E(\varepsilon; \Delta) = \exp(-\varepsilon^2/2\Delta^2)/\sqrt{2\pi\Delta^2}$ or $= [\vartheta(\varepsilon + \Delta) - \vartheta(\varepsilon - \Delta)]/2\Delta$ according to numerical convenience. For the sake of definiteness we choose our parameters in (3) to fit the band-structure and FS of Bi2212, namely $t = 200$ meV, $t' = -50$ meV, and $\mu = -180$ meV, corresponding to a hole doping $x \simeq 0.17$ with respect to half filling. The parameters appearing in the effective interaction (1) were taken as $V_{c,s} = 400$ meV, $\kappa_{c,s} = 0.01$ and $\gamma_{c,s} = 0.01$ meV$^{-1}$.

The quasiparticle spectra are characterized by a coherent quasiparticle peak at an energy $\varepsilon \simeq \xi_k$ and by shadow peaks at energies $\varepsilon \simeq \xi_{k-Q_i}$, produced by the interaction with charge and spin fluctuations. The shadow peaks do not generally correspond to new poles in the electron Green function and are essentially incoherent, although they follow the dispersion of the shadow bands. Their intensity varies strongly with $k$ and increases when $\xi_{k-Q_i}$ approaches the value $\xi_k$. In particular at the hot spots, where $\xi_k \simeq \xi_{k-Q_i} \simeq 0$ and the non-FL inverse scattering time $1/\tau_k \sim \sqrt{\xi_k}$ [11], there is a suppression of the coherent spectral weight at the Fermi energy.

We also study the $k$-distribution of low-lying spectral weight $w_k = \tilde{A}(k, \varepsilon = 0)$. The transfer of the spectral weight from the main FS to the different branches of the shadow FS at $\xi_{k-Q_i} \simeq 0$ produces features which are characteristic of the interaction with charge and spin fluctuations and of their interplay. In particular the symmetric suppression of spectral weight at the $M$ points of the Brillouin zone, which would be due to spin fluctuations alone, is modulated by charge fluctuations (Fig. 2). This is also the case for the (weak) hole pockets produced by spin fluctuations around the points $(\pm \pi/2, \pm \pi/2)$. The interference with the branches of the shadow FS due to charge fluctuations enhances these pockets around $\pm(\pi/2, \pi/2)$ and suppresses them around $\pm(\pi/2, -\pi/2)$ (Fig. 2). Experimental results on this issue are controversial. Strong shadow peaks in the diagonal directions, giving rise to hole pockets in the FS, have been reported in the literature [40,42], where other experiments found only weak (or even absent) features [3].

We point out that, because of the transfer of spectral weight to the shadow FS, the
experimentally observed FS may be rather different from the theoretical FS, determined as
\[ \xi_k + \text{Re}\Sigma(k, \varepsilon = 0) = 0. \] The observed evolution of the FS could, indeed, be associated with the change in the distribution of the low-laying spectral weight, without the topological change in the quasiparticle FS which was proposed in [43].

III. SUPERCONDUCTIVITY

In the stripe-QCP scenario the dynamical precursors of the ICDW mediate an attractive interaction in the Cooper channel [44]. As a matter of simplification we solve the BCS-like equation

\[ \Delta_k = - \int \int_{-\pi}^{+\pi} \frac{dk_x' dk_y'}{4\pi^2} \frac{\tanh(\epsilon_{k'}/2T)}{2\epsilon_{k'}} \times \left[ \frac{V_s}{\kappa_s^2 + \eta_{k-k'-Q_s}} - \frac{V_c}{\kappa_c^2 + \eta_{k-k'-Q_c}} \right] \Delta_{k'} \] (5)

where \( \epsilon_k = \sqrt{\xi_k^2 + \Delta_k^2} \), \( \Delta_k \) is the gap parameter, and both the charge- and spin-induced static effective interactions in the Cooper channel have been considered, corresponding to an interaction in the particle-hole channel [44]. A constructive interference between the small-\( q \) attraction associated to charge fluctuations and the large-\( q \) repulsion associated to spin fluctuations yields both a high critical temperature and a gap parameter \( \Delta_k \) with \( d \)-wave symmetry (Fig. 3).

The variation of the critical temperature with doping follows the variation of the relevant energy scale in each region of the phase diagram. In the overdoped (quantum disordered) region \( \kappa_c^2 \sim x - x_c \) (at low temperatures) and \( T_c \) decreases rapidly with increasing doping. In the quantum critical region around optimal doping \( \kappa_c^2 \sim T \) and \( T_c \) is almost doping independent.

In the underdoped region new scales of energy appear. At \( T_0(x) \), which correspond to the mean-field ICDW critical temperature, precursors of the stripe phase show up in the reduction of spectral weight near the Fermi energy, i.e. a reduction of the static spin susceptibility [3]. At the same time the damping of the AFM fluctuations is reduced and
(almost) propagating spin waves appear, with a natural crossover from a dynamical index \( z = 2 \) to \( z = 1 \). At lower temperatures, since Boebinger’s experiment \cite{27} suggests that superconductivity is hindering the formation of a static CDW, we have to assume that the onset of (local) superconductivity introduces a cut-off for critical fluctuations. Thus, near the onset of the stripe phase \( T_{CDW}(x) \lesssim T^*(x) \), we take a mass term \( \kappa_c^2 \simeq \max[\Delta_{max}(T), T - T_{CDW}(x)] \), where \( \Delta_{max} \) is the maximum over \( k \) of the (local) superconducting gap \( |\Delta_k| \). When introduced into (5), this dependence allows the (local) gap to survive up to a temperature as high as \( T^* \), even if the phase coherence, which is necessary for bulk superconductivity, develops at a lower temperature \( T_c \). This produces a long pseudogap tail in the underdoped region, which, despite the crudeness of our approximations, displays the behavior of the analogous quantity as measured in ARPES experiments in underdoped Bi2212 \cite{1}.

**IV. CONCLUSIONS**

In this paper we briefly recapitulated the stripe-QCP scenario and presented some of its consequences in the normal and SC states. Within this scenario, the occurrence of a charge-ordering instability, only hindered by the setting in of a SC phase (in this sense it would be more appropriate to speak about a “missed-QCP” scenario) provides the underlying mechanism ruling the physics of the SC cuprates. In particular, it gives rise to the formation of the observed stripe textures in these materials, to the non-Fermi liquid properties of the normal phase, to the main features found in ARPES experiments, and to the strong pairing interaction. The most natural location for this QCP is optimum doping where the strongest violation to the FL behavior and the highest critical temperature occur. Indeed the physical properties governed by the proximity to a QCP account for the ubiquitous *universal* behavior, observed near optimal doping in all SC copper oxides. This rationale is missing in the theory of the AFM QCP or in the theory of the QCP associated to the coexistence of antiferromagnetism and superconductivity, which would also be located near the AF phase at low doping or near the spin glass transition.
We conclude by remarking that the scenario of the stripe QCP near optimum doping, hidden by the occurrence of the SC phase, shares a common origin (the Coulomb-frustrated phase separation) with the scenario proposed by Emery, Kivelson and coworkers, but relies on a distinct mechanism. In this latter proposal, the anomalous normal properties stem from the marked one-dimensional character of the metallic stripe phase, and the pairing arises from the in-and-out pair hopping from the 1D stripes into the spin-gapped AF background. A related description of the stripe phase in terms of purely one-dimensional strings has been also put forward by Zaanen [45]. In our picture the non-FL character of the metal arises from the singular scattering by critical fluctuations near the QCP for the onset of the stripe phase. The fluctuations of the stripe texture also provide a strong pairing potential accounting for high critical temperatures. We believe that this more two-dimensional physical description in terms fluctuations of the stripe texture is closer to the reality, at least for the optimal and overdoped systems, where the substantial metallic character of the systems is difficult to reconcile with the formation of strongly one-dimensional long-living stripe structures.

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REFERENCES

[1] J. G. Bednorz and K. A. M"uller, Z. Phys. B 64, 189 (1986).

[2] H. Takagi et al., Phys. Rev. B 40, 2254 (1989).

[3] D. C. Johnston, Phys. Rev. Lett. 62, 957 (1993).

[4] T. Ito, K. Takenaka and S. Uchida, Phys. Rev. Lett. 70, 3995 (1993).

[5] D. S. Marshall et al., Phys. Rev. Lett. 76, 4841 (1996).

[6] J. M. Harris et al., Phys. Rev. Lett. 79, 143 (1997).

[7] H. Ding et al., Nature 382, 51 (1996).

[8] C. Berthier et al., J. de Physique I, December 1996.

[9] J. Rossat-Mignod et al., Physica B 186-189, 1 (1993).

[10] Pengcheng Dai, H. A. Mook and F Doğan, Phys. Rev. Lett. 80, 1738 (1998).

[11] T. E. Mason et al., Physica B 199-200, 284 (1994); T. E. Mason et al., Phys. Rev. Lett. 77, 1604 (1996).

[12] S. Petit et al., Physica B 234-236, 800 (1997); P. Bourges et al., Phys. Rev. B 53, 876 (1996).

[13] J. W. Loram et al., Phys. Rev. Lett. 71, 1740 (1993).

[14] P. W. Anderson, Science 235, 1196 (1987); Phys. Rev. Lett. 64, 1839 (1990); ibid. 65, 2306 (1990).

[15] J. Feldman, H. Kn"orrer, D. Lehmann and E. Trubowitz in Constructive Physics, edited by V. Rivesseau, Springer Lecture in Physics (Springer, New York, 1995).

[16] C. Castellani, C. Di Castro and W. Metzner, Phys. Rev. Lett. 72, 316 (1994).

[17] C. Castellani and C. Di Castro, Physica C 235-240, 99 (1994).
[18] S. Sachdev and J. Ye, *Phys. Rev. Lett.* **69**, 2411 (1992).

[19] P. Montoux, A. V. Balatsky and D. Pines, *Phys. Rev. B* **46**, 14803 (1992); A. Sokol and D. Pines, *Phys. Rev. Lett.* **71**, 2813 (1993); P. Montoux and D. Pines, *Phys. Rev. B* **50**, 16015 (1994).

[20] C. M. Varma, *Physica C* **263**, 39 (1996) and references therein.

[21] R. B. Laughlin, cond-mat/9709195.

[22] C. Castellani, C. Di Castro and M. Grilli, *Phys. Rev. Lett.* **75**, 4650 (1995).

[23] C. Castellani, C. Di Castro, and M. Grilli, *Z. Phys. B* **103**, 137 (1997).

[24] J. M. Tranquada *et al.*, *Phys. Rev. B* **56**, 7689 (1996).

[25] V. J. Emery, S. A. Kivelson and H. Q. Lin, *Phys. Rev. Lett.* **64**, 475 (1990); M. Marder, N. Papanicolau and G. C. Psaltakis, *Phys. Rev.*, **41**, 6920 (1990); M. Grilli *et al.*, *Phys. Rev. Lett.* **67**, 259 (1991); N. Cancrini *et al.*, *Europhys. Lett.* **14**, 597 (1991); R. Raimondi *et al.*, *Phys. Rev. B* **47**, 3331 (1993); S. Caprara, C. Di Castro and M. Grilli, *Phys. Rev. B* **51**, 9286 (1995); F. Bucci *et al.*, *Phys. Rev. B* **52**, 6880 (1995); M. Grilli and C. Castellani, *Phys. Rev. B* **50**, 16880 (1994).

[26] V. J. Emery and S. A. Kivelson, *Physica C* **209**, 597 (1993).

[27] G. S. Boebinger *et al.*, *Phys. Rev. Lett.* **77**, 5417 (1996).

[28] J. H. Chou *et al.*, *Phys. Rev. B* **46**, 3179 (1992).

[29] J. Budnik, this Conference.

[30] Shou-Cheng Zhang, *Science* **275**, 1089 (1997).

[31] J. M. Tranquada *et al.*, *Nature* **375**, 561 (1995).

[32] A. Ino *et al.*, *Phys. Rev. Lett.* **79**, 2101 (1997).
[33] A. Bianconi, Physica C 235-240, 269 (1994); A. Bianconi et al., Phys. Rev. Lett. 76, 3412 (1996); A. Bianconi et al., Phys. Rev. 54, 12018 (1996).

[34] A. Bianconi et al., Phys. Rev. B 54, 4310 (1996).

[35] J. M. Tranquada, D. J. Buttrey and V. Sachan, Phys. Rev. B 54, 12318 (1996).

[36] G. Seibold et al., cond-mat/9803184.

[37] F. Becca et al., Phys. Rev. B 56, 12443 (1996).

[38] A. J. Millis, H. Monien and D. Pines, Phys. Rev. B 42, 167 (1990).

[39] H. A. Mook, this Conference.

[40] N. L. Saini et al., Phys. Rev. Lett. 79, 3467 (1997).

[41] R. Hlubina and T. M. Rice, Phys. Rev. B 51, 9253 (1995).

[42] S. La Rosa et al., Solid State Commun. 104, 459 (1997).

[43] A. V. Chubukov, D. K. Morr and A. Shakhnovich, Phil. Mag. B 74, 563 (1996).

[44] A. Perali et al., Phys. Rev. B 54, 16216 (1996).

[45] J. Zaanen, this Conference.
FIG. 1. Generic phase diagram for a SC copper oxide. The meaning of the temperature scales is explained in the text.
FIG. 2. Left: $k$-space distribution of the low-laying spectral weight in the case of an effective electron-electron interaction mediated by both charge and spin fluctuations. The values of the parameters are given in the text.
FIG. 3. Gap parameter on the FS as a function of the polar angle from the point \((\pi, \pi)\).