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A non-BCS mechanism for superconductivity in underdoped cuprates via attraction between spin vortices

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Abstract – We propose a non-BCS mechanism for superconductivity in hole-underdoped cuprates based on a gauge approach to the t - J model. The gluing force is an attraction between spin vortices centered on the empty sites of two opposite Néel sublattices, leading to pairing of charge carriers. In the presence of these pairs, a gauge force coming from the single occupancy constraint induces, in turn, the pairing of the spin carriers. The combination of the charge and spin pairs gives rise to a finite density of incoherent hole pairs, leading to a finite Nernst signal as precursor to superconductivity. The true superconducting transition occurs at an even lower temperature, via a 3D XY-type transition. The main features of this non-BCS description of superconductivity are consistent with the experimental results in underdoped cuprates, especially the contour plot of the Nernst signal.



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In this letter we propose a new mechanism of superconductivity in hole-underdoped high- T_c cuprates, using the spin-charge gauge approach to the 2D t - J model, describing the CuO planes [1]. In this approach the disturbance of hole doping on the antiferromagnetic (AF) background is systematically considered, giving rise to spin vortices dressing the charge excitation (fermionic spinless holon). At the same time, due to these vortices the spin excitation (bosonic spin-1/2 spinon) acquires a finite gap, leading to a short-range (SR) AF order. The interplay of that SR order with the dissipative motion of charge carrier results in a metal-insulator crossover, a pronounced phenomenon in the underdoped cuprates. A number of peculiar features of cuprates in the normal state can be well explained within this scheme [1]. Here this approach is generalized to consider the superconducting state.

The gluing force of the superconducting mechanism is an attraction between spin vortices on two opposite Néel sublattices, centered around the empty sites (holes), and we propose a three-step scenario: At the highest crossover temperature, denoted as T_{ph} , a finite density of incoherent

holon pairs are formed. We propose to identify this temperature with the experimentally observed (upper) pseudogap (PG) temperature, where the in-plane resistivity deviates from the linear behavior. However, the holon pairing alone is not enough for superconductivity to appear. Due to the no-double occupation constraint, there is a gauge interaction between holon and spinon, through which the spin vortex attraction induces in turn the formation of spin-singlet (RVB) spinon pairs with a reduction of the spinon gap. At the intermediate crossover temperature, denoted as T_{ps} , a finite density of incoherent spinon RVB pairs are formed, which, combined with the holon pairs, gives rise to a gas of incoherent preformed hole pairs. We propose to identify this temperature with the experimental crossover corresponding to the appearance of the diamagnetic and Nernst signal. Finally, at an even lower temperature, the superconducting transition temperature T_c , both holon pairs and RVB pairs, hence also the hole pairs, become coherent. The proposed superconducting mechanism is not of the BCS-type, and it involves a gain in kinetic energy (for spinons) coming from

the J -term of the spin interactions. The main features of this non-BCS description of superconductivity are consistent with the experimental results in underdoped cuprates, especially the contour plot of the diamagnetic and Nernst signal [2,3].

The spin-charge gauge approach. – This approach relies on the following key ideas [1]: 1) We decompose the hole operator of the t - J model as $c_\alpha = h^* z_\alpha$, where h is a spinless fermionic holon, carrying charge, while z_α is a spin-1/2 bosonic spinon, carrying spin, together with an emergent slave-particle gauge field (A_μ), minimally coupled to holon and spinon, taking care of the redundant $U(1)$ degrees of freedom coming from the spin-charge decomposition. With this choice the no-double occupation constraint is automatically satisfied via the Pauli principle. 2) In 2D (and 1D) one can add a “statistical” spin flux ($e^{i\Phi^s}$) to z and a “statistical” charge flux ($e^{i\Phi^h}$) to h , provided they “compensate” each other in an appropriate way so that the product $e^{-i\Phi^h} h^* e^{i\Phi^s} z$ is still a fermion. The introduction of these fluxes in the Lagrangian formalism is materialized via the Chern-Simons gauge fields. We find the optimal charge and spin statistical fluxes in the mean-field approximation (MFA) [1]. The effect of the optimal spin flux is to attach a spin-vortex to the holon, with opposite chirality on the two Néel sublattices, the rigidity holding up a vortex being provided by the AF background. These vortices take into account the long-range distortion of the AF background caused by the insertion of a dopant hole, as first discussed in [4]; they are naturally associated with the semionic nature of the spin flux (see [1]) and are reminiscent of Laughlin’s vortices of ref. [5]. Neglecting A -fluctuations, the leading terms of the Hamiltonian can be written as

$$\begin{aligned}
H = & \sum_{\langle ij \rangle} (-t) \text{AM}_{ij} h_i^* h_j e^{i(\Phi_j^h - \Phi_i^h)} + \text{h.c.} \\
& + J(1 - h_i^* h_i - h_j^* h_j)(1 - |\text{AM}_{ij}|^2) \\
& + J h_i^* h_j^* h_j h_i |\text{RVB}_{ij}|^2, \quad (1)
\end{aligned}$$

where $\text{AM}_{ij} = z_i^\dagger e^{-i\Phi_i^s} e^{i\Phi_j^s} z_j$ is a kind of Affleck-Marston spinon parameter [6] and $\text{RVB}_{ij} = \sum_{\alpha\beta} \epsilon_{\alpha\beta} z_{i\alpha} z_{j\beta}$ is an RVB spinon singlet order parameter. 3) We use the following improved MFA: in the first term of (1) we make the MFA $\langle \text{AM}_{ij} \rangle \approx 1$, while in the second term we replace the hole density by its average, and in the normal state we neglect the third term because of being higher order in doping (δ). Notice that such MFA correctly reproduces the critical exponents of the 1D t - J model [7], when dimensionally reduced (the spin-vortices becomes kink strings in 1D). A long-wavelength treatment of the second term in (1) leads to a CP^1 spinon nonlinear σ -model with an additional term coming from the spin flux,

$$\tilde{J}(\nabla\Phi^s)^2 z^\dagger z, \quad (2)$$

where $\tilde{J} = J(1 - 2\delta)$ and $\partial_\mu \Phi^s(x) = \epsilon_{\mu\nu} \partial_\nu \sum_j (-1)^{|j|} \times \Delta^{-1}(x - j) h_j^* h_j$ with Δ the 2D lattice Laplacian. 4) In

a quenched treatment of spin vortices, self-consistently justified if their density is not too small, we derive the MF expectation value $\langle (\nabla\Phi^s)^2 \rangle = m_0^2 \approx 0.5\delta |\log \delta|$, which opens a mass gap for the spinon, consistent with AF correlation length at small δ extracted from the neutron experiments [8]. This gap is also crucial for eliminating an overcounting of low-energy degrees of freedom often encountered in slave-particle approaches, giving rise to problems in the computation of thermodynamic quantities [9]. In fact, because of the spinon gap, the low- T thermodynamics in this approach is essentially dominated by the gapless holons, while the contribution of transverse and scalar gauge fluctuations to the free energy almost canceling each other [10]. 5) In the parameter region corresponding to the PG “phase” of the cuprates¹ the optimal charge flux is π per plaquette and via Hofstadter mechanism it converts the spinless holons h into Dirac fermions with a small Fermi Surface (FS) $\epsilon_F \sim t\delta$. Their dispersion is defined in the Magnetic Brillouin Zone (MBZ), which we choose as a union of two square regions, denoted as R(ight) and L(eft) centered at $\vec{Q}^R = (\pi/2, \pi/2)$ and $\vec{Q}^L = (-\pi/2, \pi/2)$, respectively; these momenta are also the centers of the holon FS. Increasing doping or temperature one reaches the crossover line $T^* \approx t/(8\pi) |\log \delta|$ entering the “strange metal phase” (SM) of the cuprates (see footnote ¹) where the optimal charge flux per plaquette is 0 instead of π and we recover a “large” FS for the charge excitations with $\epsilon_F \sim t(1 - \delta)$. 6) Holons and spinons are coupled by the dressed gauge field A giving rise to metal-insulator crossover and overdamped resonances for holes and magnons with strongly T -dependent lifetime [1].

Let us now turn to superconductivity (SC). The SC order parameter is assumed to be RVB-like: $\Delta^c = \langle \sum_{\alpha\beta} \epsilon_{\alpha\beta} c_{\alpha i} c_{\beta j} \rangle$. In our MF treatment, neglecting gauge fluctuations it is given by the product of $\langle \sum_{\alpha\beta} \epsilon_{\alpha\beta} z_{\alpha i} z_{\beta j} \rangle$ and $\langle h_i^* h_j^* \rangle$. Hence both expectation values should be non-vanishing to have SC in MFA.

Holon pairing. – In our approach the only quartic term in (1) is the RVB-like term which is repulsive. However, spin vortices centered on holons have opposite vorticity on the two Néel sublattices and that produces a long-range attraction, previously neglected in the MFA; this will be our key attractive force. Physically it is due to the distortion of the AF background caused by the holes. This effect in the simplest form was first realized by Trugman [11] in the early days of high- T_c studies, who correctly pointed out that putting two holes next to each other on two Néel sublattices would save energy $2J$. We include this effect in MFA by introducing a term coming from the average of $z^\dagger z$ in (2):

$$\tilde{J} \langle z^* z \rangle \sum_{i,j} (-1)^{|i|+|j|} \Delta^{-1}(i-j) h_i^* h_i h_j^* h_j, \quad (3)$$

¹In the T - δ phase diagram we identify the crossover from PG to SM with the inflection point in resistivity and the broad peak in the specific heat coefficient γ .

where Δ is the 2D lattice Laplacian. In the static approximation for holons (3) describes a 2D Coulomb gas with coupling constant $g = \tilde{J}\langle z^\dagger z \rangle$, where $\langle z^\dagger z \rangle \sim (\Lambda^2 + m_0^2)^{1/2} - m_0$ with $\Lambda \approx 1$ as a UV cutoff, and charges ± 1 depending on the Néel sublattice. For 2D Coulomb gases with the above parameters a pairing appears for a temperature $T_{ph} \approx g/2\pi$, which turns out to be inside the SM “phase”². Hence the whole PG “phase” lies below T_{ph} . However, we will discuss only the SC arising from the PG phase, anticipating that extrapolation to SM phase will introduce only quantitative changes.

To implement the above ideas it is useful to distinguish holons on the two Néel sublattices, denoting them as a and b , and the two square regions R and L of the MBZ. The corresponding holons are called a^L, b^L and a^R, b^R , respectively. We also measure the momentum from the left and right centers of the two regions \vec{Q}^R, \vec{Q}^L . Since not all vortices form pairs, a finite screening effect persists and the gas of vortices still have a finite correlation length, which we denote as $\xi \approx (Jk_F)^{-1/2}$ [12]. We keep track of the screening effect by replacing (in the long wavelength limit) Δ^{-1} in (3) by an effective potential between vortices on different Néel sublattices, whose Fourier transform extends around 0 in a region of order k_F and there it is given by $V_{eff}(q) \approx g/(q^2 + \xi^{-2})$. The coupling between the L and R regions occurs only through V_{eff} , but since $V_{eff}(\vec{Q}^R + \vec{Q}^L) \ll V_{eff}(\vec{0})$ we neglect this coupling.

We define:

$$\Delta_\alpha^h(\vec{p}) = \int d^2q V_{eff}(\vec{q} - \vec{p}) \langle b_{-\vec{q}}^\alpha a_{\vec{q}}^\alpha \rangle$$

with $\alpha = R, L$. As in refs. [13,14], the d -wave pairing symmetry is composed of two p -wave pairing within the left and right Dirac cones³ in the form

$$\Delta_\alpha^h(\vec{k}) = \begin{cases} \Delta^h(k)(k_x - k_y), & \text{if } \alpha = R, \\ \Delta^h(k)(-k_x - k_y), & \text{if } \alpha = L. \end{cases} \quad (4)$$

We adopt the BCS approximation and the energy spectrum takes the following four-branch form

$$E_\alpha(\vec{k}) = \pm \sqrt{(\mu \pm |2t\vec{k}|)^2 + |\Delta_\alpha^h(\vec{k})|^2}, \quad \alpha = R, L. \quad (5)$$

Neglecting branches not crossing FS and following the procedure developed for a spin-wave attraction mechanism in refs. [13,14], we obtain for the solution of the gap equation on the FS

$$\Delta^h \equiv \Delta^h(k_F) \approx g\xi \exp \left\{ -\frac{c}{g\xi^2 k_F} \right\}, \quad (6)$$

with a constant c . As common for non-weakly coupled attractive Fermi systems, the MF temperature at which

²The expression found for T_{ph} (and T^*) cannot be extrapolated to very small doping, because of the failure of the quenched approximation for spin-vortices producing the spinon gap.

³For the reconstructed hole only 4 nodes appear in the full BZ, due to the matrix effects discussed in PG [1].

Δ^h becomes non-vanishing should be identified with the pairing temperature T_{ph} . One reinserts gauge fluctuations and recovers gauge invariance in the nodal approximation by a standard recipe as discussed in [15]: one introduces explicitly the coordinate-dependent argument of Δ^h , $\arg \Delta^h$. The first component of the nodal holon is multiplied by $e^{-i\frac{1}{2}\arg \Delta^h}$, while the second component by $e^{i\frac{1}{2}\arg \Delta^h}$. The resulting field is a slave-particle gauge-invariant, hence physical, “nodon”. One then reinserts appropriately the gauge-invariant vector field $a_\mu = A_\mu - \frac{1}{2}\partial_\mu \arg \Delta^h$ by the Peierls substitution. It is now easy to derive the low-energy effective action in the nodal approximation for a , as a variant of QED₃.

Spinon RVB pairing. – This can be materialized by the RVB term in (1) combined with the gauge attraction. The RVB term is repulsive and in MFA of holons it is irrelevant in absence of the holon pairing, but it becomes relevant as soon as holon pairing appears. In fact, as shown below, the gauge attraction then favors an RVB condensation of short-range spinon pairs. Introducing a complex RVB-Hubbard-Stratonovich gauge field $\Delta_{\langle ij \rangle}^s$ and treating the holon pair in MFA, the RVB-term in (1) becomes

$$\sum_{\langle ij \rangle} \frac{\Delta_{\langle ij \rangle}^s \Delta_{\langle ij \rangle}^{s*}}{2J|\langle h_i h_j \rangle|^2} + \Delta_{\langle ij \rangle}^{s*} \epsilon_{\alpha\beta} z_{\alpha i} z_{\beta j} + \text{h.c.} \quad (7)$$

In the continuum limit following [15] we present $\Delta_{\langle ij \rangle}^s$ as a product of a space-independent, but direction dependent modulus factor times a space-dependent phase.

Neglecting gauge and phase fluctuations and assuming rotational invariance, from (7) we derive the modified four-branches spinon dispersion:

$$\omega(\vec{k}) = \pm \sqrt{(m_s^2 - |\Delta^s|^2) + (\tilde{J}|\vec{k}| \pm |\Delta^s|)^2}, \quad (8)$$

where $m_s = \tilde{J}m_0$. The positive branches of the dispersion (8) are similar to those found in a plasma of relativistic fermions [16]. It suggests the following interpretation: if $|\Delta^s| \neq 0$ the spinon system contains a gas of RVB spinon pairs, an analogue of Coulomb neutral pairs in the relativistic plasma, either in the plasma phase, if $\langle \Delta^s \rangle = 0$ or in a condensate, if $\langle \Delta^s \rangle \neq 0$. For a finite density of spinon pairs there are two (positive energy) excitations, with different energies, but the same spin and momenta. They are given, *e.g.*, by creating a spinon up and by destructing a spinon down in one of the RVB pairs. Notice that the minimum at $\tilde{J}|\vec{k}| = |\Delta^s|$ in the lower branch is like the roton minimum in superfluid helium and has an energy lower than m_s ; it implies a backflow of the gas of spinon-pairs dressing the “bare spinon”. Hence RVB condensation would lower the spinon kinetic energy. However, to make it occur one needs the gauge contribution.

For $\langle \Delta^s \rangle \neq 0$ the global slave particle symmetry is broken from $U(1)$ to \mathbf{Z}_2 , due to the condensation of the charge $2e$ holon and RVB pairs. The Anderson-Higgs

mechanism then implies a gap $\sim |\Delta^s|$ for the gauge field A_μ . To exhibit it explicitly one calculates from (7) the gauge effective action, obtained by integrating out the spinon. The result up to quartic terms for the Lagrangian density is

$$L_{eff}(a, \partial\chi) = \frac{1}{3\pi\sqrt{(m_s^2 - |\Delta^s|^2)}} [(\partial_\mu a_\nu - \partial_\nu a_\mu)^2 + |\Delta^s|^2(2(a_0 - \partial_0\chi)^2 + (\vec{a} - \vec{\nabla}\chi)^2)] \quad (9)$$

with $\chi = \frac{1}{2}(\arg \Delta^h - \arg \Delta^s)$ a phase field, slave-particle gauge-invariant, hence physical, whose gradient describes the potential of standard magnetic vortices. From previous results one can derive the gap equation for $|\Delta^s|$ in the continuum MFA, hence neglecting χ , at $T=0$, (setting $\Lambda=1$),

$$\frac{\tilde{J}}{m_s^2} - \frac{1}{J|\langle h_i h_j \rangle|^2} = \int d\omega d\vec{k} \frac{4\tilde{J}^2 k^2}{(\omega^2 + m_s^2 + \tilde{J}^2 k^2)^2 - 4\tilde{J}^2 k^2 |\Delta^s|^2}. \quad (10)$$

In the l.h.s. of eq. (10) the first term originates from the gauge action due to the lowering of the spinon mass ($m_s \rightarrow (m_s^2 - |\Delta^s|^2)^{1/2}$), while the second term comes from the original repulsive Heisenberg term. The r.h.s. is due to spinons; the contribution of the spectrum of gauge bosons is neglected as subdominant. Whereas in the slave-boson approach the RVB pairs are made of fermions and the Heisenberg term is attractive, so the pair-formation is BCS-like, in our approach the RVB pairs are made of bosons, and in our chosen representation the Heisenberg term is repulsive, while the pair formation arises from the decrease in the free energy of spinons, via the lowering of their mass gap, induced by holon-pairing through the gauge field. Notice that the leading part of the original Heisenberg term is used to provide the AF action for the spinons, using the identity (holding for bosonic spinons) $|AM|^2 + |RVB|^2 = 1$, see eq. (1). Only the subleading term proportional to the holon-pair density is used to obtain the formation of a finite density of RVB-pairs in (10), so the derived superconductivity appears a little bit in the spirit of Laughlin's Gossamer superconductivity [17]. One easily realizes that in (10) a non-vanishing solution for $|\Delta^s|$ is possible only if the second term is sufficiently small, *i.e.*, for sufficiently large MF holon pair density. Extension to finite T is straightforward including also the contribution from the gauge bosons (with an approximate form of the effective action due to holons) and roughly estimating $|\langle h_i h_j \rangle| \sim \Delta^h k_F J^{-1}$ up to rescaling. The temperature at which $|\Delta^s|$ becomes non-vanishing is T_{ps} , not yet the true condensation temperature, T_c . Notice that from (10) it follows necessarily $T_{ps} < T_{ph}$. The δ behavior obtained from (6) and the finite T version of (10) for the crossover temperatures T_{ph} and T_{ps} , as well as the contour plot for different Δ^s are shown in fig. 1.

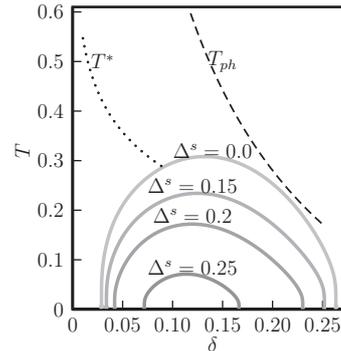


Fig. 1: The T - δ phase diagram of the mean-field gap equation of spinon for different values of MF spinon pairing Δ^s (gray lines) which could be compared with different levels of the Nernst signal; $\Delta^s = 0$ is T_{ps} . (The curves at high dopings are not quantitatively reliable as they do not take into account the crossover to SM.) The dashed line is T_{ph} , the ‘‘upper PG crossover temperature’’. The dotted line is the crossover temperature between the PG and SM phases, T^* (see also footnote ²). The temperature and Δ^s are in units of J .

SC transition. – The SC transition appears as a XY-type transition for magnetic vortices. In fact in the gauged XY-model (10) if the coefficient, $|\Delta^s|^2$, of the Anderson-Higgs mass term for a is sufficiently small the angular field χ fluctuates so strongly that it does not produce a mass gap for a_μ and $\langle e^{i\chi} \rangle = 0$. This is the Coulomb phase of the gauged XY (or Stueckelberg) model, where a plasma of magnetic vortices-antivortices appears. In the presence of a temperature gradient a perpendicular magnetic field would induce an unbalance between vortices and antivortices, giving rise to a Nernst signal. Therefore we conjecture that this phase of model (10) corresponds to the region in the phase diagram characterized by a non-SC Nernst signal and a comparison of the experimental phase diagram in [2,3] and the one derived in our model (see fig. 1) supports this idea.

For a sufficiently large coefficient $|\Delta^s|^2$, on the other hand, we are in the broken symmetry phase; the fluctuations of χ are exponentially suppressed and $\langle e^{i\chi} \rangle \neq 0$ at $T=0$ or there is quasi-condensation at $T>0$ and the gauge field is gapped. One can prove that, due to the fluctuations of the field $\arg \Delta^h$, in our approach a gapless gauge field is inconsistent with the coherence of holon pairs in PG, *i.e.*, coherent holon pairs *cannot coexist* with incoherent spinon pairs. On the other hand, due to the QED-like structure of holons-gauge action, the gauge field cannot be gapped by condensation of holon pairs alone; only the condensation of RVB spinon pairs at the same time can provide a gap to gauge fluctuations. Thus as soon as $e^{i\chi}$ (quasi-)condenses the same occurs to $\langle h_i h_j \rangle$ so that SC emerges, since the SC order parameter is $\Delta^c = \Delta^s / \langle h_i h_j \rangle$. It follows that $T_c < T_{ps}$ and the transition is of XY-type being triggered by the behavior of the XY field χ .

A few comments on comparison of the present proposal with other models on superconductivity in cuprates are in order. It is clear that our proposal differs in an essential way from the traditional BCS-Eliashberg approach [18], no matter whether the electron-phonon interaction or the antiferromagnetic fluctuations serve as the pairing glue. The SC transition occurs here in a similar way as in the preformed pairs formalism [19]. From the physical point of view, this approach is an implementation of the basic idea advocated by P.W. Anderson, attributing superconductivity to the strong correlation effects in doped Mott insulators [20–22]. It shares some similarities with other formalisms exploring the same physical idea, with, however, some substantial differences. Both in the standard slave-boson [23] and in the bosonic-RVB phase-string [24,25] approaches the Nernst effect and SC occur due to Bose-Einstein condensation (BEC) of bosonic holons. Since BEC persists for arbitrary small density in these approaches both Nernst effect and SC at $T=0$ occur as soon as the long-range AFO disappears. The same also happens in the standard “preformed pair” approaches [19], due to the persistence of condensation of pairs in the extreme BEC limit. Instead in our approach the repulsive interaction between spinons prevents the appearance of the Nernst effect below a critical doping, and the hole pairing occurs only when the holon pair density is sufficiently large to “force” the RVB spinon pairing via gauge coupling, while an even higher doping at $T=0$ is necessary to get SC. Similar “critical” dopings also appear in the phase-fluctuation approach of [26], the main physical difference with ours being in that approach nodes appear in the Nernst phase, whereas in ours a finite FS still persists and nodes appear only in the SC phase. Furthermore, in previous approaches no clear evidence of our additional crossover T^* appears, as distinct from our T_{ph} (often denoted there as T^*).

Remark: Our approach presents 3 distinct crossover lines which different authors have alternatively considered as the “pseudogap” crossover: the highest one in T , T_{ph} where holons start to pair reducing the spectral weight of the hole [27]; a second lower one, T_{ps} where incoherent hole pairs are formed, mainly affecting the magnetic properties since a finite FS still persists; and a third one, T^* , intersecting T_{ps} , where one crosses from a large to a small holon FS, with a complete suppression of the spectral weight for the holes in the antinodal region, with physical effects observable experimentally both in transport and thermodynamics [10]. The first two crossovers have mainly a magnetic origin, in the formalism described by the spin flux Φ^s , while the third one has a charge origin, due to the charge flux Φ^h , and appears only in two-dimensional bipartite lattices, see [1].

Conclusions. – We have proposed a non-perturbative pairing mechanism in high- T_c cuprates. We believe it captures the most pronounced characteristics of these compounds: a strong interplay of antiferromagnetism

and superconductivity. The same Heisenberg interaction derived from the strong on-site Coulomb repulsion is responsible for both antiferromagnetism and SC pairing: the leading term of that interaction gives rise to antiferromagnetism, while its sub-leading term providing the pairing glue due to vortex-antivortex attraction on the AF background. Compared with other proposals on pairing mechanisms, it describes this interplay in a more natural way.

Although many details of our approach are admittedly conjectural, the mechanism of SC proposed here is rather complete in its main structure and has the following appealing features: 1) It is not of simple BCS structure, in agreement with some experimental data in underdoped cuprates [28–30]. 2) SC appears only at a finite doping above the critical point where long range AF disappears. 3) It allows vortices in the normal state, as in the preformed pair scenario, supporting a Nernst signal. 4) The appearance of two positive branches in the spinon dispersion relation for a suitable spinon-antispinon attraction induces a similar structure for the magnon dispersion around the AF wave vector [27], reminiscent of the hour-glass shape of spectrum found in neutron experiments [31]. Furthermore, since the spinon gap has a maximum in δ , the energy of the magnon resonance is also expected to have a maximum. 5) In the SC state the gauge gap destroys the Reizer singularity which is responsible for the anomalous T -dependent lifetime of the magnon and electron resonances in the normal state. Therefore one expects that in the SC state, these resonances become sharper at the superconducting transition. The compositeness of excitations within the gauge approach, the holon-spinon and spinon-antispinon composites with a gauge glue coming from the single-occupancy constraint, proved to be essential in interpreting the transport [1] and thermodynamical [10] properties of cuprate superconductors, turns out to be also crucial for the superconductivity to actually occur. Furthermore, the nature of the three-step crossovers/transition needed for superconductivity, proposed in this approach: BCS-like for holon-pair formation, kinetic-energy driven for RVB spinon-pair formation and XY-like for the true superconducting transition, can be seen as specific experimental predictions of this approach. A more complete presentation of our approach to superconductivity will appear in [32].

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