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We propose a mechanism where the under-doped regime of cuprate superconductors is governed by an emergent SU(2) pseudo-spin symmetry connecting the d-wave superconducting state to a d-wave charge order. The specific SU(2) pairing fluctuations between the d-wave SC state and the charge sector lead to the formation of a new state of matter, where local patches of excitons are spontaneously generated and are responsible for the creation of Fermi arcs in the pseudo-gap phase of cuprate superconductors. In momentum space, each exciton patch is formed by a superposition of modulation wave vectors coming from the anti-nodal region. A form factor contrives the wave vectors to prefer the x and y axes, producing a typical checkerboard structure recently observed in the under-doped regime. When the amplitude of the new state vanishes, it becomes critical, and in this region a semiclassical Boltzmann calculation leads to a $T/\ln(T)$ scaling of the resistivity in three dimensions, which describes the strange-metal phase with its long-standing linear-T resistivity anomaly.

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Introduction The concept of symmetries governing the behavior of physical states is maybe the most robust in theoretical physics. From the formation of nuclei to the Higgs-Boson it has been instrumental in the determination of every emerging state in high energy physics. It would be quite remarkable that a phenomenon as complex as high temperature superconductivity is governed as well by an overall emergent symmetry. Suggestions about the existence of a pseudo-spin symmetry in the background of the physics of cuprates have been introduced since the early days of these compounds [1, 2] and has been revived over the years in different contexts. In all cases, the main simple idea is that one can rotate the d-wave superconducting (SC) state towards another state of matter quasi-degenerate in energy, like anti-ferromagnetism (AF) [3–6], a nematic state [7], or else alternating loop-currents of d-density wave [8] or π-flux phases [9, 10]. The physics is then controlled solely by the powerful constraint of the emergent symmetry which produces a vast region of the phase diagram where the fluctuations between those two states are dominant. Here, we argue that an emerging SU(2) symmetry connecting the d-wave SC state with a d-wave charge order (CO) is the main ingredient of the physics of the under-doped (UD) region, and that the very specific fluctuations associated with this symmetry are responsible for the opening of a gap in the anti-nodal (AN) zone -i.e. $(0, \pi)$ region- of the first Brillouin zone (BZ) leading to the formation of Fermi arcs in the nodal (N) zone i.e. the (π, π) region.

The SU(2) symmetry is realized explicitly in two microscopic models: the t-J model at half filling [1, 2] (for AF wave vector) and also the spin-fermion "hotspot" model [11, 12] with a linearized electron dispersion. In both models, AF interactions are at the origin of the emergent SU(2) symmetry- which is likely to be the generic case for UD cuprates. For the purpose of this study, neither the strong-coupling character of the first model, nor the vicinity of an AF quantum critical point in the second one are that crucial, but the additional symmetry induced in both. Generically when two states are degenerate - or quasi degenerate in energy, it can be accidental, or controlled by a symmetry,- here a pseudo-spin SU(2) symmetry which rotates from a d-wave SC state to a d-wave CO state [13]. The SU(2) fluctuations, typically captured within the O(4) non linear σ-model [4, 12, 14, 15], involve phase fluctuations within each state but also between the two pseudo-spin states. The presence of the SU(2) symmetry in the background of the UD region implies that at some intermediate energy scale, the two pseudo-spin states are indistinguishable. When both, d-wave SC and CO pseudo-spin states are quasi degenerate due to thermal fluctuations of the order of the pseudo-gap (PG) temperature T^* (see Fig. 1), the SU(2) pairing fluctuations trigger an instability towards a less symmetric state.

The main result of this paper is to show that the SU(2) pairing fluctuations drive a *new kind* of instability in the form of *local* patches of excitons (particle-hole excitations) which possess a checkerboard real space structure very similar to the modulations recently observed in the charge sector [16]. These solutions can be understood as local defects, or patches of the size of four to ten lattice sites, which proliferate at temperature up to T^*. They form a resonant "soup" of excitons-the Resonant Excitonic State (RES), that is responsible for the formation of Fermi arcs below T^*. Translation symmetry is broken, but in contrast to the standard charge ordering, it is broken *locally*, which does not imply global periodicity. When the RES becomes critical [17], we find that it induces a region where the resistivity vs. temperature behavior has a $T/\ln(T)$ scaling while the electron self-energy is, up to logarithms, linear in frequency, providing the expected characteristics of the strange metal (SM) phase [10, 18–21].

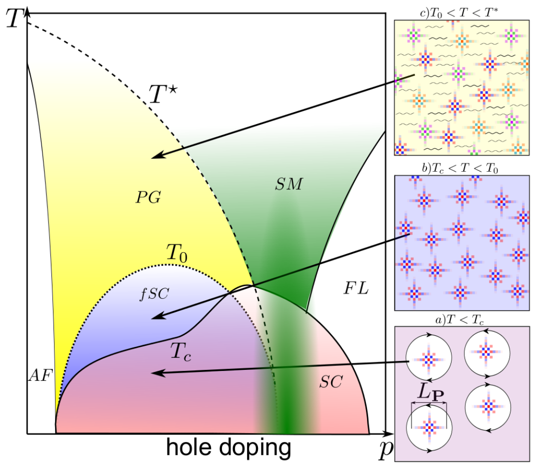

FIG. 1. Schematic phase diagram of cuprate superconductors as a function of hole doping p and temperature T. The strange metal phase is depicted in green and the Fermi liquid (FL) in white. *Insets* : a) For $T < T_c$, the excitonic patches co-exist with superconductivity and are especially visible inside the vortex-antivortex pairs. $L_{\mathbf{P}}$ is the typical size of an excitonic patch. b) For $T_c < T < T_0$, one observes the proliferation of patches and the superconducting fluctuations (fSC)- fluctuating up to a U(1) subset of the SU(2) symmetry- disappear at T_0. c) For $T_0 < T < T^*$, the patches modulations are incoherent and electron scattering through the "exciton soup" leads to dissipation which accounts for the increase of the Fermi arcs with T.

Model Our starting point is an effective action $S[\psi] = S_0[\psi] + S_1[\psi]$, where the electrons interact with pairing modes subjected to the SU(2) symmetry.

$$S_0[\psi] = -\sum_{k,\sigma} \overline{\psi}_{\sigma\mathbf{k}} G_0^{-1}(k) \psi_{\sigma\mathbf{k}}, \tag{1a}$$

$$S_1[\psi] = \sum_{k,q} \left[\Delta_{k,q} \overline{\psi}_{\uparrow\mathbf{k+q}} \overline{\psi}_{\downarrow-\mathbf{k}} + h.c. \right]. \tag{1b}$$

Here, ψ represents spin σ fermions with bare propagator $G_0^{-1}(k) = i\epsilon_n - \xi_{\mathbf{k}}$, where $\xi_{\mathbf{k}}$ represents the dispersion with subtracted chemical potential and $S_1[\psi]$ accounts for the electron-SC interaction mediated through bosonic field $\Delta_{k,q} \sim \left\langle \psi_{\mathbf{k+q}/2,\sigma} \psi_{-\mathbf{k+q}/2,-\sigma} \right\rangle$.

This bosonic field corresponds to SU(2) pairing fluctuations of a small momentum q, with the form typical of the O(4) non linear σ-model describing the thermal fluctuations between d-wave SC and CO states [11, 12, 14, 15]:

$$\pi_{k,k',q} = \langle \bar{\Delta}_{k,q} \Delta_{k',q} \rangle = \frac{\bar{\pi}_0 \left(\delta_{k,-k'} - \delta_{k,k'} \right)}{\omega_n^2 + \bar{J}_1 (\mathbf{v_k} \cdot \mathbf{q})^2 + \bar{a}_{0,\mathbf{k}}}, \tag{2}$$

where $\bar{\pi}_0$, $\bar{J}_1$, $\bar{a}_{0,\mathbf{k}}$ are non-universal parameters and $\mathbf{v_k}$ the Fermi velocity. Integrating out the bosonic field $\Delta_{k,q}$ yields a new effective two body electron-electron interaction of the form:

$$S_{\text{fin}}[\psi] = \sum_{kk'q,\sigma} \pi_{k,k',q} \; \overline{\psi}_{\sigma,\mathbf{k}} \; \psi_{\sigma,\mathbf{k'}} \; \overline{\psi}_{-\sigma,-\mathbf{k+q}} \; \psi_{-\sigma,-\mathbf{k'+q}}, \tag{3}$$

which can now be decoupled in the charge channel. Upon introducing the collective field

$$\chi_{k,k'} = T \sum_{q,\omega_n} \pi_{k,k',q} \, \langle \overline{\psi}_{\uparrow-\mathbf{k+q}} \psi_{\uparrow-\mathbf{k'+q}} \rangle, \tag{4}$$

representing a generalized charge exciton at wave vectors k and k', the self-consistent mean-field equation finally yields

$$\chi_{k,k'} = -T \sum_{q,\omega_n} \frac{\pi_{k,k',q} \; \chi_{q-k,q-k'}}{(i\omega_n - \xi_{\mathbf{q-k}})(i\omega_n - \xi_{\mathbf{q-k'}}) - |\chi_{q-k,q-k'}|^2}. \tag{5}$$

Eq. (5) is solved numerically to find the stability region for the excitonic field (particle-hole pair) at wave vectors $\mathbf{k}$ and $\mathbf{k'}$. We simplify the computation by neglecting the implicit frequency and momentum dependence of χ under the integral in order to perform the Matsubara sum at $T = 0$ exactly. It turns out that the shape of the numerical solution of χ is not much sensitive to the model parameters. For the numerical solution shown in Fig. 2b) we further restrict to a subset of coupling vectors by setting $\mathbf{k'} = -\mathbf{k}$, which, at the FS, is equivalent to the $2\mathbf{p}_\mathrm{F}$ vector. The influence of the curvature to the mass is modeled by a contribution to $\bar{a}_0$ (Eqn.(2)) of the form $\sim (\xi_{\mathbf{k+Q_0}} + \xi_{\mathbf{k}})^2$, where $\mathbf{Q_0}$ is the diagonal wave vector (depicted in blue in Fig. 2a)) [12, 22]. The SU(2) symmetry is realized when $\xi_{\mathbf{k+Q_0}} = -\xi_{\mathbf{k}}$, which favors the AN zone.

Description of the RES The solution of Eqn. (5) leads to the formation of the RES, where a wide range of $2\mathbf{p}_\mathrm{F}$ vectors give quasi-degenerate solutions as depicted in Fig.2 b). One can think of the RES as a sum of spatial modulations at vectors following the distribution $\mathbf{P} = 2\mathbf{p}_\mathrm{F}$ shown in Fig. 2d), and conjugated to a short range form factor $F_{\mathbf{k}}$:

$$\chi^{PG}_{\mathbf{r},\mathbf{r'}} = \sum_{\mathbf{P},\mathbf{k},\sigma} e^{\left[-i\mathbf{P}\cdot\frac{\mathbf{r}+\mathbf{r'}}{2}\right]} e^{\left[i\tilde{\mathbf{k}}\cdot(\mathbf{r}-\mathbf{r'})\right]} \left\langle \overline{\psi}_{\mathbf{k},\sigma} \psi_{\mathbf{k-P},\sigma} \right\rangle, \tag{6}$$

with $\tilde{\mathbf{k}} = 2\mathbf{k}-\mathbf{P}$ and $\left\langle \overline{\psi}_{\mathbf{k},\sigma} \psi_{\mathbf{k-P},\sigma} \right\rangle = \chi^{\mathbf{P}} F_{\mathbf{k}}$, and the modulation vector $\mathbf{P}$ can take indifferently any of the $2\mathbf{p}_\mathrm{F}$ wave vectors depicted in Fig.2d). The form factor $F_{\mathbf{k}}$ measures the phase space allocated to each $\mathbf{P}$ modulation, with $\left|\tilde{\mathbf{k}}\right| \leq |F_{\mathbf{k}}|$, as shown in Fig. 2c). Note that this object has a natural *local* structure due to the summation over multiple $\mathbf{P}$-wave vectors, which makes it similar to a soliton solution of the nonlinear Eq. 5. There is no global breaking of translation invariance, but rather a collection of excitonic patches breaking *locally* the translation invariance.

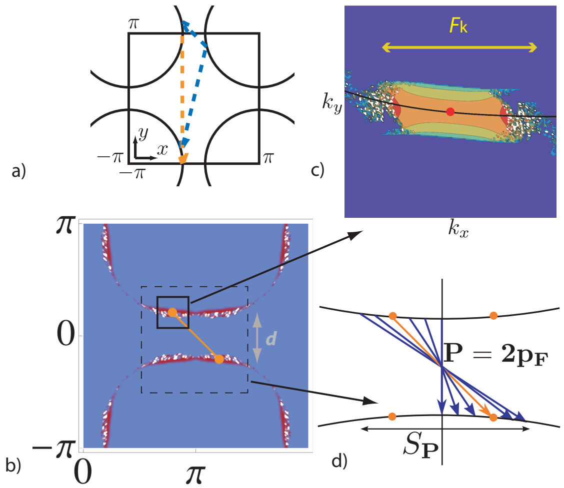

FIG. 2. (Color online) a) Schematic representation of the Fermi surface and the first BZ boundary of hole-doped cuprates with charge ordering (blue) and Umklapp (yellow) vector. b) Color plot of $|\chi_{\mathbf{k},\mathbf{k}'}|$ in the AN zone $(\pi, 0)$ of the first BZ, calculated for $\mathbf{k}' = -\mathbf{k}$. The amplitude of the electron-hole pair is maximum in the AN region along the Fermi surface, selecting a set of degenerate solutions at $2\mathbf{p}_\mathbf{F}$ wave vectors. It is suppressed in the nodal region (π, π). c) Density plot of the excitonic amplitude $|\chi_{\mathbf{k},\mathbf{k}'}|$ in the inset region of panel b), for one chosen $2\mathbf{p}_\mathbf{F}$ wave vector (depicted in yellow in d)), such that $\mathbf{k}' = \mathbf{k} + 2\mathbf{p}_\mathbf{F}$. The intensity is non zero over a width $F_\mathbf{k}$. d) Representation of the quasi-degenerate $2\mathbf{p}_\mathbf{F}$ wave vectors of the electron-hole pairs on opposed Fermi surfaces in the AN region, that give rise to the RES state. The momentum spread is $S_\mathbf{P}$.

Remarkably, each patch resembles very much a superconductor, but made of excitons (electron-hole pairs) instead of Cooper pairs (electron-electron pairs) [23]. The typical size of the pair - also called "coherence length" in the SC analogy, is typically small with $\xi \sim \hbar v_F/(\pi\chi_{AN})$ (with χ_{AN} the typical amplitude of the AN gap), which leads to the preferential formation of pair on nearest neighbor Cu-bonds. The second typical length corresponds to the size of the patches, which we call $L_\mathbf{P}$. It is controlled by the spread $S_\mathbf{P}$ of the $\mathbf{P}$ wave vectors as shown in Fig. 2b). The typical patch size of the RES $L_\mathbf{P} \sim 2\pi\hbar/S_\mathbf{P}$ is of order of a few (4 to 10) lattice sites, in good agreement with recent STM experiments [16]. Each excitonic patch also possesses an internal structure in momentum space, coming from the summation over the modulation vectors $\mathbf{P}$ in the equation (6). As shown in Fig. 2c), the modulation spreading is on the y-direction for the $(0, \pi)$- region and respectively in the x-direction in the $(\pi, 0)$- region. This leads to the typical checkerboard form, with two average modulation wave vectors around $\mathbf{Q}_x = \pm 2\pi/a(d, 0)$ and $\mathbf{Q}_y = \pm 2\pi/a(0, d)$ with $d \sim 0.3$ (see Fig.2b)), that are expected to be visible through local probes like STM [16]. The detailed structure of the modulation inside the excitonic pacthes falls beyond the scope of this manuscript and will be developped in a forthcoming publication [24].

Discussion At this stage we propose a real space picture which describes how the proliferation of patches is responsible for the opening of the PG (see Fig.1). Our theory belongs to the class of "one-gap" scenarios, for which the coherent SC state is destroyed at T_c by fluctuations [25–30], but the class of fluctuations at play here have SU(2)-character, as opposed to only the U(1)-type preformed pair scenario. Proliferation of excitonic patches destabilizes the AN part of the Fermi surface, leading to the formation of Fermi arcs whose lengths is increasing with T up to T^*, as reported in Angle Resolved Photo Emission (ARPES) (see e .g. [28, 31, 32]).

At zero temperature ($T = 0$) the whole system is in the SC state. As the temperature is raised below T_c ($0 < T < T_c$) as depicted in Fig. 1a), phase fluctuations destroy the coherent SC state. These fluctuations manifest themselves as local defects of the SC density, that are similar to vortex-antivortex pairs in classical superconductors, but with the difference that in our case, the normal phase is made of local excitonic patches which can be seen inside the vortices. An important point is the one of global phase coherence. Below T_c the global phase of modulations inside each excitonic patch can get locked with the SC one, so that the intrinsic checkerboard modulation of the patches acquires a global phase coherence. This phenomenon enhances the coherent charge density wave (CDW)-like signal experimentally observed at zero magnetic field [33].

At $T = T_c$, the gap around the node disappears, causing the loss of SC coherence, and at an intermediate temperature T_0, such that $T_c < T_0 < T^*$ the standard SC fluctuations are lost, as observed for example, by Josephson effect [34], Nernst effect [35, 36] or by study of the resistivity [37]. The transition towards the SC state can thus be described by a phase correlation length $\tilde{\lambda}_c$, which diverges at T_c and becomes very small at T^*. $\tilde{\lambda}_c^{-1}$ can be understood as the typical scale of all the phase fluctuations in the system (in the present SU(2) scenario there are three types of phase fluctuations: SC, CO, and the angle between these two).

For $T_c < T < T_0$ (Fig.1b)), the global phase coherence of the Cooper pairs and modulated excitonic patches is lost. The SC superfluid density n_s vanishes at T_0. In the regime, $T_0 < T < T^*$ shown in Fig.1c), the RES manifests itself as an incoherent set of excitonic patches. Electronic scattering through the RES induces a finite lifetime, which accounts for the size of the Fermi arcs increasing with T [28]. Above T^*, the RES disappears. Note that the incoherent and local character is likely to make this state robust to the presence of impurities [37].

Resistivity above T^* In the remaining of this paper we explore the consequences of the RES for the phase diagram of cuprates (see Fig. 1) when this mode becomes critical. We will calculate the resistivity ρ in $d = 3$ in the absence of a gap and show that it differs from the usual Fermi liquid T^2 scaling with a typical $T/\log T$ behavior. At quadratic order in the excitonic fluctuations, we obtain the following effective interaction

$$S_{\rm crit}[\psi] = \sum_{kk'qP,\sigma} \Phi_q^{\mathbf{P}}\, \overline{\psi}_{\sigma,\mathbf{k}}\, \psi_{\sigma,\mathbf{k}+\mathbf{P}+\mathbf{q}}\, \overline{\psi}_{-\sigma,\mathbf{k}'}\, \psi_{-\sigma,\mathbf{k}'-\mathbf{P}-\mathbf{q}}, \quad (7)$$

see Fig. 3a,b), with $\Phi_q^{\mathbf{P}} = \langle \chi_{-P-q}\chi_{P+q}\rangle$. The form of above interaction corresponds to a coupling with a collection of bosons (see also Eqn.(6)). The renormalized bosonic propagator follows from Dyson's equation $[\Phi_{\mathbf{q}}^{\mathbf{P}}(\Omega)]^{-1} = q^2 + m - \Pi_{\mathbf{q}}^{\mathbf{P}}(\Omega)$. Therein, the bare propagator is assumed to have Ornstein-Zernike form. The retarded bosonic polarization, $\Pi_{\mathbf{q}}^{\mathbf{P}} = \Pi'_{\mathbf{q}} + i\Pi''_{\mathbf{q}}$ in Fig. 3c), evaluated for $\mathbf{P} = 2\mathbf{p}_{\rm F}$, yields $\Pi'_{\mathbf{q}}(\Omega) = c\Big[(\Omega + q_{\parallel})\ln|\Omega + q_{\parallel}| - (\Omega - q_{\parallel})\ln|\Omega - q_{\parallel}|\Big]$ and $\Pi''_{\mathbf{q}}(\Omega) = \pi c\Big[(\Omega + q_{\parallel})\theta(-\Omega - q_{\parallel}) + (\Omega - q_{\parallel})\theta(\Omega - q_{\parallel})\Big]$. With Ω we denote real frequencies and c is a non universal factor depending on the details of the dispersion.

Next, we calculate the electronic self-energy depicted in Fig. 3c). Note that the self energy requires a summation over all ordering vectors $\mathbf{P}$. Up to logarithms, each $\mathbf{P}$-wave gives the same contribution. In the quantum critical regime, we have, the scaling behavior $\Pi'_{\mathbf{q}}(\Omega) \sim 2cq_{\parallel}\ln|\Omega|$, and $\Pi''_{\mathbf{q}}(\Omega) \sim \pi c\Omega$. We use this scaling law to evaluate the self-energy of an electron scattering through a single bosonic mode written in Matsubara form as $\Phi_{\mathbf{P}}^{-1}(i\omega_n) = \gamma|\omega_n| - v_{\parallel}q_{\parallel}\ln|\omega_n| + v_{\perp}q_{\perp}^2/2$. The evaluation is performed in $d = 3$ and at the first order in the leading singularity we obtain $\Sigma(i\epsilon_n) = i\epsilon_n/\left(4\pi v_{\parallel}v_{\perp}\ln|\epsilon_n|\right)$. We note that -with logarithmic corrections, this form is typical of a marginal Fermi liquid [18] and can account for the properties of the strange metal phase depicted in Fig. 1.

We turn now to the discussion of the relaxation time for electron-electron scattering process from a semiclassical Boltzmann treatment. The Boltzmann equation for the non equilibrium electron distribution $f_{\mathbf{k}}$ writes [38, 39]

$$\left(\frac{\partial f_{\mathbf{k}}}{\partial t}\right)_{\rm collisions} = -e\mathbf{E}\cdot\nabla_{\mathbf{k}} f_{\mathbf{k}} = -I_{ei}[f_{\mathbf{k}}] - I_{ee}[f_{\mathbf{k}}], \quad (8)$$

where e is the elementary charge, $\mathbf{E}$ a static electric field and I_{ei} respectively I_{ee} are the electron-impurity respectively electron-electron collision integrals. The electron-electron collision integral is obtained from Fermi's golden rule

$$I_{ee}[f_{\mathbf{k}}] = \frac{1}{V}\sum_{\mathbf{q}}\int_{-\infty}^{\infty} d\Omega\, {\rm Im}\Phi_{\mathbf{q}}(\Omega)\,\delta(\epsilon_{\mathbf{k}} - \epsilon_{\mathbf{k}+\mathbf{P}-\mathbf{q}} - \Omega)\times \Big[f_{\mathbf{k}}(1 - f_{\mathbf{k}+\mathbf{P}-\mathbf{q}})(1 + n_{\rm B}(\Omega)) - (1 - f_{\mathbf{k}})f_{\mathbf{k}+\mathbf{P}-\mathbf{q}}n_{\rm B}(\Omega)\Big], \quad (9)$$

with $n_{\rm B}(x) = (\exp(x/T) - 1)^{-1}$ and we drop the contribution from I_{ei}. Relaxation-time approximation amounts to set $f_{\mathbf{k}} \simeq f_{0,\mathbf{k}} - g_{\mathbf{k}}f_{0,\mathbf{k}}(1 - f_{0,\mathbf{k}})$ where f_0 is the equilibrium distribution and $g_{\mathbf{k}} = \tau e\mathbf{E}\cdot\mathbf{v}_{\mathbf{k}}/T$. In this approximation Eq. (9) becomes

$$I_{ee}[f_{\mathbf{k}}] = \frac{1}{V}\sum_{\mathbf{q}}\int_{-\infty}^{\infty} d\Omega\, {\rm Im}\Phi_{\mathbf{q}}(\Omega)\, n_{\rm B}(\Omega)\, f_{0,\mathbf{k}+\mathbf{P}-\mathbf{q}}(1 - f_{0,\mathbf{k}}) \times (g_{\mathbf{k}+\mathbf{P}-\mathbf{q}} - g_{\mathbf{k}})\,\delta(\epsilon_{\mathbf{k}} - \epsilon_{\mathbf{k}+\mathbf{P}-\mathbf{q}} - \Omega). \quad (10)$$

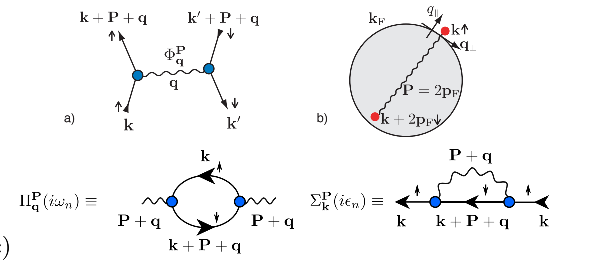

FIG. 3. (Color online) a) Graphical representation of the interaction in Eq. (7). The wavy line represents the bosonic propagator $\Phi_{\mathbf{q}}^{\mathbf{P}}$ at criticality for $|\mathbf{q}| \ll |\mathbf{P}|$. b) RES scattering between two electrons close to the FS at $\mathbf{k}$ and $\mathbf{k}+2\mathbf{p}_{\rm F}$ according to Eq. (7). c) Diagrammatic representation of the one-loop bosonic polarization $\Pi_{\mathbf{q}}$ and the fermionic self-energy $\Sigma_{\mathbf{q}}$ for the RES mode.

We see from Eq. (10) that this theory has a non-vanishing imbalance velocity factor, since for $\mathbf{q} = 0$, $(g_{\mathbf{k}+\mathbf{P}} - g_{\mathbf{k}}) \neq 0$. This implies that no additional T dependence arises from the angular part of the integral. To make the connection of the scattering time τ and resistivity ρ, we write the electrical current density as $\mathbf{J} = -2e\langle\mathbf{v}\rangle$ and note its connection to ρ via $\mathbf{J} = \rho^{-1}\mathbf{E}$. For small electric fields, $\rho \sim \tau^{-1}$ and solving above Boltzmann equation for τ yields $\tau^{-1} \sim T/\ln(T)$, such that $\rho \sim T/\ln(T)$.

Conclusion We have proposed a new mechanism for gapping out the AN region of the Fermi surface, leading to the formation of Fermi arcs below T^*. In our model, an SU(2) symmetry governs the physics of the UD region of the phase diagram, and SU(2) pairing fluctuations allow for the emergence of new local excitations in the form of excitonic patches which possess intrinsic checkerboard modulations. The proliferation of these patches leads to the formation of a metastable state -the RES, below T^* which is responsible for the formation of the PG. One striking feature of our theory, is that it reconciles the one-gap vs. two gap scenarios : fluctuations destroy the SC phase at T_c, creating Fermi arcs, as in the one-gap scenario[10, 28, 30, 40–42], but on the other hand the AN and nodal regions behave very differently upon raising the temperature (with excitonic patches forming in the AN region), which is reminiscent of two-gaps. Note that typical two gaps scenarios produce Fermi pockets instead of Fermi arcs in the nodal region [43, 44].

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[13] The SU(2) operators rotate the SC order to a CO taken as a Quadrupolar Density Wave with diagonal wave vector $\mathbf{Q_0}$. The pseudo-spin operators are

$$\eta^+ = \sum_{\mathbf{k}} c^\dagger_{\mathbf{k},\uparrow} c^\dagger_{\bar{\mathbf{k}}\downarrow}, \qquad \eta^- = (\eta^+)^\dagger, \tag{11}$$

$$\eta_z = \frac{1}{2}\sum_{\mathbf{k}} \left(c^\dagger_{\mathbf{k}\uparrow} c_{\mathbf{k}\uparrow} + c^\dagger_{\bar{\mathbf{k}},\downarrow} c_{\bar{\mathbf{k}}\downarrow} - 1 \right), \tag{12}$$

where $c^{(\dagger)}_{\mathbf{k}\sigma}$ are annihilation (creation) operators of electrons carrying momentum $\mathbf{k}$ and spin σ. These operators act on the $l = 1$ triplet representation Δ_m, $(m = -1, 0, 1)$, with $\Delta_{-1} = \frac{1}{\sqrt{2}} \sum_{\mathbf{k}} \bar{d}_{\mathbf{k}} c_{\mathbf{k}\downarrow} c_{-\mathbf{k}\uparrow}$, $\Delta_0 = \frac{1}{2}\sum_{\mathbf{k}\sigma} \bar{d}_{\mathbf{k}} c^\dagger_{\bar{\mathbf{k}}\sigma} c_{\mathbf{k}\sigma}$, and $\Delta_1 = -\frac{1}{\sqrt{2}} \sum_{\mathbf{k}} \bar{d}_{\mathbf{k}} c^\dagger_{\mathbf{k}\uparrow} c^\dagger_{-\mathbf{k}\downarrow}$, through the standard SU(2) relations $[\eta^{\pm}, \Delta_m] = \sqrt{l\,(l+1) - m\,(m \pm 1)}\Delta_{m\pm1}$, and $[\eta_z, \Delta_m] = m\Delta_m$, where $d_{\mathbf{k}}$ is the standard $d_{\mathbf{k}} = \cos(k_x) - \cos(k_y)$ with $\bar{d}_{\mathbf{k}} = (d_{\mathbf{k}} + d_{\bar{\mathbf{k}}})/2$, $\bar{\mathbf{k}} = -\mathbf{k} + \mathbf{Q_0}$ and $\overline{(-\mathbf{k})} = -(\bar{\mathbf{k}})$.

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