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Possible universal cause of high-T_c superconductivity in different metals

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Using the theory of the high temperature superconductivity based on the idea of the fermion condensation quantum phase transition (FCQPT), we show that neither the d-wave pairing symmetry, nor the pseudogap phenomenon, nor the presence of the $Cu–O_2$ planes are of decisive importance for the existence of the high-T_c superconductivity. We analyze recent experimental data on this type of superconductivity in different materials and show that these facts can be understood within the theory of superconductivity based on the FCQPT. The latter can be considered as a universal cause of the high-T_c superconductivity. The main features of a room-temperature superconductor are discussed.

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It has been a long-standing problem to develop a theory describing the high temperature superconductivity (HTS) observed in copper oxide based compounds. These compounds are extremely complex materials having a great number of competing degrees of freedom, which produce a great variety of physical properties. In turn, these properties can compete and coexist with the superconductivity hiding the understanding of the universal cause of the superconductivity. As a result, it was suggested that the unique superconducting properties in these compounds are determined by the presence of the $Cu–O_2$ planes, by the d-wave pairing symmetry, and by the existence of the pseudogap phenomena in optimal doped and underdoped cuprates (see e.g. [1 – 3]).

However, recent studies of quasiparticle tunneling spectra of cuprates have revealed that the pairing symmetry may change from d to s-wave symmetry, depending on the hole, or electron, doping level [4]. Then the high temperature s-wave superconductivity has been observed in electron doped infinite layer cuprates with a sharp transition at $T = 43$ K, without a pseudogap [5]. Therefore, we can conclude that the d-wave symmetry of pairs and the pseudogap phenomena are not necessary parts of HTS. Recent studies of HTS in chemically doped A_3C_{60} [6] have shown that the presence of the $Cu–O_2$ planes is also not a necessary condition for the existence of HTS. Moreover, it was recently demonstrated that strongly overdoped cuprates Tl-2201 obey the Wiedemann-Franz law perfectly [7]. This fact, manifesting that excitations, which carry heat, also carry charge e, shows that these quasiparticles have the same properties as Landau quasiparticles. This imposes serious constrains upon the possible theories of HTS. For instance, this fact demonstrates that there is no spin-charge separation [7] as was suggested in one group of proposals [8, 9]. On the other hand, it is well known that the superconducting transition temperature $T_c^{\alpha\gamma}(x)$ of the oxides appears to pass through a bell-shaped curve as a function of the hole (or electron) mobile charge carrier density x (see e.g. [3]). Here α denotes the material, say C_{60} or Tl-2201, etc., and γ denotes hole or electron doping. It is essential to note that the shape of the functions $T_c^{\alpha\gamma}(x)$ is similar in different samples, so that there exist empirical formulas, which connect the transition temperature $T_c^{\alpha\gamma}(x)$ with the carrier concentration x. For example, in case of Tl-2201, see e.g. [7], $T_c(x) = T_c^M(1 - 82.6(p - 0.16)^2)$. Here T_c^M is the maximum value of the transition temperature. Thus, we can use a simple approximation

$$T_c^{\alpha\gamma}(x) = T_1^{\alpha} T_2^{\gamma}(x_1 - x)x, \qquad (1)$$

where the coefficients T_1^{α} and T_2^{γ} define the transition temperature T_c for a given hole (or electron) metal, with x obviously going continuously to zero at the insulator-metal transition. It directly follows from Eq. (1) that

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the transition temperature reaches its maximum value T_c^M at the optimal doping level $x_{\rm opt}$

$$T_c^M = T_c^{\alpha\gamma}(x_{\rm opt}) = T_1^\alpha T_2^\gamma \left(\frac{x_1}{2}\right)^2. \qquad (2)$$

The general shape of the function $T_c(x)$ points to the fact that the generic properties of HTS are defined by the 2D charge (electron or hole) strongly correlated liquid rather then by solids which hold this liquid. As to the solids, they arrange the presence of the pseudogap phenomena, s or d-wave pairing symmetry, the electron-phonon coupling constant defining T_c, the variation region of x, and so on.

It is well-known that consideration of strongly correlated liquids is close to the problem of systems with a big coupling constant, which is the most important in many-body physics. One solution of this problem has been offered by Landau theory of normal Fermi liquids, which introduce the notion of quasiparticles and parameters characterizing the effective interaction between them [10]. As a result, the Landau theory has removed high energy degrees of freedom and kept a sufficiently large number of relevant low energy degrees of freedom to treat the liquid's low energy properties. Usually, it is assumed that the breakdown of the Landau theory is defined by the Pomeranchuk stability conditions and occurs when the Landau amplitudes being negative reach critical value. The new phase, in which the stability conditions are restored can in principle be again described in the framework of the theory. However, it was demonstrated recently [11] that the Pomeranchuk's conditions do not cover all possible limitations: the overlooked one is connected to the situation when, at temperature $T = 0$, the effective mass can become infinitely big. It has been demonstrated that such a situation, leading to profound consequences, can take place when the corresponding amplitude, being positive, reaches a critical value, producing a completely new class of strongly correlated Fermi liquids with the fermion condensate (FC) [11, 12]. This liquid is separated from a normal Fermi liquid by the fermion condensation quantum phase transition (FCQPT) [13]. In such a case we are dealing with the strong coupling limit where an absolutely reliable answer, based on pure theoretical first principle ground, cannot be given. Therefore, the only way to verify whether the FC occurs is to consider experimental facts, which can signal the existence of such a state. It seems to us that there are strong indications of the existence FC in high-T_c superconductors.

In this Letter we show that the new experimental facts mentioned above are the desirable evidences and that they can be explained within the theory of HTS based on the FCQPT [14]. We also describe the main features of corresponding room-temperature superconductors.

Let us start with a brief consideration of the general properties of a two-dimensional electron liquid in the superconducting state, when the system has undergone FCQPT [13, 14]. At $T = 0$, the ground state energy $E_{gs}[\kappa(\mathbf{p}), n(\mathbf{p})]$ is a functional of the order parameter of the superconducting state $\kappa(\mathbf{p})$ and of the quasiparticle occupation numbers $n(\mathbf{p})$ and is determined by the well-known equation of the weak-coupling theory of superconductivity (see e.g. [15])

$$E_{gs} = E[n(\mathbf{p})] + \int \lambda_0 V(\mathbf{p}_1, \mathbf{p}_2)\kappa(\mathbf{p}_1)\kappa^*(\mathbf{p}_2)\frac{d\mathbf{p}_1 d\mathbf{p}_2}{(2\pi)^4}. \qquad (3)$$

Here $E[n(\mathbf{p})]$ is the ground-state energy of a normal Fermi liquid, $n(\mathbf{p}) = v^2(\mathbf{p})$ and $\kappa(\mathbf{p}) = v(\mathbf{p})\sqrt{1 - v^2(\mathbf{p})}$. It is assumed that the pairing interaction $\lambda_0 V(\mathbf{p}_1, \mathbf{p}_2)$ is weak. By minimizing E_{gs} with respect to $\kappa(\mathbf{p})$ we obtain the equation connecting the single-particle energy $\varepsilon(\mathbf{p})$ to $\Delta(\mathbf{p})$

$$\varepsilon(\mathbf{p}) - \mu = \Delta(\mathbf{p})\frac{1 - 2v^2(\mathbf{p})}{2\kappa(\mathbf{p})}. \qquad (4)$$

The single-particle energy $\varepsilon(\mathbf{p})$ is determined by the Landau equation, $\varepsilon(\mathbf{p}) = \delta E[n(\mathbf{p})]/\delta n(\mathbf{p})$ [10], and μ is the chemical potential. The equation for the superconducting gap $\Delta(\mathbf{p})$ takes form

$$\Delta(\mathbf{p}) = -\int \lambda_0 V(\mathbf{p}, \mathbf{p}_1)\kappa(\mathbf{p}_1)\frac{d\mathbf{p}_1}{4\pi^2}. \qquad (5)$$

If the coupling constant $\lambda_0 \to 0$, the maximum value Δ_1 of the superconducting gap tends to zero, $\Delta_1 \to 0$, and Eq. (4) reduces to that proposed in [11]

$$\begin{gathered}\varepsilon(\mathbf{p}) - \mu = 0, \text{ if } \ \kappa(\mathbf{p}) \neq 0, (0 < n(\mathbf{p}) < 1);\\ p_i \leq p \leq p_f \in L_{FC}.\end{gathered} \qquad (6)$$

At $T = 0$, Eq. (6) defines a new state of a Fermi liquid with FC, for which the modulus of the order parameter $|\kappa(\mathbf{p})|$ has finite values in L_{FC} range of momenta $p_i \leq p \leq p_f$ occupied by FC, and $\Delta_1 \to 0$ in L_{FC} [11, 13]. Such a state can be considered as superconducting, with an infinitely small value of Δ_1 so that the entropy of this state is equal to zero. This state, created by the quantum phase transition, disappears at $T > 0$ [13].

It follows from Eq. (6) that the system brakes into two quasiparticle subsystems: the first one in the L_{FC} range is occuied by the quasiparticles with the effective mass $M^*_{FC} \propto 1/\Delta_1$, while the second by quasiparticles with finite mass M^*_L and momenta $p < p_i$. It is seen from Eq. (6) that at the point of FCQPT $p_f \to p_i \to p_F$, the effective mass is as large as $1/\Delta_1$. Equation (6) acquires nontrivial solutions at $r_s = r_{FC}$ and FCQPT takes place if the Landau amplitudes depending on the density are positive and sufficiently large, so that the potential energy prevails over the kinetic energy and the transformation of the Fermi step function $n(p) = \theta(p_F - p)$ into the smooth function defined by Eq. (6) becomes possible [11].

Looking now for a simple situation where we can solve Eqs. (4) and (5) analytically, we assume that $\lambda_0 \neq 0$ and is small, so that we can employ the standard theory of superconductivity. In that case Δ_1 becomes finite, leading to finite value of the effective mass M^*_{FC} in L_{FC}, which can be obtained from Eq. (4) [13]

$$M^*_{FC} \simeq p_F \frac{p_f - p_i}{2\Delta_1}, \quad (7)$$

while the effective mass M^*_L is disturbed weakly. Here p_F is the Fermi momentum. It follows from Eq. (7) that the quasiparticle dispersion can be presented by two straight lines characterized by the effective masses M^*_{FC} and M^*_L respectively. These lines intersect near the binding energy E_0 of electrons, which defines an intrinsic energy scale of the system:

$$E_0 = \varepsilon(\mathbf{p}_f) - \varepsilon(\mathbf{p}_i) \simeq \frac{(p_f - p_i)p_F}{M^*_{FC}} \simeq 2\Delta_1. \quad (8)$$

Equations (7) and (8) lead to the following general result for the maximum value Δ_1 coming from the contribution of FC [14]

$$\Delta_1 \simeq \beta\varepsilon_F \frac{p_f - p_i}{p_F} \ln(1 + \sqrt{2}), \quad (9)$$

where $\varepsilon_F = p_F^2/2M^*_L$ is the Fermi energy and β is the dimensionless coupling constant. If the FC contribution to Δ_1 becomes comparatively small then Δ_1 is given by the well known relation, being proportional to the exponent of $(-1/\beta)$, $\Delta_1 \propto \exp(-1/\beta)$.

In fact, as it is seen from Eqs. (4) and (5), a Fermi liquid with FC is absorbed by the superconducting phase transition and never exhibits the dispersionless plateau associated with $M^*_{FC} \to \infty$. As a result, a Fermi liquid beyond the point of FCQPT can be described by two types of quasiparticles characterized by two finite effective masses M^*_{FC} and M^*_L respectively, and by the intrinsic energy scale E_0 [13, 14]. It is reasonably reliable to suggest that we have returned to the Landau theory by integrating out high energy degrees of freedom and introducing the quasiparticles. The sole difference between the Landau Fermi liquid and Fermi liquid undergone FCQPT is that we have to increase the number of relevant low energy degrees of freedom by adding both, new type of quasiparticles with the effective mass M^*_{FC}, given by Eq. (7), and the energy scale E_0 given by Eq. (8). We have also to bear in mind that the properties of these new quasiparticles of the Fermi liquid with the FC cannot be separated from the properties of the superconducting state, as follows from Eqs. (7), (8) and (9). We may say that the quasiparticle system in the range L_{FC} becomes very "soft" and is to be considered as a strongly correlated liquid. On the other hand, the system's properties and dynamics are dominated by a strong collective effect having its origin in FCQPT and determined by the macroscopic number of quasiparticles in the range L_{FC}. Such a system cannot be disturbed by the scattering of individual quasiparticles and has the features of a quantum protectorate [13, 16]. On the other hand, as soon as the energy scale $E_0 \to 0$, the system is driven back into the normal Landau Fermi liquid [14].

Recent studies of photoemission spectra discovered an energy scale in the spectrum of low-energy electrons in cuprates, which manifests itself as a kink in the single-particle spectra [17, 18]. The spectra in the energy range (-200—0) meV can be described by two straight lines intersecting at the binding energy $E_0 \sim (50 - 70)$ meV [18]. Then, in underdoped copper oxides, there exists the pseudogap phenomenon and the shape of single-particle excitations strongly deviates from that of normal Fermi liquid (see, e.g. [1]). In the highly overdoped regime slight deviations from the normal Landau Fermi liquid are observed [7, 19]. All these peculiar properties are naturally explained within a model proposed in [13, 14, 20] and allow one to relate the doping level, or the charge carriers density x, regarded as the density of holes (or electrons) per unit area, to the density of the Fermi liquid with FC. We assume that x_{FC} corresponds to the highly overdoped regime at which FCQPT takes place, and introduce the effective coupling constant $g_{\text{eff}} \sim (x_{FC} - x)/x_{FC}$. According to the model, the doping level x at $x \leq x_{FC}$ in metals is related to $(p_f - p_i)$ in the following way:

$$g_{\text{eff}} \sim \frac{(x_{FC} - x)}{x_{FC}} \sim \frac{p_f - p_i}{p_F} \sim \frac{p_f^2 - p_i^2}{p_F^2}. \quad (10)$$

Consider now two possible types of the superconducting gap $\Delta(\mathbf{p})$ given by Eq. (5) and defined by interaction $\lambda_0 V(\mathbf{p}, \mathbf{p}_1)$. If this interaction is dominated by a phonon-mediated attraction, the even solution of Eq. (5) with an s-wave, or an $s{+}d$ mixed waves, will have the lowest energy. If the pairing interaction $\lambda_0 V(\mathbf{p}_1, \mathbf{p}_2)$ is a combination of attractive and repulsive interaction, a d-wave odd superconductivity can take place(see e.g.

[21]). But both an s-wave even and d-wave odd symmetries lead to approximately the same values of the gap Δ_1 in Eq. (9) [14]. Therefore, the non-universal pairing symmetries in HTS are likely the results of the pairing interaction, and the d-wave pairing symmetry cannot be considered as absolutely necessary for HTS existence, in accord with experimental findings [5]. If only d-wave pairing would exists, the transition from superconducting gap to pseudogap can take place, so that the superconductivity will be destroyed at the temperature T_c, with the superconducting gap being smoothly transformed into the pseudogap, which closes at some temperature $T^* > T_c$ [20, 22]. In the case of s-wave pairing we can expect that the pseudogap phenomena do not exists, in accordance with the experimental observation (see [5] and references therein).

We turn now to consideration of the maximum value of the superconducting gap Δ_1 as a function of the mobile charge carriers density x. Being rewritten in terms of x and x_{FC} related to the variables p_i and p_f by Eq. (10), Eq. (9) takes the form

$$\Delta_1 \propto \beta(x_{FC} - x)x. \tag{11}$$

Here we take into account that the Fermi level $\varepsilon_F \propto p_F^2$, the density $x \propto p_F^2$, and thus, $\varepsilon_F \propto x$. We can reliably assume that $T_c \propto \Delta_1$ because the empirically obtained simple bell shaped curve $T_c(x)$ should reflect only a smooth dependence. Then, instead of Eq. (11) we have

$$T_c^{\alpha\gamma} \propto \beta^\alpha \beta^\gamma (x_{FC} - x)x. \tag{12}$$

In Eq. (12), we made the natural change $\beta = \beta^\alpha\beta^\gamma$ since the coupling constant β is fixed by the properties of metal in question. Now it is seen that Eq. (12) coincides with Eq. (1) producing the universal optimal doping level $x_{\rm opt} = x_{FC}/2 = x_1/2$ in line with the experimental facts.

Let us make here few remarks. In Ref.[23], HTS was reported in such fullerides as C_{60}, $CHBr_3$, and $CHCl_3$. In the case of hole metals, the T_c^M ratios were measured for $C_{60}/CHBr_3$-$C_{60}/CHCl_3$-C_{60}, and the same ratios were also found for the respective electron metals [23]. It follows from Eq. (12) that among these hole doped fullerides, the T_c^M ratios for $C_{60}/CHBr_3$-$C_{60}/CHCl_3$-C_{60} have to be the same as in the case of the respective electron doped fullerides because the factor β^γ, related to hole or electron doping, drops out of the ratios. We believe that our calculations present evidences to the correctness of these experimental facts. Then, it was also shown that a Fermi liquid with FC, which exhibits strong deviations from the Landau Fermi liquid behavior, is driven into the Landau Fermi liquid by applying a magnetic field B. This field-induced Landau Fermi liquid behavior provides an $A + BT^2$ dependence of the resistivity ρ. A re-entrance into the strongly correlated regime, when ρ is a linear function of the temperature T, is observed if the magnetic field B decreases to some critical value B_{cr} [24]. We expect that the similar behavior can be observed in case of strongly overdoped high-temperature compounds at the temperatures $T \to 0$. If the superconductivity is depressed by the critical field B_c, the T^2 behavior of the resistivity ρ can be restored at fields $B_{cr} > B_c$. We assume that this behavior was observed in overdoped Tl-2201 compounds at millikelvin temperatures [7, 25].

As an example of the implementation of the previous analysis let us consider the main features of a room-temperature superconductor. As follows from Eq. (9), $\Delta_1 \sim \beta\varepsilon_F \propto \beta/r_s^2$. Noting that FCQPT takes place in three-dimensional systems at $r_s \sim 20$ and in quasi two-dimensional systems at $r_s \sim 8$ [26], we can expect that Δ_1 of 3D system comprises 10% of the corresponding maximum value of 2D superconducting gap, reaching values as high as 60 meV for underdoped crystals with $T_c = 70$ [27]. On the other hand, it is seen from Eq. (9) that Δ_1 can be even larger, $\Delta_1 \sim 75$ meV, and one can expect $T_c \sim 300$ K in the case of the s wave pairing as it follows from the simple relation $2T_c \simeq \Delta_1$. In fact, we can safely take $\varepsilon_F \sim 300$ meV, $\beta \sim 0.5$ and $(p_f - p_i)/p_F \sim 0.5$. Thus, we can conclude that a possible room-temperature superconductor has to be a quasi two-dimensional structure and has to be an s-wave superconductor in order to get rid of the pseudogap phenomena, which tremendously reduces the transition temperature. The density x of the mobile charge carriers must satisfy the condition $x \leq x_{FC}$ and be flexible to reach the optimal doping level. It is worth noting that the coupling constant β is to be sufficiently big because FC giving rise to the order parameter $\kappa(\mathbf{p})$ does not produce by itself the gap Δ. For instance, the coupling constant can be enhanced by intercalation or by some kind of a disorder. It is pertinent to note, that FCQPT can take place in three-dimensional metals at the usual metallic densities, as in heavy-fermion metals, when the effective mass is sufficiently large [24]. In that case, the potential energy can easily prevail over the kinetic energy leading to FCQPT at sufficiently high densities. Then, we can expect the appearance of HTS if the coupling constant is large to ensure high values of the superconducting gap.

In summary, we have shown by a simple self-consistent analysis that the general features of the shape of the critical temperature T_c as a function of the density x of the mobile carriers in metals and the value of the optimal doping $x_{\rm opt}$ can be understood within the framework of the theory of high-T_c superconductivity based on the FCQPT. We have demonstrated that neither the d-wave pairing symmetry, nor the pseudogap phenomenon, nor the presence of the Cu–O_2 planes are of importance for the existence of the high-T_c superconductivity. Our theory explains the experimental observation that a strongly correlated Fermi liquid in heavily overdoped cuprates transforms into a normal Landau liquid. The main features of a room-temperature superconductor have also been outlined.

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