On the superconductivity of 2D system with arbitrary carrier
density in external magnetic field

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Abstract

Selfconsistent equations which describe the order parameter and chemical potential behaviour in 2D metallic system as functions of external magnetic field, $H$, temperature, $T$, and carrier density, $n$, are obtained. It is shown that for the case of the local pairs (low carrier density and negative chemical potential $\mu$) the derivative $dH_c(T)/dT$ at $T = T_c$ is essentially less then that for the system with Cooper pairs (high carrier density and positive chemical potential $\mu \gg T_c$). It is found that in magnetic field satisfying the quantum limit criterion the system is characterized by non-trivial inhomogeneous order parameter which can exist at rather high temperatures.

Introduction

Despite of the great deal of efforts directed toward solving the problem, the nature of high temperature superconductivity (HTSC) retains to be unknown. However, those efforts have not been done in vain. At present one can observe a common consent about the most characteristic properties of HTSC [1, 2]. Such properties are: a) quasi–two–dimensional behaviour of the conductivity in the normal phase; b) relatively low (at least considerably lower than in ordinary metals) density of carriers. There are also many other distinctive properties which may appear to be very important for understanding the nature of HTSC. But it is unlikely to describe all the characteristic features of HTSC in the framework of the only theoretical approach. So, very often the role of one of them is studied and then the total picture is created by means of summing up all the results in some ”artistic” way.

One of the very important and intensively studied questions in the theory of superconductivity (SC) is the question about a dependence of SC properties of a system on the carrier density. The first who emphasized on the relevance of such a problem was Legget [3]. Later this problem was also discussed in Refs. [4, 5, 6]. In Ref. [7] the existence of a crossover driven by the carrier density was clearly established. In that paper the authors have considered a three dimensional model with a local attractive interaction and shown that the gradual transition from low to high carrier density is accompanied by a crossover from local pair SC to the SC with Cooper pairs. As for HTSC, an intermediate case seems to be realized in it.

In this paper we continue the analysis of the papers [3] and [8].
2 Model and general discussion

As in the papers \[6, 8\], we shall focus our attention only on the qualitative side of the problem. So, for our purposes, we can take the simplest Hamiltonian which describes a system of charge carriers (fermions) with a local attraction between them:

\[
\hat{H} = \int d^2r \hat{H}(r),
\]

where \( A_j \) is a vector potential corresponding to a magnetic field, \( H \), (note that in two dimensions magnetic field is pseudoscalar function), \( e \) and \( m \) are charge and mass of the carriers, respectively, \( g > 0 \) is an attractive coupling between these carriers and \( \mu \) is the chemical potential.

Introducing Nambu notations for fermion field \( \Psi = (\Psi^\dagger, \Psi) \) \[9\], we can rewrite (2) in a more convenient form:

\[
\hat{H}(r) = -\Psi^\dagger \tau_3 \left( \frac{D^2}{2m} + \mu \right) \Psi + g \Psi^\dagger \Psi \Psi^\dagger \tau_+ \Psi,
\]

where \( \tau_3, \tau_+ \equiv (\tau_1 + i\tau_2)/2, \tau_- \equiv (\tau_1 - i\tau_2)/2 \) are Pauli matrices and the covariant derivative has the form \( D_j \equiv \partial_j - ie\tau_3 A_j/c \). Then the partition function is expressed through the Hamiltonian as:

\[
Z \equiv \exp(-\Omega/T) = \text{Tr} \exp(-\hat{H}/T),
\]

Since the thermodynamical potential, \( \Omega = \Omega(V, T, \mu) \), is the function of the chemical potential (besides the volume and the temperature) and we are interested in a dependence of all physical values on the density of the carriers, it is necessary to write down the second equation which links the density, \( n \), and the chemical potential, \( \mu \):

\[
n = \frac{1}{V} \frac{\partial \Omega}{\partial \mu}.
\]

The expression for the partition function with the Hamiltonian (3) can be represented in a path integral form:

\[
Z = \int [d\Psi^\dagger d\Psi] \exp \left\{ - \frac{\beta}{\hbar} \int_0^\beta d\tau \int d^2r \left( \Psi^\dagger \partial_\tau \Psi + \hat{H}(r) \right) \right\}, \quad \beta \equiv \frac{1}{T},
\]

where Grassmann variables \( \Psi(\tau; r) \) satisfy the antiperiodic boundary condition \( \Psi(0; r) = -\Psi(\beta; r) \). After usual introduction of an auxiliary scalar field by means of the Habbard–Stratanovich trick:

\[
Z = \int [d\Phi^\dagger d\Psi d\Phi^*] \exp \left\{ - \frac{\beta}{\hbar} \int_0^\beta d\tau \int d^2r \left( \left| \Phi \right|^2 + \right.ight.

+ \Psi^\dagger \left[ \partial_\tau - \tau_3 \left( \frac{D^2}{2m} + \mu \right) + \tau_+ \Phi + \tau_- \Phi^* \right] \Psi \left. \right\},
\]

where \( A_j \) is a vector potential corresponding to a magnetic field, \( H \), (note that in two dimensions magnetic field is pseudoscalar function), \( e \) and \( m \) are charge and mass of the carriers, respectively, \( g > 0 \) is an attractive coupling between these carriers and \( \mu \) is the chemical potential.
we can perform (at least formally) the integration over Grassmann variables and represent the result through the effective action:

$$S_{\text{eff}}(\Phi, \Phi^*) = -Tr Ln G^{-1} + \frac{1}{g} \int \int d^2 r |\Phi|^2,$$

depending only on the scalar field $\Phi$. However, the next integration over scalar field we can perform only approximately, for example, using so called saddle point formalism. In our problem the "saddle point" is defined by the equation:

$$\frac{\delta S_{\text{eff}}(\Phi, \Phi^*)}{\delta \Phi^*(\tau; r)} = \text{tr} \left[ G(\tau, \tau; r, r) \tau_+ \right] + \frac{\Phi}{g} = 0,$$

where Green’s function, $G$, of interacting fermions is defined as the solution of the equation:

$$\left[ -\partial_{\tau_1} + \tau_3 \left( \frac{D^2}{2m} + \mu \right) - \tau_- \Phi - \tau_+ \Phi^* \right] G(\tau_1, \tau_2; r_1, r_2) = \delta(\tau_1 - \tau_2) \delta(r_1 - r_2),$$

with the boundary condition $G(\tau_1 + \beta, \tau_2; r_1, r_2) = -G(\tau_1, \tau_2; r_1, r_2)$.

Let $\bar{\Phi}$ be a solution to the equation (9), then the partition function in the next to the leading approximation takes a form of the Gauss type path integral:

$$Z = \exp(-\bar{S}_{\text{eff}}) \int [d\Phi d\Phi^*] \exp \left[ -\int d\tau_1 \int d\tau_2 \right.$$

$$\left. \int d^2 r_1 \int d^2 r_2 \bar{\Phi}^*(\tau_1; r_1) \Gamma^{-1}(\tau_1, \tau_2; r_1, r_2) \bar{\Phi}(\tau_2; r_2) \right],$$

where $\bar{S}_{\text{eff}} \equiv S_{\text{eff}}(\bar{\Phi}, \bar{\Phi}^*)$. New field $\bar{\Phi}$ describes fluctuations and $\Gamma$ is its propagator:

$$\Gamma^{-1}(\tau_1, \tau_2; r_1, r_2) = \frac{1}{g} \delta(\tau_1 - \tau_2) \delta(r_1 - r_2) +$$

$$+ \text{tr} \left[ G(\tau_1, \tau_2; r_1, r_2) \tau_+ G(\tau_2, \tau_1; r_2, r_1) \tau_- \right].$$

For the partition function in this approximation we have

$$Z = \exp(-\bar{S}_{\text{eff}} - Tr Ln \Gamma^{-1}).$$

So, as follows from Eqs.(4) and (5), the expression for the carrier density takes the form of a sum of two different terms:

$$n = \frac{T}{V} Tr [\tau_3 G] - \frac{T}{V} \frac{\partial}{\partial \mu} (Tr Ln \Gamma^{-1}).$$

The first term is expressed through the fermion propagator and the second through the propagator of scalar fluctuations. Having such a representation, we shall refer to the first term as fermion part in density of the carriers and to the second as boson one. The fact that the ratio of fermions and composite bosons (fluctuations can also be interpreted as a field of two-fermion composite particles) is determined by the dynamics is the main virtue of this model.

At this point we have, in principle, a closed selfconsistent system of two equations which completely describe a dependence of the order parameter and the chemical potential as functions of the temperature and carrier density.
3 On the critical line equation

As was indicated in the previous section Eqs. (9) and (14) completely describe the behaviour of the order parameter and the chemical potential of the system as a function of "external" parameters $T$, $B$ and $n$. But analyzing these equations in general case is an unsolvable problem. So we restrict ourselves only to analyzing the behaviour of the system near the critical line (related to the second type phase transition). Such a choice simplifies the problem considerably since when the system is in the nearcritical region we have a natural small value (order parameter) and as a consequence we can apply perturbation theory in this value.

In order to simplify the problem further we assume that the solution to the equation (9) does not depend on "time" coordinate $\tau$. In this case, as follows from Eq. (10), Green's function depends only on the difference of "time" variables ($\tau_1 - \tau_2$). So, after taking into account the boundary conditions, Green’s function can be expanded into the Fourier series:

$$G(\tau_1 - \tau_2; \mathbf{r}_1, \mathbf{r}_2) = T \sum_{n=-\infty}^{\infty} G_n(\mathbf{r}_1, \mathbf{r}_2) \exp[-i\omega_n(\tau_1 - \tau_2)], \quad \omega_n = \pi T(2n + 1).$$

(15)

Solving the equation for the Green's function (10) in linear approximation in order parameter and substituting it into (9), we come to the following integral equation:

$$\bar{\Phi}(\mathbf{r}) = \int d^2 \mathbf{r}_1 K_1(\mathbf{r}, \mathbf{r}_1) \bar{\Phi}(\mathbf{r}_1),$$

(16)

$$K_1(\mathbf{r}, \mathbf{r}_1) = -gT \sum_{n=-\infty}^{\infty} tr \left[ G_n^{(0)}(\mathbf{r}, \mathbf{r}_1) \tau_- G_n^{(0)}(\mathbf{r}_1, \mathbf{r}) \tau_+ \right],$$

(17)

where $G_n^{(0)}(\mathbf{r}, \mathbf{r}_1)$ denotes the Green’s function at $\Phi = 0$ which can be easily found using the Schwinger proper time method [11]. Here we write down this function without deriving (for details see appendix A in Ref. [11]):

$$G_n^{(0)}(\mathbf{r}_1, \mathbf{r}_2) = \exp\left(-\frac{i}{2l^2} \tau_3(x_1y_2 - y_1x_2)\right) G_n^{(hom)}(\mathbf{r}_1 - \mathbf{r}_2),$$

(18)

$$G_n^{(hom)}(\mathbf{r}) = \frac{1}{2\pi l^2} \exp\left(-\frac{\mathbf{r}^2}{4l^2}\right) \sum_{j=0}^{\infty} L_j \left(\frac{\mathbf{r}^2}{2l^2}\right) \frac{1}{i\omega_n - \tau_3[\omega_H(j + 1/2) - \mu]].$$

(19)

Strictly speaking, Eq. (17) which has the form of a spectral problem is valid only on the critical line where the order parameter equals zero. The condition indicating the existence of a solution to this equation determines the critical line in the problem. If we were going to leave the critical line, we would need to take into account the next nonlinear term. In this paper we are interested only in the behaviour on the critical line.

After substituting the explicit expression for the Green’s function into (17) we get:

$$\bar{\Phi}(\mathbf{r}) = \int d^2 \mathbf{r}_1 K_1^{(hom)}(\mathbf{r} - \mathbf{r}_1) \exp\left(i\frac{l}{2}(x_1y_1 - y_1x_1)\right) \bar{\Phi}(\mathbf{r}_1),$$

(20)

where $l$ is the magnetic length and the homogeneous kernel, $K_1^{(hom)}$, determined by

$$K_1^{(hom)}(\mathbf{r}, \mathbf{r}_1) = -gT \sum_{n=-\infty}^{\infty} tr \left[ G_n^{(hom)}(\mathbf{r}) \tau_- G_n^{(hom)}(\mathbf{r}) \tau_+ \right].$$

(21)
As was shown in [12, 13], Eq. (20) has the exact solution:

\[ \Phi(r) = \Delta \exp \left( -\frac{r^2}{2l^2} \right), \]  

(22)

where in the simplest case \( \Delta \) is constant. After substituting (22) into (20) we come to the condition determining the critical line:

\[
1 + \frac{g T l^2}{(2\pi)^3} \sum_{n=-\infty}^{\infty} \int \int \int d^2 K d^2 k \exp \left( -\frac{K^2 l^2}{2} \right) \cdot \text{tr} \left[ G_n^{(hom)} \left( \frac{K}{2} + k \right) \tau_- G_n^{(hom)} \left( -\frac{K}{2} + k \right) \tau_+ \right] = 0.
\]

(23)

The second equation which links the density and the chemical potential is the following:

\[
n = \frac{T}{(2\pi)^2} \sum_{n=-\infty}^{\infty} \int d^2 k \text{tr} \left[ \tau_3 G^{(hom)}(k) \right] - \frac{T}{V} \frac{\partial}{\partial \mu} (Tr L n \Gamma_0^{-1}).
\]

(24)

where \( \Gamma_0 \) is defined by the expression similar to (12) but with Green’s functions at \( \Phi = 0 \).

Eqs. (23) and (24) selfconsistently describe the critical parameters \( \mu(T) \) and \( H_{c2}(T) \) at every given carrier density. Below we shall consider two limiting cases of low and high magnetic fields.

### 4 Low magnetic field

This case is realized when the cyclotron frequency is much lower then the temperature, \( \omega_H \ll T_c \). Expanding Eqs. (23) and (13) in power series in \( l^{-1} \) and keeping only the lowest term in magnetic field, we get the equation:

\[
\frac{4\pi}{gm} = \int_0^W \frac{du}{u - \mu} \tanh \frac{u - \mu}{2T} - \frac{4T}{ml^2} \sum_{n=0}^{\infty} \int_{-\mu}^{W-\mu} \frac{du}{u} \frac{d^2(u + \mu)}{\omega_n^2 + u^2},
\]

(25)

where \( W \) is the band width.

The critical temperature, \( T_c \), when the field is switched off is defined by the equations:

\[
\frac{4\pi}{gm} = \int_0^W \frac{du}{u - \mu_c} \tanh \frac{u - \mu_c}{2T_c}, \quad \mu_c = \mu(T_c),
\]

(26)

which have been analyzed in detail in Ref.[8]. Substituting the expression for the coupling (24) into (25) and taking the limit \( W \to \infty \), we come to:

\[
\int_0^\infty du \left[ \frac{\tanh(u - \mu)/2T}{u - \mu} - \frac{\tanh(u - \mu_c)/2T_c}{u - \mu_c} \right] = \frac{2T}{ml^2} \sum_{n=0}^{\infty} \left[ \frac{1}{\omega_n^2} + \frac{\pi \mu}{2\omega_n^3} + \frac{\mu \arctan(\mu/\omega_n)}{\omega_n^3} \right],
\]

(27)
Approaching the critical line at low but nonzero magnetic field, we should come to the value of the temperature close to \( T_c \). So, expanding the left hand side of the Eq.(27) in \((T - T_c)\) we can find the slope of the critical line at \( T_c \):

\[
\frac{e}{4mc} \left( \frac{dH_{c_2}}{dT} \right)_{T_c} = \frac{1 + \tanh(\mu_c/2T_c) - (\mu_c/T_c)(\partial \mu_c/\partial T_c) \tanh(\mu_c/2T_c)}{1 + [7\zeta(3)/2\pi^2](\mu_c/T_c) + (|\mu_c|/2T_c) \int_0^\infty (du/u^3)(u - \tanh u)}.
\]

Even though the slope of the critical line is given by the explicit expression (28) this does not solve the problem completely because, as was indicated above, we should add also the second equation (24) linking the chemical potential with the carrier density. From that equation we should also find the derivative \( \partial \mu_c/\partial T_c \equiv (\partial \mu/\partial T)_{T_c} \) which is calculated at constant \( n \).

Since the analysis of the second equation for arbitrary \( \mu \) is quite a difficult problem we shall consider only limiting cases.

1. Local pairs. At first we consider the case of low density of carriers and strong interaction between fermions, \( \epsilon_F \ll |\epsilon_b| \) where \( |\epsilon_b| = 2W \exp(-4\pi/mg) \) is the energy of bound state in two particle problem and \( \epsilon_F \equiv n\pi/m \). Analyzing the propagator of the scalar field describing fluctuations we come to the conclusion that in the lowest approximation in \( l^{-1} \) the propagator should be taken at zero magnetic field (for details see [11]). Assuming that \( \mu_c < 0 \) and \( |\mu_c| \gg T_c \) (the equivalence of these conditions and the inequality stated above can be easily established), it is not difficult to find the expression for carrier density in this limiting case:

\[
n = -\frac{2mT_c}{\pi} \ln \left[ 1 - \exp \left( \frac{2\mu_c - \epsilon_b}{T_c} \right) \right] + O \left( \exp \left( -\frac{|\mu_c|}{T_c} \right) \right),
\]

where we omitted exponentially small terms.

In principle, Eqs.(28) and (29) completely describe the critical line slope at \( T_c \), if the critical temperature at zero field is nonzero and the density of carriers is a finite constant. However, it can be easily shown that equations (28) and (29) are consistent only at \( T_c = 0 \). This fact is in full agreement with the general statements about the role of fluctuations in low dimensional models with short–range interaction where, as is known, there is no room for a nontrivial order parameter [14]. This problematic situation is not actual in real HTSC since they are not strictly two–dimensional, they are only quasi–two–dimensional. So, we come to the place where we are not able to use a strictly two–dimensional model. In order to avoid this obstacle we assume that the model will be treated correctly if we take into account quasi–two–dimensional character of the model only into the equation for carrier density which is extremely sensitive to the number of dimensions and leave the equation (28) without changes.

Since in three dimensions the density is measured in different units, our enlargement of the phase volume of the system should be accompanied by a redefinition of the “two–dimensional” density in the form: \( n^{2D} = n^{3D} K_z^{-1} \) where \( n^{3D} \) is the density found in the 3D anisotropic model and \( K_z \) is some characteristic momentum. Chosing \( K_z^2/4M \simeq T \) where \( M \) is an effective mass in \( z \)–direction, we get:

\[
n = \frac{2mT}{\pi^2} \int_0^\infty du \frac{u^{1/2}}{\exp[u + (\epsilon - 2\mu)/T - 1]}.
\]
Now we find the slope of the critical line without problems:

\[-\frac{e}{4mc} \left( \frac{dH_{c2}}{dT} \right)_{T_c} \simeq \left( \frac{\mu_c}{T_c} \right)^2 \exp \left( -\frac{|\mu_c|}{2T_c} \right), \quad |\mu_c| = \frac{|\epsilon_b|}{2} \gg T_c \sim \epsilon_F.\]  

(31)

Obtained expression shows that the slope of the critical line in the case of composite bosons is very small what is in qualitative agreement with the result of the paper [13].

2. Cooper pairs. When the density of carriers is high or the attraction between fermions is relatively weak, i.e. \( \epsilon_F \gg \epsilon_b \), we come to the SC of Cooper pairs. In terms of chemical potential this condition is equivalent to the inequality \( \mu_c \gg T_c \). Since the fluctuations are suppressed in any three–dimensional model and the mean field approximation is reliable when the density of carriers is high enough, we can assume that the same statement is correct in our "quasi–two–dimensional" model (though it cannot be shown in a direct way). With such an assumption the bosonic term in Eq.(24) can be omitted and we get:

\[ n = \frac{mT}{\pi} \ln \left( 1 + \exp \frac{\mu_c}{T} \right) \sim \frac{\mu_c m}{T}, \]  

(32)

and for the slope of the critical line:

\[-\frac{e}{4mc} \left( \frac{dH_{c2}}{dT} \right)_{T_c} \simeq \frac{2\pi^2}{7} \frac{T_c}{\zeta(7) \mu_c}, \quad \mu_c \simeq \epsilon_F \gg T_c = \frac{\gamma}{\pi} \sqrt{2|\epsilon_b|\epsilon_F}, \]  

(33)

where \( \gamma \) is the Euler constant.

So, the results received in this section show that the behaviour of the critical line, in particular its slope at \( T_c \) essentially depends on the density of carriers in the system. At low density the slope is small and with increasing the density the slope becomes bigger (all other parameters are kept constant). Note that after achieving some maximum value the slope turns to decreasing since in the regime of Cooper pairs it decreases with increasing of the carrier density.

5 Strong magnetic field

When a magnetic field is so strong that it satisfies the condition, \( \omega_H \gg T_c \), we come to the so called quantum limit. It is obvious that in a very strong magnetic field, only the lowest Landau level in fermion spectrum plays an important role. So, omitting all those terms in fermion Green's function which correspond to the higher Landau levels, we get:

\[ G^{(hom)}_{n}(k) = \frac{2 \exp(-k^2l^2)}{i\omega_n - \tau_3[i\omega_H/2 - \mu]}. \]  

(34)

This, in its turn, means that the density of states at the lowest level should be higher than the density of carriers in the system. It is also clear that the higher density of carriers is in the system, the stronger magnetic field should be applied in order to achieve the quantum limit. The second condition when the approximation of the lowest Landau level is reliable is connected with the value of coupling constant. When the attraction between fermions is
very strong the substitution of the approximate expression (34) instead of an exact Green’s function will not be satisfactory since any strong interaction could considerably change the spectrum.

At last, all analysis in this section is performed in the framework of the mean field approximation. We believe that the consideration taking into account fluctuations in quantum limit would have been worthwhile problem, but as is easily seen from all our consideration, at first one need to find some simple formalism for treating quasi–two–dimensional models which is free from the obstacles incorporated into all two–dimensional models. It could be more realistic to consider an anisotropic 3D model, however as we can judge such a problem is much more difficult than ours and could be solved only numerically.

Since the expression for Green’s function is so simple, we come at once to the equation describing the critical line:

$$T = \frac{\omega_H - 2\epsilon_F}{2 \ln(\omega_H/|\epsilon_b|) \ln(\omega_H/\epsilon_F - 1)}.$$  \hspace{1cm} (35)

where we assumed that $n < eH/\pi c$ (or $\epsilon_F < \omega_H$). The obtained result is similar to the result of paper [10].

6 Conclusion

In this paper we derived the system of selfconsistent equations describing second type phase transition in the simplest ”quasi–two–dimensional” system with local attraction between fermions in an external magnetic field. We have shown that when the density of carriers in the system is relatively low the equation, linking the chemical potential and the carrier density, is nontrivial ($\mu \neq \epsilon_F$) and, as a result, this is reflected in the behaviour of the critical line (its slope at $T_c$ is small). In the opposite case of high density of carriers the system can be described in the framework of the mean field approximation (as in BCS model).

Our results are in qualitative agreement with the results obtained in [13] for composite Bose particles though analytical behaviour is different.

The research was supported in part by the grant No.K5O100 of the International Science Foundation and by the International Soros Science Education Program (ISSEP) through the grant No.PSU052143.

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