Field-Theoretical description of the crossover between BCS and BEC in a non-Fermi superconductor

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Abstract

Using the field-theoretical methods we studied the evolution from BCS description of a non-Fermi superconductor to that of Bose-Einstein condensation (BEC) in one loop approximation. We showed that the repulsive interaction between composite bosons is determined by the exponent $\alpha$ of the Anderson propagator in a two dimensional model. For $\alpha \neq 0$ the crossover is also continous and for $\alpha = 0$ we obtain the case of the Fermi liquid.
INTRODUCTION

The problem of the crossover from BCS superconducting state to a Bose-Einstein condensate (BEC) of local pairs\[1\]–\[3\] became very important in the context of high temperature superconductors (HTSC). While at the present time there is no quantitative microscopic theory for the occurrence of the superconducting state in the doped antiferromagnetic materials, it is generally accepted that the superconducting state can be described in terms of a pairing picture. The short coherence length ($\xi \sim 10-20$ A) increased the interest for the problem\[4\]–\[14\] because it showed that the BCS equations of highly overlapping pairs, or the description in terms of composite bosons cannot describe the whole regime between weak and strong coupling. The mean field method developed by different authors\[3\],\[4\],\[6\],\[14\] and solved analytically in two and three dimension, and the Ginzburg-Landau description\[7\],\[8\],\[13\] showed that the evolution between the two limits is continuous, no singularities during this evolution appearing. The zero temperature coherence length in the framework of field-theoretical method has been in Ref.\[12\]. The problem of the BCS-BEC crossover in arbitrary dimension $d$ using the field-theoretical method has been extensively discussed in Refs.\[15\]–\[18\] where the chemical potential, the number of condensed pairs and the repulsive interaction between pairs have been calculating using the analogy with the field-theoretical description of superfluidity.

In this paper we apply this method to study the crossover problem for a non-Fermi superconductor described by the Anderson model\[19\],\[20\] (See also Refs.\[20\]–\[28\]), to study the crossover between weak coupling and strong coupling. The paper is organized as follows. In Section II we study the weak coupling model for a $d=2$ non-Fermi superconductor. Section III contains the strong coupling limit. To make the paper self-contained we present in Appendix the Lagrangian formalism for the superfluid phase following Refs.\[15\],\[16\]. The results are discussed in Section IV.

WEAK COUPLING LIMIT

The BCS-like model for the non-Fermi system is described by the Lagrangian

$$\mathcal{L} = \psi_\uparrow^\dagger G_0^{-1} \psi_\uparrow + \psi_\downarrow^\dagger (G_0^{-1})^* \psi_\downarrow - \lambda_0 \psi_\uparrow^\dagger \psi_\downarrow^\dagger \psi_\downarrow \psi_\uparrow$$

(1)

where the normal state is described by the Green function

$$G_0(p, \omega) = \frac{g(\alpha)}{\omega^\alpha (\omega - \xi(p) + i\delta)^{1-\alpha}}$$

(2)

where $\omega_c \leq \omega \leq \omega_c$, $g(\alpha) = \pi \alpha / (2 \sin(\pi \alpha / 2))$ and $\lambda_0 < 0$ is the coupling constant describing the attraction between electrons. The Green function given by Eq. (2) contain a cut-off $\omega_c$. 

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and the exponent $\alpha$. This form has been proposed first by Anderson to describe the 2D non-Fermi properties of the superconducting state.

If we introduce the two-component fermionic field

$$\Psi = \begin{pmatrix} \psi_\uparrow \\ \psi_\downarrow \end{pmatrix} \quad \Psi^\dagger = \begin{pmatrix} \psi_\uparrow^\dagger \\ \psi_\downarrow^\dagger \end{pmatrix}$$

the non-interacting part of the Lagrangian (1) is

$$L_0 = \Psi^\dagger \begin{pmatrix} G_0^{-1} & 0 \\ 0 & (G_0^{-1})^* \end{pmatrix} \Psi$$

In order to calculate the partition function

$$Z = \int D\Psi^\dagger D\Psi exp\left[ i \int_x \mathcal{L} \right]$$

we will transform the interaction contribution from the Lagrangian (1) as

$$\exp\left[ -i \lambda_0 \int_x \psi_\uparrow^\dagger \psi_\downarrow^\dagger \psi_\downarrow \psi_\uparrow \right] = \int D\Delta^\dagger D\Delta \exp\left[ -i \int_x \left( \Delta^\dagger \psi_\downarrow \psi_\uparrow + \psi_\uparrow^\dagger \psi_\downarrow^\dagger \Delta - \frac{1}{\lambda_0} \Delta^\dagger \Delta \right) \right]$$

where $f_x = \int dt d^4x$ as in Ref. 16–18, and $\Delta = \lambda_0 \psi_\downarrow \psi_\uparrow$ is a bosonic field. The partition function defined by Eq. (5) will be expressed using Eq. (3) in a bilinear form as

$$Z = \int D\Psi^\dagger D\Psi \int D\Delta^\dagger D\Delta \exp\left( \frac{i}{\lambda_0} \int_x \Delta^\dagger \Delta \right) \exp\left[ i \int_x \Psi^\dagger \begin{pmatrix} G_0^{-1} & -\Delta \\ -\Delta^\dagger & (G_0^{-1})^* \end{pmatrix} \Psi \right]$$

Performing the integral over the Grassmann fields the partition function becomes

$$Z = \int D\Delta^\dagger D\Delta \exp\left( i S_{eff}[\Delta^\dagger, \Delta] + \frac{1}{\lambda_0} \int_x \Delta^\dagger \Delta \right)$$

where $S_{eff}[\Delta^\dagger, \Delta]$ is the one loop effective action, which can be written as

$$S_{eff}[\Delta^\dagger, \Delta] = -i Tr \ln \left( \begin{pmatrix} f(\alpha)(p_0 - \xi(p))^{1-\alpha} & -\Delta \\ -\Delta^\dagger & f(\alpha)(p_0 + \xi(p))^{1-\alpha} \end{pmatrix} \right)$$

where $f(\alpha) = \omega_c^\alpha g^{-1}(\alpha)$ and the trace $Tr$ has been used according to the meaning from Ref. 16.

In the mean field approximation the integral from Eq. (8) can be performed using the solution given by the saddle point and for $T \neq 0$ the critical temperature $T_c$ will be obtained as

$$T_c(\alpha) = \omega_D \left[ \frac{D(\alpha)}{C(\alpha)} \right]^{1/2\alpha} \left[ 1 - \frac{1}{A(\alpha) D(\alpha)} \frac{1}{\lambda_0} \left( \frac{\omega_c}{\omega_D} \right)^{2\alpha} \right]^{1/2\alpha}$$
where \( A(\alpha) = g^2(2^2 \alpha \sin(\pi(1-\alpha))/\pi, C(\alpha) = \Gamma^2(\alpha)[1 - 2^{1-2\alpha}]\zeta(\alpha), D(\alpha) = \Gamma(1 - 2\alpha)\Gamma(1 - \alpha)/(2\alpha \Gamma(1 - \alpha)), \Gamma(x) \) being the Euler’s gamma function. This expression is valid only in the limit \( 0 < \alpha < 0.5 \) and a positive critical temperature implies for the coupling constant the condition \( |\lambda_0| > \lambda_c \), with \( \lambda_c = (\omega_c/\omega_D)^\alpha/(A(\alpha)D(\alpha)) \). We have to mention that the critical temperature obtained in Eq. (10), calculated also in Ref. [29] is different from the one obtained in Ref. [22,24,25], and it is easy to show that it gives the exact BCS result in the limit \( \alpha \to 0 \).

If we consider the effective action as

\[
S_{eff}[\Delta^\dagger, \Delta] = -iTr \ln \begin{pmatrix} f(\alpha)(p_0 - \xi(p))^{1-\alpha} & 0 \\ 0 & f(\alpha)(p_0 + \xi(p))^{1-\alpha} \end{pmatrix} - iTr \ln \left[ 1 - \left| \frac{\Delta}{f(\alpha)(p_0^2 - \xi^2(p))^{1-\alpha}} \right|^2 \right]
\]

and the system as space time independent the partition function can be written as

\[
Z = Z_0 \exp \left[ \frac{i}{\lambda_0} \bar{\Delta}^\dagger \bar{\Delta} \right]
\]

\( Z_0 \) containing the non-interacting contribution, and we get for the renormalized coupling constant \( \lambda \) the expression

\[
\frac{1}{\lambda} = \frac{1}{\lambda_0} + \frac{i}{f^2(\alpha)} \int \frac{d^2p}{(2\pi)^2} \int \frac{dp_0}{2\pi} \frac{1}{(p_0^2 - \xi^2(p))^{1-\alpha}}
\]

Using the integral

\[
\int_{k_0} \frac{1}{(k_0^2 - E^2 + i\eta)^l} = i(-1)^l \sqrt{\pi} \frac{\Gamma(l - 1/2)}{\Gamma(l)} \frac{1}{E^{2l-1}}
\]

we calculated \( \lambda \) as

\[
\frac{1}{\lambda} = \frac{1}{\lambda_0} + \frac{1}{\lambda_1}
\]

where

\[
\lambda_1 = - \frac{4\pi \alpha g^{-2}(\alpha)}{\cos(\pi(\alpha - 1)) B(1/2, 1/2 - \alpha)} \left( \frac{\omega_c}{\omega_D} \right)^{2\alpha}
\]

\( B(x, y) \) being the Euler beta function \( B(x, y) = \Gamma(x)\Gamma(y)/\Gamma(x + y) \). The expression given by Eq. (15) is positive for \( \alpha < 1/2 \). The new coupling constant \( \lambda \) has to be also negative in order to have superconductivity (\( \lambda < 0 \)) and this condition is satisfied if \( |\lambda_0| < \lambda_1 \). If we consider also the condition \( \lambda_c < |\lambda_0| \) we get the general condition for the bare coupling constant \( \lambda_0 \), \( \lambda_c < |\lambda_0| < \lambda_1 \), which is satisfied for \( 0 < \alpha < 0.5 \).

We mention that for the weak coupling limit \( \lambda_0 \to 0^- \) the BCS limit studied in Ref. [24] is reobtained, but we also showed that the critical constant calculated from the critical temperature is smaller than \( \lambda_1(\alpha) \), which also satisfies condition \( \lambda_1(\alpha \to 0) = 0 \).

In the limit \( \lambda_0 \to -\infty \), called the strong coupling limit, we expect an important effect of the non-Fermi character of the electrons in the coupling constant.
STRONG COUPLING LIMIT

In this limit we consider $\Delta(x) = \bar{\Delta} + \tilde{\Delta}(x)$ and consider the action $S_{eff}[\tilde{\Delta}^\dagger, \tilde{\Delta}]$ obtained from Eq. (9) as

$$S_{eff}[\tilde{\Delta}^\dagger, \tilde{\Delta}] = -i Tr \ln \left[ 1 + \hat{G}_0 \hat{\Delta} \right]$$

(16)

where

$$\hat{G}_0^{-1} = \left( \begin{array}{cc} f(\alpha)(p_0 - \xi(p))^{1-\alpha} & -\bar{\Delta} \\ -\bar{\Delta}^\dagger & f(\alpha)(p_0 + \xi(p))^{1-\alpha} \end{array} \right)$$

(17)

$$\hat{\Delta} = \left( \begin{array}{cc} 0 & -\bar{\Delta} \\ -\bar{\Delta}^\dagger & 0 \end{array} \right)$$

(18)

which can be written as

$$S_{eff}[\tilde{\Delta}^\dagger, \tilde{\Delta}] = -i Tr \sum_{l=1}^{\infty} \frac{1}{l} \left[ \hat{G}_0 \hat{\Delta} \right]^l$$

(19)

with

$$\hat{G}_0(p_0, p) = \frac{1}{f^2(\alpha)(p_0^2 - \xi^2(p))^{1-\alpha} - |\Delta|^2} \left( \begin{array}{cc} 0 & -\bar{\Delta} \\ -\bar{\Delta}^\dagger & 0 \end{array} \right)$$

(20)

We are interested in quadratic terms in $\tilde{\Delta}$ and we will take the approximation

$$S_{eff}[\tilde{\Delta}^\dagger, \tilde{\Delta}] = S_{eff}^{(2)}(0) + S_{eff}^{(2)}(\mathbf{q})$$

(21)

which contains the quadratic contributions. The first term in Eq. (21) has the form

$$S_{eff}^{(2)}(0) = \frac{1}{2} i Tr \frac{1}{f^2(\alpha)(p_0^2 - \xi^2(p))^{1-\alpha} - |\Delta|^2} \left( \bar{\Delta}^2 \bar{\Delta}^\dagger + \bar{\Delta}^\dagger \bar{\Delta} \bar{\Delta}^\dagger + 2 |\Delta|^2 |\bar{\Delta}|^2 \right)$$

(22)

$$+ \frac{1}{2} i Tr \frac{1}{f^2(\alpha)(p_0^2 - \xi^2(p))^{1-\alpha} - |\Delta|^2 |\bar{\Delta}|^2}$$

(23)

which will be approximated, taking in the dominator $\bar{\Delta} \approx 0$ as

$$S_{eff}^{(2)}(0) \approx \frac{1}{2} i Tr \frac{1}{f^4(\alpha)(p_0^2 - \xi^2(p))^{2(1-\alpha)}} \left[ \bar{\Delta}^2 \bar{\Delta}^\dagger + \bar{\Delta}^\dagger \bar{\Delta} \bar{\Delta}^\dagger + 2 |\Delta|^2 |\bar{\Delta}|^2 \right]$$

(24)

the last term giving no contribution to the renormalized coupling constant. Following the same approximation we calculated $S_{eff}^{(2)}(\mathbf{q})$ as

$$S_{eff}^{(2)}(\mathbf{q}) = \frac{1}{2} i Tr \frac{1}{f^2(\alpha)(p_0 - \xi(p))^{1-\alpha}(p_0 + \mathbf{q}_0 + \xi(p + \mathbf{q}))^{1-\alpha}} \bar{\Delta} \bar{\Delta}^\dagger$$

(25)
From Eqs. (24) and (25) we have
\[
\mathcal{L}^{(2)}(0) = -\frac{B(1/2, 3/2 - 2\alpha)}{4\pi f^4(\alpha)}(2m)^{3-4\alpha} \\
\times \int \frac{d^2p}{(2\pi)^2 (p^2 + m\varepsilon_a)^{3-4\alpha}} \left[ \bar{\Delta}^2 \bar{\Delta}^\dagger \bar{\Delta}^\dagger + \bar{\Delta} l^2 \bar{\Delta} \bar{\Delta} + 2|\bar{\Delta}|^2 |\bar{\Delta}|^2 \right]
\] (26)

and
\[
\mathcal{L}^{(2)}(q) = -\frac{\sin(\pi(1 - \alpha))}{4\pi f^2(\alpha)} B(\alpha, \alpha) m^{1-2\alpha} \\
+ \frac{m}{16\pi^2 f^2(\alpha)} \frac{2^{3-4\alpha}}{1 - 2\alpha} \frac{B(1/2, 3/2 - 2\alpha)}{\varepsilon_a^{2-4\alpha}} \frac{1}{(p^2 + m\varepsilon_a + q_0 m + q^2/4)^{1-2\alpha}} \\
+ \frac{m}{16\pi^2 f^2(\alpha)} \frac{\sin(\pi(1 - \alpha))}{2\alpha} \frac{B(\alpha, \alpha)}{(\varepsilon_a + q_0 + q^2/4m)^{2\alpha}} \bar{\Delta} \bar{\Delta}^\dagger \\
+ \frac{m}{16\pi^2 f^2(\alpha)} \frac{\sin(\pi(1 - \alpha))}{2\alpha} \frac{B(\alpha, \alpha)}{(\varepsilon_a - q_0 + q^2/4m)^{2\alpha}} \bar{\Delta}^\dagger \bar{\Delta}
\] (27)

The integrals from Eqs. (26) and (27) can be performed using the formula
\[
\int \frac{1}{(p^2 + A^2)^N} = \frac{\Gamma(N - d/2)}{(4\pi)^d/2\Gamma(N)} (A^2)^{N-d/2}
\]
and we obtain
\[
\mathcal{L}^{(2)} = -\frac{m}{16\pi^2 f^2(\alpha)} \frac{2^{3-4\alpha}}{1 - 2\alpha} \frac{B(1/2, 3/2 - 2\alpha)}{\varepsilon_a^{2-4\alpha}} \left[ \bar{\Delta}^2 \bar{\Delta}^\dagger \bar{\Delta}^\dagger + \bar{\Delta} l^2 \bar{\Delta} \bar{\Delta} + 2|\bar{\Delta}|^2 |\bar{\Delta}|^2 \right]
\]
\[
+ \frac{m}{16\pi^2 f^2(\alpha)} \frac{\sin(\pi(1 - \alpha))}{2\alpha} \frac{B(\alpha, \alpha)}{(\varepsilon_a + q_0 + q^2/4m)^{2\alpha}} \bar{\Delta} \bar{\Delta}^\dagger \\
+ \frac{m}{16\pi^2 f^2(\alpha)} \frac{\sin(\pi(1 - \alpha))}{2\alpha} \frac{B(\alpha, \alpha)}{(\varepsilon_a - q_0 + q^2/4m)^{2\alpha}} \bar{\Delta}^\dagger \bar{\Delta}
\] (28)

Using the approximation
\[
(\varepsilon_a \pm q_0 + \frac{q^2}{4m})^{2\alpha} \approx (\varepsilon_a)^{2\alpha} + 2\alpha(\varepsilon_a)^{2\alpha-1} \left( \pm q_0 + \frac{q^2}{4m} \right)
\]
and using the notation
\[
\bar{\Psi} = \begin{pmatrix} \bar{\Delta} \\ \bar{\Delta}^\dagger \end{pmatrix}
\] (29)

we obtain from Eq. (28)
\[
\mathcal{L}^{(2)} = \frac{m}{16\pi^2 f^2(\alpha)} \sin(\pi(1 - \alpha)) B(\alpha, \alpha) \varepsilon_a^{2\alpha-1} \frac{1}{2} \bar{\Psi}^\dagger M \bar{\Psi}
\] (30)

where
\[
M = \begin{pmatrix} q_0 - \frac{q^2}{2m} & -\mu_0 \\
-\mu_0 & -q_0 - \frac{q^2}{2m} & -\mu_0 \end{pmatrix}
\] (31)
\( m_b = 2m \) being the boson mass and \( \mu_0 \) the chemical potential

\[
\mu_0 = \frac{1}{f^2(\alpha)} \frac{2^{2-4\alpha} B(1/2, 3/2 - 2\alpha)}{(1 - 2\alpha) \sin(\pi(1 - \alpha)) B(\alpha, \alpha)} |\tilde{\Delta}|^2 \varepsilon_a^{2\alpha - 1} \tag{32}
\]

The velocity \( c_0 \) of the sound mode is

\[
c_0^2 = \frac{\mu_0}{m_b} = \frac{1}{f^2(\alpha)} \frac{2^{2-4\alpha} B(1/2, 3/2 - 2\alpha)}{(1 - 2\alpha) \sin(\pi(1 - \alpha)) B(\alpha, \alpha) m} |\tilde{\Delta}|^2 \varepsilon_a^{2\alpha - 1} \tag{33}
\]

and the repulsive interaction \( \lambda_{0b} \) between pairs is

\[
\lambda_{0b}(\alpha) = \frac{\pi^2}{m (1 - 2\alpha) |\sin(\pi(1 - \alpha)) B(\alpha, \alpha)|^2} 2^{4-4\alpha} B(1/2, 3/2 - 2\alpha) \tag{34}
\]

We mention that \( \lim_{\alpha \to 0} \lambda_{0b}(\alpha) = 2\pi/m \) a result identical to the result obtained in Ref. 16 for the two dimensional case.

**I. RESULTS AND DISCUSSIONS**

Using the field-theoretical methods we studied the crossover between BCS and BEC in a non-Fermi liquid. The weak coupling case leads to the same results as in the mean field like models\(^{25 - 28}\). In the strong coupling limit we showed that the pairs form a Bose gas with a repulsive coupling constant which is controled by \( \alpha \).

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**APPENDIX:**

In this appendix we briefly present the field-theoretical formulation of the Bogoliubov theory for the interacting Bose gas. The system of an interacting Bose gas is described by the Lagrangean

\[
\mathcal{L} = \frac{1}{2} \left\{ \Phi^* [i \partial_0 - \hat{\varepsilon} + \mu_0] \Phi - \lambda_0 |\Phi|^4 + c.c \right\} \tag{A1}
\]

where \( \Phi \) is the operator corresponding to the complex scalar field, \( \hat{\varepsilon} \) the kinetic energy operator, \( \mu_0 \) the chemical potential and \( \lambda_0 \) the bare repulsive coupling constant. The model described by Lagrangean (A1) has a global symmetry to the transformation

\[
\Phi(x) \to e^{ia} \Phi(x) \tag{A2}
\]
where $a$ is a constant defining the transformation. At $T = 0$ this symmetry is spontaneously broken and it is associated with the occurrence of the superfluid phase. According to the Goldstone theorem the dynamical restoring of the symmetry implies the occurrence of the Goldstone bosons and it can be included in the theory taking
\[
\Phi(x) = e^{i\theta(x)} \left( \Phi_0 + \tilde{\Phi}(x) \right) \tag{A3}
\]
However, in the simple mode which is important for the pair description we will neglect the Goldstone mode and will take $\theta(x) = 0$. The expression for the Lagrangean (A1) becomes
\[
L = \frac{1}{2} \left\{ \tilde{\Phi}^\dagger(x) \left[ i\partial_0 - \varepsilon + \mu_0 \right] \tilde{\Phi} + \tilde{\Phi}(x) \left[ -i\partial_0 - \varepsilon + \mu_0 \right] \tilde{\Phi}^\dagger(x) 
- \lambda_0 \left( \Phi_0^2 \tilde{\Phi}^\dagger(x)^2 + \tilde{\Phi}^2(x)\Phi_0^2 \right) 
- \lambda_0 \left( \Phi_0^2 \tilde{\Phi}^2(x) + \tilde{\Phi}^\dagger(x)^2\Phi_0^2 \right) \right\} \tag{A4}
\]
In the $p$-representation using
\[
\hat{\Phi} = \begin{pmatrix} \tilde{\Phi}(x) \\ \Phi_0 \end{pmatrix} \tag{A5}
\]
Eq. (A4) becomes
\[
\mathcal{L} = \frac{1}{2} \hat{\Phi}^\dagger \hat{M} \hat{\Phi}(x) \tag{A6}
\]
where
\[
\hat{M} = \begin{pmatrix} p_0 - \varepsilon(p) - U(x) - 4\lambda_0|\Phi_0|^2 & -2\lambda_0|\Phi_0|^2 \\ -2\lambda_0|\Phi_0|^2 & -p_0 - \varepsilon(p) - U(x) - 4\lambda_0|\Phi_0|^2 \end{pmatrix} \tag{A7}
\]
and
\[
U(x) = \partial_0 \theta(x) + \frac{1}{2m} [\nabla \theta(x)]^2 \tag{A8}
\]
Eq. (A7) has been obtained neglecting the terms containing $\nabla^2 \theta(x)$, which is irrelevant in the low momentum approximation, and the term $j \nabla \theta(x)$, where $j$ is the current associated with $x$-variation of the phase. For the elementary excitations spectrum we neglect $U(x)$ (which is equivalent to neglect the $x$-dependence of $\theta(x)$) and from the condition
\[
det \hat{M} = 0 \tag{A9}
\]
we get
\[
E^2(p) = \varepsilon^2(p) + 2\mu_0\varepsilon(p) = \varepsilon^2(p) + 4\lambda_0|\Phi_0|^2\varepsilon(p) \tag{A10}
\]
The spectrum expressed by (A10) can be approximated, in the limit $p \to 0$, as
\[
E \cong u_0|p| \tag{A11}
\]
with $u_0 = \sqrt{\mu_0/m}$, is gapless and it remains gapless to all orders in perturbation theory. For large momentum from (A10) we get
\[
E(p) \cong \varepsilon(p) + 2\lambda_0|\Phi_0|^2 \tag{A12}
\]
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