Effect of disorder in BCS–BEC crossover

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Abstract

In this paper, we have investigated the effect of weak random disorder in the Bardeen–Cooper–Schrieffer to Bose–Einstein condensation (BCS–BEC) crossover region. The disorder is included in the mean field formalism through Nozières–Smith–Rink (NSR) theory of superconducting fluctuations. A self-consistent numerical solution of the coupled equations involving the superfluid gap parameter and density as a function of the disorder strength, albeit unaffected in the BCS phase, yields a depleted order parameter in the BEC regime and an interesting nonmonotonic behaviour of the condensate fraction in the vicinity of the unitary region, and a gradual depletion thereafter, as the pairing interaction is continuously tuned across the BCS–BEC crossover. The unitary regime thus demonstrates a robust paradigm of superfluidity even when the disorder is introduced. To support the above feature and shed light on a lingering controversial issue, we have computed the behaviour of the sound mode across the crossover that distinctly reveals a suppression of the sound velocity. We also find the Landau critical velocity that shows similar nonmonotonicity as that of the condensate fraction data, thereby supporting a stable superfluid scenario in the unitary limit.

(Some figures may appear in colour only in the online journal)

1. Introduction

Atomic gases at very low temperature are a unique system where one can observe the continuous evolution of a fermionic system (BCS type) to a bosonic system (BEC type) by changing the inter-atomic interaction by means of Fano–Feshbach resonance \cite{1}. The experimental advances to address this transition have introduced the possibility for studying the so-called BCS–BEC crossover more closely \cite{2, 3}. The physics of the crossover focuses on the change of $s$-wave scattering length ($1/a$) from attractive to repulsive (from $-\infty$ to $+\infty$) by tuning an external magnetic field. In the intermediate region where $1/a \rightarrow 0$, a new region emerges, where dimensionless coherence length becomes in the order of unity, which is the focal point of crossover physics.

The situation becomes more interesting, if one assumes the existence of random disorder in an otherwise very clean system. It is well known that every wavefunction is spatially localized when the disorder is introduced in the one-dimensional case, which is called Anderson localization \cite{4}. However, it is immensely difficult to directly observe the localization in electronic systems on crystal lattices, so that one has to take the indirect route of conductivity measurement to observe the effects of localization. It is an intriguing question how the Anderson localization modifies the BCS superconductivity. It has been found by Anderson \cite{5} that the order parameter is unaffected when the disorder is not too strong to give rise to the Anderson localization. Here, the time reversal pairs, rather than the opposite momentum pairs, then form the Cooper pairs.

The cold atomic system offers a great amount of controllability and allows one to observe the macroscopic wavefunction. Therefore, it was conceived as a very useful candidate to visualize the disorder effects directly. Recently, the ultracold Bose gas ($^{87}\text{Rb}$ and $^{39}\text{K}$) enabled us to see the localization directly \cite{6, 7}. The latest experiments are conducted in three dimensions for both noninteracting atomic Fermi gas of $^{40}\text{K}$ \cite{8} and Bose gas of $^{87}\text{Rb}$ \cite{9}. These experiments have widened the possibility of studying the crossover in light of disorder \cite{10} experimentally.

The static disorder in Fermi and Bose systems is not a new issue. A considerable amount of attention has been paid to
disorders in superconductors [11–13] and in Bose gas [14–25]. Of late, the interest at unitarity is also gaining pace [26–31], but it has still failed to address all the questions associated with it. It is exciting to envision the three-dimensional phase diagram involving temperature, interaction and disorder through the evolution from BCS to BEC superfluid. At the beginning, it can be considered that the random potentials are independent of the hyperfine states of the atoms, but then they can be extended to the correlated disorder problem in the crossover region.

In this paper, we present our investigation on BCS–BEC crossover with weak uncorrelated disorder at zero temperature. Precisely, we show (i) monotonic depletion of the order parameter, (ii) the nonmonotonic nature of the condensate fraction and (iii) suppression of sound velocity as a function of disorder. Furthermore, we add a study on Landau critical velocity to attempt a qualitative understanding of the nonmonotonic behaviour of the condensate fraction. Although the order parameter is mainly for academic interest, the other three quantities are also experimentally viable [32–34]. In this study, we observe that superfluidity is more robust in the unitary regime, which is also consistent with the behaviour of the coherence length in the crossover regime. The robust nature of superfluidity has already been pointed out in the context of vortex core structure [35], Josephson current [36] and collective modes [37]. We also observe a progressive decay of the sound velocity on the BEC side with increasing disorder possibly occurring due to the enhancement of impurity scattering [38]. Here we include a systematic study of the physical observables and their response to the random disorder.

The recent experiments on ultracold Bose and Fermi gas with disorder were carried out with optical speckle [6, 8] and quasi-periodic optical lattices [7]. Although they pose interesting physics to study, here we only consider quenched delta-correlated disorder which remains the subject of interest at zero temperature [26] and at finite temperatures [31] in previous studies. One can visualize the situation of the few heavy atoms (say 40K) in a homogeneous bath of light \(^6\)Li atoms [31]. Since the number of heavy atoms is very limited, one can call it a quasi-homogeneous system. To include the disorder effects in the mean field approach, we follow the NSR theory [39] of superconducting fluctuations extended to the broken symmetry state [40]. Furthermore, we take advantage of the developed technique to study the condensate fraction of clean Fermi gas with Gaussian fluctuations [41, 42].

We arrange our study in the following way: in section 2, we present the basic formalism. Section 3 is dedicated to discussing our results in three parts. The first part contains the order parameter, the second one is for the condensate fraction and the third leads to the sound mode and Landau critical velocity. Finally, in section 4 we draw our conclusions.

2. Formalism

To make our presentation self-contained, we briefly summarize the mathematical formalism presented in [26]. To describe the effect of impurity in Fermi superfluid in the crossover from the BCS to BEC regime, one needs to start from the real space Hamiltonian in three dimensions for an s-wave superfluid:

\[
\mathcal{H}(x) = \sum_{\sigma} \Phi^\dagger_\sigma(x) \left[ -\frac{\nabla^2}{2m} - \mu + V_d(x) \right] \Phi_\sigma(x) + \int dx^2 \mathcal{V}(x, x') \Phi^\dagger_\sigma(x') \Phi^\dagger_\sigma(x) \Phi_\sigma(x) \Phi_\sigma(x'),
\]

where \( \Phi^\dagger_\sigma(x) \) and \( \Phi_\sigma(x) \) represent the creation and annihilation of fermions with the mass \( m \) and spin state \( \sigma \) at \( x \). \( V_d(x) \) signifies the (weak) random potential and \( \mu \) is the chemical potential. We set \( \hbar = 1 \), where \( \hbar \) is the Planck constant. The s-wave fermionic interaction is defined as \( \mathcal{V}(x, x') = -g \delta(x - x') \). The disorder potential is modelled as \( V_d(x) = \sum_i g_d \delta(x - x_i) \), where \( g_d \) is the fermion-impurity coupling constant and \( x_i \) are the static positions of the quenched disorder. We assume that it exhibits white noise correlation, that is, \( \langle V_d(q) V_d(q') \rangle = \kappa \beta \delta_{qq'} \), where \( \beta \) is the inverse temperature, \( v_m \) is the bosonic Matsubara frequency and \( \kappa = n_i g_d^2 \), which describes the strength of the impurity potential with \( n_i \) being the concentration of the impurities.

The partition function corresponding to equation (1) can be written in the path integral formulation as

\[
Z = \int D[\Phi, \Phi^\dagger] \exp \left[ -\mathcal{S}(\Phi, \Phi^\dagger) \right],
\]

where \( \mathcal{S}(\Phi, \Phi^\dagger) = \int dx \partial_\tau \Phi^\dagger \partial_\tau \Phi + \mathcal{H} \). By introducing the pairing field \( \Delta(x, \tau) \) and applying the Grassman identity \( \int D[\Delta, \Delta^\dagger] \exp \left[ -\frac{1}{g} \int dx \int dx' \partial_\tau \Delta \partial_\tau \Delta^\dagger \right] = 1 \), equation (2) can be given as

\[
Z = \int D[\Phi, \Phi^\dagger] \int D[\Delta, \Delta^\dagger] \exp \left[ -\mathcal{S}_{\text{eff}} \right],
\]

where \( \mathcal{S}_{\text{eff}} = \mathcal{S}(\Phi, \Phi^\dagger) + 1/g \int dx \int dx' \partial_\tau \Delta \partial_\tau \Delta^\dagger \). Following Hubbard–Stratonovich transformation, equation (3) can be written in terms of the inverse Nambu propagator as

\[
Z_{\text{eff}} = \int D[\Delta, \Delta^\dagger] \exp \left[ -\frac{1}{g} \int dx \int dx' \partial_\tau \Delta \partial_\tau \Delta^\dagger \right] \times \int D[\Phi, \Phi^\dagger] \exp \left[ -\frac{1}{g} \int dx \int dx' \partial_\tau \Phi^\dagger \partial_\tau \Phi \right],
\]

where the inverse Nambu propagator \( G^{-1}(x, \tau) \) is defined as

\[
\begin{pmatrix}
-\partial_\tau + \frac{\nabla^2}{2m} + \mu - V_d \\
-\partial_\tau - \frac{\nabla^2}{2m} - \mu + V_d
\end{pmatrix}
\]

(5)

After integrating out the fermionic fields from equation (4), we are left with the effective action as

\[
\mathcal{S}_{\text{eff}} = \int dx \int_0^\beta d\tau \left[ \frac{\Delta^2}{g} - \frac{1}{\beta} \text{Tr} \ln[\beta G^{-1}(x, \tau)] \right],
\]

(6)

where \( r = (x, \tau) \). It is important to mention that the main contribution in the partition function comes from a small fluctuation, \( \delta \Delta(x, \tau) = \Delta(x, \tau) - \Delta \), where \( \Delta \) is the homogeneous BCS pairing field. Green’s function in equation (5) can be written as a sum of Green’s function in the absence of disorder \( \left( G_0^{-1} = -\partial_\tau + (\nabla^2/2m + \mu) \sigma_z + \Delta \sigma_i \right) \).
and a self-energy contribution \( \Sigma = -V_p \sigma_x + \delta \Delta \sigma_x + \delta \Delta \sigma_\mu \) which contains the disorder as well as the small fluctuations of the BCS pairing fields. \( \mathcal{I} \) denotes the identity matrix, and \( \sigma_i \) are the Pauli matrices and ladder matrices \((i \in \{x, y, z, +, -\})\).

By expanding the inverse Nambu propagator up to the second order, one can write the effective action \( (S_{BCS}) \) in equation (6) as a sum of bosonic action \( (S_B) \) and fermionic action \( (S_F) \). Also, it contains an additional term which emerges from the linear order of self-energy expansion \( (G_0^\Sigma) \). It is possible to set the linear order to zero if we consider that \( S_F \) is an extremum of \( S_{BCS} \) after performing all the fermionic Matsubara frequency sums. The constrained condition leads to the BCS gap equation which after appropriate regularization through the s-wave scattering length reads

\[
-\frac{m}{4\pi a} = \sum_k \left[ \frac{1}{2E_k} \right] - \frac{1}{2E_k}.
\]

Equation (7) suggests that the BCS gap equation does not have any contribution from the disorder potential explicitly.

Now to construct the density equation with a usual prescription of statistical mechanics, the thermodynamic potential \( \Omega \) should be differentiated with respect to the chemical potential \( \mu \). \( \Omega \) can be written as a sum over fermionic \( (\Omega_F) \) and bosonic \( (\Omega_B) \) thermodynamic potentials, which implies that

\[
n = n_F + n_B = -\frac{\partial}{\partial \mu} (\Omega_F + \Omega_B) = -\frac{1}{\beta} \frac{\partial}{\partial \mu} (S_F + S_B).
\]

(8)

The well-known BCS density equation can be restored from equation (8) if we consider only \( n_F \), then it yields the familiar \( \sum_k (1 - \frac{\xi_k}{E_k}) \). However, the presence of disorder and fluctuation leads to \( n_B \neq 0 \). Hence, the final mean field density equation will be

\[
n = \sum_k \left( 1 - \frac{\xi_k}{E_k} \right) = \frac{\partial \Omega_B}{\partial \mu}.
\]

(9)

The bosonic thermodynamic potential consists of two parts. One comes from the thermal contribution and the other is due to disorder. Since we are interested in zero temperature, we neglect the thermal contribution from here on. Henceforth, the disorder-induced thermodynamic potential can be written as

\[
\Omega_{B_d} = -\frac{k}{\beta} \sum_{q, \nu \omega \sigma} \mathcal{N}^\dagger \mathcal{M}^{-1} \mathcal{N},
\]

(10)

where \( \mathcal{N} \) is a doublet which couples disorder with fluctuation. After performing the fermionic Matsubara frequency summation over \( \mathcal{N} \),

\[
\mathcal{N}_1 = \sum_k \frac{2}{E_k \xi_k + E_k + E_k - \xi_k}.
\]

(11)

The inverse fluctuation propagator matrix \( \mathcal{M} \) is a \( 2 \times 2 \) symmetric matrix whose elements are given by

\[
\mathcal{M}_{11} = \frac{1}{g} + \sum_k \left[ \frac{\xi_k^2}{\nu_m - E_k - E_k + \xi_k} - \frac{\nu_m^2}{\nu_m - E_k + E_k - \xi_k} \right],
\]

\[
\mathcal{M}_{12} = \sum_k \left[ \frac{1}{\nu_m - E_k + E_k + \xi_k + \nu_k} - \frac{1}{\nu_m - E_k - E_k + \xi_k + \nu_k} \right] \times \left[ \frac{1}{\nu_m - E_k + E_k + \xi_k - \nu_k} - \frac{1}{\nu_m - E_k - E_k + \xi_k - \nu_k} \right].
\]

(12)

\[\text{Figure 1.}\] The order parameters \( \Delta \) normalized by Fermi energy as a function of \((k_F a)^{-1}\) for various disorder strengths. The shaded area represents the crossover region. In the inset, the behaviour of chemical potential is depicted.

3. Results

Order parameter

Equations (7) and (9) are now ready to be solved self-consistently. Our analysis is valid only for the weak disorder. Considering [31], we can safely assume that the disorder is weak if the dimensionless disorder strength \( \eta = \kappa m^2/k_F \leq 5 \) is satisfied. Figure 1 demonstrates that in the BCS limit \((1/k_F a \rightarrow -\infty)\) the behaviour of the order parameter for different values of disorder strength follows the mean field expression of \( \Delta/k_F \) which is a very small quantity as \( \Delta \rightarrow 0 \) for \((k_F a)^{-1} \rightarrow 0 \). On the BEC side, the correction comes through the effective chemical potential of the composite bosons. However, the BEC chemical potential is dominated by the binding energy and it turns out quite large compared to the effective chemical potential. The external potential \( V_d(x) \) induced by the disorder usually has no direct influence on the internal degrees of freedom, e.g. interaction amongst fermions. It is thus no wonder that the binding energy of the composite bosons exhibits no pronounced change as a function of \((k_F a)^{-1}\) and so does the chemical potential.

Condensate fraction

Although \( \Delta \) and \( \mu \) are the first quantities to study the crossover through the mean field theory, they are mostly of academic
Here we follow the similar mean field description of various disorder strengths. The shaded area and the vertical orange line at \((k_F a)^{-1} = 0\) represent the crossover region and the unitary limit, respectively.

Figure 2. The condensate fraction \(n_c\) as a function of \((k_F a)^{-1}\) for various disorder strengths. The shaded area and the vertical orange line at \((k_F a)^{-1} = 0\) represent the crossover region and the unitary limit, respectively.

Interest. From now on, we shall focus on the quantities which are more prone to the experimental observation. Our first choice is the condensate fraction which is also one of our main results (depicted in figure 2). In a clean Fermi gas, it is possible to work out the condensate fraction through mean field theory [44] which shows good agreement with the experiment [33]. Here we follow the similar mean field description

\[
    n_c = \sum_k \left[ \frac{\Delta(\eta)}{2E_k(\eta)} \right]^2, \tag{13}
\]

where \(n_c\) is the condensate fraction. The only difference with the clean system calculation in equation (13) is that we used the disorder-induced values for \(\Delta\) and \(\mu\). We observe quite remarkable behaviour for \(n_c\). As one expects that the condensate fraction decays similar to that of the clean limit when \(1/(k_F a) \to -\infty\), since \(\Delta\) does not change with the variation of \(\eta\) in the BCS regime. If one extends the self-energy up to the second order in the condensate fraction calculation, one can observe the effect of disorder in the BCS limit as well, where \(n_c - n_{c0}\) gets saturated at some finite value instead of exponential decay [26], with \(n_{c0}\) denoting the condensate fraction in the clean limit. On the BEC side, the disorder destroys part of the condensate and turns it into a normal fluid. The condensate fraction approaches roughly \(\eta/\sqrt{k_F a}\) as obtained from the study of hard sphere Bose gas in random disorder [43].

The nonmonotonic behaviour in the crossover region (grey area in figure 2) is the most intriguing point. In the study of quantized vortices, one sees that the accumulation of a number of vortices becomes maximum in the crossover region [45]. Later on, in a theoretical study of a single vortex it was also observed that the circulation current is maximum in the crossover regime [35]. A similar feature was also reported in the Josephson current study [36]. This unique behaviour was qualitatively attributed to the maximization of Landau critical velocity in this region. In general, all these observations actually point to the robustness of the superfluidity at the unitarity. But the normal mean field does not show any precise nonmonotonic behaviour for the condensate fraction. Here with a weak disorder, we are able to generate a picture which affirms the belief that in the presence of less impurity, the condensate fraction is less affected across the crossover, implying a comparatively high yield of superfluidity in this region.

One should also note the position of the extrema. With the increase in disorder, we observe that the extrema slowly move towards \(1/a \to 0^+\) from \(+\infty\), but never cross it as we are exhausted with the weak impurity limit. A further increase of disorder will break down the weakness condition. This result agrees qualitatively with [26, 31]; however in those cases it has been suggested that the most robust region of superfluidity emerges when \(1/a \to 0^+\). But there is no good justification for that. Here it is pointed out that this behaviour is consistent with that of the critical velocity discussed in the following section.

Sound mode and critical velocity

Another important and experimentally relevant [34] quantity is the lowest modes of the collective excitation of the condensate. The nature of sound velocity, in the presence of disorder, in a Bose gas has been studied quite extensively, but it lacks a real consensus so that a considerable amount of ambiguity still exists. Using a perturbative method, the sound velocity is enhanced in the presence of uncorrelated disorder and very weak interaction [16, 46–48], whereas within a non-perturbative self-consistent approach and a spatially correlated weak disorder case, depression in sound velocity has been reported [15, 49–51]. In a more recent study, it has been shown that there exists no generic behaviour of sound in the presence of disorder using a perturbative approach [52]. Hence, in the crossover region, the behaviour of sound is expected to be quite interesting.

From a technical point of view in order to obtain the sound velocity, one needs to carry out an analytic continuation of the Matsubara frequency. Hence, the fluctuation propagator matrix \(M\) is expanded to the second order for both momentum and frequency. The determinant of which leads to [53]

\[
    M_{12}(q, v)M_{12}(q, -v) - M_{12}^2(q, v) = A(\Delta, \mu, k)q^2 + B(\Delta, \mu, k)v^2 + \cdots = 0, \tag{14}
\]

for \(v = v_j q_j\), where \(v_j\) represents the sound velocity, \(A\) and \(B\) are the functions of \(\Delta\) and \(\mu\), respectively, and can be evaluated by summing them over \(k\). Hence, the disorder enters here through the modified \(\Delta\) and \(\mu\) (as we restricted ourselves to observe the disorder effect in the zeroth order of the sound mode). In the BEC limit, it has been shown that the sound velocity is \(v_s^2 = A^2/(8m|\mu|)\) [37], so that the suppression of the order parameter should directly result in diminishing sound velocity.

Near the unitarity, the sound velocity is directly connected to the chemical potential potential through \(v_s^2 = 2\mu/(3m)\). In effect, the sound velocities for various disorders are merged near the unitarity as the chemical potential exhibits almost no change regardless of the addition of the disorder. However, in the sound mode one can point out a spread of around 15% near
unitarity. For that purpose in table 1 we present the chemical potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the chemical potential results in a visible spread in sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation. The deviation in the potential data and compare the numerically calculated sound velocity with the mean field estimation.
words the resonant gas screens the disorder better than a BCS or a BEC superfluid. This may lead to a pronounced behaviour of the condensate fraction in the vicinity of the unitary regime.

4. Conclusion

In conclusion, we have studied several important physical quantities such as gap parameter, condensate fraction, sound velocity, etc to address the issue of BCS to BEC crossover in a disordered environment. To this end, we have included weak disorder via the Gaussian fluctuation as prescribed earlier [26, 31], and hence solved the coupled BCS mean field equations self-consistently. This enables us to obtain the two basic mean field parameters, $\Delta$ and $\mu$. We thus show that the order parameter gets depleted leaving the chemical potential unchanged as we go from a BCS to a BEC regime. Afterwards, the condensate fraction has been calculated using the well-known mean field description with the disorder affected $\Delta$ and $\mu$. As the disorder strength increases, we observe a pronounced maximum developed in the crossover region justifying the expectation of robust superfluid in this region. We have tried to connect this nonmonotonic nature qualitatively to the Landau critical velocity, which also shows a sharp maximum near unitarity. With the increase of impurity, this peak slightly moves towards the unitary point. A similar feature is observed for the condensate fraction. In addition, the depression of the sound velocity has also been addressed, which might be related to the enhanced scattering from the random scatterers employed in this model.

To be precise, the nature of $\Delta$, $n_\alpha$, $v_\alpha$ and $v_s$ has been reported here, when subjected to a weak random impurity. We hope this study has shed some light on the physics of the BCS–BEC crossover in disordered systems. An interesting future perspective can be the generalization of the theory for arbitrarily strong disorder and interaction. We also hope that these results will be observed and verified in an experiment soon.

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