Interior gap superfluidity in a two-component Fermi gas of atoms

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A new superfluid phase in Fermi matter, termed as “interior gap” (IG) or “breached pair”, has been recently predicted by Liu and Wilczek [Phys.Rev.Lett. 90, 047002 (2003)]. This results from pairing between fermions of two species having essentially different Fermi surfaces. Using a nonperturbative variational approach, we analyze the features, such as energy gap, momentum distributions, and elementary excitations associated with the predicted phase. We discuss possible realization of this phase in two-component Fermi gases in an optical trap.

PACS numbers: 03.75.kk, 74.20.-z, 32.80.Lg

Since the first realization of Fermi degeneracy in $^{40}$K gas by Jin’s group [1], several other groups have also achieved quantum degeneracy in atomic Fermi gases. A recent experiment by Regal et al. [2], has reported condensation of fermionic atom-pairs in trapped $^{40}$K gas. Recently, several groups have demonstrated formation of cold molecules [3, 4] and molecular Bose-Einstein condensates [5] from trapped Fermi gases of atoms near the Feshbach resonance. O’Hara et al. [11] demonstrated strongly interacting degenerate Fermi gas of $^6$Li atoms. Studies on trapped degenerate Fermi gases are important in the context of diverse fields, such as nuclear physics, astrophysics, strong-coupling superconductivity and superfluidity, [12] and so on.

A new form of Fermi matter, associated with what is termed as “interior gap” (IG) superfluidity [12], has recently come into fore. Liu and Wilczek [13] have shown that, in two species of fermions differing in Fermi momenta, an attractive interaction can lead to pairing within the interior of the larger Fermi surface, while the exterior remains gapless. This pairing phenomenon gives rise to IG superfluidity having superfluid and normal exterior remains gapless. This pairing phenomenon gives rise to IG superfluidity having superfluid and normal Fermi liquid components simultaneously. Many years ago, an analogous state of two-species fermions was predicted by Sarma [14], which, however, corresponded to a metastable phase. Recently, it was shown in the context of color superconductivity that, such a phase is the only stable phase when the relative density of the two species [15, 16] is kept fixed. IG superfluidity arises naturally in finite density quark matter [15, 17, 18, 19].

In this paper, we analyze the energy gap, momentum distributions and quasiparticle excitations in an IG superfluid system. For possible realizations of this phase, optically trapped two-component atomic Fermi gases may be useful. It has already been suggested [13] that a mixture of $^6$Li and $^{40}$K Fermi gases should be used to create this new phase. Let us consider a mixture of two hyperfine spin components, for instance, $|1\rangle = |F = 1/2, m_F = 1/2\rangle$ and $|2\rangle = |F = 1/2, m_F = -1/2\rangle$ of $^6$Li atoms, in an optical trap. The relative number density of the two components can be controlled by a rf field [11, 21, 21]. The lifetime of each spin component is long enough to carry out an experiment at a fixed relative density which is a necessary condition for IG superfluidity [13]. It may be noted that controlled mixtures, other than 50:50, of fermion gas have already been created [21].

To discuss IG superfluidity, let us consider a system of interacting fermions of two species in a harmonic trap, described by the Hamiltonian $\hat{H} = \hat{H}_0 + \hat{H}_I$ where

$$\hat{H}_0 = \sum_{i=1,2} \int \, dr \hat{\Psi}_i^\dagger(r) \left[ -\frac{\hbar^2}{2m_i} \nabla^2 + V_{ho}(r) - \mu_i(r) \right] \hat{\Psi}_i(r),$$

where $V_{ho}(r)$ is the harmonic trapping potential and

$$\hat{H}_I = \frac{1}{2} \int \, dr' dr \hat{\Psi}_1^\dagger(r) \hat{\Psi}_2^\dagger(r') \hat{\Psi}_1(r) \hat{\Psi}_2(r').$$

Here, $\hat{\Psi}_i(r)$ represents fermionic field operator, $\mu_i(r)$ denotes the local chemical potential and $\hat{V}(\mathbf{r} - \mathbf{r'})$ is the two-body interaction potential. We now introduce the annihilation operator $\hat{c}_i^\dagger(\mathbf{k})$ which, for a uniform gas, is related to field operator by $\hat{\Psi}_i(r) = (1/\sqrt{V}) \sum_k \exp(-i\mathbf{k}\cdot\mathbf{r}) \hat{c}_i^\dagger(\mathbf{k})$, where $V$ is the volume of the system. For a uniform gas with number density $n$, the mean field energy in the weak interaction regime is $U_{mf} = gn$, where $g = 4\pi\hbar^2a_s/(2m)$, with $m = m_1m_2/(m_1 + m_2)$ being the reduced mass, and $a_s$ the s-wave scattering length. In the strong interaction regime, it is unitarity-limited and takes the form $U_{mf} = \beta E_F$, [11, 23, 24], where $\beta$ is a constant. In this regime, one can define an effective scattering length from the relations $U_{mf} = gn = \beta E_F$ yielding $a_{eff} \propto 1/k_F$. Under local density approximation (LDA) in a harmonic trap of potential $V_{ho}$, the equilibrium conditions are given by

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\[ \frac{\hbar^2}{2m_i} |6\pi^2 n_i(r)|^{2/3} + g|\epsilon_{ij}|n_j(r) + V_{ho}(r) = \mu_i(r), \]  

where \( \epsilon_{ij} \) is the Levi-Civita tensor. LDA is particularly good when the local Fermi energy is larger than the average level spacing of the trap and the coherence length of the fermion pair is shorter than the average trap size \( \delta \). The Fermi momentum of the \( i \)th component is then related to the respective peak density \( n_i(r_0) \) at trap center by \( k_F = [6\pi^2 n_i(r_0)]^{1/3} \). In what follows, we thus consider a uniform gas and derive the gap equation.

To consider the pairing state, we take the variational ansatz as given by \( \Delta \) where we have defined the gap \( \Delta = -\bar{\epsilon} \).

The superscript “prime” refers to operators in the Hilbert space. In Eq. (5), the function \( \bar{\epsilon}(k) \) will be related to the distribution functions, for the two species. Using the Bogoliubov transformation, it is easy to calculate the thermodynamic potential given by

\[ \Omega = \exp \left[ \frac{1}{2} \int \left( \epsilon_i' k^1 f(k) \epsilon_j' (-k^1) \right) \epsilon_{ij} dk - H.c. \right] \left| \Omega \right| \]

where \( \left| \Omega \right| \) represents the vacuum, annihilated by \( \epsilon_i' \)’s. To include the effects of temperature and density, we write down the state at finite temperature and density \( \rho \), taking a thermal Bogoliubov transformation over the state \( \Omega \) using thermofield dynamics as described in Refs. 27, 28.

The superscript “prime” refers to operators in \( \left| \Omega \right| \) basis, i.e., \( \epsilon_i' \left| \Omega \right| = 0 \). The tilde operators are the ones in the extended Hilbert space. In Eq. (6), the function \( \theta_i(k, \beta, \mu) \) will be related to the distribution functions.

Using the Bogoliubov transformation, it is easy to calculate the thermodynamic potential given by

\[ \Omega = \left< \Omega(\beta, \mu) | \mathcal{H} - \mu_j \psi_j' \psi_j \right| \Omega(\beta, \mu) \right> \left| \Omega \right| - \frac{1}{\beta} s; \]

where \( \mathcal{H} \) is the Hamiltonian density for constant potential. In the above \( \left| \Omega \right| \), \( s = -\frac{\hbar^2}{4\pi}\sum_k \int dk \left( \sin^2 \theta_i \ln \sin^2 \theta_i + \cos^2 \theta_i \ln \cos^2 \theta_i \right) \) is the entropy density.

Minimizing the thermodynamic potential \( \Omega \) with respect to condensate function \( f(k) \), we get\n
\[ \tan 2f(k) = \frac{-2gI_D}{(\epsilon_1 - \epsilon_2 + \epsilon_3 - \epsilon_4 + \epsilon_5 + \epsilon_6)} \equiv \frac{\Delta}{(\bar{\epsilon} - \bar{\nu})}, \]

where we have defined the gap \( \Delta = -gI_D \). The order parameter here is \( \left< \psi_i(x) \psi_j(x) \right> = -I_D \) where \( I_D = \frac{-\hbar^2}{4\pi}\sum_k \int 2\sin 2\theta_i \left( 1 - \sin^2 \theta_i - \sin^2 \theta_j \right) dk \). Further, \( \epsilon = (\epsilon_1 + \epsilon_2)/2, \bar{\nu} = (\nu_1 + \nu_2)/2, \) and \( \nu_i = \nu_i - g\rho_j \epsilon_j \). Thus, it may be noted that the pair condensate functions depend upon the average energy and the average chemical potential of the fermions that condense. Finally, the minimization of the thermodynamic potential with respect to the thermal functions \( \theta_i(k) \) yields \( \sin^2 \theta_i = 1/(\exp(\beta\omega_i) + 1) \). The quasi particle energies \( \omega_i \)’s are given by \( \omega_1 = \omega + \delta \xi, \omega_2 = \omega - \delta \xi \). Here, \( \omega = \sqrt{\Delta^2 + \xi^2}, \xi = \bar{\epsilon} - \bar{\nu} \) and \( \delta \xi = (\epsilon_1 - \epsilon_2)/2 - (\nu_1 - \nu_2)/2 \). Note that in the degenerate mass and chemical potential case we shall have the same quasi particle energies but if the masses and/or the chemical potentials are different there is possibility of having gapless modes \( \omega_1 = 0 (\omega_2 = 0) \) when \( \omega = -\delta \xi \). So, although we shall have nonzero order parameter \( \Delta \) there will be fermionic zero modes or gapless modes \( \xi \). Indeed, we shall show below that such a situation is possible depending upon the mismatch in the fermion number densities and the magnitude of \( \Delta \).

Substituting the solutions for the ansatz functions in definition for \( I_D \), we can have the gap equation which, however, is divergent. The origin of this divergence lies in the point-like four fermion interaction which needs to be regularized. We tackle this problem by defining the regularized coupling as was done in Refs. 23, 24 to access the strong coupling regime \( 24, 30 \) by subtracting out the zero temperature and zero density contribution. The regularized gap equation is then given by

\[ -\frac{1}{g_R} = \frac{1}{(2\pi)^3} \int dk \left[ \frac{1 - \theta(-\omega_1) - \theta(-\omega_2)}{2\sqrt{\Delta^2 + \xi^2}} - \frac{1}{2\xi} \right] \]

where \( g_R = 2\pi\hbar^2a_{eff}/\bar{m} \). Unless otherwise stated explicitly, we take \( \hbar = v_F = 1 \) in our calculations, where \( v_F \) is the Fermi velocity. Substituting the expressions for the optimized condensate functions and the distribution functions, finite part of the thermodynamic potential Eq. (4), in the zero temperature limit, reduces to

\[ \Omega = \frac{1}{(2\pi)^3} \int dk \left[ \xi - \sqrt{\Delta^2 + \xi^2} \right] - \frac{\Delta^2}{g_R} - g_R\rho_1\rho_2 \]

\[ + \frac{1}{(2\pi)^3} \int dk \left( (\omega + \delta \xi)\theta(-\omega_1) + (\omega - \delta \xi)\theta(-\omega_2) \right) \]
where $\delta^2 k^2_F = k^2_F - k^2_{F_1}$. For fixed average chemical potential and for nonzero $\delta \xi$ there are, in general, two solutions for the gap as shown in Fig. 1. Clearly, when $\delta \xi = 0$, there is only one solution which for the parameters chosen turns out to be $\Delta = 0.71|a_{\text{ren eff}}|^{-1}$. As $\delta \xi$ is increased, the second integral in Eq. (11) starts contributing and the curve can cross in general the RHS=1 line at two places. The larger $\Delta$ is the conventional BCS solution. The smaller value of $\Delta$ will have breached gap character $[15]$ and, in this case, the number densities of the two species will be different. In fact, the number densities of the two species are given as, for $\delta \xi < 0$,

$$\rho_1 = \frac{1}{(2\pi)^3} \int \left[ \theta(-\omega_1) + \frac{1}{2} \left( 1 - \frac{\xi}{\omega} \right) (1 - \theta(-\omega_1)) \right] \, dk, \quad (10)$$

$$\rho_2 = \frac{1}{(2\pi)^3} \int \frac{1}{2} \left( 1 - \frac{\xi}{\omega} \right) (1 - \theta(-\omega_1)) \, dk. \quad (11)$$

Clearly, without any mismatch in the chemical potential of the two species, $\omega_1 = \omega > 0$ and hence the two densities are equal as the theta functions will not contribute and the distribution function becomes the usual BCS distribution function. In presence of mismatch in densities, the $\theta$ function contributes in the range of momenta between $k_{\text{min}}$ and $k_{\text{max}}$. In that region, the particle distribution function becomes unity for species one and zero for the other species. This is what is seen in Figs. 2(c) and 2(d). It is easy to see that in the breached gap phase, the difference in number densities is $\delta \rho \equiv \rho_1 - \rho_2 = (k^3_{\text{max}} - k^3_{\text{min}})/(6\pi^2)$

For systems where chemical potential is kept fixed, the breached pair phase will always have higher free energy as compared to the BCS phase. However, for systems where the relative densities of the condensing fermions can be fixed, the breached pairing state will be the only solution provided the ratio between the densities is different from unity. In Fig 2(b), almost linear behavior of the gap with the density ratio may be noted. Figures 2(c) and (d) show that in a strip of momenta between $k_{\text{min}}$ and $k_{\text{max}}$, the momentum distribution resembles that of normal fermionic distribution. Below $k_{\text{min}}$ and beyond $k_{\text{max}}$, the distribution is of the BCS type. The quasi-particle excitation energy $\omega_1$ becomes zero at the points $k_{\text{min}}$ and $k_{\text{max}}$ showing gapless behavior. For the parameters chosen in Figs. 3(c) and (d), $k_{\text{min}}|a_{\text{eff}}| = 1.42$ and $k_{\text{max}}|a_{\text{eff}}| = 1.58$, i.e., $k_{\text{min}}/k_0 = 0.95$ and $k_{\text{max}}/k_0 = 1.05$, $k_0$ being the average Fermi momentum of the two species.

We thus show that the ansatz state, as defined in Eq. (1), shows breached pair superfluidity with the excitation energies becoming zero at $k_{\text{min}}$ and $k_{\text{max}}$ and that when relative number density is fixed this state is the state having lower free energy. In this region, only one of the species is completely occupied and the other is completely empty. Thus, it is a more general state as compared to Ref. [13] where the excitation energy was zero at $k_{\text{min}}$ which was equal to $p_d$ of Ref. [13] beyond which one species was completely occupied and the other was completely empty.

We next discuss the possibility of probing the gap by stimulated Bragg scattering (SBS) $[31, 32]$ in analogy with Raman scattering in anisotropic superconductors $[32]$. There are many theoretical proposals $[34, 35]$ for...
The laser photon with momentum $k_1$ is scattered into a photon with momentum $k_1 - k_3$ to the atoms along the $x$ axis, while the laser photon with momentum $k_3$ is scattered into a photon with momentum $k_2$ transferring a momentum $\mathbf{q} = k_3 - k_2$ to the atoms. The frequencies of the three laser beams are chosen such that the energy transfer $\delta = (\omega_1 - \omega_3) = (\omega_3 - \omega_2) > 0$. A pair of atoms which may form a Cooper pair having mutually opposite momenta are scattered equally. The one-dimensional (along the $x$ direction) momentum and density distribution of the gas can be determined from the analysis of time of flight images. The spectrum (number of scattered atoms versus $\delta$ for different $\mathbf{q}$ values) may reveal the existence of the gap and the detailed comparison of the spectra for different relative densities of the two components may provide a proof of IG superfluidity.

FIG. 3: A possible scheme for measuring dynamic structure function by SBS. The laser photon with momentum $k_1$ is scattered into a photon with momentum $k_1 - k_3$ to the atoms along the $x$ axis, while the laser photon with momentum $k_3$ is scattered into a photon with momentum $k_2$ transferring a momentum $\mathbf{q} = k_3 - k_2$ to the atoms. The frequencies of the three laser beams are chosen such that the energy transfer $\delta = (\omega_1 - \omega_3) = (\omega_3 - \omega_2) > 0$. A pair of atoms which may form a Cooper pair having mutually opposite momenta are scattered equally. The one-dimensional (along the $x$ direction) momentum and density distribution of the gas can be determined from the analysis of time of flight images. The spectrum (number of scattered atoms versus $\delta$ for different $\mathbf{q}$ values) may reveal the existence of the gap and the detailed comparison of the spectra for different relative densities of the two components may provide a proof of IG superfluidity.

probing atomic Cooper pairs by resonant \footnote{34} and off-resonant \footnote{35} laser light. It is well established that, in the case of a Fermi superfluid, there will be no spectrum of density fluctuation unless energy transferred to the superfluid due to scattering exceeds $2\Delta$ \footnote{33}. The quantum Monte Carlo simulation \footnote{36} suggests that, in the large $a_s$ limit, $\Delta \sim 0.81(3/5)c_F \simeq 0.49c_F$. In the recent experiments \footnote{4} \footnote{10} \footnote{11} on two-component $^6$Li atoms in the unitarity-limited regime $(a_s k_F >> 1)$, the typical value of the Fermi velocity $v_F = h k_F / m \sim 15$ cm/s. For two-component $^6$Li atoms, if the Bragg pulses are tuned near an excited level (say, $2\Sigma_3/2$, wavelength $\lambda = 670.776$ nm), the momentum transfer $\mathbf{q} = k_1 - k_3 \simeq 2k_L \hat{x}$, where $k_L = 2\pi / \lambda$. This momentum transfer raises the velocity of the scattered atom by $2(h k_L / m) \simeq 20$ cm/sec which exceeds $v_F$. Therefore, the scattered atoms can be distinguished in time of flight images as discussed in the caption of Fig. 3.

In conclusion, we have analyzed momentum distributions and energy gap for an IG superfluid. Two-component Fermi gases of atoms in optical traps seem to be promising systems for experimental exploration for such superfluid state. Fermi atoms in a two-band optical lattice may also be considered for studying breached pair superfluidity.

The authors are thankful to I. Shovkovy for useful communications and discussions; and to M. Randera for helpful discussions and bringing to our attention Ref. \footnote{33}.

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