Nonuniform mixed-parity superfluid state in Fermi gases

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(Dated: May 10, 2018)

We study the effects of dipole interaction on the superfluidity in a homogeneous Fermi gas with population imbalance. We show that the Larkin-Ovchinnikov-Fulde-Ferrell phase is replaced by another nonuniform superfluid phase, in which the order parameter has a nonzero triplet component induced by the dipole interaction.

PACS numbers: 74.20.Rp, 74.20.-z, 03.75.Ss

Bardeen-Cooper-Schrieffer (BCS) theory is the standard model of superconductivity and superfluidity in fermionic systems. The central concept of the BCS model is that of a Cooper pair composed of fermions with opposite momenta and spins. If there is a mismatch between the Fermi surfaces of spin-up and spin-down fermions, e.g. because of the Zeeman splitting in magnetic field, then the formation of the Cooper pairs costs energy, which results in the suppression of the critical temperature. It was shown by Larkin and Ovchinnikov \cite{1} and Fulde and Ferrell \cite{2} (LOFF) that the pair-breaking effect. It was shown by Larkin and Ovchinnikov \cite{1} and Fulde and Ferrell \cite{2} (LOFF) that the pair-breaking effect of the Fermi surface splitting can be reduced if the Cooper pairs have a nonzero center-of-mass momentum. The resulting nonuniform superconducting state turns out to be more favorable at low temperatures than the uniform BCS state.

Due to the sensitivity of the nonuniform state to disorder and the orbital effects, the attempts to find it in superconductors have been unsuccessful (at least until the recent observation of the features consistent with the LOFF state in the phase diagram of the heavy-fermion compound CeCoIn$_5$ \cite{3}). The recent surge of interest to the nonuniform superconducting and superfluid states, see e.g. Refs. \cite{4,5,6,7,8,9,10}, has been stimulated by the experimental progress in ultracold atomic Fermi gases, such as $^6$Li and $^{40}$K. The Fermi-surface splitting in these systems is due to the populations of atoms in different hyperfine states being unequal \cite{11,12}. When a pairing interaction between the fermions is turned on, the system becomes formally equivalent to a neutral and perfectly clean superconductor in a Zeeman field.

In this Letter we consider the effects of the long-range and anisotropic interaction between atomic dipole moments in a Fermi gas with population imbalance. The magnetic dipole interaction is quite small for alkali atoms but can be considerably enhanced to become experimentally observable for atoms with large dipole moments \cite{12}. Another possibility is to control the magnitude of electric dipole moments by applying external electric field, as suggested in Ref. \cite{13}. In either case, we assume that the dipole interaction represents a correction to the dominant isotropic $s$-wave pairing. We will show that increasing the population imbalance produces a novel nonuniform mixed-parity (NMP) superfluid state, which has higher critical temperature than the LOFF state.

We consider a homogeneous Fermi gas consisting of two different species of atoms of equal mass $m$. For instance, in $^6$Li one can have a mixture of the hyperfine states $|+\rangle \equiv |m_e = 1/2, m_i = 1\rangle$ and $|-\rangle \equiv |m_e = 1/2, m_i = 0\rangle$ \cite{14}, where $m_e$ and $m_i$ are the electron and nuclear spin projections respectively. The free-particle Hamiltonian has the form $H_0 = \sum_{k}(\xi_k - \hbar \sigma_3 \alpha^2)\hat{c}_{k\alpha} \hat{c}_{k\beta} \delta$, where $\xi_k = (k^2 - k_F^2)/2m$, $k_F$ is the Fermi momentum, $\hbar$ is the Fermi surface splitting due to unequal particle concentrations in the states labelled by $\alpha, \beta = \pm$, and $\sigma_3$ is the Pauli matrix (we set $\hbar = k_B = 1$). We assume that the atomic dipoles are fully polarized, e.g. by an external magnetic field along the $z$-axis: $\mu = \mu_d \hat{z}$. The interaction between two atoms at distance $R$ is

$$U(R) = -\lambda_0 \delta(R) + \mu_d^2 \frac{1 - 3R^2}{R^3},$$

where the first term describes a short-range attraction with the coupling constant $\lambda_0 = 4\pi|\alpha|/m > 0$, and the second term is the dipole interaction. The $s$-wave scattering length $a < 0$ depends on the field, diverging in the strong-coupling regime near the Feshbach resonance. The effects of the dipole interaction \cite{11} have been extensively studied for Bose gases, see e.g. Ref. \cite{15} and the references therein.

In the BCS regime the pairing Hamiltonian, which takes into account the possibility of the Cooper pairs having a nonzero momentum $q$, can be written in the following form:

$$H_{int} = \frac{1}{2} \sum_{k,k',q} V(k,k') \Gamma_{\alpha\beta\gamma\delta}$$

$$\times \hat{c}_{k+\hat{q},\alpha} \hat{c}_{-\hat{q},\beta} \hat{c}_{-k',\gamma} \hat{c}_{k',\delta},$$

where $V(k,k') = -\lambda_0 + 4\pi \mu_d^2 (k_0 - k'_0)^2 /|k - k'|^2$ is the Fourier transform of Eq. \cite{11,16}. The matrix $\Gamma$ has the following nonzero elements: $\Gamma_{+++} = \Gamma_{+-+} = \Gamma_{++-} = \Gamma_{---} = 1$, which is due to the fact that the pairing interaction depends only on the total particle densities $n(r) = n_+(r) + n_-(r)$.\

\textbf{References}

\cite{1-15}
The symmetry factors \( \phi \) for \( \hat{s} \) satisfy the orthonormality condition:
\[
\langle s | t \rangle = \sum_i \lambda_i \phi_i(k) \phi_i^*(k').
\]
The symmetry factors \( \phi_i(k) \) are the eigenfunctions and the coupling constants \( -\lambda_i \) the eigenvalues of the operator \( V \) defined by the kernel \( V(k_F, k_F') \). The symmetry factors are even (odd) for the singlet (triplet) pairing and satisfy the orthonormality condition: \( \langle \phi_i^*(k) \phi_j(k) \rangle = \delta_{ij} \), where the angular brackets denote the average over the spherical Fermi surface.

The effect of the dipole interaction on superfluidity is two-fold: in addition to breaking the rotational symmetry of the system and introducing gap anisotropy, it also leads to the possibility of triplet pairing. We include one singlet and one triplet pair quantum state with the highest critical temperatures. One can show that the lowest eigenvalues of \( V \) in both cases are achieved for the eigenfunctions that do not depend on the azimuthal angle: \( \phi(k) = f(\cos \theta), \) where the function \( f(s) \) satisfies the integral equation \( \int_1^\infty ds' K(s, s') f(s') = -\lambda f(s), \) with the kernel \( K(s, s') = -\lambda_0/2 + 2\pi \mu_d^2 |s - s'|. \) In the singlet case the solution is \( f(s) \propto \cos(\kappa_\sigma |s|), \) and in the triplet case \( f(s) \propto \sin(\kappa_\sigma |s|), \) where \( \kappa = \sqrt{2\pi \mu_d^2}. \) For the singlet pairing, the normalized symmetry factor has the form
\[
\phi_\sigma(k) = \sqrt{\frac{2}{1 + \frac{\Delta_\sigma^2}{\kappa_\sigma^2}}} \cos(\kappa_\sigma \cos \theta), \tag{3}
\]
where \( \kappa_\sigma \) satisfies the equation
\[
\kappa_\sigma \tan \kappa_\sigma = \frac{2\pi \mu_d^2}{\lambda_0 - 2\pi \mu_d^2}, \tag{4}
\]
from which one obtains \( \lambda_0 = 2\pi \mu_d^2/\kappa_\sigma^2. \) The limit of weak dipole interaction, \( \mu_d^2 \ll \lambda_0, \) we find \( \phi_\sigma(k) \simeq 1 + \kappa_\sigma^2 (1 - 3s^2)/6, \) and \( \lambda_\sigma \simeq \lambda_0 - 2\pi \mu_d^2, \) i.e. the dipole interaction leads to a downward renormalization of the coupling constant in the singlet channel. For the triplet pairing, one can show that the lowest eigenvalue corresponds to \( \kappa = \pi/2, \) so that
\[
\phi_t(k) = \sqrt{2} \sin \left(\frac{\pi}{2} \cos \theta\right), \tag{5}
\]
and \( \lambda_t = (8/\pi) \mu_d^2. \) Note that the coupling constant and the symmetry factor in the triplet channel are not sensitive to the value of \( \lambda_0 \) and therefore are the same as in the pure dipolar case.

The singlet and triplet contributions to the pairing Hamiltonian can now be explicitly separated:
\[
H_{int} = -\lambda_s \sum_q B_s^\dagger(q) B_s(q) - \lambda_t \sum_q B_t^\dagger(q) B_t(q), \tag{6}
\]
where \( B_s^\dagger(q) = (1/2) \sum_k \phi_s^e(k) (i\sigma_2)_{\alpha\beta} c^\dagger_{k+q/2,\alpha} c^\dagger_{k-q/2,\beta} \) and \( B_t^\dagger(q) = (1/2) \sum_k \phi_t^e(k) ((i\sigma_2)_{\alpha\beta} c^\dagger_{k+q/2,\alpha} c^\dagger_{k-q/2,\beta} \) are the pair creation operators. The representation \( \phi_\sigma(k) \) emphasizes the full invariance of the pairing interaction with respect to rotations in the “pseudospin” space spanned by the states \(|+\rangle \) and \(|-\rangle \). Decoupling \( H_{int} \) yields the gap function
\[
\Delta_{\alpha\beta}(k, \mathbf{r}) = (i\sigma_2)_{\alpha\beta} \psi^e(r) \phi_\sigma(k) + (i\sigma_2)_{\alpha\beta} \mathbf{d}(r) \phi_t(k), \tag{7}
\]
where the complex scalar \( \psi \) and the complex vector \( \mathbf{d} \) are the order parameters of the singlet and the triplet pairing respectively.

In the absence of population imbalance \( h = 0, \) and the critical temperatures for the singlet and triplet pairing are given by the standard BCS expressions: \( T_s = T_{c0}(s) = (2e^2/\pi) \omega_c e^{-1/N_F \lambda}, \) where \( s = s, t \) and \( N_F \) is the density of states at the Fermi surface per one hyperfine state (all three components of \( \mathbf{d} \) have the same critical temperature). Of the most interest to us is the limit of weak dipole interaction, in which \( T_s \gg T_t \).

The phase diagram at \( h \neq 0 \) can be obtained from the free energy, which is represented as an expansion in powers of the order parameter components: \( F = F_0 + F_2[\psi, \mathbf{d}] + ... \) (\( F_0 \) is the free energy in the normal state). To find the critical temperature \( T_c(h) \) of the second-order phase transition into the superfluid state, or conversely the critical Fermi-surface splitting \( h_c(T) \), it is sufficient to keep only the quadratic terms in \( F \). Using Eq. (6) we find that the contributions to \( F_2 \) from \( \psi \) and \( d_z \), which describe the pairing between fermions of different species, are decoupled from \( d_x \pm (d_x \pm id_y)/\sqrt{2}, \) which correspond to the intra-species pairing:
\[
F_2 = \sum_q \left[ \left( \psi_q^* d_{z,q}^d \right) \left( \begin{array}{cc} A_{ss} & A_{st} \\ A_{ts} & A_{tt} \end{array} \right) \left( \begin{array}{c} \psi_q \\ d_{z,q} \end{array} \right) + A_+ |d_{+q}|^2 + A_- |d_{-q}|^2 \right]. \tag{8}
\]
The coefficients in this expression have the following form: \( A_{ab} = A_{ba} = N_F [\delta_{ab} \ln(T/T_s) + I_{ab}], \) \( A_\pm = N_F [\ln(T/T_t) + I_{tt}(h = 0) \pm \delta I], \) \( I_{ab} = \langle \phi_a(k) \phi_b(k) \rangle \Re \Psi \left( \frac{1}{2} + \frac{iW}{4\pi T} \right) - \delta_{ab} \Psi \left( \frac{1}{2} \right), \tag{9}
\]
where \( \phi_{a,t}(k) \) are defined by Eqs. (3) and (5). \( \Psi(x) \) is the digamma function, \( W = v_F k - 2h, v_F \) is the Fermi velocity, \( \delta I \propto N'_F h \) is proportional to the difference between the densities of states at the Fermi levels for
the two fermionic species. Since at \( h \neq 0 \) the symmetry in the "pseudospin" space is reduced to rotations about the z-axis, different components of \( \mathbf{d} \) appear at different temperatures.

The critical temperatures for \( d_+ \) and \( d_- \) are found from the equations \( A_+ = 0 \) and \( A_- = 0 \) respectively. If one neglects the band asymmetry \( N_k^p \) then the phase transition is not affected by the population imbalance, otherwise \( T_c(h) = T_c \mp O(\delta I) \), so that a nonunitary triplet state with \( d_+ = 0, d_- \neq 0 \) is realized.

The effect of the Fermi-surface splitting on the other two components of the order parameter, \( \psi \) and \( d_z \), is more interesting. Because of the presence of the off-diagonal matrix elements in \( \mathcal{A} \), the singlet and triplet pairing channels can be mixed at \( q \neq 0 \) and \( h \neq 0 \), producing the NMP state, in which both \( \psi \) and \( d_z \) are nonzero. The critical temperature is obtained from the equation \( \det \mathcal{A}(q) = 0 \), after maximizing with respect to \( q \). The calculation is facilitated by the observation that the functions \( \beta \) can be expressed in terms of two dimensionless variables \( Q = v_F q / 2h \) and \( z = h / 2\pi T \). At any given \( 0 \leq z \leq \infty \), we find

\[
T_c(h) = T_s \max_q \mathcal{I}(Q, z), \tag{10}
\]

where

\[
\mathcal{I} = -\frac{1}{2} (I_{ss} + I_{tt} + r) + \frac{1}{2} \sqrt{(I_{ss} - I_{tt})^2 + 4I_{zt}^2}.
\]

The critical band splitting is given by \( h_c(T) = 2\pi z T_c \). If the maximum of \( \mathcal{I} \) is achieved at \( Q = Q_c \), then the wave vector of the superfluid instability is \( q_c = 2h_c Q_c / v_F \). The parameter \( r = \ln(T_s / T_1) > 0 \) characterizes the relative strength of pairing in the singlet and triplet channels.

In the absence of dipole interaction, \( r = \infty \), the pairing is isotropic, and \( \mathcal{I}(Q, z) = \Psi(1/2) - (\text{Re} \Psi(1/2 - iz + izkQ)) \). At \( z \leq z_{\text{LOFF}} \approx 0.30 \), this function has a maximum at \( Q = 0 \), therefore the phase transition occurs into the uniform superfluid state. At \( z > z_{\text{LOFF}} \), i.e. at \( T < T_{\text{LOFF}} \), the maximum of \( \mathcal{I} \) is achieved at \( Q \neq 0 \), which corresponds to the LOFF state.

Since the dipole interaction lifts the spherical degeneracy of the critical temperature, the cases \( q \parallel \hat{z} \) and \( q \parallel q \) should be considered separately. In our numerical analysis we use the following values of the parameters, appropriate for the weak dipole interaction in the BCS limit: \( (N_p \lambda_0)^{-1} = 1.0 \) and \( \mu_2^p / \lambda_0 = 0.1 \), which gives \( T_t \approx 0.02 T_s \) and \( r \approx 3.93 \).

For \( q = q \hat{z} \), we find that at \( z \leq z_{\text{NMP}} \approx 0.23 \) the maximum of \( \mathcal{I} \) is at \( Q = 0 \), therefore \( I_{zt} = 0 \) and the phase transition occurs into the uniform singlet state. In contrast, at \( z > z_{\text{NMP}} \), i.e. at \( T < T_{\text{NMP}} \approx 0.68 T_s \), the critical temperature has two degenerate maxima at \( Q = \pm Q_c \hat{z} \neq 0 \), which corresponds to the phase transition into the state

\[
\begin{pmatrix}
\psi \\
d_z
\end{pmatrix} = \eta_1 \begin{pmatrix} 1 \\ \rho \end{pmatrix} e^{izqz} + \eta_2 \begin{pmatrix} 1 \\ -\rho \end{pmatrix} e^{-izqz}. \tag{11}
\]

Here \( \rho = -I_{zt} / [\ln(T_c / T_t) + I_{zt}] \), and the weights \( \eta_{1,2} \) of the two plane-wave components are determined by minimizing the full nonlinear free energy \( \mathcal{F} \), see below. For \( q \parallel \hat{z} \), there is no singlet-triplet mixing terms, and at \( T < T_{\text{LOFF}} \) one obtains the LOFF state with \( \psi \) modulated in the equatorial plane and \( d_z = 0 \). The transition line for this state is determined by the maximum of \( \mathcal{I}(Q, z) = -I_{ss} \), which yields the critical temperature slightly higher than in the isotropic LOFF case but still lower than in the NMP case.

Thus we come to the conclusion that it is the NMP state \( \text{NMP} \) that has the highest critical temperature below \( T_{\text{NMP}} \), see Fig. 1. One can show that as \( r \) varies from \( +\infty \) (no dipole interaction, \( T_t = 0 \)) to \( 0 \) (strong dipole interaction, \( T_t = T_s \)), \( T_{\text{NMP}} \) moves from \( T_{\text{LOFF}} \) towards \( T_s \). The parameter \( \rho \) is the measure of the triplet component admixture, which is zero at \( T = T_{\text{NMP}} \) and increases as temperature decreases [for the parameter values used above, \( \rho(T = 0) \approx 0.23 \)].

In order to determine the spatial structure of the NMP phase just below the critical temperature, one needs to evaluate the fourth-order terms in the free energy density:

\[
F_4 = \beta_1(|\eta_1|^4 + |\eta_2|^4) + \beta_2 |\eta_1|^2 |\eta_2|^2. \tag{12}
\]

The coefficients \( \beta_{1,2} \) are functions of temperature. The above expression is positive definite if \( \beta_1 > 0, \beta_2 > -2 \beta_1 \). At \( \beta_2 > 0 \) the minimum is achieved in the state \( (\eta_1, \eta_2) \sim (1, 0) \) or \( (0, 1) \), while at \( \beta_2 < 0 \) one has \( (\eta_1, \eta_2) \sim (1, e^{i\varphi}) \), where \( \varphi \) is an arbitrary phase. Near \( T_{\text{NMP}} \) one can set \( \rho = q_c = 0 \) in the expressions for \( \beta_{1,2} \), which gives \( \beta_1 = \beta_2 / 4 = -\langle N_F \phi^4 \rangle / 32\pi^2 T_{\text{NMP}}^2 \text{Re} \Psi''(1/2 - iz_{\text{NMP}}) > 0 \). Therefore, the NMP phase transition in the vicinity of \( T_{\text{NMP}} \) is of the second order, with \( (\psi, d_z) \sim (1, \rho) e^{i\eta z} \).
The determination of the full phase diagram, including the spatial structure of the order parameter at all $T < T_{NMP}$, is much more difficult, because of the competition between the NMP, LOFF, and uniform BCS phases. We leave this problem for future studies.

The origin of the NMP instability can be elucidated using the Ginzburg-Landau expansion of the free energy at small $h$ and $q = q \hat{z}$ in the vicinity of $T_{c}$. In addition to the usual uniform and gradient terms, the energy density contains mixed-parity terms, which are linear in gradients [19]:

$$F_{2} = \psi^{*} \hat{A}_{\psi} \psi + d_{s}^{*} \hat{A}_{d} d_{s} + i \tilde{K} h (\psi^{*} \nabla_{s} d_{s} + d_{s}^{*} \nabla_{s} \psi), \tag{13}$$

where $\hat{A}_{\psi} = N_{F} (T - T_{a}) / T_{a} - K_{a} \nabla_{s}^{2}$, $K_{a} = 7 \zeta (3) N_{F} (\phi_{a} v_{z}^{2}) / 16 \pi^{2} T_{a}^{2}$, and $\tilde{K} = 7 \zeta (3) N_{F} (\phi_{a} \phi_{b} v_{z}^{2}) / 4 \pi^{2} T_{a}^{2}$. Because of the last term, which favors a spatial modulation of the order parameter, at $h > h_{\text{NMP}} = \sqrt{N_{F} (T_{c} - T_{a}) / T_{a}^{2} \tilde{K}^{2}}$ the superfluid phase transition is of the NMP type. In contrast, the onset of the singlet LOFF instability is marked by the coefficient $K_{s}$ changing sign at $h_{\text{LOFF}} > h_{\text{NMP}}$.

To summarize, we found that in a homogeneous Fermi gas with s-wave attraction, the atomic dipole interaction produces triplet pairs that can mix with the singlet order parameter in a spatially modulated superfluid state. If the densities of atoms in different hyperfine states are unequal then the nonuniform mixed-parity state has higher critical temperature than the LOFF state, even at weak dipole interaction. To relate these findings with the cold atom experiments the effects of the trapping potential should be included. Another interesting open question concerns the fate of the NMP state in the strong-coupling regime near the BCS-BEC crossover.

Finally, we would like to remark that the possibility of a nonuniform singlet-triplet mixing is not restricted to polarized Fermi gases. For example, in a fully isotropic system the pairing interaction $V (k, k')$ in Eq. (2) depends only on the angle between $k$ and $k'$ and can therefore be factorized in terms of the spherical harmonics $Y_{lm} (k)$. In the singlet s-wave channel the order parameter is a scalar $\psi$, with $\phi_{s} (k) = Y_{00} (k) = 1$, while in the triplet p-wave channel the order parameter is now represented by three vectors $d_{m}$, with $m = 0, \pm 1$, corresponding to $\phi_{t,m} (k) = Y_{1m} (k)$. When expressed in terms of the Cartesian components of $k$, the triplet order parameter becomes a $3 \times 3$ complex matrix $A_{\alpha \beta}$. In the presence of magnetic field $h$, there is a contribution to the free energy density of the form $i \hbar \psi^{*} \nabla_{\alpha} A_{\alpha \beta} - c.c.$, which is invariant under all required symmetry operations: orbital and spin rotations, time reversal, and inversion. Similarly to the last term in Eq. (13) this can lead to the NMP instability at sufficiently high $h$.

This work was supported by the Natural Sciences and Engineering Research Council of Canada.

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