Induced p-wave superfluidity in imbalanced Fermi gases in a synthetic gauge field

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Abstract

We study the pairing formation and the appearance of induced spin-triplet p-wave superfluidity in dilute three-dimensional imbalanced Fermi gases in the presence of a uniform non-Abelian gauge field. This gauge field generates a synthetic Rashba-type spin–orbit interaction which has remarkable consequences in the induced p-wave pairing gaps. Without the synthetic gauge field, the p-wave pairing occurs in one of the components due to the induced (second-order) interaction via an exchange of density fluctuations in the other component. We show that this p-wave superfluid gap induced by density fluctuations is greatly enhanced due to the Rashba-type spin–orbit coupling.

Keywords: imbalanced Fermi gases, induced superfluidity, spin-triplet p-wave pairing gaps

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

1. Introduction

The tremendous improvement in the techniques of dealing with ultra cold atoms in recent years has paved the way for studying many-body quantum phenomena in an unprecedented manner [1–4]. The possibility of observing exotic states of matter, which may have analogies with condensed matter, quark matter and neutron star physics, such as topological phase transitions (TPT), Majorana fermions and color superconductivity, has greatly motivated the investigation of pairing and condensation of ultra cold Fermi systems under the influence of external electric and magnetic fields [5, 6]. While the laser field allows the spin–orbit coupling, the Zeeman magnetic field leads to an imbalance between the spin up and spin down chemical potential of the two fermionic species.

The first experimental realization and investigation of a spin–orbit coupled Fermi gas has been reported in [8]. A very recent and interesting experimental paper [9] reports the demonstration that singlet and triplet states are coupled with spin–orbit coupling, which is directly related to this work.

Motivated by these experimental realizations and other related works [10–12], as well as theoretical investigations in Fermi gases with spin–orbit coupling [13, 14], we study the ground state of dilute (spin 1/2) imbalanced Fermi gases. We investigate the manifestation of two possible induced p-wave pairing gaps, a ‘direct’ one, induced by a Rashba-type spin–orbit coupling (RTSOC) generated by a synthetic gauge field [15–17], and that induced by density fluctuations.

From the cold atom side, pairing formation in non-conventional systems (e.g. the one formed by a two-species imbalanced configuration) is of great interest in the investigation of mixtures of alkali atoms as, for example, lithium-potassium mixtures [18]. A multi-pairing system turns out to be very interesting when compared to the usual configuration, since the ground state now results from a competition not only of Fermi surface, chemical potential and mass mismatches, but also from the several pairing gaps that can simultaneously be present [19–22].

However, these recent investigations in imbalanced (non-conventional) systems without RTOC did not consider interactions between the same component. In other words, only s-wave (inter-species) pairing gaps have been taken into account. In this paper we consider pairing between the same
imbalanced Fermi gas with a RTSOC, described by the Hamiltonian:

$$H = H_0 + H_{SO} + H_{int},$$

where $H_0$ is the kinetic term, $H_{SO}$ is a Rashba-type spin–orbit interaction, generated by a synthetic gauge field, and $H_{int}$ is the term with a short-range $s$-wave interaction between the two fermionic species.

$$H_0 = \sum_{k,\sigma} \xi_{k,\sigma} c_{k,\sigma}^\dagger c_{k,\sigma},$$

$$H_{SO} = \sum_k \Delta(k) (e^{-i\lambda \hat{\mathbf{k}} \cdot \hat{\mathbf{S}}_F} - 1) + h.c.,$$

$$H_{int} = \sum_{k,k',\sigma} g(k, k') c_{k,\sigma}^\dagger c_{k',\sigma}^\dagger c_{-k',\sigma} c_{-k,\sigma},$$

species from the point of view of induced interactions that can emerge in the two situations we mentioned earlier.

In an imbalanced configuration, standard (BCS) $s$-wave pairing is energetically unfavorable to occur [23, 24]. In this adverse scenario, other kinds of pairing are expected to manifest [20, 21]. Intra-species $p$-wave pairing gaps induced by density fluctuations are our main interest in this work, since those gaps survive in the limit of vanishing RTSOC and chemical potential imbalance. Previous studies (without RTSOC) found that while the $p$-wave energy gain is parametrically smaller in weak coupling, in asymptotic regions of electrical harmonics. The $p$-wave triplet refers to angular momentum $l = 1$. The spherical harmonics $Y_{l,m}$ with $l = 1$ can be expressed as a linear function of the vector $\mathbf{k}$ as [27],

$$Y_{1,0}(k) \sim k_x + ik_y, \quad Y_{1,1}(k) \sim k_x - ik_y, \quad \text{and} \quad Y_{1,-1}(k) \sim k_y,$$

where $Y_{1,\pm 1}(k)$ represent a state with two nodal points, while $Y_{1,0}(k)$ a state with a nodal plane. The states with order parameter $\Delta_{so}(k) \sim Y_{1,\pm 1}(k)$ break time-reversal symmetry. As we will see below, the two induced $p$-wave triplet pairing gaps we find here are of this type.

It is instructive to diagonalize first $H_0 + H_{SO}$, which gives

$$H_0 + H_{SO} = \sum_{k,j=\pm} \xi_{k,j} a_{k,j}^\dagger a_{k,j}.$$
In equation (3) $a^{\dagger}_{k^+_i}(a_{k^-_i})$ is the creation(annihilation) operator for the state with helicity $(\pm_3)$, $\xi_{k^+_i} = \xi_k \pm H_k$, where $\xi_k = k^2/(2m) - \mu$, and $H_k \equiv \sqrt{\hbar^2 + \lambda^2 k^2}$.

In the next section we obtain the p-wave gap induced by a RTSOC, and calculate that induced by density fluctuations. Finally we discuss the relative importance of these two p-wave gaps in an expression which contain both contributions.

3. The triplet pairing gaps

3.1. P-wave triplet pairing induced by a RTSOC

Since we are mainly interested to investigate induced (both a RTSOC and density fluctuations) p-wave triplet pairings, we take a constant (i.e. $k$-independent) interaction strength $g < 0$, which is equivalent to consider only spin-singlet pairing. We treat the $H_{int}$ term using a BCS decoupling, and then write it in the basis that diagonalizes $H_0 + H_{SO}$ [28, 29],

$$H_{int} = \frac{1}{g} \sum_k \Delta_{s-}(k)a_{k^+_i}^\dagger a_{k^-_i} + \sum_{k,j=\pm_3} \Delta_{p}(k, j)a_{k^+_j}^\dagger a_{k^-_j} + h.c. \quad (4)$$

To regulate the divergence associated with the contact interaction term in $H_{int}$, we use [30]

$$\frac{1}{g} = \frac{m}{4\pi a_s} = \int \frac{d^3k}{(2\pi)^3}k^2 \quad (5)$$

Besides the regularization of the ultraviolet divergence present in the gap equation, this equation relates the strength of the contact interaction with the three dimensional scattering length $a_s$, which is more physically relevant, since it permits to make contact with current experiments.

In equation (4) we have defined $\Delta_{s-}(k) = \Delta_i(k)$, and [29]

$$\Delta_{s+}(k) = \frac{(k_x + ik_y)}{k_\perp} \Delta_p$$

$$\Delta_{s-}(k) = \frac{(k_x - ik_y)}{k_\perp} \Delta_p \quad (6)$$

where $k_\perp = \sqrt{k_x^2 + k_y^2}$, and

$$\left( \frac{\Delta_i(k)}{\Delta_p(k)} \right) = \frac{1}{2\sqrt{\lambda^2 k^2 + \hbar^2}} \left( \frac{2\hbar}{\lambda k} \right) \Delta. \quad (7)$$

$\Delta = -g \sum_k \langle c_{-k^-_i}c_{k^-_i} \rangle$ is the s-wave energy gap. Notice that $\Delta_i(k)^2 + 4\Delta_p(k)^2 = \Delta^2$. This shows that in an imbalanced ($h \neq 0$) Fermi system with RTSOC, the original s-wave order parameter ($\Delta$) has contributions from pairing with both the same and different helicity states [31].

The pairing gaps $\Delta_{s-}(k)$ and $\Delta_{s+}(k)$ are of the p-wave triplet type ($\sim \gamma_{1, \pm_3}(k))$ as we mentioned earlier. However, they vanish in the limit $\lambda \to 0$. This is the reason we have named them as ‘induced by RTSOC’ p-wave pairing gaps. The excitation spectra read $E_{k^\pm_3} = \sqrt{\xi_k^2 + |\Delta|^2 + H_k^2 \pm 2E_0}$, where $E_0 = \sqrt{\hbar^2(\xi_k^2 + |\Delta|^2) + \lambda^2 k^2 \xi_k^2}$. The vanishing of the excitation energies (i.e. $E_{k^\pm_3} = 0$) implies the equation

$$\left( E_{k^\pm_3} - H_k \right)^2 + (2k\lambda)^2 = 0, \quad (8)$$

where $E_k = \sqrt{\xi_k^2 + \Delta^2}$. Notice that the above equation will be satisfied only for $k = 0$ and $E_{k=0} = H_{k=0}$. This gives an equation for a critical field $h_c = \sqrt{\mu^2 + \Delta^2}$ at which a quantum phase transition to a topological superfluid state occurs [32, 33]. For a fixed value of $\lambda$ and in the limit $h \gg h_c$, it is found that the pairing gap which enters in equation (7) behaves as $\Delta \sim C/h^2$ (where $C$ is a constant depending on $\lambda$) both in 2D [32] and 3D [33]. Thus, in a highly imbalanced system, the ‘direct’ triplet p-wave pairing gaps decrease as $\Delta_p \sim C/h^3$. Therefore, in such an imbalanced configuration, when the imbalance tends to lead the system to the normal state, the direct triplet p-wave pairing gaps are vanishingly small. This is natural to expect, since the spin-triplet p-wave pairing induced by RTSOC occurs due to the direct short-range s-wave inter-component interaction [34].

As we will see next, there is a second-order interaction, that also generates a p-wave pairing gap, which is induced by density fluctuations. This investigation has been carried out previously without RTSOC in, for instance, [25] and [35]. It was found that the p-wave pairing gaps induced by density fluctuations are exponentially suppressed by both the strength of the coupling constant and asymmetry between the up and down species. We show that they are significantly enhanced when the effects of RTSOC are taken into account. More importantly, we discuss the limits the p-wave pairing gaps induced by density fluctuations overcome from first-order contributions.

3.2. P-wave triplet pairing induced by density fluctuations

We start the calculation of the induced interaction for the majority species, i.e. for the atoms with + helicity. The induced interaction was obtained originally by Gorkov and Melik-Barkhudarov (GMB) in the BCS limit by second-order perturbation theory [36]. For a (back-to-back) scattering process depicted in figure 2, with $p_1 + p_2 \to p_3 + p_4$, the induced interaction to lowest order in the s-wave channel is
given by [37],
\[ U_{\text{ind}}(p_1, p_2) = -g^2 \chi_{ph}(p_1 - p_2), \]
where \( p_i = (k_i, \alpha_i) \) is a vector in the space of wave-vector \( k \) and fermion Matsubara frequency \( \omega_i = (2l + 1)\pi/\beta \). The polarization function \( \chi_{ph}(p') \) is given by
\[ \chi_{ph}(p') = \frac{1}{\beta N} \sum_p G_{0\bar{0}}(p) G_{0\bar{w}}(p + p') \]
\[ = \int \frac{d^3k}{(2\pi)^3} \frac{\delta_{k_\perp} - \delta_{k_{\perp}+q_{\perp}}}{i\Delta_k + \xi_{k_{\perp}} - \xi_{k_{\perp}+q_{\perp}}}, \]
where \( p' = (q, \Omega) \), \( \Omega = 2\pi/\beta \) is the Matsubara frequency of a boson. The Matsubara Green’s function of a non-interacting Fermi gas is given by \( G_{0\bar{0}}(p) = 1/(i\omega - \xi_{k_{\perp}}) \). Notice that this polarization function is being calculated in the basis of the helicity states. Making the integrations above assuming that \( h \ll 2k_F \), we get
\[ \chi(|\bar{q}|) = -N_{\bar{F}\bar{L}} L(x^-) - N_{\bar{F}} \left[L(z) + F(x^-) - F(z)\right], \]
where \( N_{\bar{F}\bar{L}} = \frac{m k_F^2}{2\pi^2} \) is the density of states at the ‘deformed’ Fermi surface of the + species, \( N_{\bar{F}} = \frac{m k_F^2}{2\pi^2} \), \( x^+ = \frac{q}{2k_F} \), \( z = \frac{q}{2k_F} \), \( k_{\perp}^F = k_{\perp}^+ - k_{\perp}^- \), with \( k_{\perp} = m\lambda \), and \( q = |\bar{q}| \). \( L(x) \) is the static Lindhard function,
\[ L(x) = \frac{1}{2} - \frac{1}{4x} (1 - x^2) \ln \left[ \frac{1 - x}{1 + x} \right], \]
and we have also defined
\[ F(x) = -\frac{1}{x} \ln \left[ \frac{1 - x}{1 + x} \right] + \ln \left[ 1 - \frac{1}{x^2} \right]. \]

The Fermi momenta of the (free) + - atoms are found by \( \xi_{k_{\perp}+q_{\perp}} = 0 \), which read [29]
\[ k_{F,0} = \sqrt{2m(\mu + m\lambda^2)} \pm \frac{2}{\lambda} \frac{\hbar^2 + m\lambda^2(2m\lambda^2 + 2\mu)}{2m}, \]
Expressing the coupling \( g \) in terms of the s-wave scattering length to lowest order,
\[ U_{\text{ind}} = \left( \frac{4\pi a_s}{m} \right)^2 \left[ N_{\bar{F}\bar{L}}(x^-) + N_{\bar{F}}(L(z) + F(x^-) - F(z)) \right]. \]

In the scattering process conservation of momentum implies \( \bar{k}_1 + \bar{k}_2 = \bar{k}_3 + \bar{k}_4 \) which is set to zero. \( q \) is equal to the magnitude of \( \bar{k}_1 + \bar{k}_2 = \bar{k}_3 - \bar{k}_4 \), then
\[ q = \sqrt{\bar{k}_1^2 + \bar{k}_2^2} = \sqrt{\bar{k}_1^2 + \bar{k}_3^2 + 2\bar{k}_1 \bar{k}_3} = \sqrt{\bar{k}_1^2 + \bar{k}_3^2 + 2\bar{k}_1 ||\bar{k}_3|| \cos \theta}. \]
Since both particles are at Fermi surface of the atoms with + helicity, \( |\bar{k}_1| = |\bar{k}_3| = k_F^+ \), thus, \( q = k_F^+ \sqrt{2(1 + \cos \theta)} \). Then we have
\[ x^- = \frac{q}{2k_F^+}, \quad z = \frac{q}{2k_F^+}, \quad y = \frac{\sqrt{2(1 + \cos \theta)}}{2k_F^+}, \quad \text{where} \quad y = \frac{k_F^+}{k_F^+}, \quad \text{and} \quad z = \frac{k_F^+}{k_F^+} \]
\[ \text{ Taking the projection onto the Legendre polynomial } R_{n=1}(\cos(\theta)) \ [25, 35, 38] \]
\[ U_p = \frac{1}{2} \int_0^\pi \cos(\theta) \sin(\theta) d\theta \ U_{\text{ind}}(x), \]
we find
\[ U_p = -\left( \frac{4\pi a_s}{m} \right)^2 G_1(h, \alpha), \]
\[ \text{where} \quad G_1(h, \alpha) = N_{\bar{F}\bar{L}} L_1(\eta) + N_{\bar{F}}[L_1(\eta) + F_1(\eta) - F_1(\eta)], \]
with the following definitions
\[ L_1(\eta) = \frac{5\pi^2 - 2}{15x^4} \ln |1 - x^2| - \frac{x^2 + 5}{30x} \ln \left[ 1 - x \left( \frac{1 - x}{1 + x} \right) \right] - \frac{x^2 + 2}{15x^2}, \]
and
\[ F_1(\xi) = 2 \ln \left[ \frac{1 - x}{1 + x} \right] - \frac{x}{x^2} - \frac{1}{x} \ln \left[ 1 - x^2 \right]. \]
It is worth to mention that induced interactions have been used to obtain the transition temperature (or tricritical point) beyond mean-field both in three [37, 39, 40], and two dimensions [41–43]. We can express \( y \) and \( \eta \) in terms of the modified by RTSOC Fermi vectors \( k_{F,0} \)
\[ y = \sqrt{\frac{1 + 2a^2 + \sqrt{\hbar^2 + 2a^2(2a^2 + 2)}}{1 + 2a^2 - \sqrt{\hbar^2 + 2a^2(2a^2 + 2)}}}, \]
\[ \eta = \frac{1}{\alpha} \sqrt{1 + 2a^2 + \sqrt{\hbar^2 + 2a^2(2a^2 + 2)}}. \]
Here we have defined \( \alpha = \lambda/\nu_F \), where \( \nu_F \) is the Fermi velocity, and \( \hbar = \hbar/\mu \). The Fermi vector \( k_{F,0} \) and \( y \) are both real for \( h < \mu \), as is the case without RTSOC. An analogy with the mean-field analysis [25] leads to the (second-order) p-wave superfluid triplet intra-species pairing amplitude,
\[ \tilde{\Delta}_{++} \sim E_F \exp \left[ -\frac{\pi^2}{4a^2 \tilde{k}_F k_{F,0} f_1(h, \alpha)} \right], \]
where \( E_F \) is the Fermi energy of the fermions with + helicity, and \( f_1(h, \alpha) = L_1(\eta) + N_{\bar{F}}[L_1(\eta) + F_1(\eta) - F_1(\eta)] \).

Leading order in \( \alpha \), we can approximate \( k_F k_{F,0} f_1(h, \alpha) \sim (k_F^+ k_{F,0} - k_F^- k_{F,0}) L_1(\eta) = k_F^+ g_2(h, \alpha)L_1(\eta) \), where we have defined \( g_2(h, \alpha) = g_2(h, \alpha) = g_2(h, \alpha) \), with
\[ g_2(h, \alpha) = \alpha \sqrt{1 + 4a^2 - \hbar^2} \]
and
\[ g_5(h, \alpha) = \alpha \sqrt{1 + 4a^2 + \sqrt{\hbar^2 + 2a^2(2a^2 + 2)}}. \]
Finally we
and $a_F$ in the BCS range considered and has a $k_F a_F \gg k_F \gg k_F a_F$ increases with $\Delta$. The results of [25] are readily obtained in the limit where $a_F$ and $a_F$ are given by equation (22). The fact that we have calculated the $p$-wave pairing gaps induced by density fluctuations in the helicity basis allows us to write the total triplet $p$-wave pairing gaps in a unique expression,

$$
\Delta_{++}^{\text{total}}(k) = \Delta_{++}(k) + \Delta_{++}(k),
$$

where $\Delta_{++}(k)$ and $\Delta_{++}(k)$ are given by equations (7) and (23), respectively, and

$$
\Delta_{--}^{\text{total}}(k) = \Delta_{--}(k) + \Delta_{--}(k),
$$

where $\Delta_{--}$ is given by equation (24).

Now we discuss the limiting cases of the dominant gaps (interactions) namely, the induced by RTSOC or the induced by density fluctuations ones. In the limit of very small spin–orbit coupling ($\alpha \approx 0$), no matter the value of the $s$-wave interaction $g$, the second-order $p$-wave induced by density fluctuations will dominate. However, for small $g$ and strong spin–orbit coupling i.e. $\alpha \gg 1$ (or $\lambda \gg v_F$), the RTSOC-induced $p$-wave pairing gaps dominate.

It is worth to point out that in the deep BCS limit the observation of the enhancements calculated here would be almost impossible to observe experimentally and experiments would have to be performed closer to the Feshbach resonance, where the value of the pairing gap will become a noticeable fraction of the Fermi energy or so. However, in that case it is very likely that the simple analysis we performed in this work would be inaccurate and one might expect quite significant corrections, in particular from the frequency and momentum dependence of the induced interaction. Such a calculation is rather involved and is out of the scope of the present work. Such a calculation has been performed recently in the case where no spin–orbit interaction was present. By implementing the Eliashberg formalism with full frequency and momentum dependence, Bulgac and Yoon [46] observed a significant renormalization of the pairing correlations, when compared with a simple mean-field approximation, especially while approaching the Feshbach resonance. Thus, it is natural to expect that the same mechanism can have a similar effect in the presence of spin–orbit interaction, as considered here.

4. Conclusions

In summary, we have calculated the triplet pairing gap in an imbalanced Fermi gas induced by both a RTSOC, and that from a second-order effect on the particle interaction $g$, due to calculations show that the ground state is minimized with the states $\sim(k_\perp \pm ik_\parallel)$ [44, 45]. Then we obtain

$$
\tilde{\Delta}_{++}(k) \sim (k_\perp + ik_\parallel)\Delta_{++},
$$

$$
\tilde{\Delta}_{--}(k) \sim (k_\perp - ik_\parallel)\Delta_{--},
$$

where $\tilde{\Delta}_{+}$ is given by equation (22). The absolute value of the $p$-wave pairing gaps induced by density fluctuations in the helicity basis allows us to write the total triplet $p$-wave pairing gaps in a unique expression,

$$
\Delta_{++}^{\text{total}}(k) = \Delta_{++}(k) + \Delta_{++}(k),
$$

where $\Delta_{++}(k)$ and $\Delta_{++}(k)$ are given by equations (7) and (23), respectively, and

$$
\Delta_{--}^{\text{total}}(k) = \Delta_{--}(k) + \Delta_{--}(k),
$$

where $\Delta_{--}$ is given by equation (24).

Since the lower branch is being emptied by increasing $h_\perp$, the induced pairing between atoms of the negative helicity is strongly suppressed. It is given by $\tilde{\Delta}_{--} \sim E_F \exp[-(\pi/2 k_F a_F)^2 g_\perp(h, \alpha) L_1(y_2)]$, where $y_2 = \sqrt{1 + 2\alpha^2 + \sqrt{\hbar^2 + 2\alpha^2(2\alpha^2 + 2)}} + \alpha$.

The results of [25] are readily obtained in the limit $\alpha = 0$, since $k_F g_\perp(h, \alpha = 0) = k_F^2 k_F^2$, $y(\alpha = 0) = k_F^2/k_F^2$, and $y_2(\alpha = 0) = y(\alpha = 0)^{-1} = k_F^2/k_F^2$.

$$
\tilde{\Delta}_{++} \sim E_F \exp\left[-\left(\frac{\pi}{2 k_F a_F}\right)^2 \frac{1}{g_\perp(h, \alpha) L_1(y_2)}\right].
$$

Figure 3. Behavior of the upper branch helicity induced $p$-wave pairing gap over the Fermi energy, $\Delta \equiv \tilde{\Delta}_{++}/E_F$, as a function of the non-dimensional parameters $k_F a_F$ and $\alpha$.

With the induced interaction between intra-species atoms, equation (9), we can make a BCS-type calculation and obtain the second-order triplet pairing gaps (not only their amplitudes, as before), $\tilde{\Delta}_{++}(k)$ and $\tilde{\Delta}_{--}(k)$. Energetic
density fluctuations. We have shown that the first-order contribution dominates in the limit of strong RTSOC. However, the latter prevails in the limit of zero RTSOC, in which case the former induced p-wave pairing gap (which comes from the first-order ‘direct’ interaction) vanishes.

We expect that with the current experimental techniques it will be possible soon to detect (intra-species) p-wave gaps in highly imbalanced systems. Then, the detection of the enhancement of the p-wave gaps could be done turning on a RTSOC and comparing the resulting images with the ones without the RTSOC synthetic gauge field. Since a ‘weakly-paired’ p-wave superfluid with \( p_+ + iq \) symmetry in two dimensions is of special relevance because its vortices support zero-energy Majorana fermions and exhibit non-Abelian statistics [47], we hope our work can stimulate new investigations in this promising subject using our results.

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