Pair Breaking in Rotating Fermi Gases

Michael Urban\textsuperscript{1} and Peter Schuck\textsuperscript{1,2}

\textsuperscript{1}Institut de Physique Nucléaire, CNRS-IN2P3 and Université Paris-Sud, 91406 Orsay Cedex, France
\textsuperscript{2}Laboratoire de Physique et Modélisation des Milieux Condensés, CNRS and Université Joseph Fourier, Maison des Magistères, BP 166, 38042 Grenoble Cedex, France

We study the pair-breaking effect of rotation on a cold Fermi gas in the BCS-BEC crossover region. In the framework of BCS theory, which is supposed to be qualitatively correct at zero temperature, we find that in a trap rotating around a symmetry axis, three regions have to be distinguished: (A) a region near the rotational axis where the superfluid stays at rest and where no pairs are broken, (B) a region where the pairs are progressively broken with increasing distance from the rotational axis, resulting in an increasing rotational current, and (C) a normal-fluid region where all pairs are broken and which rotates like a rigid body. Due to region B, density and current do not exhibit any discontinuities.

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The surprising properties of superfluids become most evident if one looks at rotating systems. But the rotation does not only reveal the superfluidity, it can also destroy it. To give an example, in nuclear physics, the strong reduction of the nuclear moment of inertia compared to its rigid-body value is a direct consequence of superfluidity due to pairing correlations. But with increasing angular momentum, the pairing correlations are progressively destroyed and the moment of inertia increases to its rigid-body value. This pair-breaking effect of rotation was studied many years ago \textsuperscript{1}. In trapped atomic Fermi gases, the picture is somewhat different, since, contrary to the situation in atomic nuclei, the coherence length is much smaller than the system size. It is therefore possible to create quantized vortices or even vortex lattices \textsuperscript{2}, which allow the system to stay superfluid while rotating.

However, in a recent paper by Bausmehrt, Recati, and Stringari \textsuperscript{3} it has been argued that it may be possible to put a trapped Fermi gas adiabatically into rotation without creating vortices. In that paper, the destruction of superfluidity by rotation is described in a way which is very different from the nuclear physics case: Instead of decreasing the value of the pairing gap with increasing angular velocity, the authors assume that the system separates into a paired and an unpaired phase, while the properties of the paired phase itself are not affected by the rotation. The authors consider the unitary limit, where the energy densities of the paired and unpaired phases is determined by energy minimization: Near the rotational axis, the system prefers to stay superfluid, i.e., to stay at rest, since the paired phase has a lower energy density than the unpaired one. But beyond a certain distance from the rotational axis, the centrifugal energy which the system could win if it participated in the rotation becomes equal to the energy which is needed to break the pairs. Hence, the non-rotating superfluid core is surrounded by a rotating normal-fluid phase. At the interface separating the two phases, the density and the current are discontinuous.

This picture is very intuitive, but it is lacking the microscopic understanding of the pair-breaking mechanism. In the present paper we will therefore describe the rotating Fermi gas in the framework of BCS theory. The rotation is most easily described in the rotating frame, where the hamiltonian $\hat{H}$ (minus the chemical potential $\mu$ times the particle number $\hat{N}$) is given by

$$\hat{H} - \mu \hat{N} = \int d^3r \left[ \hat{\psi}^\dagger \left( \frac{\mathbf{p}^2}{2m} + V(r) - \Omega L_z - \mu \right) \hat{\psi}(r) + g \hat{\psi}^\dagger_\uparrow(r) \hat{\psi}^\dagger_\downarrow(r) \hat{\psi}_\downarrow(r) \hat{\psi}_\uparrow(r) \right],$$

where $\hat{\psi}$ is the Fermion field operator with components for (pseudo-)spin up ($\uparrow$) and down ($\downarrow$), $m$ is the atom mass, $\mathbf{p} = -i\hbar \nabla$ and $\mathbf{L} = \mathbf{r} \times \mathbf{p}$ are momentum and angular momentum, respectively, $V(r) = m(\omega_z^2 z^2 + \omega_y^2 r^2)/2$ is the axially symmetric trap potential and $g < 0$ is the coupling constant. The system is supposed to rotate with angular velocity $\Omega$ around the symmetry ($z$) axis of the potential.

If the system is large enough, such that the coherence length is small compared with the oscillator length associated with the trap potential, we can make use of the Thomas-Fermi (TF) or local-density approximation (LDA), which amounts to treating the system at each point $\mathbf{r}$ as uniform with a local chemical potential $\mu_\text{loc}(\mathbf{r}) = \mu - V(\mathbf{r})$. Then $\mathbf{p}$ becomes a number instead of an operator, and the “cranking” term $\Omega L_z$ can conveniently be written as $\Omega L_z = \mathbf{v}(\mathbf{r}) \cdot \mathbf{r}$, where $\mathbf{v}(\mathbf{r}) = \Omega \times \mathbf{r}$ is the velocity field corresponding to a rigid rotation. All quantities depend only parametrically on $\mathbf{r}$ via $\mu_\text{loc}(\mathbf{r})$ and $\mathbf{v}(\mathbf{r})$.

The gap, density, and current can all be derived from the normal and anomalous Matsubara Green’s functions $G$ and $F^\dagger$. They have to satisfy the Gorkov equations, which in the presence of the cranking term $\Omega L_z$ become

$$i\hbar \omega_n - (\xi + \mathbf{v} \cdot \mathbf{p}) G + \Delta F^\dagger = \hbar,$$

$$i\hbar \omega_n + (\xi + \mathbf{v} \cdot \mathbf{p}) F^\dagger + \Delta^* G = 0.$$
where we introduced the abbreviation $\xi = \xi(r, p) = p^2/(2m) - \mu_{\text{loc}}(r)$, $\omega_n$ denotes a fermionic Matsubara frequency, and $\Delta(r)$ is the gap. Note that we are neglecting the Hartree mean field, but anyway it would not qualitatively change our results in the BCS-BEC crossover regime \[4\]. Eqs. \[2\] and \[3\] can readily be solved for $G$ and $F$. They are formally similar to those describing pairing between particles with unbalanced populations (see, e.g., Ref. \[2\]), except that here the chemical potential which includes the effect of the centrifugal force, and $p_F$ is the corresponding local Fermi momentum. For a given $z$ coordinate, the condition \[6\] is fulfilled beyond a certain distance $r_{\perp}(z)$ from the rotational axis, since the velocity increases as $v = \Omega r_{\perp}$. At smaller distances, the energies $E_{\pm}$ are always positive, i.e., the system is in the usual superfluid phase and does not participate in the rotation. Beyond $r_{\perp}$, the gap is reduced by the rotation. We will call this region, where a rotational current exists although the gap is non-zero, the partially paired phase. Finally, at a certain distance $r_{\perp,2}$, the gap vanishes and the system enters the normal phase where it rotates like a rigid body.

In the case of a system without superfluid flow (like in our axially symmetric trap, as long as there are no vortices), the gap can be assumed to be real ($\Delta = \Delta^*$). The gap equation is obtained in the usual way by summing $F$ over $\omega_n$ and integrating over $p$, with the result

$$\Delta = \frac{4\pi\hbar^2a}{m} \int \frac{d^3p}{(2\pi\hbar)^3} \left[ \frac{\Delta}{2E} \right] \left[ 1 - f(E_+) - f(E_-) \right] \frac{m\Delta}{p^2}$$

(4)

where we defined the quasiparticle energies $E_{\pm} = E \pm p \cdot v$, with $E = \sqrt{\xi^2 + \Delta^2}$, and $f(E) = 1/(e^{E/(k_BT)} + 1)$ denotes the Fermi function, $T$ being the temperature and $k_B$ the Boltzmann constant. In Eq. \[4\], the divergence of the gap equation due to the contact interaction has been regularized in the usual way by expressing the coupling constant $g$ in terms of the $s$-wave scattering length $a$ \[3\].

We are mainly interested in the BCS-BEC crossover regime, where it is known that the BCS description fails at higher temperatures, and in particular the BCS prediction for the critical temperature $T_c$ is much too high. However, at zero temperature, BCS theory gives a reasonable description for the critical temperature $T_c$, which we will therefore restrict ourselves to the zero-temperature case, in which the Fermi function reduces to a step function, $f(E) = \theta(-E)$. Hence, the factor $[1 - f(E_+) - f(E_-)]$ is equal to 1 if both $E_+$ and $E_-$ are positive and 0 otherwise (at most one of the two energies $E_+$ and $E_-$ can be negative). In other words, states with $E_{\pm} < 0$ are excluded from pairing. In order to better understand the role of these states, let us look at the occupation numbers $\rho(r, p)$, which are obtained by summing $G$ over $\omega_n$:

$$\rho(r, p) = \frac{1}{2} \left[ 1 - \frac{\xi}{E} \right] [1 - f(E_+)] + \frac{1}{2} \left[ 1 + \frac{\xi}{E} \right] f(E_-).$$

(5)

For states with both $E_+ > 0$ and $E_- > 0$, this reduces to the usual BCS expression. But if a state at momentum $p$ has $E_- < 0$, its occupation number is equal to 1. The corresponding time-reversed state with momentum $-p$ has then $E_+ < 0$ and its occupation number is equal to 0. As we will see below, this gives rise to a normal-fluid (rotational) current.

It is easy to see that the energies $E_{\pm}$ can only become negative if the velocity $v$ exceeds a critical value such that

$$p_F v > \Delta.$$  

(6)

Here we have introduced the abbreviation $p_F$ = $\sqrt{2m\mu_{\text{loc}}}$, where $\mu_{\text{loc}} = \mu_{\text{loc}} + mv^2/2$ denotes the local chemical potential which includes the effect of the centrifugal force, and $p_F$ is the corresponding local Fermi momentum. For a given $z$ coordinate, the condition \[6\] is fulfilled beyond a certain distance $r_{\perp}(z)$ from the rotational axis, since the velocity increases as $v = \Omega r_{\perp}$. At smaller distances, the energies $E_{\pm}$ are always positive, i.e., the system is in the usual superfluid phase and does not participate in the rotation. Beyond $r_{\perp}$, the gap is reduced by the rotation. We will call this region, where a rotational current exists although the gap is non-zero, the partially paired phase. Finally, at a certain distance $r_{\perp,2}$, the gap vanishes and the system enters the normal phase where it rotates like a rigid body.

If the condition \[6\] is fulfilled, i.e., for $r_{\perp} > r_{\perp,1}$, one can easily see that the energies $E_{\pm}$ can become negative if the momentum lies between two limits $p_-$ and $p_+$ which are given by

$$p_\pm^2 = p_F^2 + m^2 v^2 \pm 2m \sqrt{p_F^2 v^2 - \Delta^2},$$

(7)

The integrand of the gap equation \[4\] is only affected by the rotation if $p$ lies between $p_-$ and $p_+$. Integrating Eq. \[4\] over the angle between $p$ and $v$ and dividing both sides of the equation by $\Delta$, we obtain

$$1 = -\frac{a}{\pi\hbar m} \left[ \int_0^\infty dp \left( \frac{p^2}{E} - 2m \right) - \int_{p_-}^{p_+} dp \left( \frac{p^2}{E} - \frac{p^2}{v} \right) \right].$$

(8)

The first integral is the same as in the gap equation without rotation while the second one is the contribution of the $f(E_{\pm})$ terms due to the rotation.

In the weak-coupling limit, when $\Delta \ll \mu_{\text{loc}}$, the pair-breaking effects appear already at extremely low angular velocities $\Omega$. In this case it is possible to evaluate the integrals in Eq. \[5\] analytically, and one can show that the critical velocity for which the gap disappears is given by $v_c = (e/2)\Delta_{\text{loc}}/p_F$, where $e = 2.71 \ldots$ denotes Euler’s number. Hence, for a given $z$ coordinate, the radial coordinates $r_{\perp,1,2}$ separating the fully paired from the partially paired and the partially paired from the unpaired phase, respectively, are the solutions of the equations

$$p_F(r_{\perp,1}, z)\Omega r_{\perp,1} = \Delta_{\text{loc}} = \Delta_{\text{loc}}(r_{\perp,1}, z),$$

$$p_F(r_{\perp,2}, z)\Omega r_{\perp,2} = e\frac{e}{2}\Delta_{\text{loc}}(r_{\perp,2}, z).$$

(9)

(10)

In the crossover regime, the situation is more complicated, since the gap $\Delta$ may be comparable with $\mu_{\text{loc}}$. Therefore the integrals have to be evaluated numerically. In addition, the rotation can now be much faster and the centrifugal force can lead to a sizeable change of the density profile and it is necessary to readjust the global chemical potential $\mu$ as a function of $\Omega$ in order to keep the total number of particles fixed. The density per spin state, $\rho(r)$, is obtained by integrating the occupation numbers over $p$. Using Eq. \[5\], one obtains

$$\rho(r) = \frac{1}{4\pi^2 \hbar^2} \left[ \int_0^\infty dp \left( 1 - \frac{\xi}{E} \right) + \int_{p_-}^{p_+} dp \left( \frac{p^2}{E} - \frac{p^2}{v} \right) \right] .$$

(11)
The second term arises from the $f(E_{\pm})$ terms and exists only if the condition (6) is fulfilled, i.e., beyond $r_{\perp 1}$. Between $r_{\perp 1}$ and $r_{\perp 2}$, the density goes smoothly from its value with pairing to the value without pairing, \[ \lim_{\Delta \to 0} \rho(r) = \rho_{\pm 2}(r)/(6\pi^2\hbar^3). \] Once we have calculated the density, we can obtain the total number of particles by integrating the density over space. This allows us to determine the value of the chemical potential.

An interesting quantity is the current density, which can be obtained by multiplying the occupation numbers with $p/\hbar$ and integrating over $p$. From Eq. (5) it is clear that for $r_{\perp 2} < r_{\perp 1}$, i.e., close to the rotational axis where the condition (6) is satisfied, the current vanishes as it should in the superfluid phase. Beyond $r_{\perp 1}$, the result can be given in closed form as

\[ j = \frac{(\rho_1^2 - \Delta_2^2)/\varpi^3)^{1/2}}{6\pi^2\hbar^3} v. \quad (12) \]

One sees that in the partially paired phase the current increases with decreasing gap and it correctly approaches its rigid-body limit if one approaches the unpaired phase: \[ \lim_{\Delta \to 0} j(r) = \rho(r)v(r). \]

Let us now discuss some numerical results. We consider a system with $N = 4 \cdot 10^5$ atoms ($2 \cdot 10^5$ atoms per spin state) in two cases: (a) close to the BCS limit, with $1/(k_F a) = -2$ [$k_F = p_F(r = 0)/\hbar$], and (b) at unitarity, i.e., in the limit $a \to \infty$. We do not consider the BEC side of the cross-over, since as soon as the chemical potential becomes negative, the energies $E_{\pm}$ are always positive, i.e., the molecules in the BEC phase are never broken by the rotation. For simplicity we choose a spherically symmetric trap ($\omega_z = \omega_\perp$), but this will not qualitatively change our results. In the figures, we will use the harmonic oscillator units set by the trap potential, i.e., $\hbar\omega$ for energies and $l_\hbar = \sqrt{\hbar/(m\omega)}$ for lengths.

Let us first discuss the BCS case. In this case the pairing is so weak that it does not appreciably influence the density (upper panel of Fig. 1). It is also very fragile, i.e., the moment of inertia, which can be calculated within linear response theory [9], must be measured at extremely low angular velocity. Already for an angular velocity as small as $\Omega = 0.05\omega$, the gap (second panel of Fig. 1) is zero in a large part of the system. Because of the small angular velocity, the centrifugal force has no effect on the density, either. Looking at the gap, one can clearly see the point $r_{\perp 1}(z = 0) = 5.1 l_\hbar$ where the results for the non-rotating (dashed line) and the rotating (solid line) system start to differ, and the point $r_{\perp 2} = 6.1 l_\hbar$ where the gap goes to zero. The three regions are even more evident in the current (lower panel of Fig. 1): The current starts to be non-vanishing at $r_{\perp 1}$ and it has a kink at $r_{\perp 2}$ where it reaches the rigid-body value.

More interesting are the results in the cross-over regime, where the gap is strong enough to support a relatively fast rotation. In Fig. 2 we display the density, gap and current (from top to bottom) for a system at the unitary limit rotating with $\Omega = 0.45\Omega_0$ (solid lines; for comparison, the density and gap of the corresponding non-rotating system are shown as the dashed lines). In this case, the centrifugal force leads to an oblate deformation in the limit of the cross-over, since as soon as the chemical potential increases with decreasing gap and it correctly approaches its rigid-body limit if one approaches the unpaired phase: \[ \lim_{\Delta \to 0} \rho(r) = \rho_\pm (r)/(6\pi^2\hbar^3). \]

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of the system: The chemical potential $\mu$ and the axial size $r_{\perp \text{max}} = \sqrt{2\mu/m/\omega}$ of the system, which is determined by $r_{\perp \text{max}} = \sqrt{2\mu/m(\omega^2 - \Omega^2)}$, increases. The increase of the radial size is visible in the upper panel of Fig. 2 where the density is shown as a function of $r_{\perp}$ for $\omega = 0$. The depletion of the density in the center is a consequence of the reduced chemical potential. This is also the reason why the gap in the center decreases with the rotation. Due to the strong pairing, the gap has a direct effect on the density. This is the reason for the kink in the density profile at $r_{\perp} = r_{\perp \text{2}}$. However, we stress that the density stays continuous at $r_{\perp \text{2}}$.

The fact that, in contrast to the results of Bausmerth et al. [9], the density, the gap, and the current remain continuous functions of $r_{\perp}$ is the main statement of the present paper. In fact, if we followed the arguments given in Ref. [9], we would find a similar discontinuity as they do. The only difference with their result would be the different numerical values of the parameters $\xi_0$ and $\xi_N$ which determine the relationship between the density and the local chemical potential $\mu_{\text{loc}} = \xi_0 h^2(6\pi^2 \rho)^{2/3}/(2m)$ or $\mu_{\text{loc}} = \xi_N h^2(6\pi^2 \rho)^{2/3}/(2m)$ in the superfluid and normal phase, respectively. In BCS theory without mean-field shift, one obtains $\xi_S = 0.59$ and $\xi_N = 1$, whereas the QMC results used in Ref. [9] are $\xi_S = 0.44$ and $\xi_N = 0.56$. If one excluded the possibility of an intermediate “partially paired” phase, as in Ref. [9], the system would have to split into a fully paired superfluid and a fully unpaired normal-fluid phase, and the density would have a discontinuity across the phase boundary with a ratio $\rho_N/\rho_S = (\xi_S/\xi_N)^{3/5}$, which gives 0.73 with the BCS results and 0.85 with the QMC results for $\xi_S$ and $\xi_N$. From this we see that, even if BCS theory is not capable to give the right numbers for $\xi_S$ and $\xi_N$, the ratio is semi-quantitatively correct. Anyway, if our results for the unitary limit might not be very precise, we believe that they are qualitatively correct and that between the ordinary normal and superfluid phases there will be a region in which some pairs are broken while others stay unbroken. In particular, we checked that in the region between $r_{\perp \text{1}}$ and $r_{\perp \text{2}}$ our energy density is lower than both that of the non-rotating superfluid and the rigidly rotating unpaired gas.

We emphasize that the existence of the intermediate region is not a finite-size effect, but it survives in arbitrarily large systems. For instance, if the trap was a flat potential well instead of a harmonic oscillator, the ratio of the two radii $r_{\perp \text{2}}$ and $r_{\perp \text{1}}$ would become $r_{\perp \text{2}}/r_{\perp \text{1}} = \epsilon/2 = 1.36$ according to Eqs. (9) and (10), independently of the size of the system and of the angular velocity of the rotation. In the harmonic oscillator the intermediate region is smaller since the gap decreases with increasing $r_{\perp}$ already in the non-rotating case.

This does not mean that finite-size effects do not play any role. For instance, the abrupt decrease of $\Delta$ for $r_{\perp} \rightarrow r_{\perp \text{2}}$ is an artefact of the TF approximation, which requires that all spatial variations be slow compared with the length scale set by the coherence length. A necessary condition for this is $\Delta \gg \hbar \omega$. In a true quantum calculation, the profiles of $\Delta$, $\rho$ and $|j|$ would be rounded and no sharp interface between the different phases could be defined. In addition, beyond a certain critical angular velocity $\Omega$, the gap should completely disappear, even on the rotational axis [10].

An interesting extension of the present work is to study a system which is deformed in the $xy$ plane, i.e., in the plane perpendicular to the rotational axis. This question is very important since it is impossible to put the system into rotation without such a deformation (of course, once the system rotates, the deformation can be switched off and the conservation of angular momentum ensures that the system keeps rotating). In the deformed case, also the superfluid part of the system has a non-vanishing current, with an irrotational velocity field. Another important question concerns the collective excitations of the rotating system, in particular the radial quadrupole mode whose precession is used in current experiments for measuring the angular momentum of the system [11]. In order to stay in contact with the experiments, temperature effects should be taken into account, too.

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