The discovery of fermion superfluids near Feshbach resonances is a major development in cold atoms physics. Not only the long sought goal of achieving a fermion superfluid has been realized, the resulting condensate exhibits many interesting “universal” properties due to the interplay between unitarity scattering and Fermi statistics. Even though the studies of these superfluids have just begun, another new and exciting direction has already been spun off. This is the physics of resonances with non-zero orbital angular momentum.

While most experiments on Feshbach resonance focus on s-wave scattering, there are many Feshbach resonances with non-zero momentum accessible by the same energy tuning method using external magnetic fields. Recently, Salomon’s group at ENS has reported evidence of p-wave molecules by sweeping a Fermi gas of $^6$Li through a p-resonance. It is therefore conceivable that p-wave or even higher angular momentum fermion superfluids can be realized in the future. Typically, the width of the resonance decreases with increasing angular momentum. One therefore expects that $\ell \neq 0$ superfluids will be harder to observe than their s-wave counterpart. Whether this can be achieved within current technology remains to be seen. The fact that the width of p-resonances can be resolved in recent experiments is very encouraging. In any case, the difficulty is not an intrinsic one. We hope that the novel properties of the $\ell \neq 0$ superfluid pointed out below will motivate searches for these new superfluids.

In this paper, we consider $\ell \neq 0$ pairing near a scattering resonance where interaction between particles is strongest. Our goal is to determine the ground state structure, their signature, and the possible existence of universal behavior in these systems. As a first step, we shall focus at $T = 0$. In the case of s-wave resonance, it has been shown that mean field theory is valid at $T = 0$, despite large fluctuation effects near $T_c$. Using a model potential that reproduces the exact two-body scattering amplitude in vacuum, we have found analytic solutions of the BCS problem for all $\ell \neq 0$ pairing at $T = 0$. Our findings are: (A) For p and f-wave pairing, the pairing states are $Y_{11}$ and $Y_{22}$ respectively, which are orbital ferromagnets that break time reversal symmetry and carry macroscopic angular momenta. For d-wave pairing, there is a degeneracy between $Y_{22}$ and a so-called “cyclic” state. (B) The criterion for determining the ground state structure of a $\ell \neq 0$ superfluid is that the energy gap must have minimum angular fluctuation. (C) Unlike s-wave superfluids, where energy per particle has a universal form $E_F(1 + \beta)$, the properties of $\ell \neq 0$ superfluids are not universal and are determined by the effective range $r_F$ of two body scattering. The energy per particle is proportional to $E_F|k_Fr_F| << E_F$, where $k_F$ is the Fermi wavevector. (D) Although dipolar energy is insignificant for superfluid pairing, it breaks the rotational symmetry and orient the angular momentum of the pair. (E) Experimentally, the nature of the pairing state can be easily revealed by the angular dependence of the momentum distribution of the fermions, whose orientation can be controlled by an external magnetic field through the dipole interaction. These results are established below.

To begin, we first point out a major difference between the $\ell \neq 0$ pairing of atomic gases and the p-wave pairing of superfluid $^3$He. In the latter case, the pairing interaction is rotationally invariant in spin space. The p-wave interactions between $^3$He pairs ($\uparrow\uparrow$, $(\uparrow\downarrow + \downarrow\uparrow)$, $(\downarrow\downarrow)$) are identical. In the recent ENS experiment, three different p-resonances are found for the $^6$Li pairs in spin states $|\uparrow\uparrow\rangle$, $|\downarrow\downarrow\rangle$, and $|\uparrow\downarrow\rangle$. When the pair $|\uparrow\downarrow\rangle$ is at resonance, the pairs $|\frac{1}{2}\downarrow\frac{1}{2}\rangle$ and $|\frac{1}{2}\frac{1}{2}\rangle$ are not and their interactions can be ignored. In other words, particle interaction is highly anisotropic in the “pseudo-spin” space for atomic Fermi gases near resonance.

**Setting up the pairing problem:** We consider a two component Fermi gas (denoted as $\uparrow$ and $\downarrow$) with a Hamiltonian $\hat{H} = \hat{H}_0 + \hat{V}$. $\hat{H}_0 = \sum_{\mathbf{k},\sigma = \uparrow, \downarrow} \epsilon_{\mathbf{k}} c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}$, $\epsilon_{\mathbf{k}} = \hbar^2 k^2/2M$, $M$ is the mass of the fermion, $\hat{V} = \Omega^{-1} \sum_{\mathbf{q}} \sum_{\mathbf{k},\mathbf{k}'} V_{\mathbf{q}} \delta_{\mathbf{k} + \mathbf{k}', \mathbf{q}} c_{\mathbf{k}\uparrow}^\dagger c_{\mathbf{k}'\downarrow} c_{\mathbf{k}'\uparrow} c_{\mathbf{k}\downarrow}$, where $V_{\mathbf{q}} = \int dr e^{-i\mathbf{q}\cdot\mathbf{r}} V(r)$ is the Fourier transform of the potential $V(r)$ between unlike fermions, and $\Omega$ is the volume. Interactions between like fermions will be set to zero. If $b$ is the range of the potential $V(r)$, then for wavevector $k$ such that $kb << 1$, the two body scattering
amplitude in the \( \ell \)-th partial wave channel is

\[
f_\ell(k) = \frac{(kb)^{2\ell}}{-a_\ell^{-1} + r_\ell k^2/2 - i(bk)^{2\ell} k},
\]

where \( a_\ell \) and \( r_\ell \) are the scattering length and effective range respectively. At resonance, \( a_\ell \) diverges, while \( r_\ell \) remains of microscopic size. In the case of a square well \( (V(r) = -|V| \) for \( r < b \), and 0 otherwise), it is straightforward to show that near resonance \( r_\ell = -b/\pi \) \( (2\ell - 1)!! \).

Recall also that the bound state energy \( E_b \) will appear as a pole in \( f(k) \) when \( k \) is continued analytically to the pure imaginary axis, \( k \to i\kappa \). Near resonance, we have

\[
E_b = \frac{2\hbar^2}{M a_\ell r_\ell}.
\]

It is clear that a bound state exists only when \( a_\ell r_\ell < 0 \).

The scattering amplitude \( f_\ell(k) \) is related to \( V(r) \) through the \( T \)-matrix \( T_\ell(k,k';E) \),

\[
f_\ell(k) = -\frac{M}{4\pi\hbar^2} T_\ell(k,k;2\kappa_\ell + 0^+),
\]

where \( T_\ell(k,k';E) \) satisfies the integral equation

\[
T_\ell(k,k';E) = V_\ell(k,k') + \frac{\int dp \frac{p^2 V_\ell(k,p) T_\ell(p,p';E_+)}{2\pi^2}}{E_+ - \hbar^2 p^2/M},
\]

with \( E_+ = E + i0^+ \), where we have used the expansion \( V_{k' - k} = 4\pi \sum \mathcal{V}_\ell(k,k') \delta_{m\ell} Y_{m\ell}(k) Y_{m\ell}^*(k') \).

In standard BCS theory, the ground states is \( |G\rangle = \prod_k (u_k + \sqrt{\kappa} a_k^\dagger a_{-k}^\dagger) |0\rangle \), \( |u_k|^2 + |v_k|^2 = 1 \). The coherence factor \( u_k \) and \( v_k \) are determined by minimizing

\[
\langle H - \mu N \rangle = 2 \sum_k (\epsilon_k - \mu) |v_k|^2 - \Omega^{-1} \sum_{k,k'} V_{k' - k} u_k^* v_{k'} u_k v_k,
\]

which gives \( u_k = \sqrt{(1 + \xi_k/E_k)/2} \), \( v_k/u_k = \Delta_k/(E_k + \xi_k) \), \( E_k = \sqrt{\xi_k^2 + |\Delta_k|^2} \), \( \xi_k = \epsilon_k - \mu \), and the energy gap \( \Delta_k \) satisfies

\[
\Delta_k = -\sum_p V_{k-p} u_p v_p = -\sum_p \frac{V_{k-p} \Delta_p}{2E_p}.
\]

The chemical potential \( \mu \) is determined by the number constraint

\[
n = \frac{N}{\Omega} = \frac{1}{\Omega} \sum_k \left( 1 - \frac{\epsilon_k - \mu}{E_k} \right).
\]

Since we are interested in the resonance physics of the \( \ell \)-th partial wave, we replace \( V_{k' - k} \) simply by its \( \ell \)-th angular momentum component, i.e. \( V_{k' - k} \to 4\pi \mathcal{V}_\ell(k,k') \delta_{m\ell} Y_{m\ell}(k) Y_{m\ell}^*(k') \).

Even with this (single harmonic) simplification for \( V_{k' - k} \), there is a serious complication which does not occur in s-wave pairing -- that eq. (10) does not have a single harmonic solution, since nonlinearity will force \( \Delta_k \) to have all spherical harmonics. While one expects on physical grounds that only a few harmonics around \( \ell \) are dominant, there is no simple way to extract such dominant piece and to calculate the dominant energy contribution.

To eliminate this technical complication, we adopt the viewpoint that all microscopic potentials that produce the same low energy scattering amplitude will describe the same low energy physics for the system. We can therefore replace the actual \( V_{k,k'} \) by a model potential that produces the same scattering amplitude \( f_\ell \), but allows a much easier solution for the gap equation. The most convenient model is a generalization of the separable potential used by Nozieres and S. Schmitt-Rink \cite{Nozieres},

\[
V_\ell(k,k') = \lambda_\ell w_\ell(k) w_\ell(k'),
\]

\[
w_\ell(k) = \frac{(k/k_\ell)\ell}{[1 + (k/k_\ell)^2]^{(\ell+1)/2}},
\]

where \( k_\ell \) is a momentum cutoff. With eq. (8), eq. (4) has the solution

\[
T_\ell(k,k';E_+) = t_\ell(E_+) w_\ell(k) w_\ell(k'),
\]

\[
\frac{1}{t_\ell(E_+)} = \frac{1}{\lambda_\ell} - \frac{1}{\Omega} \sum_k \frac{w_\ell(k)^2}{E_+ - 2\kappa_k}.
\]

For \( \ell \neq 0 \), eq. (11) has a low energy expansion \cite{Rmp},

\[
\frac{1}{t_\ell(E_+)} = \frac{1}{\lambda_\ell} + \frac{1}{\Omega} \sum_k \frac{w_\ell(k)^2}{2\kappa_k} + \frac{E_+}{\Omega} \sum_k \frac{w_\ell(k)^2}{4\kappa_k^2} + O(E^2\ldots)
\]

\[+\frac{i\pi}{\Omega} \sum_k w_\ell(k)^2 \delta(E - 2\kappa_k).
\]

Substituting eq. (10) and (11) into eq. (3), and noting that the last term in eq. (12) integrates to \( ibw_\ell(k)^2(4\pi \hbar^2/M)^{-1} \), we achieve the form eq. (11) provided \( \lambda_\ell \) and \( k_\ell \) are related to the physical parameters \( a_\ell \) and \( r_\ell \) as

\[
\frac{M}{4\pi\hbar^2 a_\ell(k_\ell b)^{2\ell}} = \frac{1}{\lambda_\ell} + \frac{1}{\Omega} \sum_p \frac{w_\ell(p)^2}{2\kappa_p^2} = \frac{1}{g_\ell},
\]

\[
r_\ell = -\frac{2\pi(bk_\ell)^{2\ell}}{\Omega} \sum_p \frac{w_\ell(p)^2}{\kappa_p^2} \left( \frac{\hbar^2}{M} \right)^2 - 2(\ell + 1) \frac{a_\ell k_\ell^2}{\alpha k_\ell^2}.
\]

More explicitly, we have

\[
\frac{1}{\lambda_\ell} + \frac{J}{2\pi^2 \hbar^2} \frac{Mk_\ell}{\pi} = \frac{1}{g_\ell}, \quad r_\ell = -\frac{1}{k_\ell} \left[ \frac{4I_\ell(k_\ell b)^{2\ell}}{\pi} - \frac{2(\ell + 1)}{a_\ell k_\ell^2} \right]
\]
\[ J = \int_0^\infty \frac{q^2 dq}{(1 + q^2)^{\ell+1}}, \quad I_n = \int_0^\infty \frac{q^{2(\ell-2)n} dq}{(1 + q^2)^{\ell+1}}. \]  

(16)

Near resonance, \( a_\ell \to \infty \), the second term in \( r_\ell \) in eq. (14) and (15) can be ignored.

**Solution of gap equation:** With the potential given by (21), eq. (5) has a solution

\[ \Delta_k = w_\ell(k) \sum_m C_m Y_{\ell m}(k), \]

(17)

where the coefficients \( \{C_m\} \) satisfy the equations \(-C_m/\lambda_\ell = 4\pi \Omega^{-1} \sum_{p,m} w_\ell(p)^2 Y_{\ell m}^*(\hat{p}) Y_{\ell m}(\hat{p}) C_{m'}/(2E_p)\). Using eq. (13) to express \( \lambda_\ell \) in terms of physical parameters \( a_\ell \) and \( r_\ell \), we recast this equation as

\[ \frac{C_m}{g_\ell} = 4\pi \sum_{p,m'} w_\ell(p)^2 Y_{\ell m}^*(\hat{p}) Y_{\ell m}(\hat{p}) C_{m'}/(2E_p) - \frac{1}{2E_p}. \]

(18)

The energy of the system, from eq. (15), is

\[ \frac{\langle H \rangle}{\Omega} = \frac{1}{\Omega} \sum_k \left[ \xi_k - E_k + \frac{|\Delta_k|^2}{2E_k} \right] + \mu n. \]

(19)

It is useful to separate the angular structure and magnitude of \( \Delta_k \) by writing

\[ \Delta_k = w_\ell(k) C \Delta(\hat{k}), \quad \Delta(\hat{k}) = \sum m \alpha_m Y_{\ell m}(\hat{k}), \]

(20)

where \( C^2 = \sum_m |C_m|^2 \) and \( \alpha_m = C_m/C \). Eq. (13) can then be written as

\[ -\frac{C^2}{4\pi g_\ell} = \frac{1}{\Omega} \sum_p |\Delta_p|^2 \left( \frac{1}{2E_p} - \frac{1}{2E_p} \right). \]

(21)

Our goal is to solve from eq. (21) and (7) the quantities \( C \) (and \( \Delta(\hat{k}) \) and \( \mu \)) as a function of \( n, g_\ell \), and \( r_\ell \), and then determine which solution \( (C, \Delta(\hat{k})) \) minimizes eq. (10).

In the following, we shall present an analytic solution for eq. (21) and (7) near resonance. We have also solved these equations numerically for all \( g_\ell, r_\ell \). For all the angular momenta we have studied, the two results are identical within the region of \( g_\ell \) where the analytic solution is valid. To derive the analytic solution, we write eq. (21) and (7) in dimensionless form by expressing all energies and wave-vectors in units of \( \epsilon_\ell \) and \( \kappa_0 \), i.e. defining \( \mu \equiv \frac{p \epsilon_\ell}{C} \), \( k \equiv \frac{q \kappa_0}{C} \), \( \alpha \equiv \frac{q \epsilon_\ell}{\kappa_0} \), \( \Delta_\ell = \kappa_0 w_\ell(q) C \Delta(\hat{q}) \).

Next, we assume that the solution of eq. (21) and (7) for \( \ell \neq 0 \) satisfy \( \mu \). \( \Delta_\ell^2 \ll 1 \). We will verify later that this is indeed the case when \( k_F r_\ell / a_\ell \ll 1 \). With this assumption, we can expand eq. (21) and (7) in \( \mu \) and \( \Delta_\ell^2 \). To the lowest order in these quantity, eqs. (21), (7) and (19) (expressed in terms of un-scaled variables) become

\[ -\frac{|C|^2}{4\pi g_\ell} = \frac{1}{\Omega} \sum_p |\Delta_p|^2 \mu - \frac{1}{\Omega} \sum_p \frac{|\Delta_p|^4}{4\epsilon_\ell^4}, \]

(22)

and

\[ \langle H \rangle/\Omega = n \mu - \frac{1}{\Omega} \sum_p \frac{|\Delta_p|^2}{2\epsilon_\ell^2}. \]

(23)

In deriving eq. (24), we need to include second order terms \( \mu^2, \mu \Delta_\ell^2, \Delta_\ell^4 \) due to cancelation of lower order terms. Note that these expansions will not work for \( s \)-wave since the sums in eqs. (22) to (23) are infrared divergent.

Combining eqs. (22) and (23), we have

\[ n \mu = n \mu_0 - \frac{|C|^2}{4\pi g_\ell}, \quad n \mu_0 = \frac{1}{\Omega} \sum_p \frac{|\Delta_p|^4}{4\epsilon_\ell^4}. \]

(25)

Using eq. (25) and eqs. (13) near resonance, (and noting that the \( 1/(k_0 a_\ell) \) term in eq. (13) can be ignored near resonance), we find two equivalent forms for \( C \),

\[ \frac{C^2}{4\pi g_\ell} = -\frac{n}{2 a_\ell r_\ell} - \frac{n E_b}{2}, \quad C = \frac{8(k_0 b)^\ell}{\sqrt{3} \sqrt[4]{k_F r_\ell}}. \]

(26)

Comparing eq. (26) and (27), we have

\[ \mu = \mu_0 + E_b/2. \]

(27)

The explicit form of \( \mu_0 \) can be obtained by evaluating the integral in eq. (25) and using the expression of \( r_\ell \) in eq. (15) near resonance, which is

\[ \mu_0 = \gamma \left( \frac{2\pi^2}{3(k_0 b)^2 I_2^2} \right) E_F |k_F r_\ell|, \quad \gamma = \int d\hat{p} |\Delta(\hat{p})|^4 \]

(28)

where \( I_n \)’s are given in eq. (16). Eq. (26) to (27) together with eq. (2) give \( C \) and \( \mu \) as a function of \( n, a_\ell \) and \( r_\ell \). From these equations, it is easy to show that \( \mu, \Delta_\ell^2 \ll 1 \) since \( k_0 b r_\ell < 1 \), our initial assumption \( \mu, \Delta_\ell^2 \ll 1 \) is valid. Note that the structure of the gap \( \gamma \) shows up only in \( \mu \) and not in \( C \).

Finally, using eqs. (21), (25), and (27), we obtain the energy density as a function of \( n, a_\ell \) and \( r_\ell \),

\[ \langle H \rangle/\Omega = n \mu - n \mu_0/2 = (n/2)(\mu_0 + E_b). \]

(29)

Eqs. (29), (25), and (27) imply that (i) the ground state has minimum angular fluctuation in the gap, i.e. \( \gamma \). (ii) Unlike the \( s \)-wave case where \( \mu \) is of order \( E_F \) near resonance, eq. (28) shows that the chemical potential \( \mu \) for \( \ell \neq 0 \) at resonance is greatly reduced from \( E_F \), by a non-universal factor \( |k_F r_\ell| \), reflecting a stronger interaction energy than the \( s \)-wave case. This is due to the fact that the energy of the bound state \( E_b \) for \( \ell \neq 0 \) grows much faster than that of \( s \)-wave away from resonance, since \( E_b \propto 1/(a_\ell r_\ell) \) for \( \ell \neq 0 \) whereas \( E_b \propto -1/a_\ell^3 \) for \( \ell = 0 \). The effect of this larger interaction energy also shows up
The superfluid is an “orbital ferromagnet” since all pairs are in the same 3-state. We then have
\[ \gamma_{\ell=1} \propto 2(A \cdot A)^2 + |A \cdot A|^2, \]
\[ \gamma_{\ell=2} \propto 2(Tr A^2) + |Tr A|^2, \]
and a similar but lengthier expression for \( \ell = 3. \)

For \( \ell = 1, \) the minimum occurs at \( A^2 = 0. \) This means \( A \propto \hat{x} + i\hat{y} \) or any rotation of it, which is \( \Delta(k) = Y_{11}(k) \) along an arbitrary angular momentum quantization axis. The superfluid is an “orbital ferromagnet” since all pairs are in the same \( Y_{11} \) state. This is remarkable for it implies that by sweeping across the Feshbach resonance, a superfluid with macroscopic angular momentum and broken time reversal symmetry will result.

The problems of minimizing \( \gamma \) for \( \ell = 2, 3 \) were solved by N.D. Mermin in the context of superfluid \(^3\)He \[8\], \[11\]. For \( \ell = 2, \) there is an accidental degeneracy. Both \( \Delta(k) = Y_{22}(k) \) and the “cyclic” state \( \Delta(k) \propto k_x^2 + e^{2\pi i/3}k_y^2 + e^{4\pi i/3}k_z^2 \) minimize \( \gamma_{\ell=2} \)[8]. This degeneracy can be resolved in higher order in \( |k_F r| \) and will be discussed elsewhere. For \( \ell = 3, \) the ground state is \( \Delta(k) = Y_{32}(k) \) along an arbitrary direction \[11\]. This state is also an orbital ferromagnet, even though it is not of maximum angular momentum state. At present, there are no exact solutions for \( \ell \geq 4. \)

Numerical Results: Note that although \( \mu \sim \mathcal{E}_F |k_F r| \) near resonance (eq.\[28\]), it has to recover to \( \mu \sim \mathcal{E}_F \) on the atomic side of the resonance whether \( a_\ell \) is negative and small. We have solved eq.\[24\] and \[7\] numerically and have shown this is the case. (See figure 1). Our numerical results are in exact agreement with eq.\[24\] and \[28\] near resonance.

The effect of dipolar energy: Dipolar energy \( V_D \) breaks rotational symmetry in real space. Since electron spins are polarized by the external magnetic field \( B, \) we have \( V_D = \frac{1}{2} \int U(r-r')\psi_\alpha^\dagger(r)\psi_\beta^\dagger(r')\psi_\beta(r')\psi_\alpha(r), \) \( U(r) = \mu_B^2(1-3(B \cdot r)^2)/r^3, \) where \( \mu_B \) is the electron Bohr magneton. Since dipolar energy per particle is \( \mu_B^2/\mu_n \) and since \( \mu_B^2\mu_n/\mu = (e^2k_0/m_e c^2)(M/m_e)(3\pi^2/2)^{-1} \sim 10^{-4} \) for \( k_0 \sim 10^8 \text{cm}^{-1}, \) dipolar energy is not strong enough to affect the gap structure. On the other hand, it can orient the pairing state. A straightforward calculation shows that \( \langle V_D \rangle = (2\pi^2\mu_B^2/\Omega) \sum_{k,k'}\Delta_k^2\Delta_{k'}^2/(4\mathcal{E}_k\mathcal{E}_{k'}), \) where \( q = k'-k. \) In the case \( \ell = 1, \) \( \Delta_k \propto A \cdot k. \) We have \( \langle V_D \rangle \propto -|B \cdot A|^2. \) Since \( A \propto \hat{x} + i\hat{y}, \) we have \( \langle V_D \rangle \propto + (B \cdot \hat{z})^2. \) The angular momentum of the pair will lie in the plane perpendicular to \( B. \)

Signature of the \( \ell \neq 0 \) superfluid: Eq.\[23\] shows that the momentum distribution \( n_p \) of the fermions is \( |\Delta_\ell|^2/(2\pi^2) \propto |\Delta(p)|^2. \) A measurement of the angular dependence of \( n_p \) therefore gives \( |\Delta(p)|^2 \) directly.

We have thus established results (A) to (E) mentioned in the Introduction. This work is supported by NASA GRANT-NAG8-1765 and NSF Grant DMR-0109255.

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