The \(p\)-wave polaron

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We consider the properties of a single impurity immersed in a Fermi sea close to an interspecies \(p\)-wave Feshbach resonance. We calculate its dispersion and spectral response to a radiofrequency pulse. In the presence of a magnetic field, dipolar interactions split the resonance and lead to the appearance of two novel features with respect to the \(s\)-wave case: a third polaron branch in the excitation spectrum, in addition to the usual attractive and repulsive ones; and an anisotropic dispersion of the impurity characterized by different effective masses perpendicular and parallel to the magnetic field. The anisotropy can be tuned as a function of the field strength and the two effective masses may have opposite signs, or become smaller than the bare mass.

Understanding the physics of a single impurity in a degenerate ultracold gas has been essential in discovering the phase diagram of spin-imbalanced Fermi mixtures \([1-3]\). Close to a Feshbach resonance, the impurity state becomes truly many-body in character. A large effort has been devoted to its understanding, resulting in an impressive agreement between theory and experiment in recent years \([4]\). The emergent picture is that the quasiparticle formed by the impurity interacting with the background gas can be fermionic (a “polaron”) or bosonic (a “molecule”).

So far, the polaron has not been studied close to higher partial wave resonances. Particularly interesting are \(p\)-wave resonances, as \(p\)-wave coupled superfluids are predicted to display a richer phase diagram than their \(s\)-wave counterparts \([5-8]\). For example, the Bardeen-Cooper-Schrieffer (BCS) and Bose-Einstein condensation (BEC) regimes of these superfluids are separate phases \([9]\), and each of these can be either chiral or polar \([7]\). When confined to two dimensions, the superfluid BCS phase is even topologically nontrivial \([10]\). Therefore it becomes crucial to understand the nature of the quasiparticles which could form such states. For \(^6\)Li and \(^{40}\)K, \(p\)-wave resonances between atoms in the same hyperfine state were found to have an extremely narrow magnetic width, of order \(\lesssim 1\)G \([11, 12]\). However, \(p\)-wave resonances in Li-K mixtures \([13]\) proved to have larger magnetic widths, in the range \(1 - 10\)G, indicating a stronger open-channel character of the \(p\)-wave interaction. Given the stability of magnetic fields in state of the art experiments \((\approx \pm 1\)mG), such resonances are now finally accessible for detailed study.

Here we study the \(p\)-wave polaron problem by considering a single impurity atom, labelled \(\uparrow\), immersed in a spin-\(\uparrow\) Fermi sea, assuming that the two atomic species have the same mass and are strongly interacting in the \(p\)-wave channel, as may be achieved close to a \(p\)-wave Feshbach resonance. We will neglect any background \(s\)-wave interactions (see below).

An interesting feature of \(p\)-wave Feshbach resonances is that they are generally split into a doublet according to the projection \(m_l\) of the relative angular momentum onto the magnetic field axis. As first discussed in Ref. \([12]\), the \(m_l = \pm 1\) resonances are shifted towards higher energy from the \(m_l = 0\) by the magnetic dipole-dipole interaction between the outer shell electrons in the presence of the external magnetic field. As we shall see, an important consequence is that when the energy splitting between

\[ 1 \]

\[ 0 \]

\[ -1 \]

\[ -2 \]

\[ -2 \]

\[ -1 \]

\[ 0 \]

\[ 1 \]

\[ 2 \]

\[ 3 \]

\[ 4 \]

FIG. 1. (color online). Spectral function of the \(p\)-wave polaron. The horizontal and vertical axes are \(- (k_0^2 v_\pm 1)^{-1}\) and \(\omega/\epsilon_F\). The \(m_l = \pm 1\) and \(m_l = 0\) resonances are located respectively at \(x = 0\) and \(x = \delta\), with \(\delta = 0\) (a), 0.5 (b), 1 (c), 2 (d). Thick lines are the dressed molecules with \(m_l = 0\) (solid) and \(m_l = \pm 1\) (dashed). Here \(-k_0/k_F = 10\).
the two molecular levels is larger than the width of the $m_l = 0$ molecule-hole continuum, a new polaron branch appears lying between the attractive and repulsive polaron branches which had been previously observed close to $s$-wave Feshbach resonances [14, 15]. This is reflected in the spectral function of the impurity atom displayed in Fig. 1. Additionally, we find a regime close to the $m_l = 0$ resonance in which the effective mass of the attractive polaron becomes negative (positive) along (perpendicular to) the magnetic field axis with the opposite behaviour near the $m_l = \pm 1$ resonance. Clearly, these two new features could lead to many-body effects (e.g. collective modes, possible superfluid states) which would be dramatically different from that of the $s$-wave system.

The impurity may also bind to a particle from the Fermi sea to form a diatomic molecule, and as the interactions are increased this molecular state becomes energetically favorable. We calculate the energies of the polaron and the molecule, and show that the position of the polaron-molecule transition shifts towards the BCS side for increasingly narrow resonances.

To model the $p$-wave Feshbach resonance we use the following two-channel Hamiltonian [7, 16], and work in units where $\hbar = 1$:

$$H = \sum_{p,\sigma = \uparrow, \downarrow} \frac{\mu^2}{2m} a_{\sigma p}^+ a_{\sigma p} + \sum_{q,\mu} \left( \epsilon_{\mu} + \frac{q^2}{4m} \right) b_{\mu q}^+ b_{\mu q}$$

$$+ \sum_{p, q, \mu} \frac{g(|p|)}{\sqrt{V}} \mathbf{P}_\mu \left( b_{\mu q}^+ a_{\downarrow q}^+ + a_{\uparrow q}^+ b_{\mu q} + a_{\downarrow q}^+ a_{\uparrow q}^+ b_{\mu q} \right).$$

For convenience we define bosonic operators $b_{\mu x, y, z}$ in terms of the closed channel $l = 1$ molecule operators $b_{x, y, z}$ as $b_{\mu x, y, z} \equiv (b_{x, y, z} + b_{-x, -y, -z})/\sqrt{2}, b_{\mu x, y, z} \equiv (b_{x, y, z} + b_{-x, -y, -z})/\sqrt{2}, b_{z} \equiv b_{0}$. We specialize to the case where the $m_l = \pm 1$ resonances are degenerate [17]. $a_{\mu}$ and $b_{\mu}$ are the creation and annihilation operators of fermions with mass $m$, and $V$ is the system volume. The resonance splitting caused by dipolar anisotropy may be modelled by a positive shift of the $m_l = \pm 1$ molecule energies: $\epsilon_{x, y} = \epsilon_v + \delta_0$ with $\epsilon_{x, y} = \epsilon_{\pm 1}$ and $\epsilon_v = \epsilon_0$.

Interactions couple the closed channel molecule to a pair of atoms in the open channel, and are described by the momentum-dependent coupling constant $g(|p|)\mathbf{P}_\mu$. The coupling vanishes above a cutoff $\Lambda$ of the order of $1/R_e$, the inverse van der Waals length. In the low energy limit, the physical results should not depend on the actual shape of $g$ and we choose the cutoff function to be proportional to a step function, $g(|p|) = g(\Theta(\Lambda - p))$.

In the absence of the Fermi sea, the bare propagator of the closed channel molecule can be read off from the Hamiltonian as

$$D^{cl}_{\mu \nu}(p, \omega) = \frac{\delta_{\mu \nu}}{\omega - \epsilon_v - p^2/4m + i\eta} \equiv D^{cl}_{\mu \nu}(p, \omega)\delta_{\mu \nu}. \quad (2)$$

The renormalized molecular propagator $D^{0}_{\mu \nu}$ is dressed by the polarization bubble $\Pi^0$ as shown in Fig. 2a. In vacuum $\Pi^0_{\mu \nu} \equiv \Pi^0(\delta_{\mu \nu}$ is diagonal and takes the form

$$\Pi^0(p, \omega) = \frac{1}{3} \int \frac{d^3q}{(2\pi)^3} \frac{q^2 g^2(q)}{\omega^2 - q^2/m - p^2/4m + i0}. \quad (3)$$

Thus the propagator in vacuum is also diagonal, i.e. $D^{0}_{\mu \nu}(p, \omega) = \delta_{\mu \nu}/\Pi^0(\delta_{\mu \nu})$.

The elastic scattering between a spin-1 and a spin-1 atom in the ladder approximation is described by a $T$-matrix. The vacuum $T$-matrix in the present problem is given by (see Fig. 2b)

$$T(k, k') = \sum_{\mu} \frac{g^2(k)\lambda_{\mu}k'_{\mu}}{\sqrt{\Pi^0(\delta_{\mu \mu})}} - \frac{1}{\Pi^0(\delta_{\mu \mu})}. \quad (4)$$

Note that if the detuning $\epsilon_{\mu}$ is independent of $\mu$ the interaction is isotropic and the $T$-matrix is proportional to $\mathbf{k} \cdot \mathbf{k}'$. However, the splitting of the $p$-wave resonance by the dipole-dipole interaction in the presence of a magnetic field breaks rotational symmetry. In order to relate the parameters of the model to the physical observables, we compare with the low energy expansion of the $p$-wave scattering amplitude $f(k, k') \equiv -(m/4\pi) T(k, k')$:

$$f(k, k') \approx \sum_{\mu} \frac{k_{\mu}k'_{\mu}}{v_\mu - k_{\mu}k_{\mu} + \frac{1}{2} k_0 k^2 - i k^4}. \quad (5)$$

Here $v_\mu$ is the state dependent scattering volume and $k_0$, with dimension of momentum, is the $p$-wave analogue of the effective range [18]. Evaluating the polarization bubble and inserting in Eq. (4) yields the relationships

$$v_\mu = -\frac{mg^2}{12\pi \left( \epsilon_v - \frac{mg^2}{18\pi^2} \right)}, \quad k_0 = -\frac{24\pi}{m^2 g^2} (1 + c_2), \quad (6)$$

with $c_2 \equiv m^2 g^2 \Lambda/6\pi^2$. The vacuum molecule propagator is then

$$D^{0}_{\mu \nu}(p, \omega) = \frac{-12\pi / (mg^2) \delta_{\mu \nu}}{-v_\mu^{-1} + \frac{1}{2} k_0 (m(\omega - p^2/4) - (p^2/4 - m\omega - i0)^{3/2}}. \quad (7)$$

Note that the parameter $k_0$ is naturally large and negative, of the order of the cut-off $\Lambda$ [19, 20].

The $p$-wave resonances in the many-body system are characterized by two dimensionless parameters. Of these

FIG. 2. (a) The renormalized propagator of molecules (thick wavy line). The thin wavy line is the bare molecule while the straight lines are fermions. (b) The diagram which leads to the $T$-matrix, Eq. (4).
\( \gamma \sim (1 + c_2)k_F/k_0 \) controls interactions at the scale of the Fermi momentum \( k_F \) (i.e. many-body physics). Resonances with \( \gamma \ll 1 (\gamma \gg 1) \) are termed narrow (wide) and quantum fluctuations are suppressed for \( \gamma \ll 1 \) [7]. The second parameter \( c_2 \) controls interactions at the scale \( \Lambda \) (i.e. few-body physics) [19] and distinguishes strongly coupled \( (c_2 \gg 1) \) from weakly coupled \( (c_2 \ll 1) \) systems. Identical fermions interacting close to a strong Feshbach resonance may form trimer states [20]. However, using the method of Ref. [19], we have checked that trimers consisting of two identical fermions and an impurity of equal mass do not form in vacuum in the absence of resonance splitting.

We now turn to the question of the many-body state of a spin-\( \downarrow \) impurity immersed in a spin-\( \uparrow \) Fermi sea at zero temperature. In the s-wave case both the attractive and repulsive polaron branches are accurately described by dressing the impurity with a single particle-hole excitation [2, 21]; we will adopt the same approximation by dressing the impurity with a single particle-hole excitation. Indeed, this approximation becomes exact far away from resonance, or in the limit of narrow resonances, \( \gamma \ll 1 \) [7]. The propagator of the impurity with momentum \( p \) and energy \( \omega \) in the medium is

\[
G_\downarrow(p, \omega) = \frac{1}{\omega - p^2/2m - \Sigma(p, \omega) + i\epsilon},
\]

with \( \Sigma \) the self-energy. The energy of the polaron satisfies \( E = \text{Re}[\Sigma(p, E)] \). The self energy in the single particle-hole approximation is given by the diagram in Fig. 3:

\[
\Sigma(p, \omega) = \int \frac{d^3q}{(2\pi)^3} n_{F\uparrow}(q) T(p, \omega; q, \epsilon_q),
\]

where the Fermi function \( n_{F\uparrow}(q) \) takes the value 1 if the state with momentum \( q \) is occupied, 0 otherwise. The off-shell two-particle scattering \( T \)-matrix in the medium is

\[
T(p, \omega; q, \epsilon_q) = g^2(|p - q|/2) \sum_{\mu \nu} \left( \frac{p - q}{2} \right)_\mu \left( \frac{p - q}{2} \right)_\nu \times \left\{ \left[ D^{\text{sp}}(p + q, \omega + \epsilon_q) \right]^{-1} - \Pi(p + q, \omega + \epsilon_q) \right\}^{-1}.
\]

Here \( (p, \omega) \) \((q, \epsilon_q)\) are the \( \uparrow \) \( \downarrow \) atom momentum and energy entering the \( T \)-matrix, with \( \epsilon_q = q^2/2m \). The polarization bubble in the medium is given by

\[
\Pi_{\mu \nu}(q, \omega) = \int \frac{d^3k}{(2\pi)^3} \frac{k_\mu k_\nu g^2(k)[1 - n_{F\uparrow}(k + q/2)]}{\omega - q^2/4m - k^2/m + i\epsilon},
\]

Since \( |q| \ll \Lambda \) it is useful to write \( \Pi = \Pi^0 + (\Pi - \Pi^0) \). Then, upon proper renormalization, the coupling in Eqs. (10) and (11) reduces to the bare value.

The polarization bubble \( \Pi \) is a tensor, as collisions with particles in the Fermi sea generally do not preserve the projection of the molecule’s angular momentum. However, \( D^{\text{sp}} \) is diagonal, and the matrix inverse in Eq. (10) may be performed as follows. We write the polarization bubble as

\[
\Pi_{\mu \nu}(p + q, \omega + \epsilon_q) \equiv A\delta_{\mu \nu} + B(p + q)\mu(p + q)\nu/|p + q|^2.
\]

In the matrix inverse of Eq. (10), sums containing only \( A \) or only \( B \) terms form geometric series. Resuming first the \( A \) terms, and successively adding the \( B \) terms, one finds

\[
T(p, \omega; q, \epsilon_q) = \frac{g^2}{4} \left[ D_- + \frac{D_+^2}{|p + q|^2/B - D_-} \right],
\]

with \( D_{\pm} = \sum_{\mu} \frac{(p + q)\mu}{|p + q|^2/\pm - B} \) and \( D_x = \sum_{\mu} \frac{(p + q)\mu}{|p + q|^2/\pm - A} \).

The above formalism allows us to calculate the spectral function of the impurity atoms, defined as \( A_\downarrow(p, \omega) = -2\text{Im}[G_\downarrow(p, \omega)] \), which gives the spectral response of the impurity to a radiofrequency pulse of frequency \( \omega \). In particular, the spectral function peaks at the energy of the quasiparticle states with a finite wavefunction overlap with the bare impurity, i.e. the polarons. Our results for the spectral function close to resonance are shown in Fig. 1 for several values of the dimensionless resonance splitting \( \delta \equiv \delta_0 2\pi/(mg^2k_F^2) \). In the absence of resonance splitting, the picture is qualitatively the same as in the s-wave polaron case [2, 22, 23] with an attractive (repulsive) polaron of energy lower (higher) than the bare impurity. However, an additional branch appears for finite resonance splitting. For small resonance splitting, \( \delta_0 \lesssim \epsilon_F \), the intermediate branch disappears in the \( m_l = 0 \) molecule-hole continuum. For larger resonance splitting the intermediate branch shows up as an isolated spectral line which is a repulsive polaron close to the \( m_l = 0 \) resonance, approximately a free impurity between the resonances, and an attractive polaron close to the \( m_l = \pm 1 \) resonance.

In addition to energy, the polaron is characterized by its quasiparticle weight, the residue, given by \( Z = [1 - \partial \text{Re} \{\Sigma(0, E)\}/\partial E]^{-1} \) evaluated at the quasiparticle energy. In Fig. 4 we display the energy and residue across resonance for various values of \( k_0/k_F \). We see that the absolute values of the polaron energies decrease for increasing \( |k_0/k_F| \) as expected since particle-hole fluctuations become suppressed [7]. Indeed, in this limit the self energy (9), in the absence of resonance splitting (\( v_{\mu} = v \),

![FIG. 3. The polaron self energy.](image-url)
takes the form

$$\Sigma(0, E) = \frac{-3}{2\pi m} \int_0^{k_F} \frac{q^4 dq}{v^{-1} + \frac{1}{2} mk_0(E + \epsilon_q/2) + i0}.$$  \hspace{1cm} (14)$$

As mentioned above, the impurity may also bind a particle from the majority Fermi sea to form a bosonic quasiparticle — a molecule. In the “Cooper pair” approximation (i.e., with no particle-hole pairs), the energy $E_{\text{mol}}$ of the molecular state is given by the pole of the $T$-matrix at $p = 0$ and $\omega = E_{\text{mol}} + \epsilon_F$. The latter approximation is an upper bound to the real energy of the molecule, which becomes exact for $|k_0/k_F| \gg 1$. On the BEC side, $E_{\text{mol}}$ tends to $E_b - \epsilon_F$, with $E_b$ the energy of the molecule in vacuum. The molecule energy is included in Figs. 1 and 4. The latter illustrates that as the resonance becomes more narrow (for increasing $-k_0/k_F$) the transition from a polaronic to a molecular ground state takes place further towards the BCS limit, as is the case close to a narrow $s$-wave Feshbach resonance [14, 24–26]. Note, however, that the limit of narrow $p$-wave resonances corresponds to small densities, while the opposite is true for $s$-wave resonances.

The presence of a magnetic field which splits the resonances also breaks isotropy since it differentiates between $m_l$ levels. One consequence is the appearance of a strongly anisotropic dispersion of the impurity close to either of the resonances. At small momenta, the dispersion may be written as $E(p) = E(0) + p_\parallel^2/2m_\parallel + p_\perp^2/2m_\perp$, where $p_\parallel$ and $p_\perp$ are the projections of $p$ along the magnetic field, and on the plane perpendicular to it. We compute the effective mass tensor, showing how, on the BEC side, its magnitude is generally smaller than the bare mass of the particle. This is allowed since, in this region, the polaron is no longer the ground state. Moreover, close to the $m_l = 0$ resonance $m/m_\parallel^*$ becomes negative while $m/m_\perp^*$ is still positive, the opposite behavior occurring in the vicinity of the $m_l = \pm 1$ resonance. This is illustrated in Fig. 5.

We now make a few final remarks. In the above, any background $s$-wave scattering between the two atomic species has been ignored. This is justified provided the background scattering is small, in particular for a $p$-wave resonance close to a zero-crossing of an $s$-wave resonance. At vanishing scattering energy the $s$-wave channel dominates. However, in the present system the low energy scale is set by $\epsilon_F$. Then the total cross section for scattering of two atoms in the $s$-wave channel is $\sigma_s = 4\pi a^2$ while the $p$-wave cross section is $\sigma_p = 12\pi|f_p(k_F)|^2$ [27]. If $|vk_F^3|^{-1} \gg |k_0/k_F|$ then $p$-wave scattering is dominant for $|vk_F^3| \gtrsim |k_Fa|$ while in the opposite case we require $|k_0a| \lesssim 1$. Both requirements are achievable provided $k_0$ is not too large.

Early experimental studies of $p$-wave Feshbach molecules showed that these were short-lived, with lifetimes in the range 2 to 20 ms [11, 28–30]. This was explained in Refs. [19, 20, 31] as due to relaxation processes as well as possible recombination into trimers. However, radiofrequency spectroscopy has been able to study in detail the complete spectral response of metastable many-body states of strongly-interacting fermions with an intrinsic lifetime of only 30ms [14] including their polaronic and molecular branches. Furthermore, a recent proposal to modify the $p$-wave resonant interaction by coupling to a long-range excited state with the use of an optical Feshbach resonance predicts the possibility of suppressing three-body recombination [32], potentially allowing for longer lifetimes.

In this Letter, we studied the properties of the $p$-wave polaron. These are observable by radiofrequency spectroscopy. In particular, the anisotropic dispersion rela-
tion may be studied using angle resolved photo emission spectroscopy. Our results are also relevant to studies of resonantly enhanced atom-dimer scattering near the appearance of confinement induced $p$-wave trimers \[33, 34\] in polarized gases, in which case the dimer acts as the impurity investigated here.

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