

## Measurement of the zero crossing in a Feshbach resonance of fermionic ${}^6\text{Li}$

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We measure a zero crossing in the scattering length of a mixture of the two lowest hyperfine states of  ${}^6\text{Li}$ . To locate the zero crossing, we monitor the decrease in temperature and atom number arising from evaporation in a  $\text{CO}_2$  laser trap as a function of magnetic field  $B$ . The temperature decrease and atom loss are minimized for  $B = 52.8 \pm 0.4$  mT, consistent with no evaporation. We also present preliminary calculations using potentials that have been constrained by the measured zero crossing and locate a broad Feshbach resonance at  $\approx 86$  mT, in agreement with previous theoretical predictions. In addition, our theoretical model predicts a second and much narrower Feshbach resonance near 55 mT.

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Recently, three groups have predicted the possibility of high-temperature superfluid transitions arising from the strong pairing of a two-component Fermi gas in the vicinity of a Feshbach resonance [1–3]. Transition temperatures as large as half the Fermi temperature are expected [1–4]. This is a much larger fraction than in the standard Bardeen-Cooper-Schrieffer theory. One promising atomic system is  ${}^{40}\text{K}$ , for which a Feshbach resonance has been observed [5] and degeneracy has been achieved [6,7]. Another promising candidate is  ${}^6\text{Li}$ , which is predicted to exhibit magnetically tunable Feshbach resonances in mixtures of the two lowest hyperfine states. Such transition temperatures are within the reach of current experiments involving degenerate  ${}^6\text{Li}$  Fermi gases [8–11]. Hence, locating the resonance is of importance for studies of atomic gas analogs of high-temperature superconductivity.

As is well known, a Feshbach resonance arises in a colliding two-atom system when the energy of the incoming open elastic channel is magnetically tuned into resonance with a bound molecular state of an energetically closed channel. The tuning dependence arises from the difference in the magnetic moments of the open and closed channels [12]. In the vicinity of a Feshbach resonance, the  $s$ -wave scattering length  $a(B)$  is described approximately by [4]

$$a(B) = a_{bg} \left( 1 - \frac{\Delta B}{B - B_0} \right), \quad (1)$$

where  $B$  is the applied field,  $a_{bg}$  is the background scattering length,  $B_0$  is the field at which the resonance occurs, and  $\Delta B$  is proportional to the strength of the coupling between the open and closed channels.

For an ultracold gas which is a mixture of the two lowest hyperfine states,  $|1\rangle$  and  $|2\rangle$ , of  ${}^6\text{Li}$  (the  $|1/2, \pm 1/2\rangle$  states in the low-field, total angular-momentum  $|f, M_f\rangle$  basis),  $s$ -wave collisions only occur between atoms in these two different spin states. Hence, at sufficiently low temperatures, the scattering length of the mixture is the scattering length of this collision. The incoming channel of the Feshbach resonance in this mixture is predominantly triplet in character [13].

Thus near the resonance,  $a_{bg}$  is nearly equal to the triplet  $a^3\Sigma_u^+$  scattering length which is large and negative [14]. The vibrational level that gives rise to the resonance is due to the most weakly bound vibrational level ( $v=38$ ) of the singlet  $X^1\Sigma_g^+$  potential. The zero-field binding energy of the  $v=38$  vibrational level has not been directly measured, and the uncertainty in the binding energy can lead to significant theoretical uncertainty in the location of the Feshbach resonance and the associated zero crossing.

Figure 1 shows an analytic model [15] of the scattering length as a function of magnetic field between 0 and 150 mT (1 mT=10 G), and is based on adjusting model parameters to fit the results of previous multichannel calculations [4,13]. The scattering length is predicted to cross zero between 50 mT and 60 mT for a mixture of states  $|1\rangle$  and  $|2\rangle$  in  ${}^6\text{Li}$ , and in Fig. 1 is shown to cross the horizontal axis near 54 mT.

Previously, we have shown [10,16] that the scattering length for a mixture of the  $|1\rangle$  and  $|2\rangle$  states is very small at zero magnetic field in agreement with predictions [13]. This was accomplished by confining the mixture in an ultrastable  $\text{CO}_2$  laser trap for several hundred seconds at  $\approx 100$   $\mu\text{K}$  with negligible evaporation. At magnetic fields of 10–30 mT, we have observed large evaporation rates suitable for achieving degeneracy [10]. Further, we observed scattering cross sections corresponding to a scattering length with a magnitude of  $540a_0$  at a magnetic field of 0.8 mT in a mixture of the  $|1\rangle$  and  $|3\rangle = |f=3/2, M_f=-3/2\rangle$  hyperfine states. The measured scattering cross section is consistent with a calculation based on the predicted [14] singlet ( $45.5a_0$ ) and triplet ( $-2160a_0$ ) scattering lengths [17].

In this paper, we observe the vanishing of the elastic-scattering cross section, and hence the vanishing of the elastic-scattering length, for a mixture of the  $|1\rangle$  and  $|2\rangle$  hyperfine states of  ${}^6\text{Li}$  in a magnetic field. This provides a first step in determining the location of a Feshbach resonance in this system. To determine the location of the zero crossing, we monitor the temperature drop and atom loss after 40 s of free evaporation in a stable  $\text{CO}_2$  laser trap as a function of the applied magnetic field.

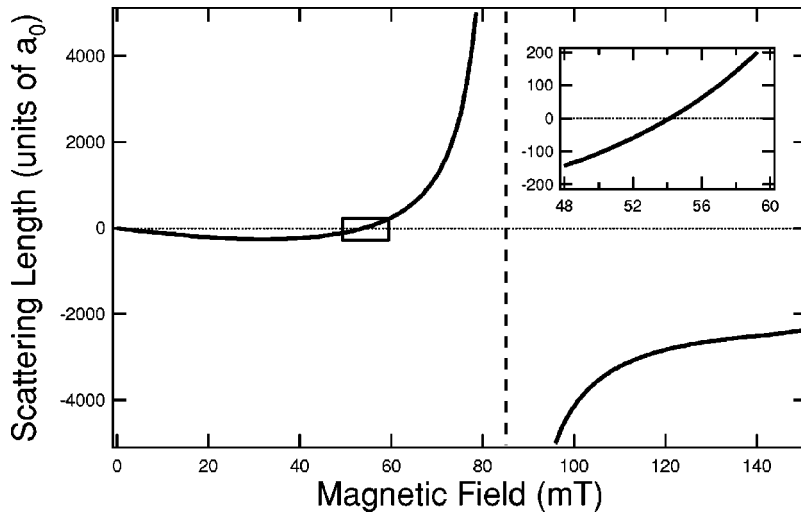


FIG. 1. Analytic model of the scattering length for a mixture of the two lowest hyperfine states of  ${}^6\text{Li}$  as a function of magnetic field showing the zero crossing (inset). ( $1a_0 = 0.052\,9177\text{ nm}$ .)

Our  ${}^6\text{Li}$  experiments employ a stable  $\text{CO}_2$  laser trap with a single focused and retroreflected beam [10]. The atoms evaporate from the trap, which has a fixed depth of  $\approx 0.75\text{ mK}$  and a geometric mean trap frequency of  $\approx 2.5\text{ kHz}$ . Since the scattering length for the  ${}^6\text{Li}$   $|1\rangle$  and  $|2\rangle$  mixture is extremely small at zero magnetic field, evaporation is turned on and off simply by applying or not applying a magnetic field.

The  $\text{CO}_2$  laser trap is continuously loaded from a  ${}^6\text{Li}$  magneto-optical trap (MOT) [10]. Typically,  $2 \times 10^6$  atoms in a nominally 50:50 mixture of the  $|1\rangle$  and  $|2\rangle$  states are initially contained in the trap at a temperature of  $140\ \mu\text{K}$ , as determined by the time-of-flight absorption imaging. Holding the atoms at zero magnetic field for an additional 40 s does not produce a measurable temperature change in the time-of-flight images, consistent with the prediction of a small scattering length at  $B=0\text{ T}$ . In free evaporation, the number of atoms decreases rapidly to  $\approx 1/3$  of the initial value as the hottest atoms are lost [10].

To measure the magnetic-field dependence of the evaporation rate, a pair of high field magnets with currents of 0–240 A are used to generate a uniform field of 0–110 mT in the trap region. The magnets have a diameter of  $\approx 20\text{ cm}$  and are located  $\approx 12.5\text{ cm}$  from the trap, producing negligible field curvature. The MOT gradient field is generated by reversing the current in one magnet for the trap loading stage. Each of the magnets is powered by a pair of current-regulated supplies which dissipate up to 10 kW per magnet.

In the experiments, the trap is initially loaded at zero magnetic field. Then a uniform magnetic field is applied for 40 s, during which time the atoms evaporate from the trap. Finally, the field is returned to 0 T and the temperature and number are measured by time-of-flight absorption imaging [10]. The final temperature obtained as a function of the magnetic field between 40 mT and 61 mT is shown in Fig. 2.

For magnetic fields far from the peak at 53 mT, the final temperature of the atoms is  $35\ \mu\text{K}$ , well below the Doppler cooling limit of  $140\ \mu\text{K}$  observed when the magnetic field is held at 0 T for 40 s. As the magnetic field approaches the peak, the observed temperature increases to  $87\ \mu\text{K}$ , which is below the loading temperature, but well above the minimum

temperature observed far from the zero crossing. The drop in the maximum temperature from  $140\ \mu\text{K}$  to  $87\ \mu\text{K}$  arises from 0.3 s during which the current-regulated supplies change current. Since the field is switched from 0 to 53 mT and back to 0 T, the atoms evaporate for an extra period of 0.6 s. By switching the magnetic field from 0 to 53 mT and back to 0 T without the 40-s holding time, we find that the final temperature is  $90\ \mu\text{K}$ . Hence, the observed reduction in temperature at the zero crossing is consistent with rapid evaporation during the switching process.

To confirm that the peak in temperature arises from a reduced cooling rate and not from heating, we also extracted the number of remaining atoms from the time-of-flight data, Fig. 3. Consistent with a reduced evaporation rate, we observe a peak in the number at 52.6 mT, nearly overlapping the peak in the temperature. Further, near the zero crossing field, we observe a  $\approx 30\%$  increase in the number of atoms with a negligible temperature change when the holding time is decreased from 40 to 20 s, indicating a negligible heating rate.

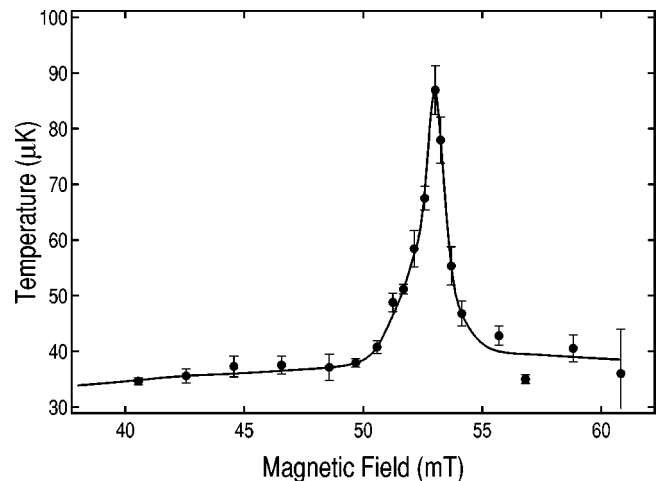


FIG. 2. Temperature after 40 s of free evaporation as a function of magnetic field. The maximum temperature occurs near the zero crossing of the scattering length where the evaporation rate vanishes. A curve has been added to guide the eye.

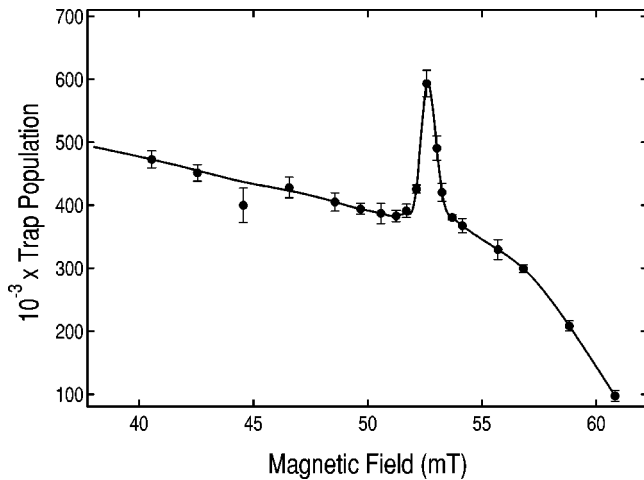


FIG. 3. Number of atoms remaining after 40 s of free evaporation as a function of magnetic field. The maximum number occurs near the zero crossing of the scattering length where the evaporation rate vanishes. A curve has been added to guide the eye.

We also observe a decrease in the remaining trap population as the magnetic field is increased, as shown in Fig. 3. This is accompanied by a slight heating, Fig. 2, and appears to arise from a broad magnetic-field-dependent loss feature that peaks near 65 mT. In this region, we observe loss and heating over time scales of 1 s in a sample at a temperature of 5  $\mu$ K and a density of  $3 \times 10^{13}$  per  $\text{cm}^3$ . This loss feature has also been observed by Dieckmann *et al.* [18]. Note that inelastic loss rates should not cause heating in the region of the zero crossing, since the elastic-scattering cross section is small and secondary scattering does not occur. However, the rapid increase in loss with increasing magnetic field may be partially responsible for the 0.4 mT shift between the peaks in the number (Fig. 3) and temperature (Fig. 2).

The magnetic field is calibrated using optical-absorption resonances to measure the splitting between the  $M_J = -3/2$  and  $M_J = -1/2$  levels of the  $2P_{3/2}$  excited electronic state. Acousto-optic (AO) modulators are used to produce two co-propagating probe beams differing by a tunable frequency up to  $\approx 1000$  MHz. One beam is  $\mathbf{z}$  polarized and interacts with the  $2S_{1/2}$ ,  $M_J = -1/2 \rightarrow 2P_{3/2}$ ,  $M_J = -1/2$  transition originating from the  $|2\rangle$  ground hyperfine state. The other beam is  $\mathbf{x}$  polarized and interacts with the  $2S_{1/2}$ ,  $M_J = -1/2 \rightarrow 2P_{3/2}$ ,  $M_J = -3/2$  transition originating from the  $|2\rangle$  ground hyperfine state. For fields near the zero crossing, note that  $|2\rangle \approx |m_s = -1/2, m_l = 0\rangle$ .

At the magnetic field corresponding to the zero crossing, the  $\mathbf{z}$  polarized beam is blocked and the dye laser frequency is adjusted to tune the  $\mathbf{x}$  polarized beam into resonance. The resonance peak is determined by spatially integrating a time-of-flight absorption image for each dye laser frequency. This method avoids degradation of the image by the diffraction arising from the small radial profile of the trapped atoms. The integrated absorption signal is insensitive to phase shifts arising from a nonzero index of refraction when the laser is detuned. The laser is then locked at the resonance frequency. Then the  $\mathbf{x}$  polarized beam is blocked and a similar method is used to tune the  $\mathbf{z}$  polarized beam into its resonance by ad-

justing an AO frequency. The excited-state frequency splitting is measured from the beat frequency of the two beams using a diode detector. At the magnetic field for the atom loss minimum, Fig. 3, we find a beat frequency of 994 MHz.

The magnetic field corresponding to the measured beat frequency is determined by calculating the expected splitting versus magnetic field, including the Zeeman, hyperfine, and fine-structure mixing contributions. From the 994 MHz splitting, the atom number peak (loss minimum) is determined to be at 52.6 mT. The minimum temperature decrease is then determined to be at 530 mT. Note that at 52.6 mT, magnetic-field-induced fine-structure mixing causes a quadratic Zeeman shift of  $\approx 12$  MHz in Li due to the small fine-structure splitting.

The uncertainty in the field determination arises primarily from the widths of the observed optical-absorption resonances, which are about 10 MHz or 0.5 mT in the vicinity of the zero crossing. The shape of each optical resonance exhibits some saturation arising from optical pumping. However, the centers are easily located with a statistical uncertainty of less than  $\pm 0.2$  mT. We report the location of the zero crossing as 52.8 mT, the average of the locations of the peaks in the temperature, Fig. 2, and number, Fig. 3. The uncertainty,  $\pm 0.4$  mT, is the root mean square or combined standard uncertainty of the systematic and statistical uncertainties. This value is dominated by the systematic error that we estimate as the shift between the temperature and number peaks (0.4 mT).

The measurement of the zero crossing can be used to further constrain the singlet potential and hence the singlet scattering length. Therefore we present preliminary results of a coupled-channel calculation that attempts to optimize the singlet potential. The interaction potentials in these calculations are constructed using the short-range singlet potential of Ref. [19], the short-range triplet potential of Ref. [20], the long-range dispersion coefficients of Ref. [21], and the methodology of Ref. [19] to connect the short- and long-range potentials. In addition, the triplet potential has been modified to be consistent with the measured binding energy of the most weakly bound state of the  $a^3\Sigma_u^+$  state of  ${}^6\text{Li}$  [14]. The singlet potential has been made to agree with the most recent prediction of the  ${}^6\text{Li}$  singlet scattering length [14,22]. Subtle mass-polarization corrections that affect the dissociation energy for  ${}^6\text{Li}$  and  ${}^7\text{Li}$  have been neglected.

The scattering length of the singlet potential can be varied within the uncertainty given by Refs. [14,22]. The corresponding uncertainty of the magnetic-field position of the zero in the scattering length does not exclude our observed zero crossing. In Fig. 4 we show the scattering length as a function of magnetic field, where the scattering length of the singlet potential has been fixed to obtain a zero crossing at 53 mT. Close examination of the zero crossing region has also revealed a narrow  $s$ -wave Feshbach resonance not found in previous calculations. This resonance is positioned at  $\approx 55$  mT between the zero crossing and the peak of the broad resonance centered at  $\approx 86$  mT. These two resonances correspond to hyperfine components of the same weakly bound singlet state. In fact, at zero field, the two states are labeled by  $F=0$  and 2, where  $F$  is the vector sum of the

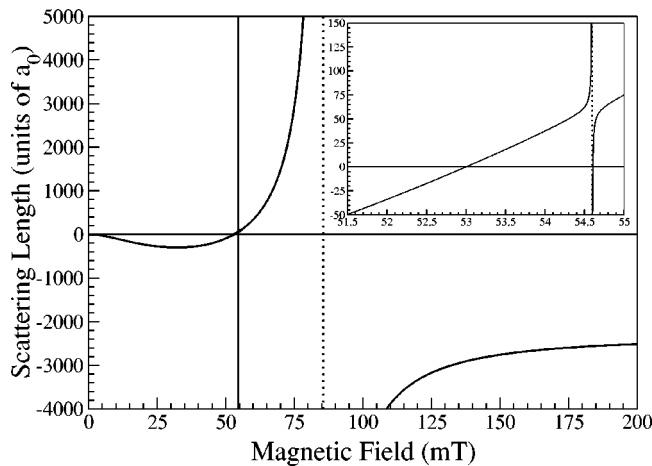


FIG. 4. Scattering length for collisions of  ${}^6\text{Li}$  atoms in the two lowest hyperfine states as a function of magnetic field. The singlet potential has been fixed to obtain a zero scattering length at 53 mT.

atomic spin  $f$  of the two atoms. A similar hyperfine structure has been observed in near-threshold singlet sodium spectroscopy [23]. Dieckmann *et al.* [18] have observed two loss features in this magnetic-field region, one broad and one narrow. The narrow feature at 55 mT is in agreement with the position of our predicted narrow resonance, but the width is an order of magnitude larger than we predict. No uncertainties are given for the positions of the resonances as these are preliminary results.

In conclusion, we have observed a zero crossing in the elastic-scattering length of a mixture of the two lowest hyperfine states of  ${}^6\text{Li}$ . The observed broad loss feature near 65 mT is not centered on the broad two-body Feshbach reso-

nance whose zero crossing is located at  $52.8 \pm 0.4$  mT and whose theoretical maximum is centered near 86 mT. We have used the measured zero crossing to further constrain the singlet potential and also predict a second, narrow Feshbach resonance at  $\approx 55$  mT. The observation of this resonance will require magnetic-field resolution beyond the capability of our present system. Further, the cross section near 86 mT is unitarity limited at the trap depths and temperatures employed in the present experiments. The evaporation rate is therefore insensitive to the scattering length, and nearly independent of magnetic field near the resonance. Since a Feshbach resonance does appear to exist in the magnetic-field region below 100 mT and the loss rates occur over time scales of 1 s at the densities and temperatures of interest, mixtures of the two lowest-lying hyperfine states of  ${}^6\text{Li}$  appear promising for studies of Fermi superfluidity.

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- [1] M. Holland, S. Kokkelmans, M.L. Chiofalo, and R. Walser, *Phys. Rev. Lett.* **87**, 120406 (2001).  
 [2] E. Timmermans, K. Furuya, P.W. Milonni, and A.K. Kerman, *Phys. Lett. A* **285**, 228 (2001).  
 [3] Y. Ohashi and A. Griffin, e-print, cond-mat/0201262.  
 [4] S. Kokkelmans, J.N. Milstein, M.L. Chiofalo, R. Walser, and M.J. Holland, *Phys. Rev. A* **65**, 053617 (2002).  
 [5] T. Loftus, C.A. Regal, C. Ticknor, J.L. Bohn, and D.S. Jin, *Phys. Rev. Lett.* **88**, 173201 (2002).  
 [6] B. DeMarco and D.S. Jin, *Science* **285**, 1703 (1999).  
 [7] G. Roati, F. Riboli, G. Modugno, and M. Inguscio, e-print cond-mat/0205015.  
 [8] A.G. Truscott, K.E. Strecker, W.I. McAlexander, G.B. Patridge, and R.G. Hulet, *Science* **291**, 2570 (2001).  
 [9] F. Schreck, L. Khaykovich, K.L. Corwin, G. Ferrari, T. Bourdel, J. Cubizolles, and C. Salomon, *Phys. Rev. Lett.* **87**, 080403 (2001).  
 [10] S.R. Granade, M.E. Gehm, K.M. O'Hara, and J.E. Thomas, *Phys. Rev. Lett.* **88**, 120405 (2002).  
 [11] Z. Hadzibabic, C.A. Stan, K. Dieckmann, S. Gupta, M.W. Zwierlein, A. Görlitz, and W. Ketterle, *Phys. Rev. Lett.* **88**, 160401 (2002).  
 [12] See, E. Tiesinga, B.J. Verhaar, and H.T.C. Stoof, *Phys. Rev. A* **47**, 4114 (1993), and references therein.  
 [13] M. Houbiers, H.T.C. Stoof, W. McAlexander, and R. Hulet, *Phys. Rev. A* **57**, R1497 (1998).  
 [14] E.R.I. Abraham, W.I. McAlexander, J.M. Gerton, R.G. Hulet, R. Cote, and A. Dalgarno, *Phys. Rev. A* **55**, R3299 (1997).  
 [15] The analytic model was provided by S. Kokkelmans.  
 [16] K.M. O'Hara, S.R. Granade, M.E. Gehm, T.A. Savard, S. Bali, C. Freed, and J.E. Thomas, *Phys. Rev. Lett.* **82**, 4204 (1999).  
 [17] K.M. O'Hara, M.E. Gehm, S.R. Granade, S. Bali, and J.E. Thomas, *Phys. Rev. Lett.* **85**, 2092 (2000).  
 [18] K. Dieckmann, C.A. Stan, S. Gupta, Z. Hadzibabic, C.H. Schunck, and W. Ketterle, e-print, cond-mat/0207046.  
 [19] R. Côté, A. Dalgarno, and M.J. Jamieson, *Phys. Rev. A* **50**, 399 (1994).  
 [20] C. Linton, F. Martin, A.J. Ross, I. Russier, P. Crozet, A. Yian-nopoulou, Li Li, and A.M. Lyyra, *J. Mol. Spectrosc.* **196**, 20 (1999).  
 [21] Z.-C. Yan, J.F. Babb, A. Dalgarno, and G.W. Drake *Phys. Rev. A* **58**, 2824 (1996).  
 [22] F.A. van Abeelen, B.J. Verhaar, and A.J. Moerdijk, *Phys. Rev. A* **55**, 4377 (1997).  
 [23] C. Samuelis, E. Tiesinga, T. Laue, M. Elbs, H. Knöckel, and E. Tiemann, *Phys. Rev. A* **63**, 012710 (2001).