Observation of Low-Field Feshbach Resonances in Collisions of Cesium Atoms

Vladan Vuletić, Andrew J. Kerman, Cheng Chin, and Steven Chu
Department of Physics, Stanford University, Stanford, California 94305-4060
(Received 6 October 1998)

We observe several Feshbach resonances in magnetic fields below 40 G for Cs atoms trapped in a 1D optical lattice. One resonance occurs in the lowest-energy ground state $F = 3$, $m_F = 3$ which is stable against inelastic binary collisions. This opens new possibilities for Bose condensation of Cs. When the elastic collision rate far exceeds the radial vibration frequency, we enter a 2D hydrodynamic regime where the thermalization rate is independent of density and temperature. Other resonances are also observed in the $F = 3$, $m_F = -3$ state through a dramatic increase in the dipolar relaxation rate. [S0031-9007(99)08457-4]

PACS numbers: 34.50.–s, 05.30.Jp, 32.80.Pj, 67.40.Hf

The collisional properties of atoms in their ground states play a crucial role for Bose-Einstein condensation (BEC) in dilute atomic gases [1]. The ratio of elastic to inelastic collision rates determines if evaporative cooling to BEC is possible [2]. Repulsive interactions lead to a stable condensate while attractive interactions permit only a metastable BEC state with a small number of atoms [3]. In most cases the energy and extension of the condensate are determined by the mean atomic interaction rather than the trap zero point energy [4]. The interaction also governs the evolution of the condensate phase and affects dephasing properties [5]. It determines the spectrum of collective excitations of the condensate and the coupling between the condensed and the noncondensed phases [6]. For condensates which are mixtures of different internal atomic states or different atomic species even the spatial pattern and the separation between the components are determined by the respective atomic interactions [7].

Since the sign and strength of the interaction determine virtually all of the condensate properties, an important goal in future BEC experiments is to develop ways to control and change these interactions, preferably in real time during a single experiment. In the low temperature limit the interaction between two ground state atoms is described completely by a single parameter, the s-wave scattering length a_s. In the vicinity of a Feshbach resonance [8] a_s varies dispersively through the whole continuum of values, thus changing not only the strength of the interactions, but also producing a region where their sign is reversed. Magnetically tunable Feshbach resonances in ultracold collisions have recently been observed in sodium through a change in the interaction energy of an optically trapped condensate [9] and in rubidium through a change in the cross section for photoassociation [10].

Of all alkalis, cesium can be optically cooled to the lowest temperatures in optical molasses [11]. However, all attempts to produce a cesium BEC using evaporative cooling in a magnetic trap have failed, due to its extraordinary inelastic collision properties [12,13]. The inelastic two-body decay from the doubly spin-polarized $F = 4$, $m_F = 4$ state was measured to be 3 orders of magnitude larger than theoretical predictions [12]. Evaporative cooling of the alternative $F = 3$, $m_F = -3$ state has stopped short of BEC for the same reason [13]. Optical traps provide a significant advantage over their magnetic counterparts with respect to inelastic binary losses, since they allow all states to be trapped, and thus provide the means to avoid exothermic two-body collisions altogether by using the lowest-energy ground state $F = 3$, $m_F = +3$. The scattering length for this state of cesium is known to be large and negative, however, which would limit the number of condensed atoms to very small values [14].

These extremely large elastic and inelastic cross sections hint at the existence of scattering resonances at or close to zero energy [15]. Recently several magnetically tunable Feshbach resonances have been predicted for the $F = 3$, $m_F = -3$ state in fields below 50 G, with one strong resonance expected at or near zero field [14]. These collisional resonances and the corresponding density-dependent level shifts would not only affect many precision experiments, such as atomic clocks [16], atomic interferometers [17], and tests of time-reversal symmetry [18], but are also of crucial importance to the prospects for Bose condensation of cesium. If a resonance were to exist in the stable ground state $F = 3$, $m_F = +3$, the field could be tuned to a region of positive scattering length. In addition, by tuning across the resonance the entire range of values for a_s could be explored, allowing continuous observation of the condensate from the strongly interacting (Thomas-Fermi) regime to the noninteracting regime, and of the collapse in the metastable regime. BEC in cesium is not only desirable in itself with respect to this atom's very unusual collisional properties but also would have practical importance as a bright source of cold atoms for future precision measurements [16–18].

In this Letter we report on several Feshbach resonances for ^{133}Cs in low magnetic fields, observed directly through changes in the collision rates of optically trapped ground state atoms. One strong resonance occurs in the $F = 3$, $m_F = +3$ state, contrary to the recent theoretical prediction [14]; additional resonances are also observed in the $F = 3$, $m_F = -3$ state. Further, we report evidence for

 0031-9007/99/82(7)/1406(4)$15.00

a 2D hydrodynamic regime where the collision rate exceeds the radial oscillation frequency of our trap by more than a factor of 20. In this regime the thermalization is mediated by collective rather than single-particle motion, and the thermalization rate is limited by the transit time of these excitations across the sample, which is close to the radial oscillation period and independent of temperature and density.

Our measurements of the elastic and inelastic collision properties are performed on a dense gas trapped in a far detuned 1D optical lattice. The trapping light with a wavelength of 1064 nm is obtained from a high power, single frequency Nd:YAG laser system we constructed. The output beam forms a vertical standing wave with a waist size of 250 μm and a running wave power of up to 17 W at the trap position. The dipole trap is loaded from a vapor cell magneto-optical trap (MOT), as described previously [19]. Up to 7×10^6 atoms are prepared in the far detuned 1D lattice, corresponding to several thousand pancake-shaped traps with peak occupation numbers of 1600 in the central region. The mean axial and radial kinetic energies are measured using a time-of-flight method with the light sheet located 12 cm below the trapping region. We determine the vibration frequencies with a parametric excitation method described previously [19]; when the trapping light intensity is modulated at twice the trap vibration frequency a resonant heating of the trapped gas is observed. Typical values for the axial and radial vibration frequencies are 90 kHz and 80 Hz, respectively. Atomic densities on the order of 10^{12} cm^{-3} are obtained from the measured atom number, temperature, and trap vibration frequencies. The trap lifetime of 1 sec is limited by the high cesium background pressure which is selected to allow fast loading of the MOT.

After loading the trap we optically pump the atoms either into the state $F = 3$, $m_F = 3$ or $F = 3$, $m_F = -3$ in a small bias field of 0.1 G. The populations in different m_F substates are determined by applying a specific microwave transition $F = 3$, m_F to $F = 4$, m_F' and subsequently monitoring the fluorescence on the $6S_{1/2}$, $F = 4$ to $6P_{3/2}$, $F = 5$ cycling transition. We achieve populations exceeding 85% in the desired $|m_F| = 3$ state, less than 10% in the $|m_F| = 2$ state, and less than 1% in each of the other m_F levels. The magnetic field is calibrated to better than 1% accuracy using the microwave transitions.

To study the elastic collisions in the $F = 3$, $m_F = 3$ state we slightly change the energy along the axial direction by either resonant parametric heating or fast degenerate Raman sideband cooling [19] and observe the subsequent collision-induced thermalization with the two radial degrees of freedom by measuring the axial temperature as a function of time after the cooling or heating has been switched off. The approach to equilibrium is very well fitted by an exponential time dependence and we define the thermalization rate as the inverse of the time constant as obtained from the fit. In Fig. 1 we plot this thermalization rate as a function of magnetic field for a temperature of 9 μK. This rate is proportional to the elastic cross section σ as long as the number of collisions required for thermalization does not depend on σ. In the limit where the de Broglie wavelength of the colliding particles is much longer than the range of the interatomic potential, the scattering of identical bosons is purely s wave with a cross section given by $\sigma = 8\pi a^2/(1 + k^2 a^2)$, where $\hbar k$ is the relative momentum of the particles [15]. In the vicinity of a Feshbach resonance the scattering length a can be approximated as $a = a_\infty(B - B_0)/(B - B_p)$, where a_∞ is the asymptotic value of the scattering length far from resonance, and B_0 and B_p are, respectively, the zero and pole field values [9]. We fit the expression $\Gamma_{\rm th} = \eta^{-1}\overline{n}\,\overline{\sigma}(a_\infty, B_0, B_p)v_{\rm rms}$ to our measured thermalization rate, where $\overline{n}$ is the mean density, v is the relative velocity of the colliding particles, and η is the mean number of collisions necessary for thermalization. The averages are performed over a thermal distribution with the measured temperature. From the fits of low density data we obtain $a_\infty = -(675 \pm 150)$ Å for the asymptotic value of the scattering length, which agrees with values obtained from a study of evaporative cooling in a magnetic trap [13]. For the zero and the pole of the scattering length we find $B_0 = 17.0(2)$ G and $B_p = 30(3)$ G. The relatively large uncertainty in B_p is due to the fact that a_∞ is extremely large. The collision cross section does not exhibit a pronounced peak at B_p since for our temperatures $ka_\infty \gg 1$ and consequently even far from resonance σ is already close to the unitarity limit of $8\pi/k^2$. The fit to the measured data also yields $\eta = 2.5$ for the number of collisions necessary for thermalization at 9 μK, in good agreement with Monte Carlo simulations for a magnetic trap [15]. The inset in Fig. 1 shows the

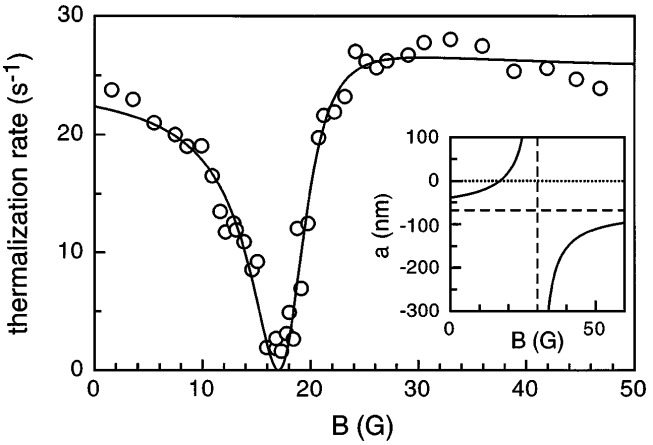

FIG. 1. The thermalization rate for the state $F = 3$, $m_F = 3$ between the axial and radial motions as a function of magnetic field for a temperature of 9 μK and a mean density of 1.5×10^{11} cm^{-3}. The solid line is a fit assuming a simple dispersive shape for the scattering length. The inset shows the scattering length as calculated from the best-fit parameters. The scattering length is positive between 17 and 30 G.

scattering length as a function of magnetic field with the parameters obtained from the thermalization rate fit. The scattering length is positive for field values between 17.0(2) and 30(3) G. Since we can experimentally control the magnetic field to a few mG accuracy, stable tuning of the scattering length is possible. We observe no loss of atoms in the resonance region. From the evolution of our atom number with the magnetic field held at the resonant value of 30 G we find an upper bound for the binary decay rate constant $G_{(3,3)} < 10^{-13}$ cm^3 s^{-1} and for the three-body decay $K_{(3,3)} < 5 \times 10^{-24}$ cm^6 s^{-1} at 10 μK.

For the opposite state $F = 3$, $m_F = -3$ the Feshbach resonances are easily detected through the accompanying exothermic two-body collisions, which lead to a loss of atoms from the trap. To measure the inelastic decay as a function of magnetic field we first optically pump the atoms to the $m_F = -3$ state. Subsequently we ramp the field in 2 ms to a certain value, hold it constant for a variable time t, and ramp it back to the initial offset value of 0.124 G. We then detect the number of atoms that have remained in $m_F = -3$. The number N of atoms per trap is fitted to $\dot{N} = -\gamma N - bN^2$, where γ is the independently measured decay rate due to background gas collisions. Since the atom number is related to the average density $\overline{n}$ by $N = \overline{n}V_{\text{eff}}$, where $V_{\text{eff}} = (4\pi k_B T/M)^{3/2}/\omega_1\omega_2\omega_3$ is the effective volume occupied by the trapped atoms, the two-body decay constant $G_{(-3,-3)}$ for the average density is related to the fitting constant b by $G_{(-3,-3)} = bV_{\text{eff}}$. Here T and $\omega_i/2\pi$ are the measured temperature and trap vibration frequencies, respectively, M is the mass of cesium, and k_B is Boltzmann's constant. We also perform an average over the Gaussian distribution of initial atom numbers per trap along the vertical direction, which amounts to a factor of 2 correction for the highest decay rates.

Figure 2 shows $G_{(-3,-3)}$ as a function of magnetic field for two temperatures, 14 and 4 μK. The error bars indicate the statistical uncertainty of the fits to the measured decay curves. The overall accuracy in G (40%) is estimated from the uncertainties in atom number (30%), temperature (20%), and vibration frequency (10%). We observe a strong Feshbach resonance near 31 G, a weaker one at 34 G, and possibly one at 22 G. We have verified that the observed decay is due to collisions among atoms in the $F = 3$, $m_F = -3$ state and not due to the small residual population in $F = 3$, $m_F = -2$. When we increase the population in the latter state to up to 50% we observe no change in $G_{(-3,-3)}$ for fields below 36 G. At 36 G we observe an additional resonance for inelastic collisions between $m_F = -3$ and $m_F = -2$ atoms, as evidenced by a linear density dependence of the decay rate both on the $m_F = -3$ and $m_F = -2$ populations. For the strong resonance at 31 G, the peak value of $G_{(-3,-3)}$ is 3 orders of magnitude higher than at 1 G. At the bias field $B_b = 0.124$ G the decay constant is below our detection sensitivity of 10^{-13} cm^3 s^{-1}.

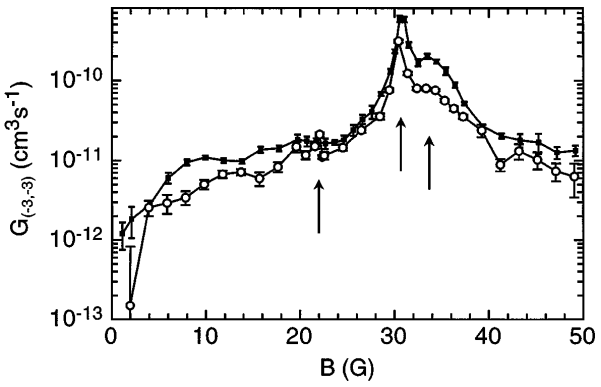

FIG. 2. The two-body loss rate coefficient for the state $F = 3$, $m_F = -3$ as a function of magnetic field for 14 μK (open circles) and 4 μK (solid squares). A strong resonance occurs at 31 G, a weaker one at 34 G, and possibly another at 22 G.

The interpretation of the above resonances in the elastic and inelastic cross sections as magnetically tunable Feshbach resonances is further supported by the fact that the spacing between the two strong resonances, their magnitude, and the overall pattern of the binary decay constant agree well with a recent prediction for cesium [14]. However, the positions of the resonances are shifted with respect to that prediction so that a strong resonance occurs in the $m_F = +3$ state.

Finally, we have investigated the temperature dependence of the peak thermalization and inelastic decay rates. For fixed trap vibration frequencies and atom numbers we use degenerate Raman sideband cooling [19] to prepare the atoms at different temperatures. We then observe the thermalization and inelastic decay rates in the unitarity limit by tuning onto the Feshbach resonances. Figure 3 shows these quantities as a function of temperature for two traps with different radial vibration frequencies. The collision rate varies as $\langle n v \sigma \rangle \propto T^{-3/2}T^{1/2}T^{-1} = T^{-2}$ in the unitarity limit and as T^{-1} for a constant cross section. Both the thermalization rate and the inelastic decay rate should exhibit the former dependence. Instead, we observe a constant thermalization rate close to the trap vibration frequency at low temperatures [Fig. 3(a)], while the peak value of the inelastic rate varies clearly as T^{-2} as expected [Fig. 3(b)].

This different behavior can be explained by the fact that at these temperatures the radial motion is in the hydrodynamic regime [20], where the elastic collision rate far exceeds the radial oscillation frequency of our trap. Thermal equilibrium cannot be established any faster than the transit time of thermal perturbations across the sample, which is on the order of the trap oscillation period. Inelastic loss from the trap, on the other hand, can occur locally in a single collision and therefore does not significantly change in the hydrodynamic regime.

In summary, we have detected one strong Feshbach resonance for the $F = 3$, $m_F = +3$ ground state of cesium through a change in the elastic cross section by more than a factor of 10. At lower temperatures the trapped gas is in the hydrodynamic regime and the thermalization

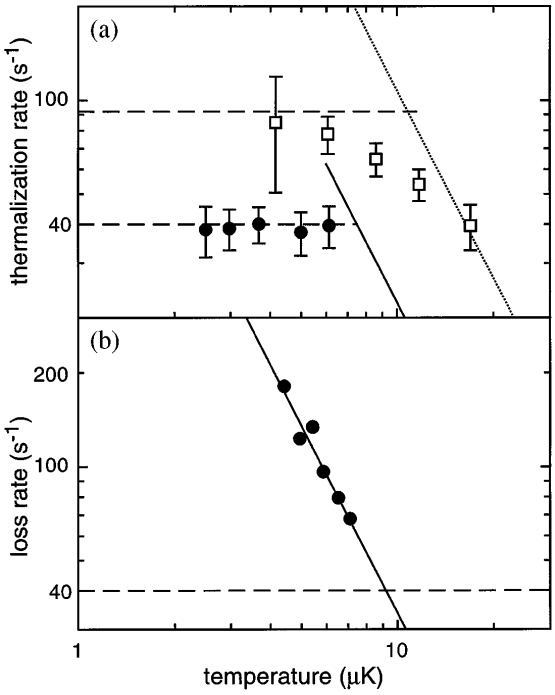

FIG. 3. (a) Thermalization rate for the state $F = 3$, $m_F = 3$ as a function of gas temperature for a trap with a radial vibration frequency of 40 Hz (solid circles) and 92 Hz (open squares). The solid and the dotted lines show the calculated collision rates. The thermalization rate is clearly limited by the radial trap vibration frequency (dashed lines), as expected for the hydrodynamic regime. (b) Peak loss rate from the state $F = 3$, $m_F = -3$ as observed at $B = 31$ G. The error bars are not visible on this scale. Since loss can occur locally in a single collision, the loss rate exceeds the trap vibration frequency of 40 Hz and scales as T^{-2}, as expected in the unitarity limit.

rate is limited by the trap vibration frequency. In the magnetically trappable $F = 3$, $m_F = -3$ state we have observed Feshbach resonances through the accompanying strong increase of the inelastic decay rate, which relates to the problems encountered in experiments aiming at BEC in this state [11,12]. The resonance in the lowest-energy ground state $F = 3$, $m_F = +3$ should allow the scattering length to be tuned to positive values, thus creating the possibility of a stable cesium condensate in an optical trap.

This work was supported in part by grants from AFOSR and the NSF. C. C. acknowledges support from the Taiwan government and V. V. from the Humboldt-Foundation.

[1] M. H. Anderson, J. R. Ensher, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Science **269**, 198 (1995); C. C. Bradley, C. A. Sackett, J. J. Tollett, and R. G. Hulet, Phys. Rev. Lett. **75**, 1687 (1995); K. B. Davis, M.-O. Mewes, M. R. Andrews, N. J. van Druten, D. S. Durfee, D. M. Kurn, and W. Ketterle, Phys. Rev. Lett. **75**, 3969 (1995).

[2] W. Ketterle and N. J. van Druten, Adv. At. Mol. Opt. Phys. **37**, 181 (1996).

[3] C. A. Sackett, H. T. C. Stoof, and R. G. Hulet, Phys. Rev. Lett. **80**, 2031 (1998); Yu. Kagan, A. E. Muryshev, and G. V. Shlyapnikov, Phys. Rev. Lett. **81**, 933 (1998).

[4] G. Baym and C. J. Pethick, Phys. Rev. Lett. **76**, 6 (1996); M.-O. Mewes, M. R. Andrews, N. J. van Druten, D. M. Kurn, D. S. Durfee, and W. Ketterle, Phys. Rev. Lett. **77**, 416 (1996).

[5] T. Wong, M. J. Collett, and D. F. Walls, Phys. Rev. A **54**, R3718 (1996).

[6] D. S. Jin, J. R. Ensher, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Phys. Rev. Lett. **77**, 420 (1996); D. S. Jin, M. R. Matthews, J. R. Ensher, C. E. Wieman, and E. A. Cornell, Phys. Rev. Lett. **78**, 764 (1997); M.-O. Mewes, M. R. Andrews, N. J. van Druten, D. M. Kurn, D. S. Durfee, C. G. Townsend, and W. Ketterle, Phys. Rev. Lett. **77**, 988 (1996); S. Stringari, Phys. Rev. Lett. **77**, 2360 (1996).

[7] C. J. Myatt, E. A. Burt, R. W. Ghrist, E. A. Cornell, and C. E. Wieman, Phys. Rev. Lett. **78**, 586 (1997).

[8] H. Feshbach, Ann. Phys. (N.Y.) **19**, 287 (1962); E. Tiesinga, A. J. Moerdijk, B. J. Verhaar, and H. T. C. Stoof, Phys. Rev. A **46**, R1167 (1992); E. Tiesinga, B. J. Verhaar, and H. T. C. Stoof, Phys. Rev. A **47**, 4114 (1993).

[9] S. Inouye, M. R. Andrews, J. Stenger, H.-J. Miesner, D. M. Stamper-Kurn, and W. Ketterle, Nature (London) **392**, 151 (1998).

[10] Ph. Courteille, R. S. Freeland, D. J. Heinzen, F. A. van Abeelen, and B. J. Verhaar, Phys. Rev. Lett. **81**, 69 (1998).

[11] D. Boiron, A. Michaud, P. Lemonde, Y. Castin, C. Salomon, S. Weyers, K. Szymaniec, L. Cognet, and A. Clairon, Phys. Rev. A **53**, R3734 (1996).

[12] J. Söding, D. Guéry-Odelin, P. Desbiolles, G. Ferrari, and J. Dalibard, Phys. Rev. Lett. **80**, 1869 (1998).

[13] D. Guéry-Odelin, J. Söding, P. Desbiolles, and J. Dalibard, Opt. Express **2**, 323 (1998).

[14] S. J. J. M. F. Kokkelmans, B. J. Verhaar, and K. Gibble, Phys. Rev. Lett. **81**, 951 (1998).

[15] M. Arndt, M. Ben Dahan, D. Guéry-Odelin, M. W. Reynolds, and J. Dalibard, Phys. Rev. Lett. **79**, 625 (1997); P. J. Leo, E. Tiesinga, and P. S. Julienne, Phys. Rev. Lett. **81**, 1389 (1998).

[16] S. Ghezali, Ph. Laurent, S. N. Lea, and A. Clairon, Europhys. Lett. **36**, 25 (1996).

[17] A. Peters, K. Y. Chung, B. Young, J. Hensley, and S. Chu, Philos. Trans. R. Soc. London A **355**, 2223 (1997).

[18] C. Chin, V. Leiber, V. Vuletic, A. J. Kerman, and S. Chu (to be published).

[19] V. Vuletic, C. Chin, A. J. Kerman, and S. Chu, Phys. Rev. Lett. **81**, 5768 (1998).

[20] D. M. Stamper-Kurn, H.-J. Miesner, S. Inouye, M. R. Andrews, and W. Ketterle, Phys. Rev. Lett. **81**, 500 (1998); A. Griffin, Wen-Chin Wu, and S. Stringari, Phys. Rev. Lett. **78**, 1838 (1997); Yu. Kagan, E. L. Surkov, and G. V. Shlyapnikov, Phys. Rev. A **55**, R18 (1997).