Adiabatic Cooling of Cesium to 700 nK in an Optical Lattice

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We localize Cs atoms in wavelength-sized potential wells of an optical lattice, and cool them to a three-dimensional temperature of 700 nK by adiabatic expansion. In the optical lattice we precool the atoms to ≈ 1 μK. We then reduce the trapping potential in a few hundred μs, causing the atomic center-of-mass distribution to expand and the temperature to decrease by an amount which agrees with a simple 3D band theory. These are the lowest 3D kinetic temperatures ever measured.

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Periodic optical potentials, or optical "lattices," are formed by the light shifts experienced by atoms in the interference pattern created by multiple laser beams. Studies of one-dimensional (1D) lattices observed atoms localized in individual potential wells with center-of-mass (c.m.) motion in the quantum regime [1–3]. Recently, optical lattices have also been demonstrated in two and three dimensions [4–6]. We report the use of a 3D optical lattice to simultaneously localize and cool Cs atoms to a temperature of ≈ 1 μK. This temperature is approximately 2 times lower than the minimum temperature of 2.5 μK [7] measured in Cs optical molasses [8]. Adiabatic expansion of the atomic c.m. distribution is used to reduce the momentum spread further, resulting in a sample of free atoms with a 3D kinetic temperature of 700 nK. This corresponds to an rms velocity less than twice the single photon recoil velocity. A simple band theory analysis predicts the adiabatic cooling limit as a function of lattice periodicity and initial temperature, in good agreement with our experiment. In an earlier experiment, Chen *et al.* reported 1D adiabatic cooling of a Li atomic beam passing through a strong standing wave, to a temperature of 12 μK (two recoil velocities) [9].

Our 3D optical lattice is formed by the interference of two pairs of linearly polarized laser beams intersecting, as shown in Fig. 1. The use of only four laser beams ensures an interference pattern which is unchanged by fluctuations in the phases of the beams, apart from an overall translation [5]. The directions of propagation and polarization of the lattice beams deviate from the geometry of Fig. 1 by less than 10 mrad, with intensity variations below $\pm 5\%$ across the relevant part of each beam profile. The lattice beams have a typical intensity of 0.5 mW/cm^2, and are tuned 25 linewidths below the $6S_{1/2}(F = 4) \rightarrow 6P_{3/2}(F' = 5)$ transition at $\lambda = 852$ nm. The resulting light field has points of pure σ^+ and σ^- polarization, each forming a centered tetragonal lattice, with a spacing between σ^+ and σ^- sites of $\lambda/2\sqrt{2}$ along $\hat{z}$ and $\lambda/\sqrt{2}$ along $\hat{x}$ and $\hat{y}$. Polarization gradient laser cooling causes Cs atoms to become tightly bound at these lattice sites [10], which are the minima of optical potential wells for the states $F = 4$, $m_F = \pm 4$. Repumping from the $6S_{1/2}(F = 3)$ hyperfine state is provided by a separate laser.

We load the lattice with Cs atoms by superimposing a magneto-optical trap (MOT) on the lattice volume. The MOT initially captures Cs atoms from a chirp-cooled atomic beam, producing a dense (10^{10} cm^{-3} [11]) sample of cold atoms in a volume ≈ 300 μm in diameter. The MOT magnetic field is then switched off, leaving an optical molasses which cools the Cs atoms to ≈ 3 μK. The molasses laser beams are extinguished, leaving the atoms in the presence of the optical lattice only. The atoms equilibrate in the lattice for 10 to 20 ms, after which adiabatic expansion is accomplished by a decrease in the lattice light intensity according to $I(t) = I(0)/(1 + \Gamma_A t)^2$, with a typical $\Gamma_A = 10^4$ s^{-1}. The expansion proceeds during an adjustable time τ, and is terminated when atoms are released

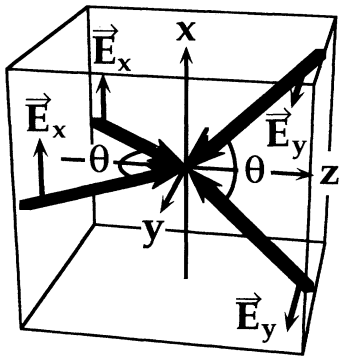

FIG. 1. Our optical lattice configuration. The lattice is formed by two pairs of linearly polarized beams. The z axis bisects the angle θ formed by each pair. The x axis is vertical and $\theta = 90°$.

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from the lattice by rapid (<1 μs) extinction of the lattice light. The 3D velocity distribution of the now free atoms is determined by measuring their spatial spreading during the 100 ms flight time to a 0.5 mm thick horizontal sheet of probe light located 5 cm below the lattice volume. The vertical velocity distribution (along $\hat{x}$) is obtained from the time dependent fluorescence from atoms falling through the probe beam, while the distributions along $\hat{y}$ and $\hat{z}$ are measured by imaging the fluorescence in the horizontal plane. We determine the rms momentum spread for each degree of freedom based on a Gaussian fit to the velocity distributions, and assign a "temperature" given by $3k_BT/2 = [\langle p_x^2\rangle + \langle p_y^2\rangle + \langle p_z^2\rangle]/2m$ as a measure of the kinetic energy of the atoms. At the lowest lattice light intensities the time-of-flight signal contains a background prior to the arrival of the cold atoms. This we ascribe to atoms that escape from the lattice during the initial equilibration phase, a phenomenon that, along with the general steady state conditions in the optical lattice, is still under study. With this background included in our fits the time-of-flight distribution is indistinguishable from a Gaussian within signal to noise [see Fig. 2(a)]. We estimate an uncertainty [12] of $\pm 3\%$ on $\langle p_x^2\rangle$, and $\pm 15\%$ on $\langle p_y^2\rangle$, $\langle p_z^2\rangle$. The larger horizontal uncertainties arise mostly from systematic error due to nonlinearity and inhomogeneity in the intensified charge coupled device imaging system. Differences between momentum spreads along the different axes are typically within estimated uncertainties.

As a first, albeit naive, model we assume that atoms localized near the bottom of deep optical potential wells have c.m. motion well approximated by a thermally excited 3D harmonic oscillator. The oscillation frequencies for atoms in the states $m_F = \pm 4$ are then

$$\omega_x = \omega_y = \frac{E_R}{\hbar}\sqrt{\frac{4\hbar|\Delta|s_0}{E_R}}, \quad \omega_z = \sqrt{\frac{88}{45}}\,\omega_x, \qquad (1)$$

where $E_R = (\hbar k)^2/2m$ is the photon recoil energy, $k = 2\pi/\lambda$ is the photon wave vector, Δ is the detuning, and $s_0 = 2\Omega_R^2/(4\Delta^2 + \Gamma^2)$ is the saturation parameter for a single lattice beam. The natural linewidth of the cooling transition is $\Gamma/2\pi = 5.2$ MHz, and Ω_R is the Rabi frequency defined so that $\Omega_R = \Gamma$ at an intensity of 2.2 mW/cm^2. The thermal excitation of each degree of freedom is described by a Boltzmann factor $f_B^i = \exp(-\hbar\omega_i/k_BT_i)$, where T_i is the temperature of that degree of freedom. The Boltzmann factor remains constant if the harmonic oscillator frequency is decreased adiabatically, so the oscillator temperature at time τ into the expansion is $T_i(\tau) = T_i(0)\omega_i(\tau)/\omega_i(0)$, approaching zero for $\tau \to \infty$. In our experiment the temperature does not go to zero, because the true periodic optical potential cannot be represented by a single harmonic oscillator. The harmonic oscillator approximation breaks down when the width of the c.m. distribution becomes comparable to the spacing between optical potential wells. As we show below, a more realistic model leads to a nonzero prediction for the final temperature.

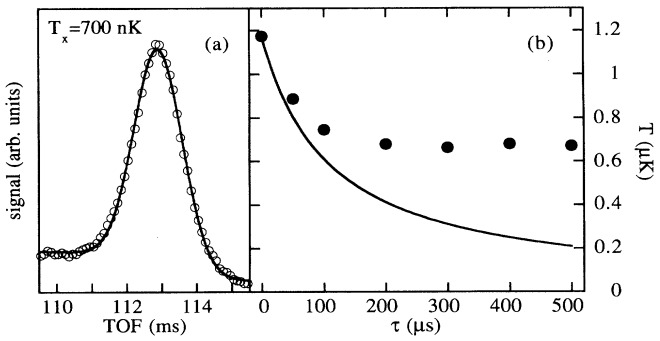

FIG. 2. (a) Typical time-of-flight spectrum. (b) Reduction in 3D temperature as the lattice light intensity is reduced at a rate $\Gamma_A \approx 1 \times 10^4$ s^{-1} (filled circles). The solid line is the reduction in temperature expected for a harmonic oscillator.

To satisfy adiabaticity one must have $|\dot{\omega}|/\omega = \varepsilon\omega$, where $\varepsilon \ll 1$ for all oscillation frequencies. The functional form given above of the decrease in lattice light intensity is chosen to satisfy adiabaticity with a time independent ε in the range 0.02 to 0.2. A lower limit on ε is set by the demand that heating be negligible during expansion. By operating at $\Delta = -25\Gamma$ (halfway to the neighboring $F' = 4$ hyperfine state), we achieve negligible heating on the time scales used.

The initial state of a trapped atom is characterized by oscillation frequencies which we calculate using Eq. (1), and Boltzmann factors determined from the observed initial momentum spreads. For conditions that produce the lowest temperatures, we find Boltzmann factors of $f_B^x \approx f_B^y \approx 0.55$ and $f_B^z \approx 0.45$, corresponding to rms c.m. spreads of order $\sqrt{\langle x^2\rangle} \approx \sqrt{\langle y^2\rangle} \approx \lambda/9$ and $\sqrt{\langle z^2\rangle} \approx \lambda/12$, and a steady state temperature of ≈ 1.2 μK. Figure 2(b) shows an example of the measured 3D temperature as defined above, as well as the prediction of the naive harmonic oscillator model, following various expansion times τ. Our data deviate noticeably from the harmonic oscillator model as early as $\tau = 100$ μs; by $\tau = 500$ μs, the measured temperature has clearly reached its final value, and differs from the harmonic oscillator prediction by more than a factor of 3.

To verify that the rate of expansion is slow enough to ensure adiabaticity, yet fast enough to avoid heating, we have varied the rate Γ_A at which the lattice light intensity is reduced. Less than a 10% change in the final temperature is observed when Γ_A is varied by an order of magnitude, from 2×10^3 to 2×10^4 s^{-1}. This robustness supports our assumptions of adiabaticity and negligible heating. Using a fixed rate $\Gamma_A \approx 8 \times 10^3$ s^{-1}, we have varied the initial conditions, determined by the initial lattice light intensity. We find that a lower final temperature results when the initial lattice light intensity

and temperature is decreased, until a limit is reached where steady state laser cooling can no longer be achieved at the initial conditions [10].

A more realistic calculation of the final temperature expected for adiabatic expansion in a periodic potential can be accomplished using a simple band theory. (A different approach was applied by Zaugg *et al.* [13] to the case of two-level atoms in a 1D lattice.) In the following we consider both a 1D lattice and a 3D cubic lattice of spherically symmetric wells to allow a simple examination of the influence of dimensionality. In our band theory, the initial condition for adiabatic expansion is taken to be an atom localized in a single optical potential well, and well approximated by a harmonic oscillator. The nth excited harmonic oscillator state H_n can be expanded in a basis of nonlocal, tight-binding Bloch states B_{nq} of band index n and quasimomentum q [14], and it can be shown that all q within a band have the same energy and population. The total population in the nth band is equal to the population of the nth harmonic oscillator state. The final state is a free atom, and is expanded in a basis of Bloch states that are simply plane waves [14]. The tight-binding and free particle band structures are shown in Fig. 3 for a 1D lattice with spatial frequency Q_0. Adiabatic expansion smoothly deforms the tight-binding bands into free particle bands. The population of each Bloch state is conserved due to adiabaticity, because a lattice of constant spatial frequency Q_0 conserves quasimomentum. After expansion the nth free particle band maps onto a pair of momentum intervals $n\hbar Q_0/2 < |p| < (n + 1)\hbar Q_0/2$, corresponding to the $(n + 1)$th Brillouin zone (BZ), within which the momentum distribution is uniform, as shown in Fig. 3(c). We assign a temperature by averaging $p^2/2m$ over this momentum distribution

$$\frac{1}{2}k_BT = \sum_n \frac{\pi_n}{\hbar Q_0}\int_{(n+1)\text{th BZ}} \frac{p^2}{2m}\,dp = E_R\left(\frac{Q_0}{k}\right)^2 \frac{1 + 4f_B + f_B^2}{12(1 - f_B)^2}, \tag{2}$$

using the initial thermal populations $\pi_n = (1 - f_B)f_B^n$. In a 3D extension of this model, the distribution of population in momentum space follows in a similar fashion from the 3D Brillouin zones. In a 3D cubic lattice the Brillouin zones are well approximated by spherical shells. Taking into account degeneracy of the spherical harmonic oscillator we find

$$\frac{3}{2}k_BT = E_R\left(\frac{Q_0}{k}\right)^2 \sum_n \left\{(1 - f_B)^3 f_B^n (n + 1)(n + 2) \times \frac{3}{40}\left(\frac{6}{\pi}\right)^{2/3}\left[(n + 1)^{5/3} - n^{5/3}\right]\right\}. \tag{3}$$

Equations (2) and (3) show that the final temperature is completely determined by Q_0 and f_B.

To compare Eqs. (2) and (3) to our experiment, one must consider which lattice spatial frequency Q_0 is appropriate. An atom moving adiabatically from a σ^+ to a σ^- site is transferred from $m_F = 4$ to $m_F = -4$ by coherent Raman coupling. For adiabatic expansion there is thus no distinction between lattice sites, and the symmetry becomes simple tetragonal with spatial frequencies $Q_x = Q_y = k\sqrt{2}$ and $Q_z = 2k\sqrt{2}$. Here we assume a simple cubic lattice with spatial frequency $Q_0 = 4k/\sqrt{5}$ (corresponding to $\theta \approx 127°$ in Fig. 1), and assign a Boltzmann factor, averaged over the 3 degrees of freedom, to each experimental data point. Figure 4 shows the temperature predicted using Eqs. (2) and (3), as

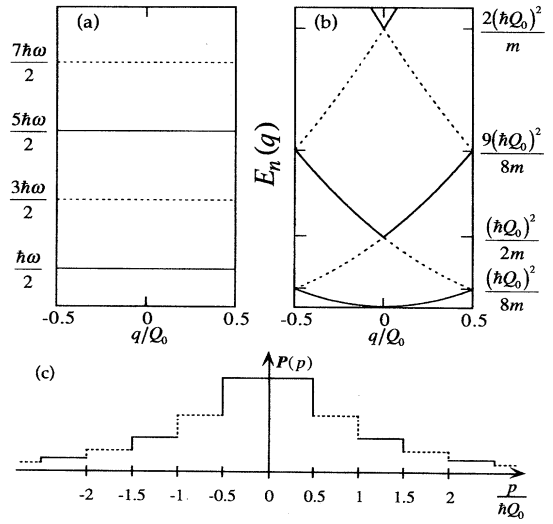

FIG. 3. (a) Tight-binding band structure. (b) Free particle band structure. (c) Final momentum distribution. Solid and dotted lines indicate odd and even numbered bands and their contributions to the final momentum distribution.

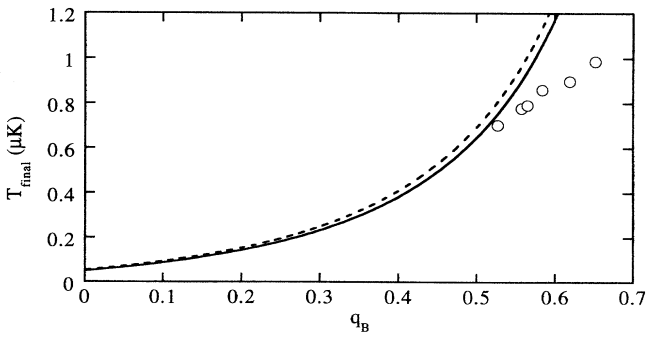

FIG. 4. Final temperature for adiabatic cooling as a function of initial Boltzmann factor. The solid line is a 3D calculation for a cubic lattice, the dashed line is a 1D calculation for a 1D lattice. In both cases the spatial frequency is $Q_0 = 4k/\sqrt{5}$. The open circles are measurements of T_x by time of flight. All measurements were taken with $\Gamma_A = 8 \times 10^{-3}\ \text{s}^{-1}$.

well as experimental data points for different Boltzmann factors, corresponding to different initial lattice light intensities. Boltzmann factors for the experimental points are computed on the basis of the vertical (along $\hat{x}$) velocity distribution only, and assume equal temperature for all degrees of freedom. We note first that there is very little difference in the temperature predicted by the 1D and 3D analysis, and second that both models reproduce the minimum observed temperatures reasonably well, given the level of approximation involved. Data points corresponding to higher initial lattice light intensity, and therefore larger initial Boltzmann factor, fall below the predicted limit for adiabatic cooling. This deviation is caused by dissipative cooling which lowers the Boltzmann factor early in the intensity reduction. A clear signature of this dissipative cooling is that at early times we measure a temperature directly proportional to the intensity, as expected for steady state cooling [7]. We also note that some discrepancy is expected, since theory predicts a slightly non-Boltzmann initial energy distribution [10], and since the estimated uncertainty in the experimentally assigned Boltzmann factors is $\pm 5\%$ [12]. Nevertheless, for the lowest initial intensities, below which there is no dissipative cooling, we find good agreement.

With a minimum temperature of 700 nK, adiabatic cooling of Cs atoms in an optical lattice offers an improvement of nearly a factor of 4 over 3D optical molasses [7]. About a factor of 2 of this improvement is due to a lower steady state temperature in the optical lattice, and the rest due to adiabatic expansion. One might suppose from Eq. (2) or (3) that a lower final temperature can be achieved by decreasing Q_0. However, if the initial potential depth and temperature are held constant, the accompanying change in f_B leads to a nearly unchanged final temperature. One route to lower final temperatures is to lower the initial temperature. This may be achieved with Raman sideband cooling [15], in a lattice detuned so far from resonance that heating and cooling rates are negligible. If atoms are cooled mostly into the ground state of the optical potential, Eq. (3) implies that adiabatic cooling in our model cubic lattice would lead to a temperature $E_R/2k_B \approx 50$ nK. Thus the band theory implies a new fundamental cooling limit. Even if atoms are cooled to the ground state (i.e., $T = 0$ K), it is impossible to release these atoms with a kinetic energy below approximately $(Q_0/k)^2 E_R/4$. An alternate way to achieve subrecoil temperatures would be to adiabatically expand the lattice, thereby reducing Q_0. In the geometry of Fig. 1 an arbitrarily small spatial frequency in the x-y plane can be achieved by choosing the proper angle θ. Anderson, Gustavson, and Kasevich have recently demonstrated a lattice configuration with $Q_0 < k$ in all directions, and achieved adiabatic cooling of Li by reduction of the lattice potential [16].

Optical lattices and adiabatic cooling can be used to create a colder atomic fountain and may provide important improvements in atom interferometers [17] and atomic clocks [18]. Attempts to achieve high phase space density, needed for Bose-Einstein condensation, might similarly benefit from adiabatic cooling. This might seem paradoxical since adiabatic expansion strictly conserves phase space density; however, localization on a microscopic scale represents phase space density which is lost when atoms are released into a macroscopic trap. This is avoided if microscopic localization is first traded for lower momentum spread by adiabatic cooling. Temperatures below those reported here may be achieved at the cost of decreased density, by releasing a small cloud of atoms into a weak macroscopic trap [19] which is turned off after a $\frac{1}{4}$ oscillation period [20].

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