Experiments with a 3D Double Optical Lattice

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(Received 24 May 2002; published 4 February 2003)

We present a setup where we trap two different cesium hyperfine ground states in two different near-resonant optical lattices with identical topographies. We demonstrate that we can change the relative spatial phase between the lattices and we measure the equilibrium temperature as a function of the relative spatial phase. This provides a topographical chart of the optical potential. We also determine the rate at which atoms are transferred between the lattices and show that the setup is a promising candidate for implementing coherent quantum state manipulation.

DOI: 10.1103/PhysRevLett.90.053001 PACS numbers: 32.80.Pj, 03.67.Lx, 32.80.Qk

The trapping of ultracold atoms in periodic light-shift potentials, optical lattices, is by now a fairly established technique [1,2]. In the past few years, the interest in optical lattices has increased considerably for two chief reasons. First, the advent of quantum degenerate gases has provided a new stage for optical lattices, where collective effects in periodic potentials can be studied (see, e.g., [3]). Moreover, optical lattices have been suggested as a promising candidate for quantum computation and quantum state manipulation (see [4] and references therein). The seminal suggestions for the feasibility of optical lattices in the latter context came in Refs. [5–7]. A key ingredient in the suggestions is two interpenetrating optical lattices, trapping two different quantum states, with as little cross talk between the lattices as possible. It is also required that the lattices have equal spatial periodicity and that the relative spatial phase is an adjustable parameter. In other words, it should be possible to translate one optical lattice, without moving the other, in order to induce controlled interactions between neighboring atoms. To construct a full scale quantum computer, additional prerequisites are that the atoms are trapped in the motional ground states of the potential wells, that the filling factor is close to unity, and that it is possible to address single qubits (i.e., lattices sites) both for setting and readout. However, in Refs. [5,7] it is shown that even without the possibility for single bit addressing, and with relaxed demands on occupation number and excitation of motional states, the double optical lattice can still be useful for quantum state manipulation.

Most suggestions to realize a double optical lattice are based on variations of the "lin $\perp$ lin configuration" [2]. This standard model for sub-Doppler laser cooling is in one dimension based on two counterpropagating laser beams with crossed linear polarization. This provides alternating lattices sites with pure circular polarization of opposite handedness (σ^+ and σ^-). At these sites atoms with opposite spins will be trapped, and, by changing the angle between the polarization vectors, the distance between the wells can be adjusted. In this Letter, we present a different approach, where a double optical lattice is produced by two different sets of laser beams with slightly different wavelengths, trapping atoms in different states. Even though the wavelengths are sufficiently far apart to minimize cross talk, they are so close that the difference in spatial periodicity is insignificant for an atomic sample with a diameter of a few hundred micrometers. Besides the possible applications in quantum state manipulation, the double optical lattice presents possibilities for studies of cold collisions and photoassociation, under slightly different conditions than in previous reports (see, e.g., [8]). By moving the lattices we can induce binary collisions with better control of interaction distance and interaction time. Moreover, many-body effects, such as three-body collisions that constitute a significant loss-channel in Bose-Einstein condensation (BEC), are essentially eliminated. Since one of our optical lattices replaces the laser field which is typically used as a repumper in standard laser cooling configurations, the double optical lattice is also of interest for studies of improved laser cooling, important, e.g., for metrology and high precision spectroscopy.

The idea behind our experiment is explained in Fig. 1. Two laser fields with frequencies ω_A and ω_B interact with two different ground states $|g_A\rangle$ and $|g_B\rangle$. These are so far apart that the optical light-shift potential for state $|g_A\rangle$ ($|g_B\rangle$) is determined by laser A (B). On the other hand, since the periodicity of an optical lattice is determined by the wavelength of the laser light, the splitting $\hbar\Delta_g$ must be small enough so that the two sets of lattice constants are practically equal. In order to perform consistent experiments with this setup, one has to be able to control the spatial position of the lattices. Position changes of the lattices will be caused by phase changes in one or several laser beams. In our experiment phase fluctuations between the two laser sources are unimportant since we choose a configuration with four laser beams for each three-dimensional lattice [2]. In such a configuration a phase fluctuation in one of the beams will result in a translation of the lattice. If the same phase fluctuation occurs in all beams, however, the position of the lattice will remain unchanged. This means that a random phase

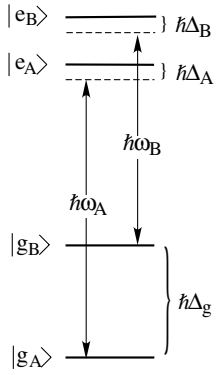

FIG. 1. A model system for a double optical lattice, operated on two transitions connecting the ground states $|g_A\rangle$ and $|g_B\rangle$ with the excited states $|e_A\rangle$ and $|e_B\rangle$. The lasers with frequencies ω_A and ω_B are detuned from resonance by Δ_A and Δ_B. The ground states are separated by $\hbar\Delta_g$.

jump in either laser will not translate the lattices. It is therefore not necessary to phase-lock the two lasers. To make our setup insensitive to fluctuations caused by vibrations in the optical components defining the beam paths, we overlap laser beam A with laser beam B before they are split up into four beams. Phase jitter occurring before the beams are overlapped will affect all lattice beams and hence will not have any effect. Mechanical vibrations in optical components after the overlap will cause identical translations of both lattices. Because the optical wavelengths are slightly different, it is possible to vary the relative spatial phase by a significant change ($\sim$1 cm) of the optical path lengths in either of the four beam pairs.

We operate our double optical lattice on the D2 line in Cs ($\lambda \approx 852$ nm). Two diode lasers are detuned below the ($F_g = 4 \rightarrow F_e = 5$) and ($F_g = 3 \rightarrow F_e = 4$) transitions, respectively. The two laser frequencies are separated by about 9 GHz, implying that the two interference patterns will not dephase significantly across the optical lattice volume (≈ 0.6 mm in diameter). The two beams are overlapped with a polarizing beam splitter cube and the combined beam is fed through an optical fiber for spatial filtering. After the fiber the beams are split into four beam pairs that form angles $\theta = 45°$ with the quantization axis (z axis). This results in two tetragonal optical lattices, with alternating σ^+/σ^- sites, that trap atoms in $m_{F_g} = \pm F_g$ (see further [2]). The lattice constants are $a_x = a_y = \lambda/\sqrt{2}$ and $a_z = \lambda/2\sqrt{2}$. For this beam geometry it is necessary to change the relative path length by 72 mm (36 mm) to achieve a change in the relative spatial phase of 2π along the x and y axis (z axis), which we define as the distance between identical double-lattice configurations (coinciding σ^+ sites of lattices A and B). The changes in path lengths are accomplished with retroreflectors and carefully aligned translation stages. Their travel range allows us to change the spatial phases in the x and y directions by approximately 2.2π and by 4.4π along z.

In a setup described in more detail in [9,10] we collect atoms in a magneto-optical trap (MOT). The atomic cloud is further cooled in an optical molasses. We then turn on the lattice beams while simultaneously switching off the molasses light. We trap 6×10^6 atoms, resulting in a dilute lattice with a peak number density of 1.7×10^{11} cm^{-3} and a filling factor of 0.5%. Since the optical lattices A and B operate on the two hyperfine ground states in Cs, the need for a special repumper beam is eliminated. Atoms will occasionally be optically pumped out from one of the optical lattices only to find themselves trapped in the other one. After letting the atoms equilibrate in the optical lattices for about 50 ms, we turn off the lattice beams and perform a time-of-flight (TOF) measurement [2] to extract the velocity distribution and relative number of atoms in the lattices. The uncertainty in the absolute temperature is about 10%, but the relative precision is much better ($\pm$ 20 nK). The TOF measurement is done with a resonant laser beam that traverses the vacuum chamber 5 cm below the lattice. The beam is spatially filtered with an optical fiber and focused into the interaction region with a cylindrical lens, creating a beam less than 50 μm thick. The TOF probe also contains two laser frequencies tuned to the ($F_g = 3 \rightarrow F_e = 4$) (probe A) resonance and the ($F_g = 4 \rightarrow F_e = 5$) (probe B) resonance, respectively. Each probe can be blocked individually. The temperature of the atoms in the $F_g = 4$ ground state (i.e., we block probe A) was measured as a function of the relative spatial phases (φ_y, φ_x, φ_z) between the two lattices. For this purpose both lattices are detuned 19Γ ($\Gamma/2\pi = 5.2$ MHz is the natural linewidth) below the ($F_g = 4 \rightarrow F_e = 5$) and ($F_g = 3 \rightarrow F_e = 4$) resonances. Figure 2 shows the result for a single scan where we varied φ_z while keeping φ_y and φ_x constant. The temperature is plotted versus the relative position of the optical lattices in terms of lattice constants. The temperature varies with Δ_z between 2.52 and 2.95 μK. The periodicity is $q_z = 0.9a_z$, which within the experimental uncertainties agrees well with the

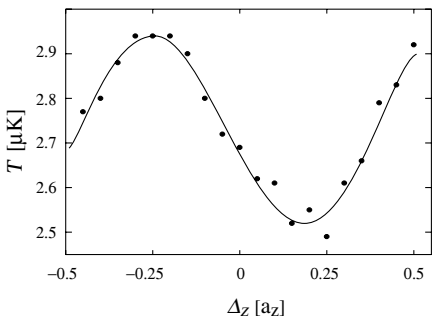

FIG. 2. Temperature as a function of relative displacement Δ_z of the two lattices. Each point is an average of five TOF measurements. Also shown is a fit of the form $T = T_0 + C\sin(\frac{2\pi}{q}z)$.

expected a_z. The temperature dependence on the spatial phase is expected. In a semiclassical model shown in Fig. 3, the atom is optically pumped from $|g_B\rangle$ to $|g_A\rangle$ and back into $|g_B\rangle$. In the extreme case of the lattices being "in phase," i.e., the σ^+ sites of both lattices overlap, nothing dramatic will happen to the atom. One optical pumping cycle changes the m_F by maximum 1 so that the atom is transferred from one trapped state to another. The situation changes when the lattices are completely "out of phase," i.e., the σ^- sites of lattice A overlap with the σ^+ sites of lattice B. The atom is now transferred from a trapping potential in lattice B to an antitrapping potential in lattice A. It slides down the potential (gaining kinetic energy) until it is optically pumped into a trapped state. This requires multiple photon scattering leading to further heating. This process is repeated when the atom is pumped back into lattice B.

In Fig. 4(a) we present a complete 2D scan where both φ_y and φ_z are varied. We obtain a temperature surface where the modulation depth along the z axis varies with φ_y. The periodicity in the y direction is $q_z = 0.9a_y$, which compares well with the predicted a_y. Our result can be compared with a theoretical plot shown in Fig. 4(b), where the irradiance $I_{\sigma+}$ of the σ^+ component in one lattice is shown in the y-z plane, which corresponds to the diabatic potential for the $m_g = +F_g$ ground state. This illustrates that we are able to map out the topography of the optical lattice potential by using the relative spatial phase between the lattices as a control parameter and the temperature as a probe.

When two optical lattices are spatially overlapped, each site will effectively contain a two-level system, formed by the ground states of lattices A and B, with dissipative couplings between the levels. Their strength is determined by the optical pumping rates which depend on irradiances and detunings. These rates are highly asymmetric, since only one of the optical lattice transitions is cycling. We have measured the optical pumping rate out of

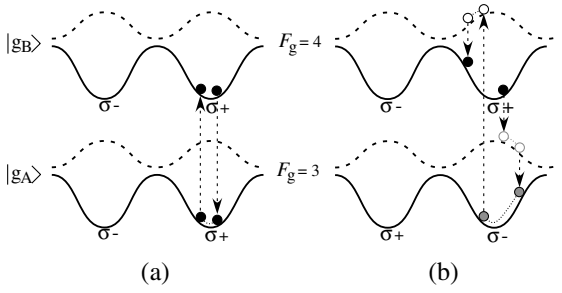

FIG. 3. Optical potential curves and atom trajectories for two different situations in a double optical lattice. Shown are only the highest/lowest adiabatic potentials for the two ground states. In (a) atoms are transferred between two trapping potentials for each depumping cycle. In (b) the atom is first transferred to an antitrapping potential before it is pumped into the trapped state.

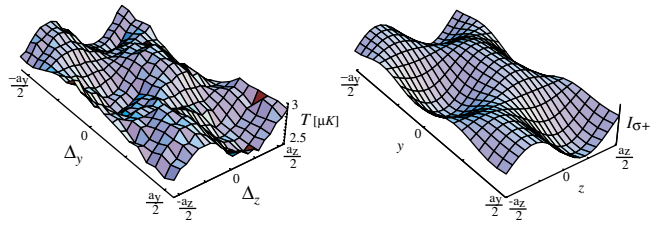

FIG. 4 (color online). (a) Measured temperature in the double optical lattice plotted as a function of the relative displacements Δ_y, Δ_z of the two lattices in the z-y plane. (b) Calculated irradiance (in relative units) of the σ^+ component as a function of position y, z in one of the lattices.

the lattice operating on the cycling ($F_g = 4 \rightarrow F_e = 5$) transition (lattice B) by turning off the lattice operating on the noncycling ($F_g = 3 \rightarrow F_e = 4$) transition (lattice A). Here the relative spatial phases were adjusted in such a way that the situation resembled the one depicted in Fig. 3(a). We recorded the relative number of remaining atoms in lattice B as a function of time by integrating the TOF signal. For detunings Δ_B between -7.7Γ and -38Γ [where $\Delta_B = \Delta_{45}$, i.e., the detuning from the ($F_g = 4 \rightarrow F_e = 5$) resonance] and irradiances between $I_B/I_0 = 0.7$ and $I_B/I_0 = 3.6$ (where $I_0 = 1.1\ \mathrm{mW/cm^2}$ is the saturation irradiance and I_B is the irradiance in a single lattice beam), we measured optical pumping times τ_B between 0.6 and 4 ms. The pumping rate $\gamma_B = 1/\tau_B$ increases with increasing irradiance and also with increasing detuning. This is expected, since in this closed transition, optical depumping occurs only due to off-resonant excitations to the $F_e = 4$ level ($\Delta_{44} = 48.1\Gamma + \Delta_{45}$). We have also measured the relative populations in the two lattices in steady state by recording relative numbers of atoms. First we use only probe B to get a relative number N_B for the atoms in lattice B. Then we use both probe beams where probe A now works as an optical pump. This yields a number proportional to the total number of atoms N_{tot}. The ratio N_B/N_{tot} was measured for the same lattice B parameters as above and with Δ_A ranging from -7.6Γ to -55Γ and I_A/I_0 ranging from 0.7 to 6.3. For these parameters, N_B/N_{tot} was between 0.75 and 0.99. The optical pumping rate out of lattice A, γ_A, was estimated by assuming a trivial rate equation at steady state:

$$\frac{dN_B}{dt} = -\gamma_B N_B + \gamma_A N_A = 0. \quad (1)$$

We found that τ_A ranges between 6 and 300 μs. Thus, τ_A and τ_B are in the order of or exceed the decoherence time in an optical lattice [11], which thus sets the upper limit to the duration of coherent quantum operations. In conjunction with appropriate design of the modulation depth, this opens a time window for a scheme in which relative displacements can be done adiabatically with respect to the oscillation frequency ν in the trapping potential, by

using electro-optic modulators (typical rise times $\sim 1\ \mu$s). Such a scheme is currently being implemented into our setup.

Most schemes proposed in [5–7] rely on the atoms being in the vibrational ground state of their potential wells. In a dissipative optical lattice, atoms can be cooled to an average vibrational state of around 1 [12]. Our double optical lattice configuration presents a possibility to improve this by using a variation of Raman sideband cooling [13–15]. A prerequisite for Raman sideband cooling is that the vibrational levels in a lattice site are well resolved. It has been shown [2,12] that this criterion is fulfilled for a standard dissipative optical lattice. In our lattice A, which operates on an open transition, a major contribution to the width of the vibrational states stems from the escape rate out of the lattice due to optical pumping. This width ($\delta\nu_A$) is extracted from our rate measurements and is between $\delta\nu_A = 0.5$ kHz and $\delta\nu_A = 25$ kHz depending on I_A and Δ_A. The oscillation frequency in a well is given by $\nu = \nu_{rec}\sqrt{2U_0/E_{rec}}$, where U_0 is the modulation depth of the optical potential and ν_{rec} and E_{rec} the recoil frequency and energy, respectively. For our parameter range we create modulation depths from $U_0 \approx 40E_{rec}$ up to $800E_{rec}$, which results in a ratio between the width and the oscillation frequency ranging from $\delta\nu_A/\nu_A = 0.03$ to $\delta\nu_A/\nu_A = 0.3$. Hence, the condition for Raman sideband cooling $\delta\nu/\nu \ll 1$ is fulfilled. Furthermore, even for the smallest modulation depth ($U_0 \approx 40E_{rec}$), the spacing between the vibrational levels is far greater than the recoil energy. The cooling scheme could be realized in the following way: the two optical lattices are overlapped according to Fig. 3(a) and an additional π-polarized (in the z direction) laser beam is introduced in the xy plane. The frequency of this Raman beam is adjusted so that it together with lattice light A creates a resonant stimulated Raman transition between the nth vibrational state in ($F_g = 4, m_F = \pm 4$) to the $(n-1)$th state in ($F_g = 3, m_F = \pm 3$). Lattice A also provides light that pumps the atoms back to lattice B with one vibrational quantum less than initially. Thus, atoms will accumulate in the $n = 0$ states in ($F_g = 4$, $m_F = \pm 4$) which are dark for the Raman resonance. Once trapped in a vibrational ground state of lattice B, atoms could be coherently transferred between the lattices or be prepared in different superpositions by applying Raman pulses. The time scale of such an operation would be in the order of $1/\Gamma$ since the Rabi frequency at the center of a well is $\Omega^2 = \frac{\Gamma^2}{2}\frac{I}{I_0}$. This is well within the decoherence time of $\sim 50\ \mu$s.

A double optical lattice could also be interesting for state manipulation. A controlled interaction could be achieved by switching off the short-lived lattice A to prevent optical pumping and quickly moving the second lattice, thus causing an interaction, and then, for instance, measure the trap loss. A concrete example is to study photoassociation and light-assisted collisions. The lattices are brought together and at different relative distances a short photoassociation pulse can be turned on. The speed at which the lattice can be displaced depends on the nature of the experiment. In the case of coherent quantum manipulations, care must be taken so that the evolution of the wave function stays coherent [7]. Collision experiments would, however, benefit from a larger filling factor than that in the present situation. The filling factor can be increased, e.g., by loading from a compressed MOT [16].

To conclude, we have successfully implemented a setup where both Cs ground states are trapped in individual optical lattices. The equilibrium temperature for atoms is a function of the relative spatial displacement of the lattices which allows us to map out the topography of the optical potential in a unique and precise way. By carefully choosing irradiance and detuning in each lattice, we can achieve average lifetimes in the lattices long enough to give us a time window for performing a host of experiments for quantum state manipulations.

This work was supported by the Swedish Natural Sciences Research Council, Carl Tryggers Foundation, and Knut & Alice Wallenbergs Foundation.

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