# Two-Dimensional Atomic Crystal Bound by Light

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We have observed the vibrational levels of rubidium atoms spatially confined in two dimensions by two crossed optical standing waves (oriented along the $x$ and $y$ axes) with mutually orthogonal linear polarization (in the $x$-$y$ plane). When the two standing waves oscillate with a 90° time phase delay, the atoms are confined by dipole forces to the intensity maxima forming a two-dimensional collimated array of linear de Broglie waveguides, spaced by half an optical wavelength. When the standing waves oscillate in phase no vibrational energy structure can be observed.

Very recently two research groups have observed one-dimensional confinement and quantized motion of atoms in an optical standing wave [1,2]. In this work we explore the potential of light forces to confine atoms with respect to all spatial coordinates and at the same time provide cooling down to temperatures in the microkelvin range. This could lead to atomic samples highly ordered in three-dimensional periodic lattices on the micron length scale. The physics describing such a sample resembles that of an ultracold and very dilute solid. In this Letter we report a two-dimensional realization of such an atomic crystal bound by light. In this 2D example the atoms are confined by dipole forces to the intensity maxima which form a two-dimensional collimated array of linear de Broglie waveguides, spaced by half an optical wavelength. Unlike laser cooled ions in a "Wigner crystal" [3], neutral atoms do not show strong Coulomb repulsion. Thus, one could study collective absorption and emission processes and collisions in systems of two or more closely spaced atoms trapped in the same de Broglie waveguide. A three-dimensional extension of our experiment would provide a cubic 3D lattice of microscopic light traps. Interesting quantum statistical phenomena might occur if several atoms are confined in the same microscopic light trap.

To achieve trapping and efficient cooling of the atoms at the same time, polarization gradients of the light are necessary. In one dimension the configuration of two counterpropagating laser beams with orthogonal polarization (i.e., the lin⊥lin scheme) is appropriate [1,2]. However, when two or more standing waves are superposed, inherently multidimensional phenomena can arise as the action of vortical radiation pressure which may prevent the localization of the atoms [4]. The presence of such complications depends on the time phase differences between the employed standing waves and the choice of polarizations.

In this Letter we demonstrate two-dimensional quantized motion and confinement of rubidium atoms in the particularly simple 2D field geometry arising in the intersection region of two standing waves oriented along the $x$ and $y$ axes and linearly polarized in the $x$-$y$ plane. As shown in Fig. 1, this light field acquires two different types of spatial polarization gradients for the two characteristic values $\phi = 0°$ and $\phi = 90°$ of the time phase difference between the two standing waves. When $\phi = 0°$ the polarization is linear everywhere, only its direction varies. In the $\phi = 90°$ case there exists a 2D array of straight lines parallel to the $z$ axis where the light exhibits circular polarization with alternating sign. There is a continuous change to linear polarization when one moves away from those locations. The two cases shown in Fig. 1 resemble the two one-dimensional optical fields which have been at the basis of the theory of polarization gradient cooling, namely, the $\sigma + \sigma -$ configuration and its counterpart, the lin⊥lin configuration [5]. Thus, we expect our two-dimensional light field to contain the physics of both elementary 1D cases depending on the choice of the time phase difference. In particular, our 2D field provides efficient sub-Doppler cooling for any value of $\phi$ [6]. At the same time it does not suffer from vortical radiation pressure [4].
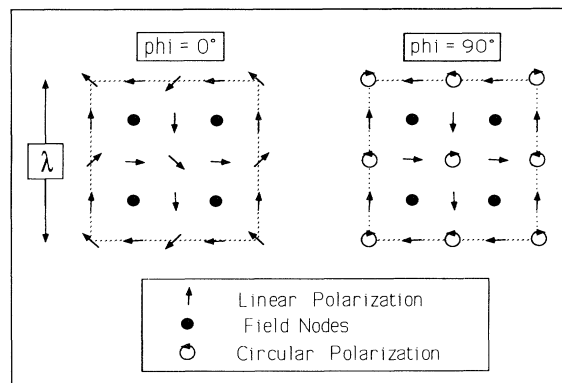


FIG. 1. The geometry of the electric field in the intersection region of two orthogonal standing waves with linear polarizations in the drawing plane. At $\phi = 90°$ the locations of circular polarization form a two-dimensional array of de Broglie waveguides.

Let us consider the case $\phi = 90°$ in some more detail for an atom with a $J \to J+1$ transition, e.g., $J = 3$. We choose the quantization axis perpendicular to the plane of the 2D field. The Zeeman levels experience different spatially varying light shifts due to the different size of the Clebsch-Gordan coefficients and the spatially varying polarization and intensity of the light field. At negative detuning $\delta$ of the light frequency with respect to the atomic resonance the outermost Zeeman levels $(m = \pm 3)$ are shifted by a maximum amount downward at those locations where the light is circularly polarized. Thus, atoms in the $m = \pm 3$ states experience deep potential minima at those locations and we expect a confinement of the $m = \pm 3$ atoms in two dimensions. In particular, because the circular polarization components of the light are strongly unbalanced at nearly any location, optical pumping should prepare the atoms predominantly into the outermost Zeeman states. Aided by efficient sub-Doppler cooling, a large fraction of atoms should get localized in a 2D array of de Broglie waveguides with $\lambda/2$ periodicity ($\lambda$ is the optical wavelength). More precisely, $m = 3$ atoms should collect at places of positive circular polarization whereas $m = -3$ atoms should collect where the light has negative circular polarization. Thus, atoms at neighboring waveguides should exhibit opposed magnetization oriented along the $z$ axis similar to an antiferromagnetic medium.

In the vicinity of the local minima the optical potential is approximately described by a two-dimensional harmonic oscillator potential. Therefore, we may expect discrete nearly equidistant vibrational quantum states in that potential. A quantitative analysis has to take into account tunneling between different potential minima, which attributes a width to each vibrational state. At higher vibrational quantum numbers these widths significantly wash out the discrete energy structure. Therefore, a band theory is more appropriate for the analysis of the problem. However, one-dimensional band calculations have shown that the widths of the low-lying states are of the order of $10^{-6}$ times the recoil energy which is far below the energy separation of the vibrational levels [7]. When the atoms are sufficiently cold we may thus treat the potential minima independently.

In the following, we consider $m = 3$ atoms, slowly moving in the vicinity of a potential minimum. We may, thus, study the atomic vibrational motion by considering the forces on a stationary two-level atom. This yields a particularly simple expression for the trapping potential:

$$U(x,y) = \frac{\hbar\delta}{2}\ln(1+S),$$

$$S(x,y) = S_0 \tfrac{1}{4}[\cos(kx)+\cos(ky)]^2, \tag{1}$$

$$S_0 = \frac{\omega_{max}^2}{2[\delta^2+(\Gamma/2)^2]}.$$

Here, $S_0$ denotes the saturation parameter of the $(J = 3,$ $m = 3) \to (J = 4, m = 4)$ transition, $\omega_{max}$ is the antinode

Rabi frequency for this transition, $\Gamma$ is the inverse lifetime of the excited state, and we have parametrized our light field such that the locations $(x,y)$ of positive circular polarization are given by $kx = \pm 2n\pi$, $ky = \pm 2n\pi$, where $n$ is an integer. We may expand Eq. (1) up to second order obtaining

$$U(x,y) = \frac{\hbar\delta}{2}\left[\ln(1+S_0)\right.$$
$$\left. - \frac{1}{2}\frac{S_0}{1+S_0}k^2(x^2+y^2)+O(4)\right]. \tag{2}$$

For negative values of the detuning $\delta$, the second term in Eq. (2) provides a two-dimensional harmonic oscillator potential yielding a frequency separation of the equidistant vibrational levels of

$$\Delta\nu_{osc} = (153 \text{ kHz})\left[\frac{|\delta|}{\Gamma}\frac{S_0}{1+S_0}\right]^{1/2}. \tag{3}$$

The value of 153 kHz corresponds to the $D_2$ resonance line of $^{85}$Rb. There are two sources of anharmonicity both leading to a lowering of the $n$th vibrational level and a splitting of its $(n+1)$-fold degeneracy which increase with the vibrational quantum number $n$. On the one hand, there are higher-order terms in Eq. (2) which, however, introduce corrections of only a few kilohertz. On the other hand, the two-level model used in Eq. (2) is correct only for an atom which is located exactly in a potential minimum. For a slightly displaced atom optical pumping partially populates the $m = 2$ state and the atom thus experiences an effective potential which is lower than predicted by Eq. (2). Both types of perturbations exhibit a spatial geometry such that the twofold degeneracy of the first excited state is preserved and the threefold degeneracy of the second state is lifted. The lifetimes of the low-lying vibrational states should be determined by optical pumping time.

We now consider the case when $\phi = 0°$ in our 2D field. The polarization remains linear in the plane of the light field (see Fig. 1). Depending on their internal state and the corresponding Clebsch-Gordan coefficient the atoms now move in potentials of different size which arise from the spatial variation of the energy density of the light. The minima of these $0°$ potentials are at the same locations, in contrast to the $90°$ case where two different optical potentials confine $m = -3$ and $m = 3$ atoms at separated locations. As in the $90°$ case optical pumping tends to prepare the atoms into the most light shifted Zeeman state; however, this pumping is less complete for $\phi = 0°$ because the local $\sigma+$ and $\sigma-$ polarization components (with respect to the chosen quantization axis $z$) are balanced everywhere. In addition, the polarization gradient cooling at $0°$ is less efficient [6]. As a result we expect population of vibrational states up to higher order arising from various potentials of different depth. The lifetimes of these vibrational states should be smaller

than in the 90° case because of the rapid redistribution of the atoms among the different potentials by optical pumping.

In our experiment a cold (4 $\mu$K) dense ($10^{10}$ atoms/cm$^3$) cloud of atoms is prepared by a magneto-optical trap operating in a rubidium-vapor cell [8]. The cloud is about 0.3 mm in diameter and contains approximately $3 \times 10^5$ atoms. The trap is active for 40 ms. During the last 0.5 ms the magnetic field is switched off and the frequency detuning of the trapping light field is increased for most efficient sub-Doppler cooling. Then, the trapping light field is switched off for a probing period of 2 ms, while the repumping laser (which counteracts hyperfine pumping) is kept active. During this time our 2D field is activated at the position of the atomic cloud.

The 2D field is produced by crossing the two branches of a Michelson interferometer which is fed by the spatially filtered output of a grating stabilized diode laser. The time phase difference between both branches can be servo controlled to any desired value by analyzing the light reflected from the interferometer. We couple 3 mW in a 6-mm-diam beam into the interferometer. This results in an antinode Rabi frequency $\omega_{max} = 5\Gamma$. The frequency of the 2D field is detuned with respect to the ($F=3$, $m=3$) $\rightarrow$ ($F=4, m=4$) transition of $^{85}$Rb by $\delta = -8\Gamma$. According to Eq. (3) the harmonic part of the frequency splitting between adjacent vibrational states for $\phi = 90°$ is

$\Delta\nu_{osc} = 161$ kHz.

A weak linearly polarized probe laser beam of 0.4 mm diameter is directed through the cloud. This beam travels within the $x$-$z$ plane and is tilted by 6° with respect to the $x$ axis. Its frequency $\nu_p$ is tuned across the frequency of the 2D field $\nu_{2D}$ during the probing period, while its absorption by the atoms is recorded. We adjust the probe polarization parallel to the plane of the 2D field to excite Raman transitions between the differently populated vibrational states. We interpret the absorption spectrum shown in Fig. 2(a) as the Stokes and anti-Stokes lines due to the Raman transitions connecting the vibrational ground state with the first and second excited states. These resonances occur at center frequencies $\nu_{2D} \pm \Delta n_1$ and $\nu_{2D} \pm \Delta n_2$ symmetrically detuned from the frequency of the 2D field $\nu_{2D}$. To evaluate $\Delta n_1$ and $\Delta n_2$ we have fitted the Raman part of the spectrum with Lorentzians located symmetrically around $\nu_{2D}$ (Fig. 3). We obtain the values $\Delta n_1 = 159$ kHz and $\Delta n_2 = 285$ kHz, respectively. The value of 159 kHz is in good agreement with the value of 161 kHz expected according to Eq. (3).

The above interpretation is based on the following considerations. The observation of well-resolved sidebands with significantly decreasing separation indicates that a large fraction of the observed transitions should start from the same state which should be the ground state if a thermal distribution of the populations is assumed. If transitions starting from highly excited states would strongly contribute to the spectrum, we would expect to observe broader and more evenly spaced sidebands because each of them would be due to a series of strongly decreasing energy separations. According to measurements similar to those described by other authors [9] our atomic cloud is prepared at a temperature of about 4 $\mu$K and due to the presence of efficient sub-Doppler cooling in the 2D field [6] this temperature should be maintained. Assuming a 4-$\mu$K thermal distribution of the vibrational populations leads to a ground-state population which is about 4 times the population in the first excited state. Is
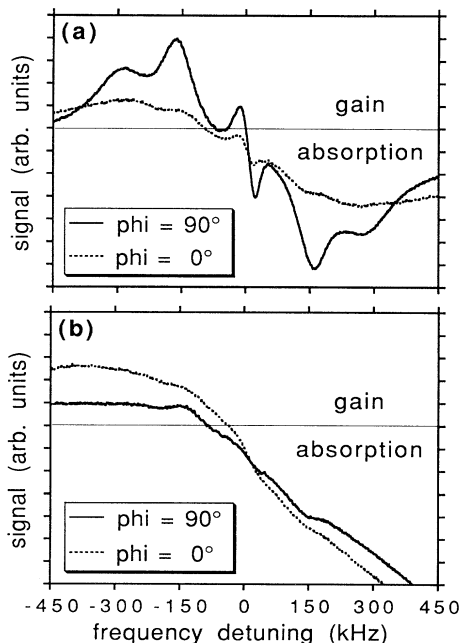


FIG. 2. Absorption spectra for $\phi=0°$, 90° and different orientations of the linear probe polarization. In (a) it is parallel and in (b) it is perpendicular to the plane of the 2D field. The frequency resolution is a few kHz and the signals are typically a few percent of the probe intensity.
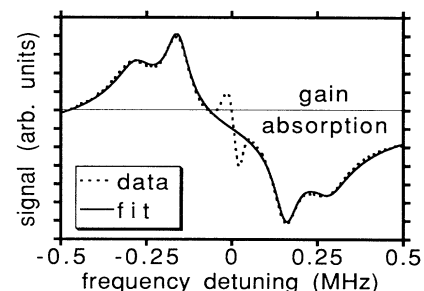


FIG. 3. The dotted curve repeats the 90° spectrum from Fig. 2(a). The solid line is a Lorentzian fit of that part of the spectrum which we interpret in terms of Raman transitions between different vibrational levels. The residual narrow structure in the center of the dotted curve results from Bragg scattering due to the periodic spatial order in our atomic sample.

this assumption compatible with the observed relative weights of the first and second sidebands? To check that, we have conducted a calculation of the transition rates from the vibrational ground state to the first two excited states. In this calculation the atoms are subjected to a potential with a harmonic part as given in (2) and an anharmonic part which is adjusted to yield the experimentally observed sideband frequencies. This potential is distorted by the interference between the probe and pump waves yielding vibrational excitation of the atom. We obtain a ratio of the transition rates from the ground state to the first $(P_1)$ and second $(P_2)$ excited states of $P_2/P_1 = 0.5$ which is in reasonable agreement with the observations. The $1/e$ radius of the ground-state probability density is only $\lambda/29$. This indicates that a large fraction of the observed atoms is localized along the axes of the light-induced waveguides. The vibrational energy levels are not well resolved for $\phi = 0°$ [Fig. 2(a)]. In particular we observe more atoms at higher frequencies than in the 90° case corresponding to sidebands of higher order. As a check of our interpretation in terms of Raman transitions we have adjusted the probe polarization perpendicular to the plane of the 2D field [Fig. 2(b)]. In this case the probe drives only $\pi$ transitions and thus cannot cooperate with the 2D field to excite a Raman transition.

The narrow structure in the center of the spectrum shown in Figs. 2(a) and 3 results from a Bragg scattering process arising from the periodic order in our atomic sample. A quantitative investigation of this Bragg resonance is beyond the scope of this Letter. However, in the following we briefly sketch a physical model for the origin of this resonance which modifies the one-dimensional considerations from Ref. [1]. In the following the terms $\sigma+$ and $\sigma-$ always refer to the quantization axis (i.e., the $z$ axis). The probe contributes equal amounts of $\sigma+$ and $\sigma-$ polarization to the total local light field. This results in a modulation of the populations of the vibrational states within each waveguide with the frequency $|v_p - v_{2D}|$. Each of the four traveling (pump) waves which produce the 2D field can be scattered by these modulated populations at each lattice site yielding a scattering component propagating on the same axis as the probe beam with the same frequency and polarization. To understand this, note that the oscillating electric field of each pump wave at a given lattice site can be split into $\sigma+$ and $\sigma-$

portions of equal size. These are scattered by the oriented atoms with different phases, thus, giving rise to a small "Faraday rotation" of the linear pump polarization within the $x$-$y$ plane. The scattering components from different lattice sites interfere constructively for the following reasons. In the case of the pump wave counterpropagating with the probe the path length difference from components scattered by successive lattice sites is $\lambda$ and the scattering phase is the same for each lattice site. For the pump waves propagating perpendicular to the probe there is an optical path length difference of $\lambda/2$ between components scattered by successive lattice sites. However, in this case we obtain a compensating phase difference of $\pi$ from the opposed magnetizations. The total scattered wave interferes with the probe beam and gives a dispersively shaped resonance. The width of this resonance in our model is determined by the inverse relaxation time for the vibrational populations. Preliminary results from a three-dimensional extension of our experiment show a central resonance with a dramatically decreased linewidth of less than 2 kHz.

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