Pis'ma v ZhETF, vol.61, iss.5, pp.362 - 367 © 1995 March 10

TRAPPING ATOMS BY RECTIFIED FORCES IN BICHROMATIC OPTICAL SUPERLATTICES

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Submitted 12 January, 1995

We consider the trapping of atoms by sub-Doppler rectified dipole forces in the bichromatic superlattice that is formed as macroscopic spatial interference pattern of two superimposed individual optical lattices differing in frequency and wavelength. A simple model is presented for a three-dimensional calculation of the rectified forces and for a semiclassical Monte-Carlo simulation of the trapping process. Our results show that the macroscopic bichromatic traps created at the superlattice sites are very promising for spatially compressing atomic ensembles released, e.g., from a conventional magneto-optical trap.

In their pioneering work [1], Kazantsev and Krasnov first described an effect underlying a new class of optical forces [2] exerted on atoms in near-resonant laser light: for atoms in bichromatic standing-wave laser fields they predicted 'rectified dipole forces' that oscillate in space on a macroscopic scale greatly exceeding the optical wavelength. They also pointed out potential applications to realize new types of atom traps. Meanwhile, rectified forces have been experimentally observed in a variety of different field configurations and rectification schemes, e.g. [3–5].

The aim of this Letter is to point out possible ways for constructing highly efficient atom traps on the basis of the rectification effect. In contrast to the work of Kazantsev and Krasnov we do not consider two-level atoms, but atoms with a ground-state substructure, for which cooling [6] and rectified forces [7,8] can occur in the sub-Doppler regime with very sharp dependence on the atomic velocity [5,8]. We show that, in simple laser beam configurations used in recent work on 'optical lattices' [9-12], bichromatic light creates an additional superlattice structure where rectified forces act together with usual polarization-gradient cooling [6] to efficiently cool and confine atoms in an array of macroscopic traps.

In our theoretical model, adopting well-known concepts in laser cooling [6], we consider atoms with a $J = 1/2 - J' = 3/2$ transition in a one-, two-, or three-dimensional bichromatic standing-wave laser field. We focus on lattices that can be decomposed into circular ($\sigma^{\pm}$) polarization components without any π–light contribution [9-12]. The bichromatic light field is characterized by the position-dependent $\sigma^{\pm}$ saturation parameters of the two frequency components $G^{\pm}_{1,2}(\mathrm{r}) = I^{\pm}_{1,2}(\mathrm{r})/I_{SAT}$ with $I_{SAT} = \hbar\omega_0^3\gamma/6\pi c^2$ and by the reduced frequency detunings $\delta_{1,2} = (\omega_{1,2} - \omega_0)/\gamma$, where ω_0 is the transition frequency and 2γ its natural linewidth. We treat the atomic motion semiclassically by considering the motion along the light-shifted ground-state potentials in the so-called sub-Doppler regime of low velocities ($v \ll \gamma/k$), where the Doppler shifts remain small compared to γ and contributions by the scattering force (Doppler cooling) can be neglected. We furthermore assume large detunings ($\delta_{1,2} \gg 1$), low optical saturation

$(G^{\pm}_{1,2}(\mathbf{r})/\delta^2_{1,2} \ll 1)$, and a negligible effect of mutual light shifts $(G^{\pm}_{1,2}(\mathbf{r})/\delta_{1,2} \ll \delta_{2,1})$. This allows us to simply add up the individual light shifts as first order perturbations in order to calculate the relevant ground-state potentials

$$U_{\pm}(\mathbf{r}) = \frac{1}{2}\hbar\gamma \left(\frac{G^{\pm}_1(\mathbf{r})}{\delta_1} + \frac{G^{\mp}_1(\mathbf{r})}{3\delta_1} + \frac{G^{\pm}_2(\mathbf{r})}{\delta_2} + \frac{G^{\mp}_2(\mathbf{r})}{3\delta_2} \right). \tag{1}$$

Here the first and the second term result from the ω_1 excitation of the $m = \pm 1/2 - m' = \pm 3/2$ and the $m = \pm 1/2 - m' = \mp 1/2$ subtransitions with Clebsch–Gordan coefficients 1 and $1/\sqrt{3}$, respectively, and the third and the forth contribution is due to the corresponding ω_2 excitation.

We now make the central assumption that the intensity to detuning ratios in the two field components are chosen to provide light shifts of about the same size, but that one frequency component is much more detuned from resonance than the other one,

$$|\delta_2| \gg |\delta_1|. \tag{2}$$

Under this condition, transitions between the two ground-state potentials are predominantly induced by the first field (ω_1) since the optical pumping rates scale with $G^{\pm}_i/\delta_i^2$. Thus neglecting the pumping effect of the second field (ω_2), the transition rates for transfering an atom out of the potential U_- into U_+ and vice versa can simply be written as

$$\Gamma_{\pm}(\mathbf{r}) = \frac{2\gamma}{9\delta_1^2} G^{\pm}_1(\mathbf{r}). \tag{3}$$

The steady-state force exerted on atoms at rest in the bichromatic field can be calculated as [6]

$$\mathbf{F}(\mathbf{r}) = -\Pi_-(\mathbf{r})\nabla U_-(\mathbf{r}) - \Pi_+(\mathbf{r})\nabla U_+(\mathbf{r}), \tag{4}$$

where

$$\Pi_{\pm}(\mathbf{r}) = \frac{\Gamma_{\pm}(\mathbf{r})}{\Gamma_+(\mathbf{r}) + \Gamma_-(\mathbf{r})} \tag{5}$$

are the steady-state occupation probabilities of the potentials $U_{\pm}(\mathbf{r})$.

In a usual, monochromatic optical lattice, where the maxima in the population of a light-shifted ground state always coincide with the potential minima, the steady-state force according to eq.(4) vanishes in an average over a unit cell of the lattice. A wavelength-averaged force occurs only as a motion-induced effect, leading to the well-known 'Sisyphus' cooling force in polarization gradients [6]. In the bichromatic field, however, optical pumping can preferentially populate locations where the potentials have a certain slope. This effect, which obviously depends in sign and magnitude on the slowly varying phase relation between the spatial modulations of optical pumping and light-shift potentials, is the basic mechanism of dipole force rectification.

As a simple 1D example, we now consider an extension of the well-known lin-$\perp$-lin polarization-gradient cooling scheme [6] in a bichromatic pair of counterpropagating laser beams with orthogonal linear polarizations. The counterpropagating traveling-wave components of same frequency have equal intensities, so that each of the two frequency components can be decomposed into a σ^+ and a σ^- standing wave with position-dependent saturation parameters

$$G^{\pm}_{1,2}(z) = g_{1,2}\left[1 \pm \cos(2k_{1,2}z)\right], \tag{6}$$

where the $g_{1,2}$ represent the position-independent saturation parameters of the single-frequency traveling wave components and $k_{1,2} = \omega_{1,2}/c$ are the two wave numbers. For this one-dimensional case the steady-state force according to Eq.(4) can be calculated to

$$F_z(z) = \frac{2}{3}\hbar\gamma \cos(2k_1 z) \left[k_1 \frac{g_1}{\delta_1} \sin(2k_1 z) + k_2 \frac{g_2}{\delta_2} \sin(2k_2 z)\right]. \quad (7)$$

The slowly varying rectified part $F_{rect}(z) = \langle F_z(z)\rangle_\lambda$ is then obtained by averaging over the optical wavelength,

$$F_{rect}(z) = -\frac{1}{3}\hbar k_2 \gamma \frac{g_2}{\delta_2} \sin(2\,\delta k\, z) \quad (8)$$

with $\delta k = k_1 - k_2$. The rectified force oscillates in space with a macroscopic period $L = \pi/|\delta k|$, and shows restoring character around $z = L/2$ with a spring constant $\kappa = 2\pi\hbar k_2 \gamma g_2/(3|\delta_2|) \times L^{-1}$.

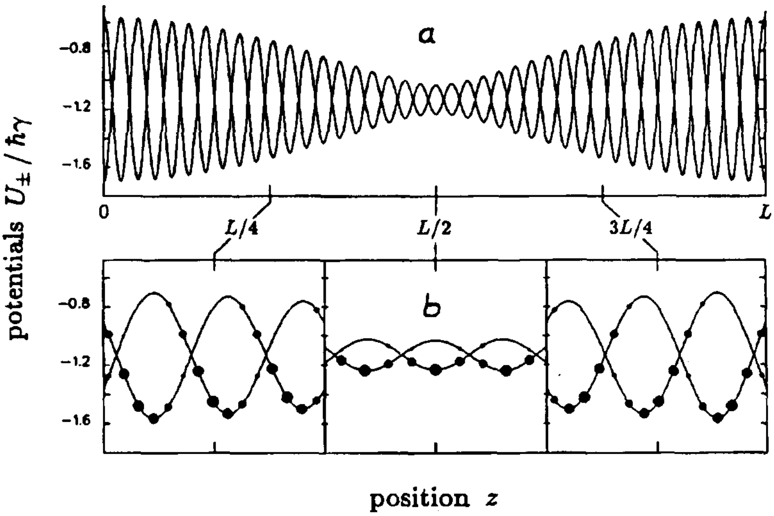

Fig.1. Bichromatic optical superlattice in 1D for $g_1/\delta_1 = -1$, $g_2/\delta_2 = -0.7$, and $k_2/k_1 = 20/21$: (a) Light-shifted ground-state potentials over a full rectification period L. (b) Dot sizes illustrating the occupation probabilities in the vicinity of $z = 0.25L$, $0.5L$, and $0.75L$, where the rectified force is positive, zero, and negative, respectively

These results are illustrated in fig.1: the potential curves $U_\pm(z)$ in (a) display the interference structure of the two different light-shift contributions, where the envelope of the modulations reflects the rectification period L. The spatial dependence of the occupation probabilities is shown in (b): for $z \approx 0.25L$ ($z \approx 0.75L$) one sees a preferential population of the negative (positive) potential slope, which leads to a positive (negative) rectified force. At $z \approx 0.5L$, i.e. in the trapping center of the superlattice, one recovers the situation of usual polarization-gradient cooling [6]: for atoms at rest, the maximum population of the light-shift potentials is found at their minima and thus, for moving atoms, the Sisyphus cooling mechanisms acts in the same way as in a monochromatic field. Here, the partially destructive interference between the two light-shift contributions is of advantage for the final cooling, because the low modulation depth of the potentials leads to a low limit temperature [6].

We have studied the dynamics and the equilibrium of the cooling and trapping process in such a superlattice site by performing a 1D semiclassical Monte-Carlo simulation based on the above equations, where the atom performs a classical motion in a light-shift potential between random quantum jumps transfering it from one ground state into the other. We also include the heating resulting from the photon momenta transfered in cycles of absorption and spontaneous emission. For the atom we take the parameters of ^{133}Cs, transition wavelength $\lambda = 852\,\mathrm{nm}$ and linewidth $2\gamma = 2\pi \times 5.3\,\mathrm{MHz}$, but for simplicity we keep the assumption of a $J = 1/2 - J' = 3/2$ transition. For the first field component we choose $g_1 = 25$ and $\delta_1 = -25$, similar to a usual monochromatic optical lattice, and in order to obtain a rectification period $L = 1\,\mathrm{mm}$ in the range of main experimental interest we take $\delta_2 = -56575$ for the second, far off-resonant field. We set $g_2 = 0.8\,|\delta_2| = 45260$ (traveling-wave intensity $\sim 500\,\mathrm{mW/mm^2}$), which produces a light-shift that is 20% less than the one induced by the first frequency component. The potential curves then look similar to the ones in fig.1(a) with the exception that now L/λ is on the order of 1000. As initial conditions for the numerical simulations, we choose parameters typical for a standard magneto-optical trap, assuming a Gaussian position distribution with a $1/e$-width of $0.5\,\mathrm{mm}$ and a temperature of $100\,\mu\mathrm{K}$.

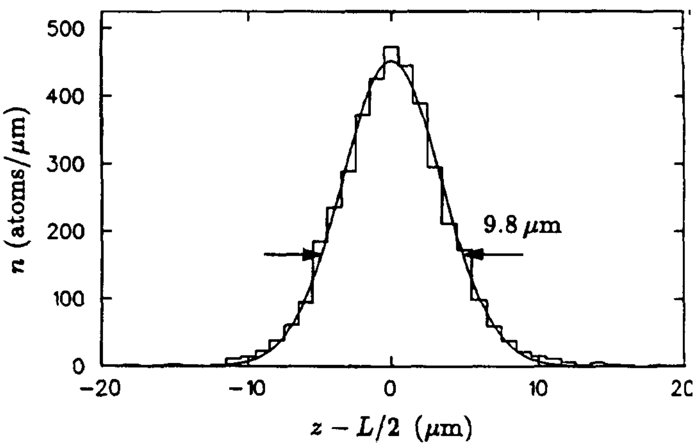

Fig.2. Result of a 1D Monte-Carlo simulation for the spatial equilibrium distribution in a bichromatic superlattice trap, performed with 4000 atoms. The solid line is a Gaussian fit to the data

Under these conditions, the results of our 1D simulation show a rapid spatial compression which completely reaches equilibrium after only 30 ms. We then find a temperature of $T = 10\,\mu\mathrm{K}$, as expected from the depth of the light-shift potentials in the trap center [6], and, most importantly, we observe a remarkably sharp spatial equilibrium distribution with a $1/e$-width of $\Delta z = 9.8\,\mu\mathrm{m}$ (see fig.2), which corresponds to a spatial compression of the initial distribution by more than a factor of 50. From these simulation results an effective spring constant of the trap can be calculated as $\kappa_{eff} = 8k_BT/\Delta z^2 = 1.2 \times 10^{-17}\,\mathrm{N/m}$, which is close to the idealized spring constant $\kappa = 2.1 \times 10^{-17}\,\mathrm{N/m}$ derived from eq.(8); the minor deviation is readily explained by the sharp velocity dependence of the

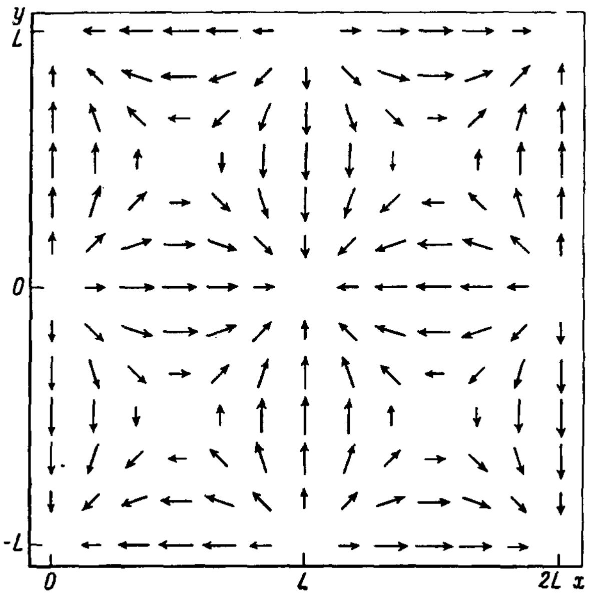

Fig.3. Rectified force field in the considered 2D configuration, illustrating a 2D bichromatic superlattice trap

rectified force [8]. This result demonstrates that the rectified force calculated as a wavelength average for atoms at rest contains not only qualitative but also important quantitative information on the macroscopic confining properties of the bichromatic light field.

In order to give an outlook on bichromatic optical superlattices in more than one dimension, let us now consider a simple and illustrative 2D example: In extension of [9] we consider two superimposed bichromatic standing waves in the $x-y$ plane, one along the x-axis with linear polarization in y-direction, and the other one along the y-axis with linear polarization in x-direction; the two relative time phases between the pairs of standing waves of same frequency are set to $\psi_{1,2}=90°$. This configuration leads to an optical lattice structure with $\sigma^{\pm}$ intensities represented by the saturation parameters

$$G^{\pm}_{1,2}(x,y) = 2g_{1,2}\left[\cos(k_{1,2}x) \pm \cos(k_{1,2}y)\right]^2. \tag{9}$$

A straightforward calculation by means of the above equations yields the rectified force field

$$\mathbf{F}^{2D}_{rect}(x,y) = -0.485\,\hbar k_2 \gamma\,\frac{g_2}{\delta_2}\begin{pmatrix}\sin(\delta k\,x)\,\cos(\delta k\,y)\\ \cos(\delta k\,x)\,\sin(\delta k\,y)\end{pmatrix}, \tag{10}$$

which is plotted in fig.3. One sees that macroscopic 2D traps exist with restoring forces being of similar strength as in the 1D case. In the same way, we have also considered a 3D bichromatic superlattice in extension of the four-beam configuration described in [12] and found macroscopic traps with restoring rectified forces acting in all directions. Since 3D optical lattice configurations [10-12] also provide very efficient polarization-gradient cooling, we expect bichromatic light to allow a similar spatial compression in each of the three dimensions as we have seen in our 1D simulation. Corresponding 3D Monte-Carlo simulations are in progress.

We finally point out a possible variation of the above scheme for alkali atoms. When tuning one frequency component close to the D_2 line and the other one to the D_1 line, a superlattice structure can be created at moderate laser power with a period determined by the fine-structure splitting of the excited state [13]. For Li, Na, K, Rb, and Cs one then obtains rectification periods $L = 1.5\,\mathrm{cm}$, $290\,\mu\mathrm{m}$, $87\,\mu\mathrm{m}$, $21\,\mu\mathrm{m}$, and $9.0\,\mu\mathrm{m}$, respectively. The mesoscopic superlattice formed for Rb and Cs atoms might be of particular interest for achieving an extremely fast local density compression into the trapping sites before substantial losses by ultracold trapped-atom collisions can occur.

Bichromatic optical superlattice traps are experimentally easy to realize and may have a bright future as powerful tool for experiments on ultracold atoms at high densities.

We gratefully acknowledge partial support by the Deutsche Forschungsgemeinschaft.

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