Trapping atoms by rectified forces in bichromatic optical superlattices

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We consider the trapping of atoms by sub-Doppler rectified dipole forces in the bichromatic superlattice which forms as macroscopic spatial interference patterns of two superimposed individual optical lattices that differ 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" which 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 field configurations and rectification schemes (e.g., Refs. 3–5).

In this letter we wish to point out possible ways of constructing highly efficient atom traps on the basis of the rectification effect. In contrast with 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 a very strong dependence on the velocity of atoms.[5,8] We will show that in simple laser beam configurations used in the recent work on optical lattices[9–12] bichromatic light creates an additional superlattice structure, in which rectified forces act together with the usual polarization-gradient cooling[6] to efficiently cool and confine atoms in an array of macroscopic traps.

In our theoretical model, which adopts 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 which 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 $G^{\pm}_{1,2}(\mathbf{r})=I^{\pm}_{1,2}(\mathbf{r})/I_{SAT}$, where $I_{SAT}=\hbar\omega_0^3\gamma/6\pi c^2$, and by the reduced frequency detunings $\delta_{1,2}=(\omega_{1,2}-\omega_0)/\gamma$, where ω_0 is the transition frequency, and 2γ is 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 the contributions from the scattering force (Doppler cooling) can be disregarded. We furthermore assume that the detunings are large ($\delta_{1,2} \gg 1$), the optical saturation is low ($G^{\pm}_{1,2}(\mathbf{r})/\delta^2_{1,2} \ll 1$), and the effect of mutual light shifts is negligible $[G^{\pm}_{1,2}(\mathbf{r})/\delta_{1,2} \ll \delta_{1,2}]$. This allows us to simply add up the individual light shifts as first-order perturbations in order to calculate the relevant ground-state potentials

$$U_{\pm}(\mathbf{r})=\frac{1}{2}\hbar\gamma\left(\frac{G_1^{\pm}(\mathbf{r})}{\delta_1}+\frac{G_1^{\mp}(\mathbf{r})}{3\delta_1}+\frac{G_2^{\pm}(\mathbf{r})}{\delta_2}+\frac{G_2^{\mp}(\mathbf{r})}{3\delta_2}\right). \tag{1}$$

Here the first term 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 fourth terms are due to the corresponding ω_2 excitation.

We now make the central assumption that the intensity of the detuning ratios in the two field components are chosen to provide light shifts of about the same size, and 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 $G_i^{\pm}/\delta_i^2$. Thus, ignoring the pumping effect of the second field (ω_2), we can write the transition rates for transferring an atom out of the potential U_- into U_+ and vice versa as follows:

$$\Gamma_{\pm}(\mathbf{r})=\frac{2\gamma}{9\delta_1^2}G_1^{\mp}(\mathbf{r}). \tag{3}$$

The steady-state force exerted on the 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 the 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 the polarization gradients.[6] In the bichromatic field, however, optical pumping can preferentially populate locations in which 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 the 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-⊥-lin polarization-gradient cooling scheme[6] in a bichromatic pair of counterpropagating laser beams with orthogonal linear polarizations. The counterpropagating traveling-wave components of the 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:

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

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 from

$$F_z(z)=\frac{2}{3}\hbar\gamma\cos(2k_1z)\left[k_1\frac{g_1}{\delta_1}\sin(2k_1z)+k_2\frac{g_2}{\delta_2}\sin(2k_2z)\right]. \tag{7}$$

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

$$F_{\mathrm{rect}}(z)=\frac{1}{3}\hbar k_2\gamma\frac{g_2}{\delta_2}\sin(2\delta kz), \tag{8}$$

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

These results are illutrated 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 $z\approx0.25L$ ($z\approx0.75L$) we see a preferential population of the negative (positive) potential slope, which leads to a positive (negative) rectified force. At $z\approx0.5L$, i.e., in the trapping center of the superlattice, the usual polarization-gradient cooling is restored:[6] for atoms at rest the maximum population of the light-shift potential is found at their minima, and thus for moving atoms the Sisyphus cooling mechanism acts in the same way as in a monochromatic field. Here, the partially destructive interference between the two light-shift contributions is beneficial 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 executes a classical motion in a light-shift potential between random quantum jumps which transfer it from one ground state to another. We also include the heating resulting from the photon momenta which have transferred in cycles of absorption and spontaneous emission. For the atom we use the parameters of ^{133}Cs, the transition wavelength $\lambda=852$ nm, and the linewidth $2\gamma=2\pi\times5.3$ MHz, but for simplicity we assume a $J=1/2-J'=3/2$ transition. For the first field

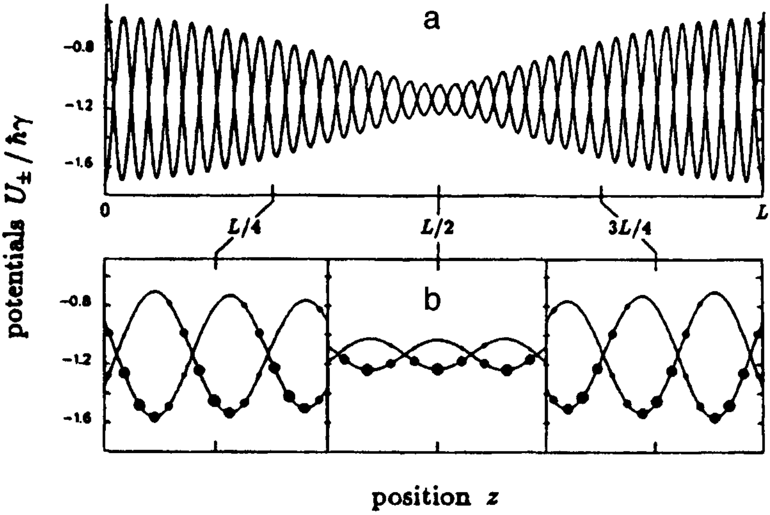

FIG. 1. Bichromatic optical superlattice in 1D for $g_1/\delta_1=-1$, g_2/δ_2=-0.7, and k_2/k_1=20/21; (a) Light-shifted ground-state potentials over a full rectification period L. (b) The dot sizes illustrate the occupation probabilities in the vicinity of $z=0.25L$, $0.5L$, and $0.75L$, where the rectified force is positive, zero, and negative, respectively.

component we choose $g_1=25$ and $\delta_1=-25$, similar to the usual monochromatic optical lattice. To obtain a rectification period $L=1$ mm in the range of main experimental interest, we take $\delta_2=-56575$ for the far off-resonant second field. We set $g_2=0.8|\delta_2|=45260$ (traveling-wave intensity ~500 mW/mm^2), which produces a light-shift that is 20% smaller than the one induced by the first frequency component. The potential curves then look similar to those 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 mm and a temperature of 100 μK.

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 $T=10$ μK, 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 $\Delta z=9.8$ μm (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 $\kappa_{\text{eff}}=8k_BT/\Delta z^2=1.2\times10^{-17}$ N/m, which is close to the idealized spring constant $\kappa=2.1\times10^{-17}$ N/m derived from Eq. (8); the minor deviation is readily explained by the sharp velocity dependence of the rectified force.[8] This result demonstrates that the rectified force calculated as a wavelength average for atoms at rest contains not only qualitative information 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

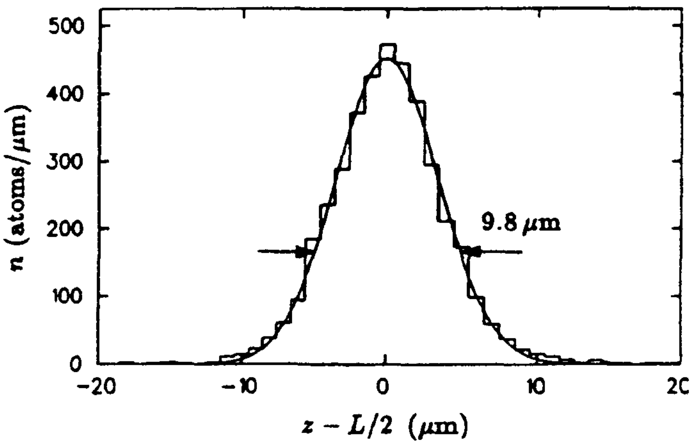

FIG. 2. The result of a 1D Monte Carlo simulation of the spatial equilibrium distribution in a bichromatic superlattice trap, performed with 4000 atoms. The solid line is a Gaussian fit to the data.

dimension, let us consider a simple and illustrative 2D example: In the extension of Ref. 9 we consider two superimposed bichromatic standing waves in the $x-y$ plane, one along the x axis with linear polarization in the y direction and the other one along the y axis with linear polarization in the x direction; the two relative time phases between the pairs of standing waves of the 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

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

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

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

which is plotted in Fig. 3. We see that there are macroscopic 2D traps with restoring forces of similar strength, as in the 1D case. In the same way, we have also considered a 3D bichromatic superlattice in the extension of the four-beam configuration described in Ref. 12, and we 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 the 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 close to the D_1 line, a superlattice structure can be created at a moderate laser power with a period determined by the fine-structure splitting of the excited state.[13] For Li, Na, K, Rb, and Cs we then obtain rectification periods $L=1.5$ cm, 290 μm, 87 μm, 21 μm, and 9.0 μm, respectively. The mesoscopic superlattice which is formed for Rb and Cs atoms might be

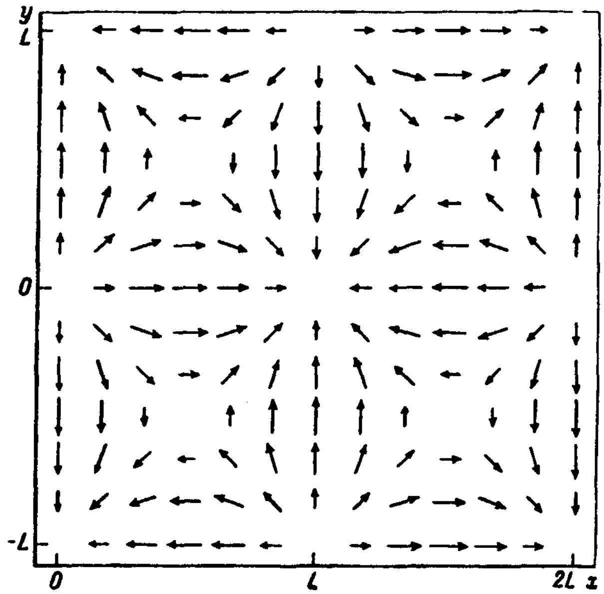

FIG. 3. Rectified force field in the considered 2D configuration, illustrating a 2D bichromatic superlattice trap.

of particular interest for achieving an extremely fast local density compression in to 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 a powerful tool for experiments on ultracold atoms at high densities.

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[1] A. P. Kazantsev and I. V. Krasnov, Pis'ma Zh. Éksp. Teor. Fiz. **46**, 264 (1987) [JETP Lett. **46**, 332 (1987)]; J. Opt. Soc. Am. B **6**, 2140 (1989).
[2] A. P. Kasantsev, *Memorial Issue on Laser Cooling and Trapping*, Laser Phys.**4**, 5 (1994).
[3] R. Grimm, Yu. B. Ovchinnikov, A. I. Sidorov, and V. S. Letokhov, Phys. Rev. Lett. **65**, 1415 (1990).
[4] J. Söding, R. Grimm, J. Kowalski *et al.*, Europhys. Lett. **20**, 101 (1992).
[5] R. Gupta, C. Xie, S. Padua *et al.*, Phys. Rev. Lett. **71**, 3087 (1993).
[6] J. Dalibard and C. Cohen-Tannoudji, J. Opt. Soc. Am. B **6**, 2023 (1989).
[7] J. Javanainen, Phys. Rev. Lett. **64**, 519 (1990).
[8] R. Grimm, J. Söding, Yu. B. Ovchinnikov, and A. I. Sidorov, Optics Commun. **98**, 54 (1993).
[9] A. Hemmerich and T. W. Hänsch, Phys. Rev. Lett. **70**, 410 (1993).
[10] G. Grynberg, B. Lounis, P. Verkerk *et al.*, Phys. Rev. Lett. **70**, 2249 (1993).
[11] A. Hemmerich and T. W. Hänsch, Europhys. Lett. **22**, 89 (1993).
[12] P. Verkerk, D. R. Meacher, A. B. Coates *et al.*, Europhys. Lett. **26**, 171 (1994).
[13] R. Grimm, Verhandl. DPG (VI) **29**, Q3A.1 (1994); Europ. Research Conf. on Electronic and Atomic Collisions, Giens, France, 10–15 Sept. 1994, Book of Abstracts, Eds. H. Hotop and M. W. Ruf, 40 (1994).

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