Ratchet for Cold Rubidium Atoms: The Asymmetric Optical Lattice

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It was suggested recently that particles evolving in a bipotential that is macroscopically flat but locally asymmetric can be set into directed motion [J. Prost *et al.*, Phys. Rev. Lett. **72**, 2652 (1994)]. We show evidence of this effect using cold rubidium atoms in a grey optical lattice. [S0031-9007(98)08213-1]

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Molecular motors that convert chemical energy into mechanical energy, for example in muscles, are a subject of considerable interest. Motivated by recent experimental observations on individual molecular motors [1], a ratchet-type model was proposed to explain the molecular dynamics [2]. This model describes the molecules as classical or quantum [3] Brownian particles evolving in a periodic sawtooth potential. Because both parity and time reversal symmetry are violated, a directional motion of particles is possible, although there is apparently no net force acting on them [4]. Model experiments performed with colloidal particles or polystyrene spheres showed good agreement with theoretical predictions [5]. More elaborate theoretical models consider the particle as existing in two states, each of them experiencing a different potential. Both potentials are *periodic* but *asymmetric*. When the transition rates between the two states differ from their values at Boltzmann equilibrium, a directional average motion is predicted [6]. In this paper we present a closely related situation of an asymmetric bipotential occurring in laser cooling and report the observation of directed motion of rubidium atoms in this optical bipotential.

Consider an atom with an electronic transition connecting a ground state having an angular momentum $F = 1$ to an excited state having the same angular momentum $F' = 1$. If this atom interacts with an electromagnetic field polarized in the xOy plane (no π component), it is well known that all the physics can be described in terms of a Λ system which includes the two ground-state Zeeman sublevels $m = -1$ and $m = 1$ and the excited sublevel $m' = 0$ [7]. This transition has been widely studied in the context of velocity selective cooling by population trapping (VSCPT) [8]. Furthermore, if the light intensity is sufficiently small or if the frequency detuning Δ from the atomic resonance is sufficiently large, it is possible to eliminate adiabatically the excited sublevel and to describe all the physical processes inside the two-level subspace $\{m = -1, m = 1\}$ [7]. The master equation then contains (i) a Hamiltonian evolution in the optical bipotential due to the light shifts and (ii) damping processes originating from absorption-spontaneous emission cycles. Because such a Λ system accommodates a state $|NC\rangle$ which is not coupled to the excited state [9], the corresponding optical potential V_{NC} is flat ($V_{NC} = 0$) and the Hamiltonian appears as the sum of the kinetic energy $p^2/2M$ (where M is the atomic mass) and of the potential energy $V_C|C\rangle\langle C|$ (with $\langle NC \mid C\rangle = 0$).

In the following, we consider a unidimensional (1D) situation with two counterpropagating beams $\mathbf{E}_1$ and $\mathbf{E}_2$ of the same amplitude E_0, same frequency ω and linear polarizations $\mathbf{e}_1 = \cos(\theta/2)\mathbf{e}_x + \sin(\theta/2)\mathbf{e}_y$ and $\mathbf{e}_2 = \cos(\theta/2)\mathbf{e}_x - \sin(\theta/2)\mathbf{e}_y$ (lin-θ-lin configuration). The optical potential associated with $|C\rangle$ is $V_C = \hbar\Delta s D(z)/2$, where $\Delta = \omega - \omega_0$ (ω_0: atomic resonance frequency), $s = \frac{\Omega^2/2}{\Delta^2+\Gamma^2/4}$, $\Omega = dE_0/\hbar$ (d being the reduced matrix element of the electronic dipole moment), Γ the natural width of the excited state, and $D(z) = 1 + \cos\theta\cos(2kz)$. In order to have a spatially modulated potential V_{NC}, a small magnetic field $\mathbf{B}_0$ along Oz can be added. The Zeeman shift then leads to

$$V_{NC} = -\hbar\Omega_0 \sin\theta \sin(2kz)/D(z)\,, \tag{1}$$

with $\Omega_0 = g\mu_B B_0$ (g: Landé factor of the ground state, μ_B: Bohr magneton). The potentials V_{NC} and V_C (the latter now differs slightly from $\Delta s D(z)/2$ because of the Zeeman effect) are shown in Fig. 1. Such a system corresponds to a 1D grey optical lattice where efficient Sisyphus cooling is achieved for $\Delta > 0$ [10], leading to trapping of a large fraction of the atoms in the asymmetric wells of V_{NC}. Consider an atom oscillating in such a well (between A and B in Fig. 1). The transition probability from $|NC\rangle$ to $|C\rangle$ due to B_0 shows a smooth variation between A and B. In particular, a transition can start from a point M located on the left-hand side of the potential minimum. Once in $|C\rangle$ ($M \to N$ in Fig. 1), the atom undergoes an oscillation that is interrupted by a spontaneous Raman process that brings the atom back into $|NC\rangle$ ($P \to Q$ in Fig. 1). The jumps between the potential curves thus tend to drag the atom to the right-hand side of Fig. 1. We observe that because of the long lifetime of atoms in $|NC\rangle$, a jump from $|NC\rangle$ to $|C\rangle$ is a rather rare event, so that the atoms undergo many oscillations in a well of the $|NC\rangle$ potential between successive jumps. (There is another reason for the jumps which is the nonadiabatic coupling—this coupling is essential for the velocity damping.) It should also be noticed that in

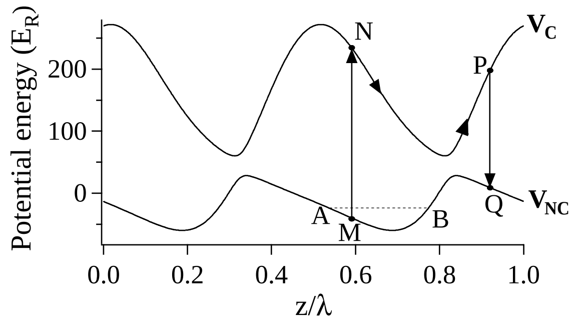

FIG. 1. Optical bipotential in the lin-θ-lin configuration with a longitudinal magnetic field $\mathbf{B}_0$. The parameters are $\theta = 45°$, $\Delta/\Gamma = 2$, $\Omega/\Gamma \simeq 0.65$, and $\Omega_0 = 60\omega_R$ [where $\hbar\omega_R = \hbar^2\omega^2/(2Mc^2)$ is the recoil energy]. The curves represent the exact optical potential, obtained by diagonalizing the total Hamiltonian (light shift and Zeeman shift). An atom oscillating between A and B can shift to the neighboring well because of a transition in V_C.

the case of thermal equilibrium where detailed balance is satisfied, the transition probability from $|NC\rangle$ to $|C\rangle$ peaks near B (Fig. 1), which inhibits the directed motion.

On inspection of Eq. (1), we see immediately that the asymmetry of V_{NC} changes sign when $\mathbf{B}_0$ is reversed. Second, we remark that the potential becomes symmetrical for $\theta = 0, \pm\pi/2$. Third, V_{NC} has opposite asymmetries for $\theta = \theta_0$ and $\theta = -\theta_0$. Implications for the occurrence and the sign of the directed motion will be presented with the experimental results.

To perform the experimental study, we start from rubidium atoms (^{87}Rb) trapped and cooled in a magneto-optical trap (MOT) [11]. At a certain time we switch off the MOT beams and inhomogeneous magnetic field and switch on the grey lattice beams operating on the blue side of the $5S_{1/2}(F = 1) \rightarrow 5P_{3/2}(F' = 1)$ transition, the static magnetic field $\mathbf{B}_0$ and repumping beams blue-detuned by 2Γ from the $5S_{1/2}(F = 2) \rightarrow 5P_{3/2}(F' = 2)$ transition (which prevent a leakage of atoms into the $F = 2$ hyperfine sublevel). The counterpropagating grey lattice beams are vertical while the repumping beams lie in the horizontal plane. The spatial displacement of cold atoms is measured using a CCD camera by direct imaging of the atomic cloud (whose initial radius is $\sim$1 mm). One technique is to use the faint image of the grey lattice itself. This requires the averaging of a significant number of images ($\sim$100) because atoms in $|NC\rangle$ scatter very few photons. An example of measurements done with this method is shown in Fig. 2a where the atomic displacement d as a function of the lattice duration τ_L is presented. After a transient period, the displacement increases linearly with the lattice duration, showing the occurrence of a directional motion with an average velocity $v = 100$ mm/sec (i.e., $v \simeq 17v_R$ where v_R is the recoil velocity $v_R = \hbar\omega/Mc$). Note that this velocity is larger by several orders of magnitude than the velocities found in macrophysics [5,6]. When the sign of B_0 is reversed, the asymmetry of V_{NC} is inverted and we

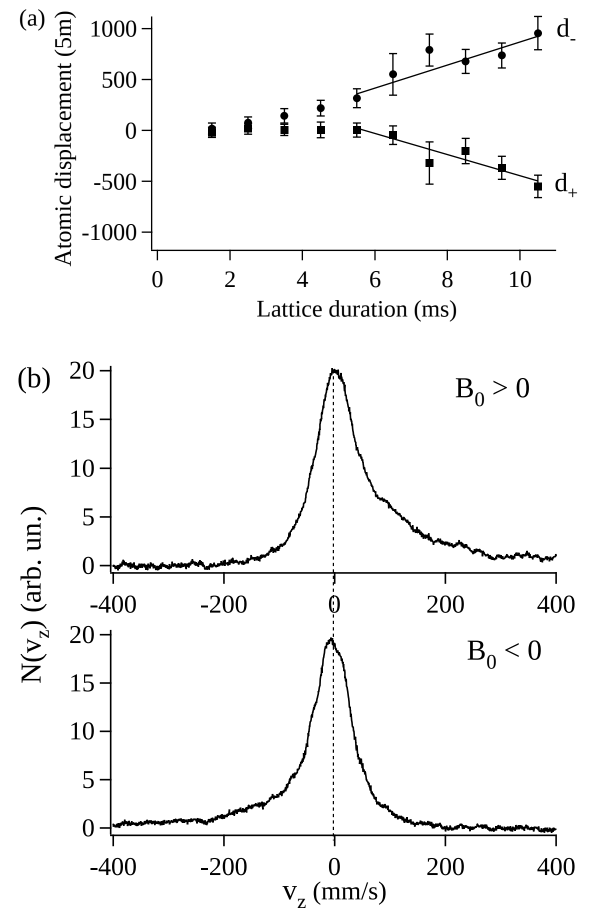

FIG. 2. (a) Atomic displacement d_+ (squares, for $B_0 > 0$) and d_- (circles, for $B_0 < 0$) as a function of the duration τ_L of the asymmetric lattice. The experimental parameters are $\theta = 45°$, $\Delta/\Gamma = 2$, $\Omega_0 = 140\omega_R$, and $s \simeq 0.3$. The displacements were measured by averaging 100 images of the cloud taken during the grey lattice phase. The axis is oriented in the direction of gravity. Gravity is the origin of the asymmetry of the displacements observed for opposite magnetic fields. (b) Velocity distribution $N(v_z)$ of the atoms for $B_0 > 0$ (top) and for $B_0 < 0$ (bottom), obtained by a time-of-flight method for $\theta = 45°$, $\Delta/\Gamma = 2$, $|\Omega_0| = 95\omega_R$, $s \simeq 0.1$, $\tau_L = 6$ ms. The asymmetry of the two velocity distributions is reversed for opposite signs of B_0. For $B_0 = 0$ the velocity distribution is symmetric and centered around $v_z = 0$ (dashed line). The effect of gravity is also visible here ($\overline{v} = 7$ mm/s for $B_0 > 0$ and $\overline{v} = -5$ mm/s for $B_0 < 0$).

observe an atomic displacement occurring in the opposite direction. However, the asymptotic velocity measured in this case is not exactly opposite to the velocity measured in the former case. Because of the additional effect of gravity, v is always larger for a motion directed downwards than for a motion directed upwards.

In fact, we can get information that goes beyond the sole determination of the average velocity. We can determine the whole velocity distribution using a ballistic method. For this purpose, we create a sheet of light 10 cm below the cloud of cold atoms. This sheet of light is resonant with the $5S_{1/2}(F = 1) \rightarrow 5P_{3/2}(F' = 2)$ transition. When the grey lattice beams and $\mathbf{B}_0$ are switched off, the rubidium atoms fall freely. The time variation of the absorption when the atoms cross the sheet of light is directly related to the velocity distribution [11]. We present in Fig. 2b the velocity distributions $N(v_z)$ measured for $\Omega_0 = \pm 95\omega_R$ and for $\theta = -45°$, $\Delta/\Gamma = 2$, $s \simeq 0.1$, $\tau_L = 6$ ms. One notices that the asymmetry of the velocity distribution is reversed for opposite signs of B_0 and is mostly located in the wings of $N(v_z)$. In fact, the core of $N(v_z)$ is associated with the atoms localized in a well of the V_{NC} potential where they undergo an oscillatory motion. On this core is superimposed the contribution from atoms traveling from one well to another because of a transition from $|NC\rangle$ to $|C\rangle$ (see Fig. 1).

The experimental results presented hereafter were also obtained by imaging, but instead of tracking the faint image of the grey lattice, we applied after a grey lattice phase τ_L a flash of two beams nearly resonant on the $5S_{1/2}\,(F = 1) \rightarrow 5P_{3/2}\,(F' = 2)$ and $5S_{1/2}\,(F = 2) \rightarrow 5P_{3/2}\,(F' = 3)$ transitions. This gives the displacement d of the atomic cloud at $t = \tau_L$. To avoid the supplementary effect of gravity, we measured the displacements d_+ and d_- for opposite values of $\mathbf{B}_0$ and plotted the variation of $\overline{d} = (d_- - d_+)/2$. We show in Fig. 3 the dependence of $\overline{d}$ versus θ for $\Omega_0 = -170\omega_R$, $\Delta/\Gamma = 2$, $s \simeq 0.2$, and $\tau_L = 6$ ms. As expected $\overline{d} = 0$ when the potentials V_C and V_{NC} become symmetric, i.e., for $\theta = 0, \pm\pi/2$. We also find that $\overline{d}$ is an odd function of θ, in agreement with the change of asymmetry of the potentials under the transformation $\theta \rightarrow -\theta$. We also studied the variation of the displacement $\overline{d}$ with the Zeeman shift Ω_0. It increases at low Ω_0, reaches a maximum for $\Omega_0 \sim \Delta s/2$ and decreases when Ω_0 further increases. These observations correspond to the fact that the potentials lose their asymmetry for $\Omega_0 \longrightarrow 0$ and $\Omega_0 \gg \Delta s/2$. [In this last case the perturbative expression of Eq. (1) is no longer valid—one has to calculate the exact potential.]

The preceding experimental observations are supported by results from a semiclassical numerical Monte Carlo simulation of the atom dynamics [12]. For example, we show in Fig. 4 the time evolution of the spatial distribution $N(z)$ of an initially well-localized atomic sample (parameters: $\theta = 45°$, $\Delta/\Gamma = 2$, $\Omega/\Gamma = 1.8$, $\Omega_0 = 250\omega_R$). The directed motion and the spreading of the density are clearly visible in these data. (We have checked that such a motion disappears when we force the transition rates to obey detailed balance.) The average velocity is $v = 1.5 v_R$. This velocity is three times smaller than the velocity found experimentally in the same conditions. In fact, in the case of a $1 \rightarrow 1$ transition it is well known that the atomic dynamics for small B_0 requires a full quantum treatment [8] and quantum corrections remain important as long as $\Omega_0 < \Delta s$ [13]. The semiclassical simulation is quantitatively correct only when $\Omega_0 \gg \Delta s$ [14]. Indeed, there might be some relation between this discrepancy and the recent observation by Reimann *et al.* [3] that quantum tunneling can significantly modify the flux in a similar problem. (An additional problem in the simulation is the fact that the distant hyperfine transitions are neglected.)

It is interesting to compare our experimental situation with previous studies on cold atoms where a directed motion is found. The simplest case corresponds to $\sigma^+ - \sigma^-$ molasses in the presence of a longitudinal magnetic field. In this case, the atoms are cooled and bunched around a velocity group that shifts linearly with B_0 [15]. A similar argument can be applied for VSCPT in the same field configuration. However, it can easily be shown that the optical potentials for this field configuration are *flat* whatever the angular momenta J and J' of the ground and excited states. The microscopic description of the directed motion is thus completely different. Another related problem was studied by Grimm *et al.* [16] and described in terms of rectified forces. This force is found for a field configuration identical to the one considered in this paper, but theory and experiments were developed for $J \rightarrow J'$ transitions and a strong saturation is generally

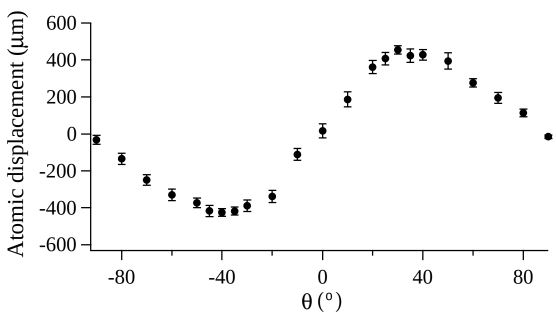

FIG. 3. Experimental results showing the atomic displacement $\overline{d} = (d_- - d_+)/2$ as a function of the angle θ between the linear polarizations of the two lattice beams, for $\Delta/\Gamma = 2$, $\Omega_0 = -170\omega_R$, $s \simeq 0.2$, $\tau_L = 6$ ms.

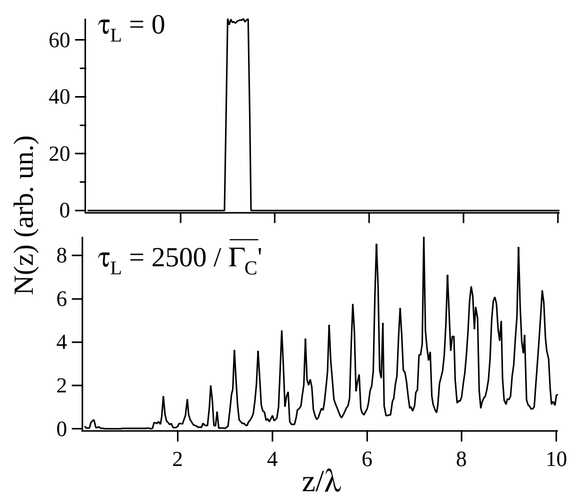

FIG. 4. Numerical results obtained by a semiclassical Monte Carlo simulation, showing the temporal evolution of the spatial distribution $N(z)$ of the atoms, for $\theta = 45°$, $\Delta/\Gamma = 2$, $\Omega/\Gamma = 1.8$, $\Omega_0 = 250\omega_R$. The unit of time is $1/\overline{\Gamma'_C}$, where $\overline{\Gamma'_C}$ is the atomic lifetime in $|C\rangle$ averaged over one spatial period.

considered. As a result, the atoms undergo frequent jumps between the various optical potentials (the distance covered by an atom in a potential curve is much smaller than the spatial period). In this case, which corresponds to the so-called "jumping regime," one can define at each point an average force (the dipole force) [7]. This force derives from a potential which is *not* periodic. From the average slope of this potential, one can determine the velocity around which the atoms are bunched. Here one expects a shift of the velocity distribution and not a modification of the wings as in the case of our experiment (Fig. 2b). The difference is that our study is performed in the "oscillating regime" where localized atoms undergo several oscillations inside a potential well of the lowest potential curve (V_{NC}) before being transferred towards the other potential (V_C). It is well known [7,12] that the dipole potential cannot be used in the oscillating regime and the dynamics must be studied with all the potential curves. (For the $1 \rightarrow 1$ transition considered here, provided the Zeeman shift $\hbar\Omega_0$ is much larger than the recoil energy $\hbar\omega_R$, the oscillating regime is found when $\Omega_0 < s\Delta^2/\Gamma$, the jumping regime occurring above this value.) To conclude this section it is interesting to note that although neither the $\sigma^+ - \sigma^-$ configuration nor the rectified force leads to a problem that is described by an asymmetric periodic bipotential, they do, however, correspond to a field configuration in agreement with the Curie symmetry law [4].

In conclusion, we have shown that cold atoms provide an interesting demonstration for a statistical physics problem that was originally considered for a completely different subject [6]. Because it is easy to tailor potentials and dissipation with cold atoms, we believe that tighter links between statistical physics and laser cooling will be developed in the future.

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