Laser Cooling in an Optical Shaker

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(Received 30 August 2004; published 20 April 2005)

We propose a novel generic approach to laser cooling based on the nonresonant interactions of atoms and molecules with optical standing waves experiencing sudden phase jumps. The technique, termed "optical shaking," combines the elements of stochastic cooling and Sisyphus cooling. An optical signal that measures the instantaneous force applied by the standing wave on the ensemble of particles is used as feedback to determine the phase jumps. This guarantees a drift towards lower energies and higher phase-space density without the loss of particles typical of evaporative cooling.

DOI: 10.1103/PhysRevLett.94.153002 PACS numbers: 33.80.Ps, 32.80.Pj

Laser cooling has led to a number of remarkable achievements, including the observation of Bose-Einstein condensation [1]. Most current schemes for laser cooling are based on the exchange of momentum between resonant light and a closed atomic system consisting of few active levels. Very few atoms obey the stringent requirements of these resonant interactions, while most other atoms and practically all molecules are not amenable to existing schemes of resonant laser cooling. Moreover, the deepest stages of cooling just prior to Bose-Einstein condensation are usually achieved via evaporative cooling, a nonoptical scheme in which evaporation of the "hotter" atoms leads to the cooling of the rest of the ensemble and a permanent loss of almost 99% of the initially trapped atoms [2]. The elusive task of cooling molecules [3] has been much more difficult. In several different recent experiments [4–9], during evaporative cooling of fermionic atoms, dimers were formed near Feshbach resonances and were observed to condense to a Bose-Einstein condensate. For polar molecules, such as NH_3 and ND_3, Meijer and collaborators have demonstrated decelerating and even stopping of molecular beams by properly switched static electric fields [10], but a general scheme is not yet available for the cooling of molecules.

We propose a new approach to laser cooling, which relies on *nonresonant* interaction of atoms or molecules with laser fields. The method combines the idea of Sisyphus cooling [11] and the concept of stochastic cooling [12,13], recently considered as promising for the cooling of neutral atoms [14]. An ensemble of precooled and trapped atoms or molecules is allowed to interact with a "shaking" optical standing wave. The very simple optical setup is depicted in Fig. 1, and consists of two nearly counterpropagating laser beams, an electro-optic modulator to introduce sudden phase changes to one of the beams, and two detectors to measure the intensity of the laser beams after they cross the interaction region.

To illustrate the principle of operation of the optical shaker, consider a collection of N atoms that interact with a nonresonant optical standing wave produced by two counterpropagating laser beams $[A\cos(kx - \omega_l t + \varphi)$ and $A\cos(kx + \omega_l t)]$. Here A is the electric field amplitude, φ is the relative phase between the beams, and $k = 2\pi/\lambda$, where λ is the wavelength of light. The spatially dependent potential energy of interaction with the standing wave is (after a constant space-averaged term has been dropped) $V(x, \varphi) = -V_0\cos(2\pi x/\Lambda + \varphi)$, where $V_0 = \alpha A^2/2$, $\Lambda = \lambda/2$, and α is the nonresonant polarizability at frequency ω_l. For two-level atoms of transition frequency ω_0, the polarizability is given by $\alpha = |d_{12}|^2/\hbar\Delta$, where $\Delta = \omega_0 - \omega_l$ is the detuning of the transition. In what follows, we consider high-field seeking atoms, namely, the case of a red-detuned optical lattice ($\Delta > 0$).

When a particle of mass m and energy E moves in this periodic potential, it slows down while climbing a potential hill, and speeds up when sliding down the hill. If the energy E exceeds the maximum value of V, the particle propagates over many periods of the potential without changing the total energy. Assume now that the standing wave makes a sudden spatial displacement, such as resulting from a sudden change of the phase φ. For a fast enough jump, the kinetic energy of the particle remains unchanged, while the potential energy may increase or decrease depending on the initial position of the particle and the value of the phase jump. In order to guarantee that the shaking decreases the potential energy (and hence the total energy) of the whole ensemble of N particles, we introduce a

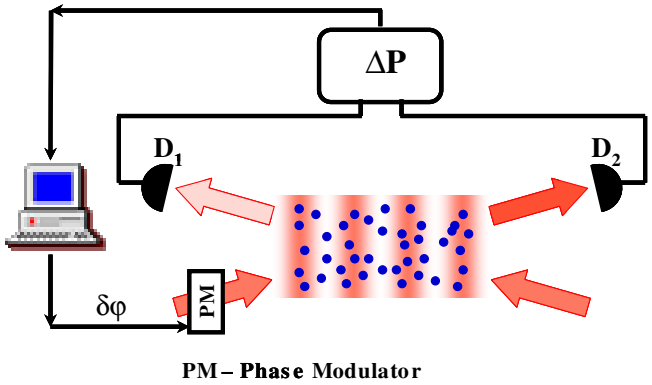

FIG. 1 (color online). Schematic setup of the optical shaker.

feedback loop that controls the value of the sudden phase jumps. The feedback signal is proportional to the total instantaneous force exerted by the standing wave upon the set of N particles. This force depends on the current value of the phase, φ, and on the spatial position of all the particles, and is given by

$$F(\varphi, t) = -\frac{2\pi V_0}{\Lambda} \sum_{j=1}^{N} \sin\left(\frac{2\pi x_j(t)}{\Lambda} + \varphi\right). \quad (1)$$

The gradient force [15] [Eq. (1)] appears because of the scattering of photons from one of the laser beams to the other one. Such a force can be measured by two detectors recording the intensity of the laser beams forming the standing wave *after* their interaction with the atomic sample. Because of the conservation of momentum, $F(\varphi, t) = \Delta P(\varphi, t)/c$, where c is speed of light, and $\Delta P = P_1 - P_2$ is the difference between the power levels measured by the detectors D_1 and D_2 (see Fig. 1). Similar arrangements have been used to study atomic wave-packet motion in optical lattices [16,17], and also for real-time feedback control of coherent oscillations of cold atoms trapped by a standing optical wave [18]. For a given set of atomic coordinates x_j, the force $F(\varphi, t)$ is a periodic function of φ. A nonzero value of the detected force indicates that the system is not at a potential minimum, and thus may be brought to the minimum by a proper sudden phase shift $\delta\varphi$, such that $F(\varphi + \delta\varphi, t) = 0$. It is easy to show with the help of Eq. (1) that

$$\tan(\delta\varphi) = -\frac{F(\varphi, t)}{F(\varphi + \pi/2, t)} = -\frac{\Delta P(\varphi, t)}{\Delta P(\varphi + \pi/2, t)}. \quad (2)$$

The solution of this equation, which corresponds to the potential minimum (not maximum), is given by

$$\delta\varphi = -\arctan\left[\frac{\Delta P(\varphi, t)}{\Delta P(\varphi + \pi/2, t)}\right] + \frac{\pi}{2}\{1 + \mathrm{sgn}[\Delta P(\varphi + \pi/2, t)]\}. \quad (3)$$

Thus, two measurements of ΔP are needed to predict the phase shift leading to a guaranteed energy extraction at each jump. Remarkably, accurate knowledge of the signal $\Delta P(\varphi, t)$ itself is not needed, and in order to find $\delta\varphi$, only the ratio of two such signals and their sign is required, greatly simplifying the calibration of the measuring system. Naturally, the two measurements should be performed in rapid succession such that the atoms do not significantly move. Immediately after a jump, the particles (on average) climb the hills of the shifted standing wave and decrease the average kinetic energy as well (*Sisyphus cooling*). To allow for efficient cooling, the system should be allowed to freely evolve and remix between any two consecutive phase jumps

Let us estimate the required phase shift $\delta\varphi$ and energy loss ΔE following a single phase jump. For a "well-mixed" ensemble of a large number $N \gg 1$ of particles, $F(\varphi, t)$ presents a random process with a zero mean value and a correlation time of the order of Λ/ν_T, where ν_T is a typical particle velocity. If the optically induced potential is rather weak ($V_0 \ll m\nu_T^2/2$), most of the particles are not trapped by the standing wave, and travel over it almost freely. The trigonometric sum defining $F(\varphi, t)$ in Eq. (1) can be then estimated as $\sum_{j=1}^{N} \sin[2\pi x_j(t)/\Lambda + \varphi] \approx \sqrt{N/2}B$, where B is a normally distributed random number with zero mean value and unity variance. The value of $F(\varphi + \pi/2, t)$ depends on the sum $\sum_{j=1}^{N} \cos[2\pi x_j(t)/\Lambda + \varphi]$, and is proportional to the potential energy of the whole ensemble at the moment of jump. The time-dependent potential energy may be treated again as a random process; however, it has a nonzero mean value depending on the strength of the potential. The corresponding sum may be approximated in the weak-field limit as $\sum_{j=1}^{N} \cos[2\pi x_j(t)/\Lambda + \varphi] \approx rN + \sqrt{N/2}A$. Here A is another random number independent of B but having the same statistical properties. The parameter $r \ll 1$ depends mainly on the particles with energy $E \sim V_0$ whose motion is "essentially" affected by the standing wave. Particles trapped near the bottom of the potential wells span only a limited range of the coordinate space where $\cos[2\pi x(t)/\Lambda + \varphi] \sim 1$. On the other hand, particles moving in the upper part of the standing wave slow down near its maxima and spend more time in the regions where $\cos[2\pi x(t)/\Lambda + \varphi] \sim -1$. The balance of these two effects contributes to the nonzero value of the parameter r. It can be shown that $r = 0.5V_0/k_BT$ if a relatively weak ($V_0/k_BT \ll 1$) one-dimensional standing wave is suddenly applied to a thermal ensemble (k_B is Boltzmann's constant). The phase shift $\delta\varphi$ needed in order to bring the whole system to the potential minimum and the energy loss per particle ΔE are given by

$$\tan(\delta\varphi) = -B/(\rho + A), \qquad \Delta E = \frac{1}{N}\frac{V_0}{r} f(\rho),$$
$$f(\rho) = \frac{\rho}{2}\left[\sqrt{(\rho + A)^2 + B^2} - (\rho + A)\right]. \quad (4)$$

Here $\rho = r\sqrt{2N}$ is an important parameter that determines the scale of the phase jumps. For $\rho \leq 1$ the phase shift $\delta\varphi$ is of the order of 1 rad, while for $\rho \gg 1$ the phase jump obeys $\delta\varphi \sim N^{-1/2}$, with the crossover occurring near $\rho \sim 1$. For a very large number of particles, the required phase jump tends to be impractically small, which will require the use of weaker laser fields to induce the optical standing wave, which in turn dictates a smaller ρ. The possibility of using variable field strength to optimize the process is discussed below. The expectation value of the energy loss per jump can be obtained by averaging ΔE over A and B. The averaged function $\langle f(\rho)\rangle$ behaves as $\rho\sqrt{\pi/8}$ for very small ρ, and tends to an asymptotic value of 0.25 for $\rho \gg 1$. This asymptotic regime can be achieved for relatively

weak fields [$r \leq (2N)^{-1/2}$]. For a thermal ensemble and $\rho \sim 1$, the averaged ΔE reaches the saturated value $\langle \Delta E \rangle = V_0/(4Nr) = k_B T/(2N)$ that is independent of V_0. This gives a rough estimate for a number of phase jumps ($\sim 2N$) needed for a substantial cooling.

Figures 2 and 3 depict numerical simulations of optical shaking in a harmonic trap. The amplitude of the light-induced potential was chosen as $V_0 = 0.1 k_B T$ (T is the initial temperature). Figure 2(a) displays the average particle energy (in units of $k_B T$) as the function of the number of controlled phase jumps. The jumps were evenly distributed in time. The time interval between them was of the order of a period of oscillations of low-energy particles trapped by the wave. For the chosen parameter values, the cooling proceeds in the regime of $\rho \geq 1$. At the initial stage, the mean particle energy loss per jump is of the order of the asymptotic estimate $\langle \Delta E \rangle = k_B T/(2N)$. The rate of cooling decreases with time as more and more particles become trapped by the standing wave, in a good agreement with Eq. (4), as the parameter r increases in the course of cooling. Moreover, the numerically calculated rate keeps the same tendency at the stages of deep cooling, where the weak-field estimates in Eq. (4) are no longer valid. The simulation was stopped when the rms of the phase jumps $\delta\varphi$ reduced to the level of about 1.5° (0.025 rad). By that time, almost 80% of the energy was extracted from the system. Figure 3 depicts the distribution of the particles in the (x, p) (dimensionless) phase space (a) at the beginning of the process, (b) after 3000 jumps, and (c) after 10^4 jumps. The last snapshot [Fig. 3(c)] was taken after adiabatic switch off of the standing wave following the end of cooling. Condensation of the particles in phase space is clearly seen as the cooling progresses. In addition, Fig. 3(b) depicts a sizable number of particles trapped in the minima of the standing wave (see visible vertical strips). Even after the adiabatic release of the trapped particles from the standing wave [Fig. 3(c)], some leftover structure may still be observed.

A simpler feedback scheme may be employed if the condition of a *guaranteed* energy loss on *each* step is replaced by a milder requirement of losing the energy *on average* in the course of many jumps. It requires only a *single* measurement defining the *sign* of ΔP (not its value), and uses a *fixed* phase shift $\pm\delta\varphi_0$ displacing the standing wave *opposite* to the direction of the instantaneous force. If small enough, such a correlated shift will *most probably* decrease the energy of the system. Using the arguments similar to the above, we find the energy loss after a single phase jump:

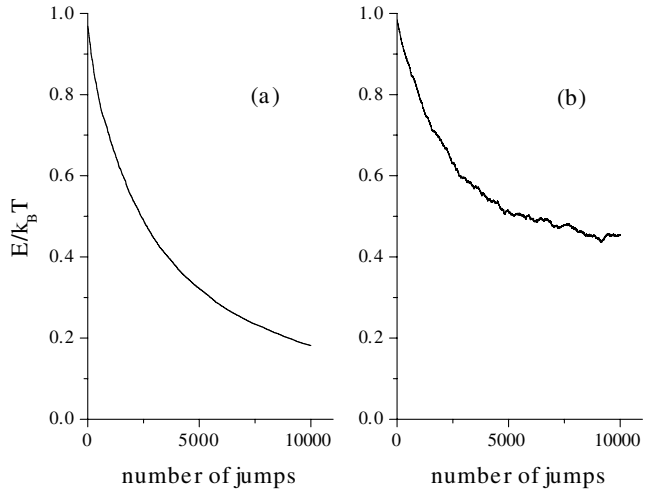

FIG. 2. Mean particle energy vs number of the controlled phase jumps for an ensemble of $N = 10^3$ particles confined in a harmonic trap. $V_0 = 0.1 k_B T$. In (a), the phase shift is defined by two measurements according to Eq. (4). Panel (b) represents results for the single-measurement feedback algorithm with $\delta\varphi_0 = \pi/10\,\mathrm{sgn}(\Delta P)$.

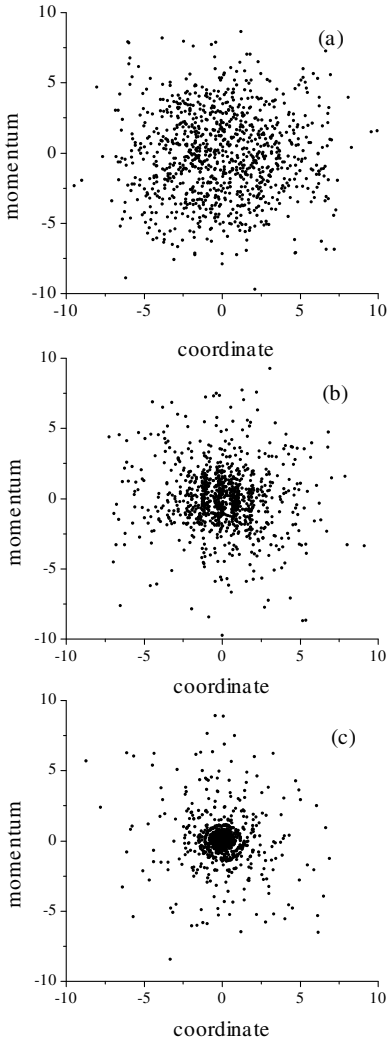

FIG. 3. Distribution of the particles in the (x, p) (dimensionless) phase space (a) at the beginning of the cooling process, (b) after 3000 jumps, and (c) after 10^4 jumps. $V_0 = 0.1 k_B T$. Particles are confined in a harmonic trap.

$$\Delta E = V_0\Big[\sin(\delta\varphi_0)\sqrt{\frac{N}{2}}|B| - [1-\cos(\delta\varphi_0)] \times \Big(rN + \sqrt{\frac{N}{2}}A\Big)\Big]. \quad (5)$$

The correlation between the force direction and $\delta\varphi$ is reflected in the modulus sign applied to the random variable B in Eq. (5). After averaging Eq. (5) over A and B, one can easily find the optimal $\delta\varphi_0$ leading to the fastest cooling on average: $\tan(\delta\varphi_0) = \sqrt{2/\pi}\rho^{-1}$. The maximal expectation value of the energy loss per jump is given by $\langle\Delta E\rangle = (V_0/2Nr)\rho(\sqrt{\rho^2 + 2/\pi} - \rho)$. For a thermal initial state with $\rho \geq 1$, this yields $\langle\Delta E\rangle \approx k_B T/(\pi N)$. Figure 2(b) depicts the results of feedback cooling of that kind applied to the same system of particles as discussed above. The preselected fixed value of the phase jumps $\delta\varphi_0 = \pm\pi/10$ is close to the optimal one at the initial temperature. Larger fluctuations are clearly observed [compare with Fig. 2(a)], but approximately 55% of the energy is extracted after 10^4 steps. Deeper cooling in this example requires adaptive lowering of $|\delta\varphi_0|$ as the process progresses. The single-measurement scheme is, probably, the most promising configuration for the first demonstration of the shaker cooling.

To estimate the parameters needed for successful operation of the optical shaker cooling, consider a standing wave that is produced by two counterpropagating Gaussian beams of power P and beam waist w_0. For the atoms in the waist region, we equate $\Delta P/c$ to the rms force and arrive at the following estimate for the relative signal: $\Delta P/P \sim 4\sqrt{2N}\alpha/\lambda\varepsilon_0 w_0^2$. For a quasiresonant transition in a two-level atom, $\alpha/4\pi\varepsilon_0 \sim (c^3/\omega_0^3)\Gamma/\Delta$ (Γ is the decay rate of the upper level). Therefore, for 10^6 trapped particles irradiated by the focused laser beams ($w_0/\lambda \sim 10$) acting on a far-detuned transition ($\Gamma/\Delta \sim 10^{-2}$), intensity stability of several percent is needed. Using a standing wave of period $\lambda/2 \sim 0.5$ μm and a feedback system with a response time of $\tau_r \sim 100$ ns allows cooling starting from the initial temperature of $T \leq 15$ mK $\sim 10^5 E_{\text{recoil}}$ (for particles having the mass of the Cs atoms). For laser power of $P \sim 2$ μW, the estimated cooling time is about several seconds, which is well within reported trapping times of many hundreds of seconds. Note also that the actual phase jump may be done on a nanosecond time scale that is much shorter than the time between jumps which in turn is determined by the need for efficient remixing.

In conclusion, we have presented optical shaking as a new approach to laser cooling by nonresonant fields. Its main advantage is the generic nature and applicability to a wide class of atoms and molecules, without specific requirements on the level structure. Moreover, no particles are lost from the trap while cooling, as is the case for evaporative cooling. It relies on a simple measurement procedure, and involves the feedback control only over the optical field, without any attempt to act selectively on different groups within a broad velocity distribution. The energy transfer from neutral particles to the laser fields is ensured by a self-acting mechanism resembling Sisyphus cooling. Classically, there is no lower limit to cooling by this method, other than practical considerations such as laser field intensity and phase fluctuations, measurement noise, feedback loop response time, etc. Quantum effects, which are not included in the present discussion, will certainly limit the shaker cooling near the recoil limit. As was pointed out above, the optical shaking is less efficient when a large fraction of the particles being cooled has already accumulated in the minima of the optical potential wells. This provides a hint that adaptive lowering of the potential during the shaking may give rise to more efficient cooling. The search for the best cooling rate, and the resulting optimization of the process, is the subject of our current research.

We acknowledge useful discussions with Y. Band, N. Davidson, and G. Raithel and appreciate help in numerical simulations from M. Finegold-Adler and M. Vilensky.

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