Atom slowing via dispersive optical interactions

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A promising technique of atom slowing is proposed. It is based upon the dispersive interaction of atoms with optical potential pulses generated by a far-off-resonance standing wave modulated in time. Each pulse reduces the velocity by a small amount. By repeating the process thousands of times, the velocity can be lowered from several hundreds of meters per second down to almost zero, over a path as short as 20 cm. In the absence of any random recoil process, the initial characteristics of the beam are preserved.

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A well-known method to slow down alkali metals or rare-gas metastable atoms— in most cases in view of trapping them—is to use a so-called *Zeeman slower* [1]. The Zeeman slower is based on the radiation pressure exerted on an atom beam by a counterpropagating light beam, the velocity-dependent frequency tuning being controlled thanks to a spatially inhomogeneous magnetic field. This slower works satisfactorily for what concerns the access to low velocities and trapping, but it has the disadvantage of producing poorly collimated beams of slow atoms (angular aperture of 0.1 rad, velocity dispersion of a few tens of percent) because of the spontaneous emission. In this article, we propose to use the dispersive interaction of atoms with comoving optical pulses generated via far-of-resonance time-modulated standing waves. In this approach, each optical pulse reduces the atomic velocity by a small amount. We show below that by repeating the process thousands of times, the velocity can be lowered from several hundreds of meters per second down to nearly zero, over a very short path of a few centimeters. Because of the absence of random recoil processes, the initial characteristics of the atomic beam as angular aperture, velocity dispersion, etc., should be preserved. One indeed shows that at the end of the slowing process the phase space density is maintained, contrary to what happens in free space propagation due to vacuum dispersion.

The approach proposed here bears slight similarities with the one known as "adiabatic slowing" [2,3]. Adiabatic slowing has been applied to a wide variety of species, such as hydrogen atoms, polar and nonpolar molecules [4], and Rydberg atoms and molecules [5]. In these methods, an external magnetic or electric field, periodic in space, is pulsed in time in such a way that the atom (or molecule) experiences a slowing gradient and nothing else. Low final velocities (a few tens of meters per second) are accessible, but at the price of rather strong fields (e.g., $B = 5.2$ T in Ref. [4]).

In the method proposed here, the nature of the force is quite different, since it derives from a special potential depending on both space and time, so-called "comoving" potential, of the form [6] $V(x,t) = S(t)\cos[2\pi x/\Lambda]$, where $S(t)$ is a limited-range ($[0,\tau_1]$) function of time, e.g., $S(t) \propto \exp(-t/\tau)$, and Λ is a spatial period. Experimental demonstration of this kind of force acting on atoms has been achieved on metastable hydrogen in a Stern-Gerlach interferometer [6]. As shown in a previous paper [7] for the case of comoving magnetic potential, for sufficiently large values of τ_1 ($\tau_1 \gg \tau$), this potential transforms the initial atomic wave function (for a given value of k) Ψ_0 into $\Psi = \exp[i\varphi_1(k,t)]\Psi_0$, where the phase shift is given by

$$\varphi_1(k,t) = -\hbar^{-1}\int_0^{(t)} dt'\, S(t')\cos\left(2\pi\frac{\hbar k}{m\Lambda}t'\right), \qquad (1)$$

where $\hbar k$ is the momentum along x, m is the atomic mass, and $(t) = \min\{t,\tau_1\}$. As a consequence, the motion of the center of a synchronized wave packet is altered: the group velocity becomes $v = \hbar k_0/m - [\partial_t\partial_k\varphi_1]_{k_0}$, k_0 being the center of the momentum distribution. This effect has been theoretically investigated at small velocities with metastable argon atoms Ar*($^3P_2$) to make negative values of v, creating *de facto* a negative-index "meta-medium" for matter waves [7]. Actually expression (1) of the phase shift is obtained using the WKB or short-wavelength approximation in which the spatial variation of the potential is assumed to be small at the de Broglie wavelength scale, i.e., $\Lambda \gg \lambda_{dB}$. This approximation, largely valid in all cases considered previously (λ_{dB} of a fraction of a nanometer, Λ of a few millimeters) becomes *a priori* questionable in the situation considered below. Obviously in a deceleration process leading in principle to the zero velocity by use of a short spatial period $\Lambda = \lambda_{opt}/2$, where $\lambda_{opt} = 811.5$ nm is an optical wavelength, there exists a lower limit v_{min} for the velocity v below which a more rigorous treatment would be needed. However, owing to the very low value of this limit ($v_{min} = 2.75$ cm/s to get $\Lambda = \lambda_{dB}$), such an effort appears to be useless.

The above expression of the phase shift [Eq. (1)] implicitly assumes that the total energy E of state Ψ takes a final value (at $t = \tau_1^+$) equal to the initial one in Ψ_0 (at $t = 0^-$). In other words the time-dependent factor $\exp(-iEt/\hbar)$ is common to wave functions Ψ_0 and Ψ, and it does not appear in the phase shift φ_1. Actually this assumption is valid for a potential pulse tending to zero "adiabatically," i.e., without any characteristic time constant smaller or comparable to τ. Indeed in such a case the only discontinuity in the wave function occurs at $t = 0$ and the time-dependent Schrödinger equation containing V holds from $t = 0^+$ to t infinite, leading to expression (1) for any positive value of t. Since $\lim_{t\to\infty}\varphi_1$ is a constant and

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$\lim_{t\to\infty}[\partial_t\partial_k\varphi_1]_{k_0} = 0$, one obtains asymptotically a "final" group velocity equal to the incoming one, namely, $\hbar k_0/m$. Since the potential is zero for $t < 0$ and t infinite, the final total energy is equal to the initial one. This does not mean that E remains constant all along the pulse duration. It does not indeed, as it can be easily verified and as it is expected with a time-dependent potential. The time dependence of the energy is entirely contained in $\varphi_1(k,t)$.

When the potential is abruptly interrupted at $t = \tau_1$, three time intervals must be considered, in a similar way as separated spatial intervals are considered in the case of a square potential: (i) $t < 0$ (free wave), (ii) $0 < t < \tau_1$ (V is present), and (iii) $t > \tau_1$ (free wave). In the third interval, the phase shift cannot be a constant, because the final energy $E(\tau_1)$ is different from E, the initial energy. Then for $t > \tau_1$, the phase shift must be a linear function of time: $\varphi(k,t) = \varphi_1(k,\tau_1) + \hbar^{-1}[E - E(\tau_1)](t - \tau_1)$. This form is readily derived from the continuity of the phase shift at $t = \tau_1$. One obtains instead of Eq. (1), a more general expression, valid at any time t:

$$\varphi(k,t) = \varphi_1 - \frac{S(\tau_1^-)}{\hbar}\Theta(t-\tau_1)(t-\tau_1)\cos\left(2\pi\frac{\hbar k\tau_1}{m\Lambda}\right), \quad (2)$$

where τ_1^- is a value of t smaller than, and arbitrarily close to, τ_1 and Θ is the Heaviside function: $\Theta(x) = 0$ for $x < 0$ and $\Theta(x) = 1$ for $x \geqslant 0$. Both spatial shift $-\partial_k\varphi|_{k_0}$ and velocity change $-\partial_t\partial_k\varphi|_{k_0}$ of the wave-packet center, derived from expression (2), are continuous at $t = \tau_1$. Note that, because of its limited range, $S(t) \propto \Theta(\tau_1 - t)\exp(-t/\tau)$. This cancels any singularity coming from $\Theta(t - \tau_1)$ in Eq. (2). This is illustrated in Fig. 1 where trajectories of the wave-packet center are shown for various values of τ_1. The conditions chosen for this simulation are typically those leading to negative refraction, namely, a magnetic field of 500 G, a velocity of 2 m/s, and a spatial period $\Lambda = 5$ mm and $\tau = 0.37$ ms. Indeed, when successive pulses are considered, it is the continuity of the velocity at the end of each pulse that will allow us to use it as the initial velocity for the next pulse (separated from the preceding one by a "blank" of duration much smaller than τ_1). Therefore the continuity condition is the key element of the comoving-potential slower operation: each pulse reduces the velocity by a small but significant amount. Repeating the process thousands of times, the velocity can be lowered from several hundreds of meters per second down to almost zero. In a simulation of a magnetic version of this slower ($\Lambda = 5$ mm, magnetic field of 800 G), the total path needed to reach the "zero" velocity is rather long (2.2 m), however, comparable to that of a standard Zeeman slower.

Other classes of methods of slowing down atoms or molecules make use of optical forces. In particular, white light [8] and bichromatic standing waves are able to generate stimulated forces largely exceeding the spontaneous contributions [9,10]. Owing to these forces, strong decelerations have been performed over amazingly short distances (a few centimeters). Deceleration of molecules can be also achieved by means of the (second order) Stark effect induced, e.g., on benzene, by an intense (1.6×10^{12} W/cm^2) far-of-resonance pulsed laser beam [11] or by the standard Stark effect induced on a molecule, such as MgO, possessing a permanent dipole

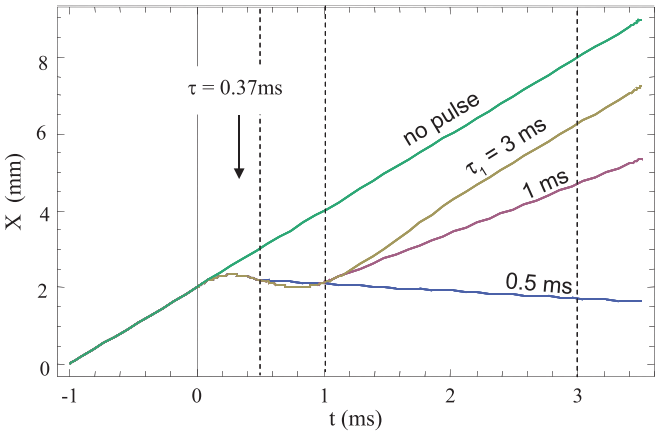

FIG. 1. (Color online) Calculated trajectories of the wave-packet center when comoving potential pulses of different durations τ_1 are present. Parameters used in this simulation are such that a negative refraction is realized (see text): magnetic field 500 G, initial velocity, 2 m/s, spatial period, $\Lambda = 5$ mm and $\tau = 0.37$ ms. At time τ_1, both position and velocity are continuous.

by a pulsed electric field inside a microwave cavity (which can be used as well as a trap) [12].

For our purpose, the use of comoving optical potentials in place of conventional magnetic or electric potentials provides us with a spatial period reduced by a huge factor ($\approx 10^3$). In principle such a potential can be generated by a linearly polarized, far-off-resonance standing light wave modulated in time and generated by an interferometer (see Fig. 2). For a Rabi frequency, Ω, and a sufficiently large detuning, $\delta\omega$, the optical potential takes the simple form [13–15]

$$V_{\rm opt}(x,t) = \frac{1}{2}S(t)\left[1 + \cos\left(4\pi\frac{x}{\lambda_{\rm opt}}\right)\right], \quad (3)$$

with, e.g., $S(t) \approx \frac{\hbar\Omega^2}{\delta\omega}\exp(-t/\tau)$ within a definite time interval $[0,\tau_1]$ and $S(t) = 0$ elsewhere. The results obtained using either the spatial dependence in $\cos^2(2\pi x/\lambda_{\rm opt})$ or in $\cos(4\pi x/\lambda_{\rm opt})$ are identical to an accuracy better than 10^{-6},

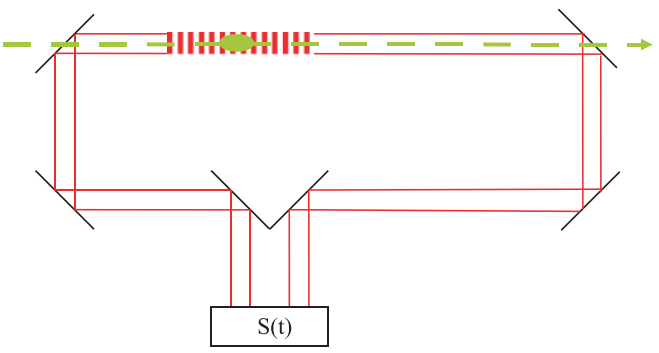

FIG. 2. (Color online) Scheme of the comoving optical potential slower. A far-off-resonance, linearly polarized standing wave ($\lambda =$ 811.5 nm) is produced by a separated-arm interferometer. Light intensity is modulated in time by a signal $S(t)$ (see text). The atom beam is introduced with a small inclination through holes drilled in the mirrors.

which means that the x-independent part of the potential has no significant effect on the evolution of the wave function. For far off resonance and linearly polarized light, the above expression of $V_{\rm opt}$ is common to all magnetic sublevels because these sublevels are shifted by the same amount [16]. As a consequence the dipolar potential is almost independent of the atomic spin orientation. It is important to note that an atom passing through this potential at some velocity v does not "feel" a unique frequency, but rather two frequencies related to the two components of the standing wave, Doppler-shifted by $\pm\Delta\omega(v) = \pm k_{\rm opt} v$. At the largest velocity considered here, 560 m/s, $\Delta\omega(560) = 2\pi\ 6.90 \times 10^8$ rad/s (690 MHz). In order to operate "far from resonance," to avoid spontaneous emission, we need that the difference between the laser frequency and the frequencies of these two Doppler-shifted resonances be large compared to $(1+s)^{1/2}\gamma = \gamma'(s)$, the power-broadened line width, γ being the natural line width and s the saturation parameter. Actually, we want to get a magnitude $S(0)$ of the potential, sufficiently high to achieve the complete slowing over a distance shorter than, e.g., 20 cm. As we wish also to use a reasonable laser power, we shall take a moderately large (negative) detuning, $\delta\omega = \omega_L - \omega_0 = -5\Delta\omega(560) = 2\pi\ 3.45 \times 10^9$ rad/s (3.45 GHz). Such a large detuning can be obtained either by means of a modulator or (for polarized atoms, as metastable argon atoms Ar* $^3P_2$ in the Zeeman state $M = +2$) with a static magnetic field of 818.1 G. For the simulation, one assumes a laser intensity of 32 mW/mm^2. Then the saturation parameter $s = I/I_{\rm sat}$, with $I_{\rm sat} = 14$ W/m^2, is equal to 2 267.3. The natural width of the transition being $\gamma = 2\pi\ 5.8 \times 10^6$ rad/s, the power-broadened width is $\gamma' = 2\pi\ 2.762 \times 10^8$ rad/s, which leads to the ratio $R = \delta\omega/\gamma' = 12.49$, large compared to 1. As the velocity v is lowered, $\Delta\omega$ decreases, tending to zero at $v = 0$ m/s. Then either the detuning is kept constant and the condition $R \gg 1$ is better and better verified or $\delta\omega$ is kept equal to $5\Delta\omega(v)$ allowing us to reduce the intensity (as v) as well as γ' (as $v^{1/2}$), but then the ratio R decreases as $v^{1/2}$, which implies a lower limit for v ($R = 1$ at $v = 3.59$ m/s).

A series of many pulses separated from each other by a small amount of time (of the order of $\langle\tau_1\rangle/100$, $\langle\tau_1\rangle$ being an average over the different pulses) is applied, each of them (numbered n) providing a small decrease $\delta v(n)$ of the velocity (a few mm/s). The duration $\tau_1(n)$ of each pulse is adjusted in such a way that the first maximum value of $|\delta v|$ is reached at the end of the pulse. In the following example, $\lambda_{\rm opt} = 811.5$ nm and $\tau = 3.7$ ns, this duration (in nanoseconds) obeys the approximate law $\tau_1(n) \approx 0.23 + 6.40\exp[-v(n)/40] + 0.60\exp[-v(n)/200]$, where $v(n)$ in meters per second is the initial velocity for the nth pulse and the path $v\tau_1$ covered by the atom during successive pulses is roughly a constant ($\approx$0.12 μm). It should be noted that, despite the pulsed character of light, the building up of the standing wave presents no difficulty: indeed at the shortest value of τ_1 (0.234 ns) for $v = 560$ m/s, the length of the interfering trains of light (Fig. 2) is 12 cm and this length increases at lower values of v. The velocity decrease per pulse $|\delta v|$ increases quasiexponentially from 0.2 mm/s at $v = 560$ m/s to 2.9 mm/s at $v \approx 0$. To get an almost complete stopping, a large number of pulses, namely, $N = 1\,987\,500$,

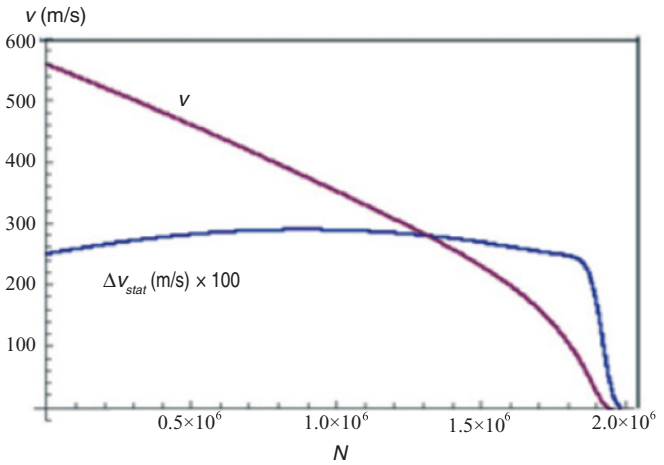

FIG. 3. (Color online) Evolution of the velocity v and the statistical velocity width $\Delta v_{\rm stat}$ (see text) multiplied by 100 as a function of the total number N of comoving optical potential pulses. The total length needed to reach the quasizero velocity (for $N = 2 \times 10^6$) is 19.3 cm. The laser intensity is 32 mW/mm^2 and the detuning is 3.45 GHz.

is needed (see Fig. 3). Nevertheless this large number does not mean a very long distance since the velocity is lowered from 560 m/s down to "zero" (within the limits defined previously) over a distance of about 19.3 cm. This is much shorter than the distance needed with the Zeeman slower or adiabatic slowing and is comparable to that used in bichromatic deceleration. It could be reduced by a factor of about 2 by using doubled laser intensity (60 mW/mm^2). The total time employed by an atom to be stopped, i.e., the total duration of the pulse sequence, is 6.07 ms. Then the complete process can be repeated at a rate of about 164 times per second.

An important characteristic of the decelerator is its effect on the atomic wave-packet spatial width. This width, $\delta x(t)$, has been calculated as a function of time using the following expression of the wave function:

$$\Phi(x,t) = \int_{-\infty}^{+\infty} dk\ \sqrt{\rho(k)}\, e^{i[kx - \frac{\hbar k^2}{2m}t + \varphi(x,t)]}, \qquad (4)$$

where $\rho(k)$ is a Gaussian distribution centered at $k_0 = 3.14 \times 10^{11}$ m^{-1} (group velocity of 560 m/s), the standard deviation of which is $\delta k = 0.002 k_0$. Under such conditions the initial width is $\delta x(0) = 0.635$ nm. In this calculation we have used the concept of "effective" time developed in a previous paper [17]. In this paper it has been shown that the comoving potentials are able to transiently narrow the wave packet, compensating for the free-propagation natural spreading. In the present case this effect is very small at the beginning of the deceleration process but becomes more and more important as the velocity decreases. The result is shown in Fig. 4 (upper part). It is seen that δx progressively deviates from the free-propagation width $\delta x_0(t)$ to rejoin the initial

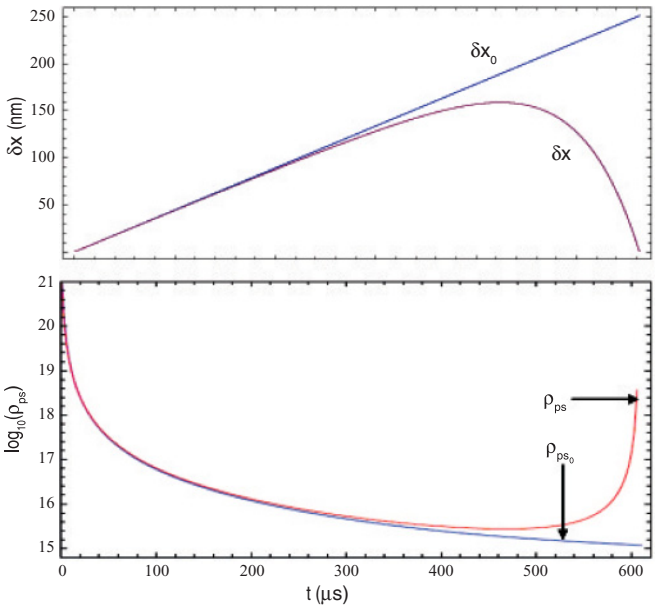

FIG. 4. (Color online) (Upper part) Evolution of the spatial width δx of the wave packet as a function of time. Upper curve: Free evolution (no potential); lower curve: the potential is present. The initial width $[\delta x_0(0) = 0.635$ nm] corresponds to a relative dispersion $\delta k/k = 0.2$ % in the momentum space. (Lower part) Evolution (in $\log_{10}$ scale) as a function of time of densities in phase space (k,x), defined as $\rho_{\mathrm{ps}_0} = (\delta k \delta x_0)^{-1}$ for free propagation and $\rho_{\mathrm{ps}} = (\delta k \delta x)^{-1}$ when the potential is present. Note that ρ_{ps_0} corresponds as well to the separable free coordinates y and z.

width at the end of the process. From Eq. (4), one readily derives the Wigner function:

$$W(x,K,t) = \frac{1}{2\pi}\int_{-\infty}^{+\infty} d\xi\, \Phi^*\left(x-\frac{\xi}{2},t\right) e^{-iK\xi}\, \Phi\left(x+\frac{\xi}{2},t\right),$$

$$W(x,K,t) = \int_{-\infty}^{+\infty} dk \sqrt{\rho(k)\rho(2K-k)} \times e^{i[2(k-K)x - \frac{2\hbar K(k-K)t}{m} + \varphi(k,t) - \varphi(2K-k,t)]}. \quad (5)$$

Then the momentum distribution is

$$w_k(k,t) = \int_{-\infty}^{+\infty} dx\, W(x,K,t) \equiv \rho(k). \quad (6)$$

It is time invariant because the potential affects the wave function *via* a pure (real) phase shift. Finally the width in k remains equal to δk_0 whereas the spatial width δx evolves as described above. Consequently the density in phase space (k,x) defined by $\rho_{\mathrm{ps}} = (\delta k \delta x)^{-1}$ first decreases then increases up to its initial value, whereas this density $\rho_{\mathrm{ps}0}$ for free propagation monotonously decreases (see Fig. 4, lower part). Note that $\rho_{\mathrm{ps}0}$ directly gives the phase-space density related to y and z coordinates, which are separable from x. We also examined this question from a statistical viewpoint, considering an ensemble of initial velocities around $v_0 = 560$ m/s [with relative width $\Delta v_{\mathrm{stat}}(0)/v_0 = \pm 0.2\%$] and positions (with width $\Delta x_{\mathrm{stat}} = \pm 0.7$ nm). As seen in Fig. 3, $\Delta v_{\mathrm{stat}}(t)$ is almost a constant, except at the very end of the process where it drops down to zero, contrarily to a strong increase observed at higher velocity, typically $v_0 > 565$ m/s, a value that can be considered as an upper bound of velocities that could be decelerated. Δx_{stat} first increases and then decreases leading to a phase-space density rather similar to ρ_{ps}.

In this article, we have proposed an approach to atom beam deceleration based on dispersive optical forces. Atom stopping should be almost achieved on short distances. The absence of spontaneous emission processes should allow preservation of the transverse coherence properties of the initial beam. It is especially applicable to narrow supersonic beams, like metastable rare-gas atom beams [18], and provides ultra-low-velocity beams for coherent atom optics and atomic interferometry. It is also a promising technique applicable to slowing down molecules since any optical pumping toward molecular levels (other than those interacting with light) is absent.

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