Mechanical effect of retroreflected frequency-chirped laser pulses on two-level atoms

G. Demeter,* G. P. Djotyan, Zs. Sörlei, and J. S. Bakos

Institute for Particle and Nuclear Physics of the Hungarian Academy of Sciences, Konkoly-Thege Miklós út 29-33,
H-1121 Budapest, Hungary

We discuss the mechanical momentum transfer to two-level atoms by a pair of short, counterpropagating, frequency-chirped laser pulses, which partially overlap each other in the atomic medium. We show that such a pulse pair can induce a much greater change of momentum than pulses that do not overlap (interact separately with the atoms). The dispersive effect on the atomic velocity distribution is shown to be much smaller for the case of overlapping pulses. Furthermore, there are regimes where the method is not sensitive to the exact values of the pulse parameters or the initial velocity distribution. The interaction can be repeated for a cumulative effect, so a sequence of such pulse pairs can be used very effectively to manipulate the mechanical motion of atoms.

I. INTRODUCTION

Using laser radiation to control and manipulate the translational state of neutral atoms is now an immensely wide field, and various schemes of laser cooling and trapping are of everyday use in the field of atomic physics. Among these schemes, those using light forces due only to stimulated processes have attracted attention for two major reasons. On the one hand, they allow one to overcome the saturation limit of spontaneous light forces; on the other, they are capable of preserving the coherence of the atomic wave packet. In particular, there have been a number of suggestions to utilize pulsed laser radiation for the manipulation of atoms—e.g., to deflect or focus atomic beams. These schemes often use laser pulses propagating in opposite directions to excite the atoms in a controlled way first and then to return them to the ground state. A momentum change of $2\hbar k$ accompanies a cycle of excitation and deexcitation, and the cycle can be repeated at a high rate, so the overall force on the atoms can greatly exceed the light pressure force mediated by excitation spontaneous-emission cycles.

Perhaps the first of these schemes was the suggestion to use a sequence of counterpropagating, frequency-chirped laser pulses that induce population inversion of the atoms in the adiabatic regime [adiabatic rapid passage (ARP)] for isotope separation in [1]. This idea was later developed further [2–4]. Another scheme has been proposed [5] and realized [6] that uses a sequence of picosecond π-pulse pairs to confine atoms or to focus atomic beams. Here two simultaneous laser pulses from opposite directions are used to push atoms towards the trap center. Atoms displaced from the central region (where the pulses overlap) will interact with the two pulses separately—they will be excited by the pulse propagating toward the central region first and deexcited by the counterpropagating pulse, thus receiving $2\hbar k$ momentum towards the center. A similar scheme used the intense standing wave that develops in the overlap regime of the pulses to focus atomic beams on the nanometer-length scale for lithographic purposes [7].

Several schemes also exist, which use cw radiation, but have a similar effect to counterpropagating laser pulses. One is the intensive force acting on atoms in a bichromatic standing wave, which can be described in terms of a sequence of π pulses from alternating directions [8–11]. Another technique is the coherent acceleration of Bose-Einstein condensates using a frequency-chirped optical lattice, which effectively induces multiple adiabatic rapid passage transitions with a momentum transfer of $2\hbar k$ in each [12,13]. A similar arrangement of chirped standing waves has also been shown to be useful for the construction of atomic mirrors and beam splitters [14]. Yet another method uses frequency-modulated standing waves, whose effect is similar to a sequence of frequency-chirped pulses, to achieve the deflection of an atomic beam [15]. The primary advantage of this latter scheme (and of the original chirped pulse technique) is its great velocity capture range, which allows its use even when there is considerable Doppler broadening of the atomic transition line, and that the scheme is, to a good measure, insensitive to variations in the parameters of the laser light.

Some interesting methods have been proposed lately which use delayed laser pulses for the manipulation of two-level atoms. The special feature of these is that they involve two laser fields that act on the same atomic transition, one of which suffers a distinct time delay with respect to the other, but still overlapping long enough to create an ambiguous situation as to which of the two fields excites the atom. (By comparison, the frequently used stimulated Raman adiabatic passage (STIRAP) scheme involves time-delayed fields that act on two separate transitions of an atom with a Λ level scheme.) One of these schemes [16] uses two counterpropagating pulses of equal amplitude, whose detunings with respect to the atomic transition are equal in magnitude, but opposite in sign. It was shown that if the pulses are intensive enough and the detunings are also appropriate, an integer times $2\hbar k$ momentum can be transferred to the atoms with a single pair of overlapping pulses. Furthermore, the interaction involves an adiabatic transfer of the atomic population and thus is robust with respect to a slight variation of the pulse parameters. Another scheme, dubbed retroreflection-induced bichromatic adiabatic passage (RIBAP), was also proposed and demonstrated [17,18] that uses two oppositely

---

*Electronic address: gdemeter@rmki.kfki.hu

FIG. 1. Timing of the pulses. The counterpropagating pulse arrives with a time delay $T=t_2-t_1$ that is comparable to the pulse length τ. The cycle may be repeated after $T_c=t_3-t_1>T,\tau$.

detuned pulses, the second of which is substantially smaller than the first one. This arrangement was shown to be useful for achieving complete transfer of the population to the excited state of the atom in a robust way.

Similar schemes have been proposed also for the manipulation of three-level Λ atoms. In [19] it was noted that robust population transfer is possible between the two metastable states of a Λ atom, even when the two fields are sufficiently strong to couple both bottom states to the excited state. In this regime, a coupling ambiguity is created and the conditions for STIRAP are violated. In [20], counterpropagating, bichromatic laser pulses with appropriate detunings and delay were proposed to drive transitions between the stable states to achieve a momentum transfer of $4n\hbar k$. Multiphoton adiabatic transitions have also been shown to be very useful for population transfer between the internal states of Rydberg atoms [21] and diatomic molecules [22].

In the present paper, we explore the possibilities of using a sequence of retroreflected, chirped laser pulses for the manipulation of the translational state of two-level atoms. A single retroreflected pulse creates a pair of counterpropagating, overlapping laser pulses, the second of which arrives with a distinct time delay T (see Fig. 1). We explore the interaction of two-level atoms with these pulse pairs. We demonstrate that the motion of the atoms can be influenced in various ways and that the mechanical effect of such pulse pairs can be much greater than that of separated pulses. We analyze the influence of spontaneous emission on the proposed manipulation scheme and show that the adverse effects are much less severe than for the case of separated pulses.

II. THEORETICAL FRAMEWORK

To describe the motion of two-level atoms under the influence of quasimonochromatic electric fields that are near-resonant with the atomic transition, we use the relevant Schrödinger equation for the wave function:

$$i\hbar\frac{\partial\Psi}{\partial t}=\left(\frac{\hat{p}^2}{2m}+\hat{H}_a-\hat{d}E\right)\Psi. \qquad (1)$$

The Hamiltonian in Eq. (1) consists of the kinetic energy term due to atomic motion ($\hat{p}$ is the center-of-mass momentum of the atom and m is its mass), the internal energy of the atom $\hat{H}_a=\hbar\omega_g|g\rangle\langle g|+\hbar\omega_e|e\rangle\langle e|$, and the dipole interaction energy with the electric field of the light, which is treated classically. Spontaneous transitions are not included in this description, so such a theory will be valid only if the laser pulses are short and the time the atoms spend in their excited state is much less than the spontaneous lifetime $\tau_{spont}=1/\Gamma$. One can, however, study the effects of spontaneous emission

by using Eq. (1) in a Monte Carlo wave function simulation [23,24].

The state of the atom is described by the wave function Ψ, which includes both the internal state of the atom and center-of-mass motion. In momentum space, we may write it as

$$\Psi=a(\kappa,t)e^{-i\omega_g t}|g,\kappa\rangle+b(\kappa,t)e^{-i\omega_e t}|e,\kappa\rangle, \qquad (2)$$

where $\hbar\kappa=p$ is the atomic momentum. The classical electric field is taken to be

$$\begin{aligned}E(x,t)=&E^+(t)\cos[kx-\omega_{ge}t-\phi^+(t)]\\&+E^-(t)\cos[-kx-\omega_{ge}t-\phi^-(t)],\end{aligned} \qquad (3)$$

where $E^\pm(t)$ and $\phi^\pm(t)$ are the slowly varying (real) amplitudes and phases of the two counterpropagating components. The sign that parametrizes the amplitude and phase shows the direction of propagation of the pulse along the x axis. Note that the optical frequency $\omega_{ge}=\omega_e-\omega_g$ of the atomic transition has been used to describe the fast temporal oscillation of the electric field; any constant detuning and frequency modulation is contained in the time-dependent phase $\phi^\pm(t)$. With the notation $\hbar\Omega^\pm(t)=d_{ge}E^\pm(t)$ for the Rabi frequencies, the equations for the probability amplitudes become

$$\begin{aligned}i\partial_t a(\kappa,t)=&\frac{\hbar\kappa^2}{2m}a(\kappa,t)-\frac{\Omega^+(t)}{2}e^{i\phi^+(t)}b(\kappa+k,t)\\&-\frac{\Omega^-(t)}{2}e^{i\phi^-(t)}b(\kappa-k,t),\end{aligned}$$

$$\begin{aligned}i\partial_t b(\kappa,t)=&\frac{\hbar\kappa^2}{2m}b(\kappa,t)-\frac{\Omega^+(t)}{2}e^{-i\phi^+(t)}a(\kappa-k,t)\\&-\frac{\Omega^-(t)}{2}e^{-i\phi^-(t)}a(\kappa+k,t).\end{aligned} \qquad (4)$$

Here d_{eg} is the dipole matrix element of the transition and is taken to be real, so $\Omega^\pm(t)$ are also real. We consider Gaussian pulses (with amplitude Ω) and a linear frequency chirp (with a chirp speed β), whose central frequency is exactly resonant with the atomic transition (i.e., there is no constant detuning). The second pulse arrives with a delay $T\sim\tau$, and the cycle may possibly be repeated after $T_c>T,\tau$ (see Fig. 1):

$$\Omega^+(t)=\Omega\sqrt{\frac{\pi}{2}}\exp\left[-\frac{t^2}{2\tau^2}\right],$$

$$\Omega^-(t)=\Omega\sqrt{\frac{\pi}{2}}\exp\left[-\frac{(t-T)^2}{2\tau^2}\right]=\Omega^+(t-T), \qquad (5)$$

$$\phi^+(t)=\frac{\beta}{2}t^2+\phi_0^+,$$

$$\phi^-(t) = \frac{\beta}{2}(t-T)^2 + \phi_0^-. \quad (6)$$

With these notations, the constant-amplitude parameter Ω gives the pulse area A through $\Omega\tau = A/\pi$; i.e., $\Omega\tau = 1$ is exactly a π pulse.

Analytical solutions of Eqs. (4) can be obtained in only a few special cases, but it is immediately evident that when the pulses do not overlap, each probability amplitude in momentum space is coupled only to one other, so the problem is basically the same as a simple two-level system without center-of-mass motion. The simplest case is when we are dealing with resonant π pulses, or chirped pulses, that induce complete population inversion between the ground and excited states. Provided that the spectrum of the pulse is much wider than the Doppler width of the initial atomic momentum distribution, a single pulse simply exchanges the populations of the ground and excited states, while shifting the momentum distribution functions corresponding to the two electronic states by $\hbar k$ in opposite directions. A second, identical pulse, counterpropagating the first, will restore the initial populations for the internal atomic states, while the momentum distributions will end up shifted by $2\hbar k$. This process is the basic idea between numerous manipulation schemes [1–5,15]. For sizable momentum transfer to the atoms, the process has to be repeated many times. This, however, means that imperfect population inversion, or spontaneous emission during the manipulation process, will lead to a quick dispersion of the initial momentum distribution—i.e., to considerable heating of the atoms.

When the pulses overlap for an appreciable amount of time, an infinite number of probability amplitudes are coupled in Eqs. (4) and the problem is much more complex. When the delay is small compared to the pulse length, we may expect the atomic motion to be similar to that in standing waves. As the delay increases, the motion is expected to deviate from this behavior and become more and more like the separated pulse case discussed in the previous paragraph. The details will depend on the pulse parameters and the delay, and it is this transition regime that is the subject of our investigations.

III. RETROREFLECTED, CHIRPED PULSES

A. Single pulse pair without relaxation

We have performed a numerical study of Eqs. (4) with various pulse parameters. The first step was to calculate the effect of a single pulse pair without taking into account relaxation by spontaneous emission. The most important features of the interaction were then extracted from the final wave function. The foremost of these are the average momentum transferred to the atoms, the spreading of the wave function in momentum space (the heating effect of the pulses), and the final population of the levels. For the initial condition, the atoms were considered to be in the ground state and we chose an infinitely sharp momentum distribution around zero—i.e., $|\psi_0\rangle = |g, \kappa = 0\rangle$. This initial condition is convenient for several reasons. On the one hand, since this state is infinitely extended in real space, the result will be
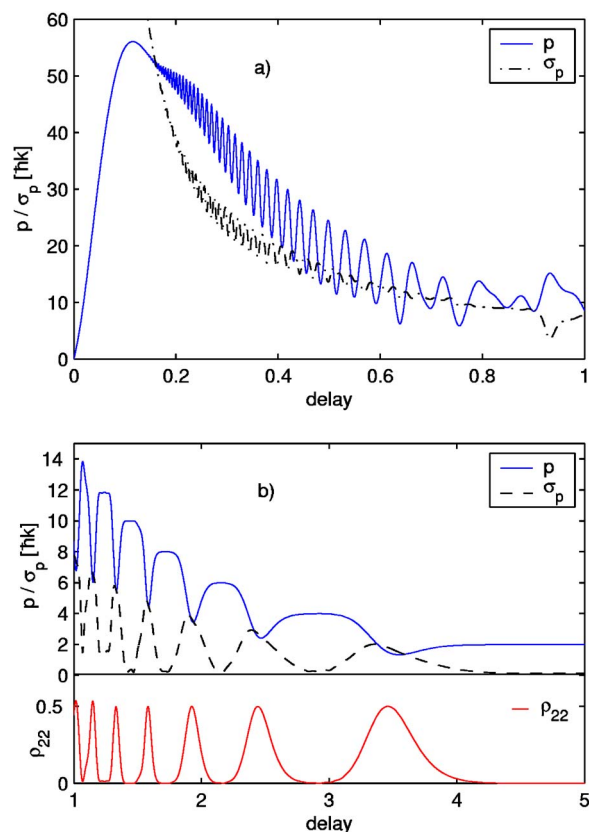

FIG. 2. (Color online) (a) Final momentum expectation value $\langle p \rangle$ and momentum spread σ_p of an atom initially in the state $|g,0\rangle$, in units of $\hbar k$, as a function of the delay $T/\tau \in [0,1]$. (b) The same curves (top) and the final value of the excitation ρ_{22} (bottom) for $T/\tau \in [1,5]$. The parameters used for the calculation were $\Omega = 100/\tau$, $\beta = 20/\tau^2$, and $\tau = 10^{-4}$ in recoil units.

insensitive to the relative phase constant $\phi_0^+ - \phi_0^-$ of the two waves and we can set both phases to 0. (Indeed, if we took a minimal uncertainty wave packet that is extended over many $\hbar k$ in momentum space as an initial condition, this would describe an atom that is localized in real space in a region $\Delta x \ll \lambda$, and thus to obtain results that are experimentally relevant, we would have to average over that phase.) Second, from the solution the results for an experimentally achievable initial momentum distribution can be calculated by simple averaging. Furthermore, since a transition between electronic states is always accompanied by a momentum change of $\hbar k$, the final momentum distribution function will also unambiguously determine the electronic state of the atom: values that correspond to an even multiple of $\hbar k$ will describe atoms that are in the ground state, and those corresponding to an odd multiple of $\hbar k$ will describe atoms that are in the excited state—in short, only the $|g, 2n\hbar k\rangle$ and $|e, (2n+1)\hbar k\rangle$ states will be populated.

Figure 2 shows the result of the interaction of a two-level atom with a single pair of chirped pulses, as a function of the time delay between them, in two distinct regimes. The curves depict the momentum transferred to the atoms in units of $\hbar k$ by a single pulse pair and the spread of the atomic wave packet in momentum space after the interaction. Additionally, Fig. 2(b) also shows the excited-state probability after

the interaction. The pulse parameters ($\Omega = 100/\tau, \beta = 20/\tau^2$) are sufficient for a complete adiabatic transfer of the atomic populations when the pulses are separated. Clearly, for no delay, the atoms do not gain any average momentum, but the spread of the wave packet becomes very large after the interaction. (The velocity spread for 0 delay is almost $200\hbar k$—not shown in the figure.) This is basically the scattering of the wave packet by an intense standing wave. At the other extreme, for large delay, where the pulses interact independently with the atoms, the momentum transfer is $2\hbar k$, the wave packet is not dispersed by the interaction, and the atoms are returned to the ground state after the second pulse leaves—this is the scenario described in [1–5,15]. In between the two extremes, the momentum transfer first increases linearly as noted in [6], then oscillates strongly with the time delay, and finally "plateaus" appear, where the momentum transfer is an integer multiple of $2\hbar k$, the velocity spread becomes negligible, and the atom is returned to the ground state after the interaction. The effect of the two overlapping pulses is exactly the same in these regions, as that of nonoverlapping pulses, but the momentum transfer is increased several times. For our pulse parameters, the plateau that corresponds to $\Delta p = 10\hbar k$ momentum transfer per pulse pair is still visible. If the pulse amplitude and chirp is changed, the curves change their position, but the overall picture remains the same. Changing either of these parameters just slightly, the peaks, valleys, and plateaus shift to slightly different delay values. If the amplitude and chirp are both increased considerably, more plateaus appear that correspond to $12\hbar k$, $14\hbar k$, etc. For each amplitude, one can determine an optimal range of the chirp where the most plateaus appear. This is a fairly large interval, but for values outside it, the conditions of momentum transfer deteriorate. On the other hand, if all parameters (amplitude, chirp, and delay) are chosen so that we are near the middle of a plateau, changing any (or all) of the parameters slightly has no effect on the momentum transfer; i.e., the process can be somewhat robust with respect to parameter changes.

The process is analogous to the scheme described in [16], where a pair of delayed pulses with opposite detuning from the atomic transition frequency were used to the same effect—i.e., to induce a momentum transfer of $2n\hbar k$ without heating the atoms. The appearance of plateaus where robust multiphoton transitions can be observed has been explained in terms of dynamical resonances that arise in connection with the topology of the dressed-state energy surfaces [25,26]. Our chirped pulse scheme has two major advantages compared to that one. First, chirped laser pulses can achieve the same momentum transfer with a pulse amplitude that is more than one order of magnitude smaller than that needed for the constant detuning scheme. (This means an over two orders of magnitude reduction in the peak intensity of the pulses.) Second, the two pulses are identical in the current scheme, so the first pulse can be simply retroreflected to obtain the second one. In contrast, pulses with opposite detunings require either a synchronization of two laser pulses with different frequencies, or tilting the laser beam with respect to the atomic beam, and, in addition, a sufficient reduction of the longitudinal velocity spread of the atoms. It is interesting to note that in the scheme described in [16], the

atoms are pushed in a "counterintuitive" direction—i.e., in the direction of propagation of the second pulse. In our case, this is not so—the atoms are pushed in the direction of the first pulse. It is also notable that just as in the case of separated chirped pulses, the outcome of the interaction is insensitive to the sign of the chirp—only the magnitude is important.

The second scheme worth mentioning, which bears a resemblance to ours, is the coherent acceleration of Bose-Einstein condensates by chirped standing-wave fields [12,13]. This method, however, has a very limited velocity capture range (less than a single-photon momentum), so it is only suitable for the manipulation of condensates or equally cold atoms. Our scheme, on the other hand, has a wide velocity capture range and is capable of uniformly accelerating an atomic ensemble where the transition line has a considerable inhomogeneous broadening. We note again that taking a δ function for the initial momentum distribution for our calculations is purely for convenience. With the pulse parameters used in our simulation, the momentum range in which the atoms receive $10\hbar k$ momentum for a delay of $T/\tau = 1.42$ is about $[-500\hbar k, 500\hbar k]$. Furthermore, the method of [12,13] requires a fairly long interaction time ($\sim$ms), while the method under discussion needs much less time to achieve the same amount of momentum transfer (from $\sim$100 ns to 1 μs). This property is important primarily if the atoms are not cold and trapped, but flying through the interaction region as part of an atomic beam.

The population transfer induced by the chirped laser pulses is also similar to the RIBAP scheme described in [17,18] where retroreflected laser pulses were used to obtain robust population transfer of the atoms, despite coupling ambiguities occurring in the simultaneous action of the two laser fields. However, in this case the second pulse was smaller in magnitude than the first one and the atoms were left in the excited state after the interaction.

Further interesting properties of the interaction are revealed by the final momentum distribution of the atoms. This is shown in Figs. 3–6, as a function of the delay. The strong white stripes at even multiples of the photon momentum in Fig. 3 correspond to the plateaus of Fig. 2(b), marking atoms in the ground state. The two states $|e, \pm \hbar k\rangle$ are populated at certain values of the delay, where there is a transition in the final distribution between two plateaus. During the transition, the atom is scattered into four states: $|g, 2n\hbar k\rangle$, $|g, 2(n+1)\hbar k\rangle$, and $|e, \pm \hbar k\rangle$. The atomic momentum distribution is plotted in Fig. 4 for three different values of the delay to illustrate such a transition—namely, between the $8\hbar k$ and $10\hbar k$ plateaus. In Figs. 4(a) and 4(c) the atom is scattered cleanly into a single-momentum state and is left in the ground state. At an intermediate value of the delay depicted in Fig. 4(b), the atom is scattered into four final states with equal probabilities.

When the delay is very small, the plateaus and the transitions between them are no longer resolved (see Fig. 5), but it is nevertheless true that the atoms split into two beams that are well separated in momentum space. Atoms in the excited state remain around $p = 0$, and atoms in the ground state obtain a fairly large momentum. Both beams have small divergence ($\sigma_p \sim 1\hbar k$). This is illustrated in Fig. 6, where the final
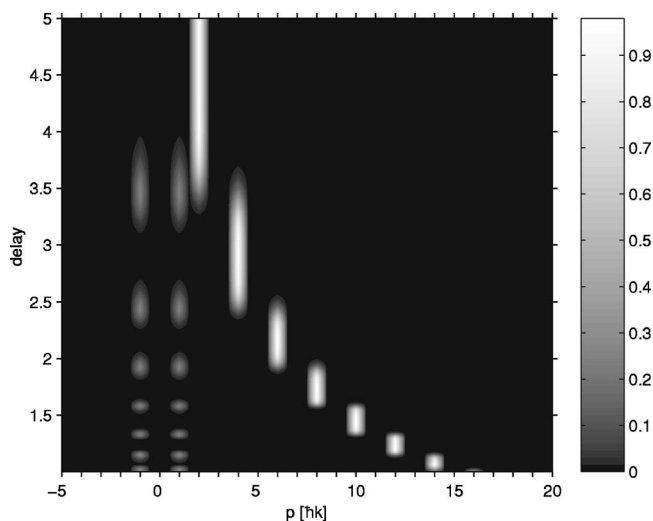
FIG. 3. Final momentum distribution of an atom after the action of two overlapping laser pulses as a function of the delay T/τ in the interval $[1,5]$. All parameters are the same as in Fig. 2(b).

momentum distribution has been plotted for a specific value of the delay. This means that retroreflected chirped pulses may also be good candidates for the realization of polarizing beam splitters.

B. Repeated interaction: Cumulative effect and spontaneous emission

For numerous applications, such as atomic beam deflection or slowing or accelerating clouds of atoms, it is desirable to transfer large momentum to the atoms without substantial heating. Counterpropagating laser pulses that first excite, then deexcite, the atoms are favorable because the atomic ensemble is accelerated as a whole, all the atoms receiving exactly the same amount of momentum. However, any method that transfers only a few $\hbar k$ momentum to the atoms in a single interaction period must be repeated many times. This is possible only if the atoms are returned to the initial internal state at the end of each interaction cycle. Imperfect population transfer will inevitably lead to heating of the atomic ensemble. For example, if two counterpropagating, separated laser pulses deliver $2\hbar k$ momentum to most of the atoms, but a few percent are left in the excited state after the second pulse, these atoms will gain $-2\hbar k$ momentum in the next interaction cycle; i.e., they will be accelerated in the opposite direction compared to the majority. The larger the number of interaction cycles used, the smaller the imperfections of population transfer that can still be tolerated. The importance of the final population of the excited state that is plotted in Fig. 2(b) is now obvious: the interaction cycle using retroreflected chirped pulses can be repeated for a cumulative effect only in the parameter regions where the final population of the excited state is close to zero. These regions coincide with the plateaus where the effect of a single interaction cycle is to transfer $2n\hbar k$ momentum to the atoms uniformly, without any heating. They are thus fine candidates for useful applications. Especially, since a single interaction cycle is several times more effective than for separate pulses,
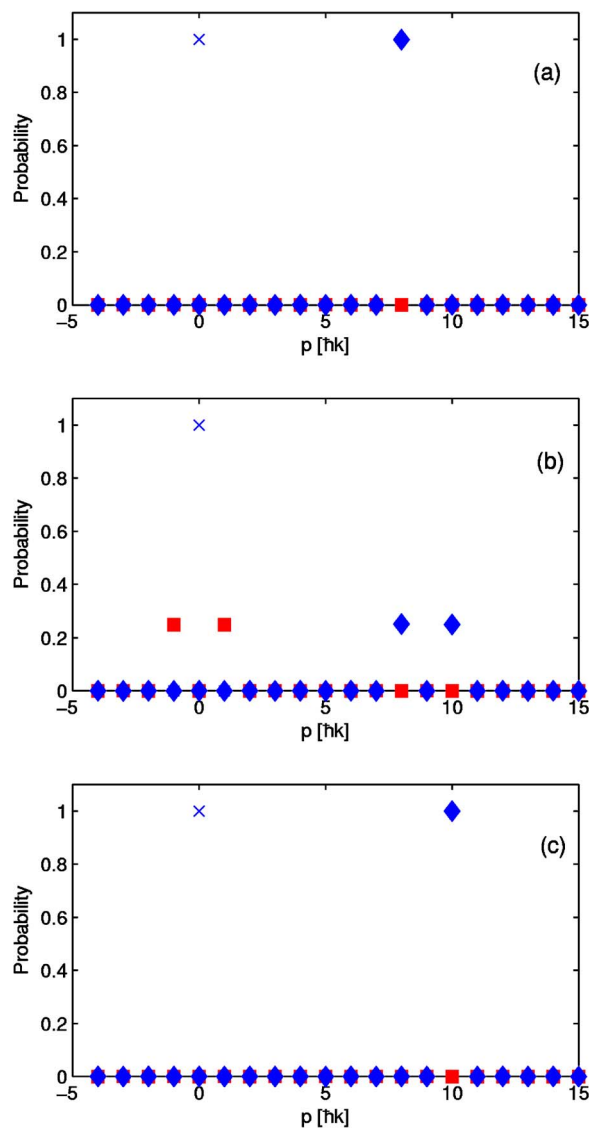

FIG. 4. (Color online) Final momentum distribution of the atom at three different values of the delay, showing a transition between two plateaus of Fig. 3. Diamonds mark ground-state atomic probabilities, squares mark excited-state atomic probabilities, and crosses mark the initial momentum state. (a) $T/\tau=1.7$, in the interval of the $8\hbar k$ plateau; (b) $T/\tau=1.582$, halfway between the two plateaus; (c) $T/\tau=1.42$ in the interval of the $10\hbar k$ plateau.

so considerably fewer repetitions are needed and thus greater imperfections of the population transfer are tolerable.

So far we have discussed only coherent processes, neglecting spontaneous relaxation of the atoms during the interaction altogether. When we consider the effects of a single pulse pair, this may well be justified, as laser pulses can have durations much smaller than the spontaneous lifetime of the atoms. When it comes to the cumulative effect of a sequence of many pulses, however, greater care must be taken, as it is very difficult to realize a whole sequence of many pulses in such a short time. In practice, the manipulation of the atoms often continues for longer than the lifetime of the excited state and relaxation has to be taken into account even if the individual pulses themselves are much shorter than τ_{spont}. In
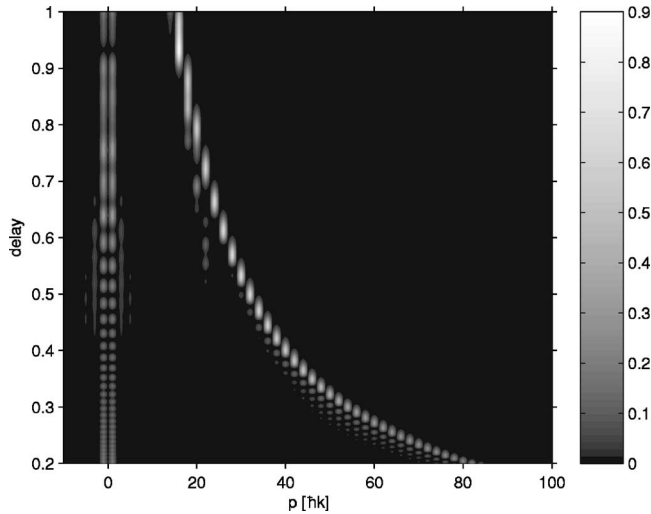
FIG. 5. Final momentum distribution of an atom after the action of two overlapping laser pulses as a function of the delay T/τ in the interval $[0.2, 1]$. All parameters are the same as in Fig. 2(a).

[27] the effect of spontaneous emission has been considered on a manipulation scheme using a sequence of nonoverlapping, counterpropagating, chirped pulses. It has been shown that if the pulses are very short, but the atoms interact with the repeated sequence of pulses longer than the spontaneous lifetime, the average force acting on the atoms converges to a stationary value. This is also true for the average heating suffered by the atomic ensemble. These values depend on the relative magnitude of the delay between the first and second pulses of a cycle, the cycle time (in the present case T and T_c), and the excited-state lifetime τ_{spont}. When $T_c \gg T$ and $T \ll \tau_{spont}$, the stationary value of the force is close to the ideal $2\hbar k/T_c$ and the heating effect is minimal, but in practice, a value of about $T_c = 4T$ has been achieved [15], for which the average force is only about half the theoretical
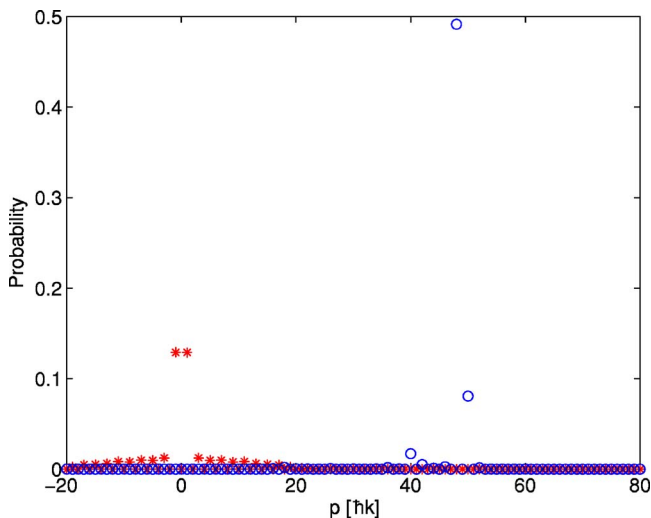

FIG. 6. (Color online) Final momentum distribution of the atoms for $T/\tau = 0.335$. Circles mark ground-state probabilities. The primary peak is at $p = 48\hbar k$. Excited-state probabilities are marked with asterisks. Atoms in the excited state are to be found predominantly in the $p = \pm 1\hbar k$ momentum states.
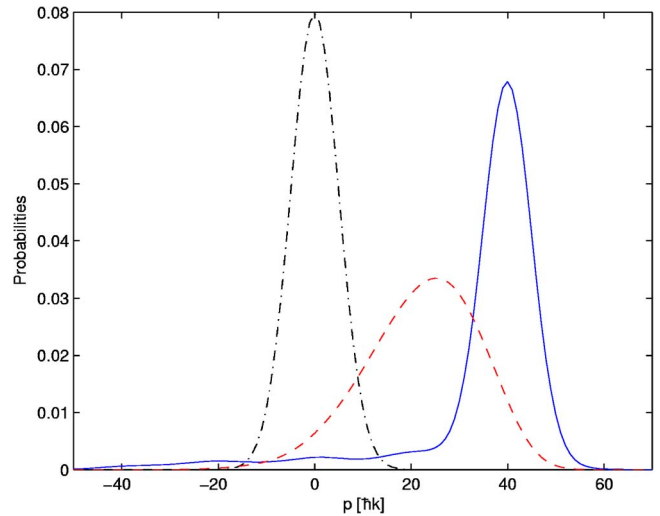

FIG. 7. (Color online) The effect of several interaction cycles on the atomic momentum distribution in the presence of spontaneous decay. The initial distribution is shown with a dash-dotted line, the distribution after the interaction with four cycles of overlapping, chirped pulses with a solid line, and the distribution after the interaction with 20 cycles of nonoverlapping chirped pulses with a dashed line. The calculations have been performed using the Monte Carlo wave function method with $N = 1000$ atoms.

maximum. (In this paper [15], it was also noted that the force unexpectedly increased when the parameters of the experiments, originally designed to demonstrate the force exerted by separate counterpropagating pulses, were changed. While no quantitative information has been given on this observation, it is possible that the effects of pulse overlap have been observed.)

Based on the results of [27], we may anticipate that the scenario described here is considerably less sensitive to the adverse effects of spontaneous emission, as the atoms spend much less time in the excited state during a single cycle and fewer cycles are needed for effective momentum transfer. To quantify the difference, we have performed a calculation of Eqs. (4) with spontaneous emission using the Monte Carlo wave function scheme [23,24]. Figure 7 shows the atomic momentum distributions obtained from this calculation. The initial distribution before the interaction (dash-dotted line) is a Gaussian distribution with $\sigma_p = 5\hbar k$. The momentum distribution after the interaction with four pairs of overlapping pulses is shown with a solid line. The pulse parameters are $\Omega = 100/\tau$, $\beta = 20/\tau^2$, the same as for Fig. 2. The delay between the pulses is $T = 1.42\tau$, the cycle time $T_c = 10\tau$, and the spontaneous lifetime is $\tau_{spont} = 1/\Gamma = 27\tau$. Note that these pulses are not extremely short compared to the spontaneous lifetime of the atoms, the full width at half maximum (FWHM) is $\sim 2.35\tau$ (in amplitude) or $\sim 1.67\tau$ (in intensity) which is only about one order of magnitude smaller than τ_{spont}. Still, the atomic momentum distribution has been translated unchanged to a fairly good measure. The maximum of the distribution is at $p = 40\hbar k$, and almost 90% of the atoms reside in the foremost peak, in the ground state, whose width is $\sigma_p = 8.4\hbar k$. A long "tail" with atoms in the excited state has also developed which contains about 10%

of the atoms and is strongly dispersed. The overall momentum transfer to the atoms is $\langle p \rangle = 35 \hbar k$. (In the absence of spontaneous emission there would have been a uniform $40 \hbar k$ translation of the momentum distribution.)

The dashed line in Fig. 7 shows the atomic momentum distribution after the interaction with 20 cycles of nonoverlapping chirped pulses of the same magnitude. The delay between the two pulses within a cycle is $T = 7.5 \tau$, and the cycle time is $T_c = 30 \tau$. The resulting momentum distribution can be seen to be much wider, $\sigma_p = 11.9 \hbar k$, and the average momentum transfer to the atoms is $\langle p \rangle = 22.1 \hbar k$. (In the absence of spontaneous emission these atoms would also have received $20 \times 2 \hbar k = 40 \hbar k$ momentum each.) From the figure it is clear that the sensitivity to the adverse effects of spontaneous emission is much smaller for the overlapping pulse case. This can be attributed to two factors. On the one hand, the atoms spend less time in their excited state during a single cycle—a pulse delay of 7.5τ between the first and second pulses of the cycle is more than a fourth of the spontaneous lifetime, so atoms have a greater chance to decay in the separated-pulse case. On the other hand, fewer cycles are needed to achieve the same amount of momentum transfer. (In the case of overlapping pulses the whole sequence lasted for $t_{max} = 40 \tau$, while for the separate pulses $t_{max} = 600 \tau$.)

The calculation presented in Fig. 7 has been performed using a one-dimensional calculation, so it shows the heating of the atomic ensemble only in the longitudinal direction (the direction of the acceleration of the atoms). Naturally, random recoils due to spontaneously emitted photons heat the ensemble also in the transverse directions. However, the effect of these random recoils is completely negligible—the heating of the atoms in the longitudinal direction is caused mostly by atoms being accelerated in the wrong direction after a spontaneous transition. Since this mechanism is effective only in the longitudinal direction, transverse heating is orders of magnitude smaller than the longitudinal one.

One further comparison is to be made at this point: namely, to the bichromatic force [8–11], which can be used very effectively to slow atoms in a beam. The bichromatic force can be much larger than radiation pressure and also possesses a great velocity capture range—both of these are similar to that achievable by our proposed scheme. The current method has two advantages over it. The first one is that the bichromatic force is velocity dependent (although in a sufficiently large velocity range, this dependence is not great). In our scheme, the force is completely independent of the atomic velocity in a large velocity range. The other advantage concerns heating of the atoms. Heating of the atoms by the bichromatic force in the longitudinal direction cannot be avoided. The directionality of that force is provided by an asymmetric choice of the beat phase of the two counterpropagating, bichromatic beams. The asymmetry, however, cannot be too great, as the series of counterpropagating π pulses that are created should not overlap. In practice, this means that with optimal parameters, the atoms are accelerated in one direction by the bichromatic beams approximately 75% of the time and in the opposite direction 25% of the time, changing direction randomly due to spontaneous events. In our scheme, the directionality of the force comes from the direction of propagation of the first pulse. Atoms

also switch their direction of acceleration due to spontaneous emission events, but the ratio of the atoms accelerating opposite to the majority can be very much smaller than $1/3$ (which is the value for the bichromatic case). Thus, when the initial atomic velocity distribution is not very wide (several times 100 m/s), our method can be more convenient. (Note that the heating effect of the bichromatic force is not always obvious in the literature, as the initial distribution of the atoms is usually a very wide thermal distribution, so slowing the atoms down to a velocity in the range of a few times 10 m/s can result in a compression of the distribution—i.e., a cooling.)

To evaluate the practical realizability of the proposed method, we first take the $2\,{}^3S_1 \rightarrow 2\,{}^3P_2$ transition of metastable He. The parameters used throughout this paper then translate to $\tau_{spont} = 100$ ns, $\tau = 3.7$ ns, $\Omega = 27$ GHz, and $\beta = 1.46$ GHz/ns. The peak intensity required is about $I \approx 1.8$ kW/cm^2, and the optical path difference between the pulses should be about 1.58 m. This intensity is much higher than, for example, the ones used in the experiments of [15] with cw lasers, but is not unreasonable to reach with modern pulsed lasers. We note that it is possible to reduce the peak Rabi frequency by a factor of 4 while doubling the pulse length; then, the atoms will still receive a momentum of $10 \hbar k$ per cycle and the peak intensity required will be only about 110 W/cm^2. (Though the atomic velocity distribution then suffers a slightly greater spreading due to spontaneous processes.) On the other hand, the method using two symmetrically detuned, overlapping pulses described in [16] would require about $\Omega = 2000/\tau$ for a similar effect, which translates into a peak intensity of 720 kW/cm^2 for the above-mentioned transition and $\tau = 3$ ns pulse duration. Clearly, this would be much harder to realize.

Another frequently used transition is the $5S_{1/2} \rightarrow 5P_{3/2}$ transition of ^{85}Rb. For this transition, our parameters correspond to $\tau_{spont} = 27$ ns, $\tau = 1$ ns, $\Omega = 100$ GHz, and $\beta = 20$ GHz/ns. This, however, raises the concern that since the ground-state hyperfine splitting of this transition is about $\Delta_{hyperfine} = 3$ GHz—i.e., the same order of magnitude as the frequency sweep covered by the chirp, $\tau \beta / 2\pi \sim 3$ GHz—the application of the two-level atom model is inappropriate. The simplest way to circumvent this would be to double the pulse length to $\tau = 2$ ns, in which case the required chirp rate becomes $\beta = 5$ GHz/ns and the two-level approximation is well justified. (Our quantities refer to angular frequency, so in terms of frequency this chirp rate is only about 0.8 GHz/ns.) The process is then slightly more susceptible to incoherent effects. The peak intensity required is $I \approx 4.6$ kW/cm^2, and the optical path difference between the pulses should be about 0.86 m. By comparison, the symmetrically detuned pulse scheme of [16] would require an intensity of 1.8 MW/cm^2 for a similar effect. Another way through the problem would be to use laser pulses, whose transform limited bandwidth $1/\tau$ is much larger than the ground-state hyperfine splitting—these are much harder to realize though with a sufficiently large intensity and chirp. Finally we mention that even in this intermediate regime (i.e., $1/\tau < \Delta_{hyperfine} \lesssim \beta \tau$), it is still possible to utilize our method, as adiabatic population transfer between selected states using chirped pulses is possible [28]. However, then the direction

of the chirp will become important too. A detailed treatment of multilevel atoms interacting with partially overlapping, chirped laser pulses will be presented in a forthcoming publication.

IV. SUMMARY

We have investigated the mechanical effect of retroreflected, frequency-chirped, overlapping laser pulses on two-level atoms. We have shown that the momentum transferred to the atoms by a single pulse pair can be several times $2\hbar k$, which is the value for two counterpropagating, but nonoverlapping chirped pulses or π pulses. In certain parameter regimes, the momentum transfer is accompanied by negligible heating of the atoms and they are returned to the ground state at the end of the interaction. This makes the process repeat-able a number of times with a cumulative effect. The proposed method has a wide velocity capture range, so an atomic ensemble with considerable Doppler width may be accelerated uniformly. We have also shown that the adverse effects of spontaneous emission are much smaller in the case of overlapping pulses. In other regimes, the atoms are split into two well-separated beams, both of which have a limited divergence. Therefore these laser pulses can be very useful in the construction of atom-optical elements such as mirrors or beam splitters.

[1] I. Nebenzahl and A. Szöke, Appl. Phys. Lett. **25**, 327 (1974).

[2] J. S. Bakos, G. P. Djotyan, G. Demeter, and Zs. Sörlei, Phys. Rev. A **53**, 2885 (1996).

[3] G. P. Djotyan, J. S. Bakos, G. Demeter, and Zs. Sörlei, J. Opt. Soc. Am. B **13**, 1697 (1996).

[4] T. Freegarde, G. Daniell, and D. Segal, Phys. Rev. A **73**, 033409 (2006).

[5] T. G. M. Freegarde, J. Walz, and T. W. Hänsch, Opt. Commun. **117**, 262 (1995).

[6] A. Goepfert, I. Bloch, D. Haubrich, F. Lison, R. Schutze, R. Wynands, and D. Meschede, Phys. Rev. A **56**, R3354 (1997).

[7] M. Mützel, D. Haubrich, and D. Meschede, Appl. Phys. B: Lasers Opt. **70**, 689 (2000).

[8] V. S. Voitsekovich *et al.*, Sov. Phys. Tech. Phys. **33**, 690 (1988).

[9] J. Söding, R. Grimm, Y. B. Ovchinnikov, P. Bouyer, and C. Salomon, Phys. Rev. Lett. **78**, 1420 (1997).

[10] M. R. Williams, F. Chi, M. T. Cashen, and H. Metcalf, Phys. Rev. A **61**, 023408 (2000).

[11] M. T. Cashen and H. Metcalf, Phys. Rev. A **63**, 025406 (2001).

[12] E. Peik, M. Ben Dahan, I. Bouchoule, Y. Castin, and C. Salomon, Phys. Rev. A **55**, 2989 (1997).

[13] R. Battesti, P. Clade, S. Guellati-Khelifa, C. Schwob, B. Gremand, F. Nez, L. Julien, and F. Biraben, Phys. Rev. Lett. **92**, 253001 (2004).

[14] V. S. Malinovsky and P. R. Berman, Phys. Rev. A **68**, 023610 (2003).

[15] M. Cashen, O. Rivoire, L. Yatsenko, and H. Metcalf, J. Opt. B: Quantum Semiclassical Opt. **4**, 75 (2002).

[16] V. I. Romanenko and L. P. Yatsenko, JETP **90**, 407 (2000).

[17] L. P. Yatsenko, B. W. Shore, N. V. Vitanov, and K. Bergmann, Phys. Rev. A **68**, 043405 (2003).

[18] A. P. Conde, L. P. Yatsenko, J. Klein, M. Oberst, and T. Halfmann, Phys. Rev. A **72**, 053808 (2005).

[19] R. Unanyan, S. Guérin, B. W. Shore, and K. Bergmann, Eur. Phys. J. D **8**, 443 (2000).

[20] V. I. Romanenko and L. P. Yatsenko, JETP **100**, 242 (2005).

[21] H. Maeda, J. H. Gurian, D. V. L. Norum, and T. F. Gallagher, Phys. Rev. Lett. **96**, 073002 (2006).

[22] G. N Gibson, Phys. Rev. A **72**, 041404(R) (2005).

[23] J. Dalibard, Y. Castin, and K. Molmer, Phys. Rev. Lett. **68**, 580 (1992).

[24] R. Dum, P. Zoller, and H. Ritsch, Phys. Rev. A **45**, 4879 (1992).

[25] S. Guérin, L. P. Yatsenko, and H. R. Jauslin, Phys. Rev. A **63**, 031403(R) (2001).

[26] L. P. Yatsenko, S. Guérin, and H. R. Jauslin, Phys. Rev. A **65**, 043407 (2002).

[27] G. Demeter, G. P. Djotyan, and J. S. Bakos, J. Opt. Soc. Am. B **15**, 16 (1998).

[28] G. P. Djotyan, J. S. Bakos, G. Demeter, and Zs. Sörlei, J. Opt. Soc. Am. B **17**, 107 (2000).