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Coherent atomic mirrors and beam splitters by adiabatic passage in multilevel systems

P. Marte, P. Zoller, and J. L. Hall*

Joint Institute for Laboratory Astrophysics, University of Colorado, Boulder, Colorado 80309-0440

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We study atomic-beam deflection by adiabatic passage between Zeeman ground levels via Raman transitions induced by counterpropagating $\sigma^{\pm}$-polarized lasers. We show that complete population transfer between the ground states can be achieved, which corresponds to the scattering of the atomic wave packet into a *single* final momentum state by absorption and induced emission of laser photons. Although the lasers can be resonant, the excited state(s) are never populated during the adiabatic transfer, which suppresses the effects of spontaneous emission and preserves the coherence of the atomic wave function. This scheme has attractive features as a beam splitter and mirror for atomic interferometry.

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Matter-wave interferometry with neutral atoms is currently a focus of research in atomic physics [1-5]. Compared to a neutron interferometer [6], an atomic interferometer (AI) promises enhanced sensitivity for high-precision experiments in both gravitational and quantum physics, and a new generation of frequency standards and rotation sensors. In addition, one of the new perspectives of atomic interferometry is the wealth of possible experiments to manipulate and probe atoms with laser radiation: this includes scattering of atomic wave packets from light waves (e.g., the Kapitza-Dirac effect) to build atomic-beam splitters and mirrors [1,2,7,8], laser-cooling techniques to prepare slow atomic beams [9], and applications to laser spectroscopy. Recently, several groups have reported the first experimental observations of atomic interference fringes by scattering atoms from mechanical gratings [3,4], or by applying a sequence of short $\pi/2$, π, and $\pi/2$ laser pulses to atoms in an atomic fountain [5] similar in concept to optical Ramsey experiments.

The key element of an AI is the atomic-beam splitter and mirror. An atomic-beam splitter separates the single-atom wave function into a macroscopic superposition state corresponding to two center-of-mass wave packets propagating in different spatial directions. An atomic mirror, on the other hand, deflects these wave packets so that the matter waves traveling along two paths of the interferometer can be brought to interference. This interference will, of course, be observed only if these scattering processes are coherent.

In this paper we discuss and analyze a scheme for coherent atomic-beam deflection from laser light waves which combines several attractive features, close to the requirements of an "ideal" atomic-beam splitter and mirror for atomic interferometry. The proposed scheme is based on the concept of "coherent adiabatic population transfer in Raman processes with *time-delayed laser pulses*," as first demonstrated in the context of molecular spectroscopy by Bergmann and co-workers [10,11] and discussed theoretically by Oreg, Hioe, and Eberly [12], Hioe and Carroll [13], and Kuklinski *et al.* [11].

We consider a stationary collimated beam of atoms propagating along the x axis of our coordinate system. The atoms are scattered from two counterpropagating σ^+ and σ^- light waves that both have frequency ω and wave vectors k directed along the $+z$ and $-z$ axis, respectively. The atomic configuration is a three-level system with two Zeeman ground states $|g_{m=\pm 1}\rangle$ coupled to an excited state $|e_{m=0}\rangle$ corresponding to a $J_g=1$ to $J_e=1$ transition as illustrated in Fig. 1. An example of this configuration is the $2s\,{}^3S_1$ to $2p\,{}^3P_1$ transition in metastable helium [14]. The incident atoms are deflected from the light waves by absorption and subsequent reemission of laser photons. For small deflection angles (when the longitudinal atomic momentum $p_z=Mv_z$ is larger than the photon recoil $\hbar k$) our problem can be reduced to solving the one-dimensional Schrödinger equation for the transverse atomic motion [1,2]; the corresponding atomic Hamiltonian is [14]

$$\begin{aligned} H_{0A}(t)=\frac{\hat{p}_z^2}{2M}+\hbar\omega_{eg}|e_0\rangle\langle e_0| & \\ -\tfrac{1}{2}\hbar[\Omega_-(t)e^{-ik\hat{z}-i\omega t}|e_0\rangle\langle g_1| & \\ +\Omega_+(t)e^{ik\hat{z}-i\omega t}|e_0\rangle\langle g_{-1}|+\text{H.c.}]\,. & \quad (1)\end{aligned}$$

The time coordinate is related to the atomic motion $x=v_xt$ along the x axis. In particular, the time dependence of the Rabi frequencies $\Omega_\pm(t)$ (see Fig. 1) corresponds to the atom moving through the laser interaction zones. We choose $\Omega_\pm$ real. ω_{eg} is the atomic transition frequency. The Hamiltonian $H_{0A}(t)$ has the property that it couples only states within the family $\{|g_-,p_z-\hbar k\rangle,|e_0,p_z\rangle,|g_+,p_z+\hbar k\rangle\}$ with p_z the transverse

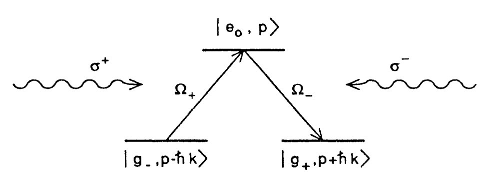

FIG. 1. Three-level system with Zeeman ground states $|g_{\pm 1}\rangle$ which are coupled by $\sigma^{\pm}$-polarized laser light to the upper state $|e_0\rangle$.

momentum. In addition to the laser-induced scattering, we allow for spontaneous decay during the interaction. Thus the total Hamiltonian is

$$H = H_{0A}(t) + H_{0F} + V_{AF} \tag{2}$$

with $H_{0A}(t)$ the atomic Hamiltonian equation (1), H_{0F} the Hamilton operator of the quantized free radiation field, and V_{AF} the dipole coupling of the vacuum modes to the atom [2]. The state vector $|\Psi(t)\rangle$ of the combined atom-radiation field consists of contributions where the atom has emitted no spontaneous photon, one photon, two photons, etc. [2]. Thus the state vector can be expanded according to

$$|\Psi(t)\rangle = |\text{vac}\otimes \int dp_z [|g_-,p_z-\hbar k\rangle a_-(p_z,t) + |e_0,p_z\rangle a_0(p_z,t)e^{-i\omega t} + |g_+,p_z+\hbar k\rangle a_+(p_z,t)] + \cdots \tag{3}$$

where $a_{\pm,0}(p_z,t)$ are the atomic vacuum amplitudes and the ellipses represents one-, two-, etc., photon contributions. Since the momentum transfer to the atom by spontaneous emission corresponds to random momentum kicks, the vacuum contribution in Eq. (3) is the (coherent) part of the state vector responsible for interference fringes in the AI. It can be shown that these amplitudes obey the Schrödinger equation [2]

$$i\frac{\partial}{\partial t}a_\pm = \left[\frac{p_z^2}{2\hbar M} + \omega_R \pm kv_z\right]a_\pm - \tfrac{1}{2}\Omega_\mp a_0(p_z,t)\,, \tag{4}$$

$$i\frac{\partial}{\partial t}a_0(p_z,t) = \left[\frac{p_z^2}{2\hbar M} - \Delta - i\tfrac{1}{2}\kappa\right]a_0 - \tfrac{1}{2}\Omega_- a_+(p_z,t) - \tfrac{1}{2}\Omega_+ a_-(p_z,t)\,, \tag{5}$$

with $\Delta = \omega - \omega_{eg}$ the detuning from resonance, and κ the spontaneous emission rate of the upper state. The ground states are shifted from the two-photon Raman resonance condition due to the Doppler detunings $\pm kv_z$ (with $v_z = p_z/M$); $\omega_R = \hbar k^2/2M$ is the recoil shift. In the following we assume that the atom is initially prepared in the $|g_0\rangle$ state with center-of-mass distribution corresponding to a well-collimated atomic beam. Note that the population transfer $\{|g_-,p_z-\hbar k\rangle \rightarrow |g_+,p_z+\hbar k\rangle\}$ corresponds to a momentum transfer $2\hbar k$, i.e., to a deflection of the atom.

Equation (4) is analogous to the equations studied in Refs. [11-13] in the context of optimizing population transfer in three-level (molecular) systems. Adopting the arguments presented in these papers in our present case, we see that an initial state $|g_-,p_z-\hbar k\rangle$ can be *adiabatically* transferred to $|g_+,p_z+\hbar k\rangle$ provided the two pulses $\Omega_\pm(t)$ are time delayed with respect to each other (but still overlapping) such that the σ^- wave, i.e., the light wave acting on the second transition in Fig. 1, *precedes* the σ^+ pulse. This population transfer (and hence momentum transfer) is illustrated in Fig. 2 where we have plotted the population of the final state $|g_+\rangle$ (i.e., $|a_+|^2$) as a function of the time delay τ between two pulses. The curves in Fig. 2 were obtained by numerically integrating the Schrödinger equation (4) for two Gaussian pulses with equal pulse durations $T_1 = T_2 = T$ (full width at half maximum) and intensities ($\Omega_+ = \Omega_-$), and the initial condition that the atoms are prepared in the $|g_-\rangle$ state. The dashed and solid curves correspond to $\Omega T = 50$, with $\kappa = 0$ and $\kappa T = 5$ ($\Omega/\kappa = 10$), respectively. In both cases we find a broad maximum for $\tau \approx -1.57T$, which is fairly insensitive to the presence of spontaneous decay even if the interaction time is long compared with the lifetime of the upper state ($\kappa T = 5$, solid curve). For $\tau \geq 0$, on the other hand, the upper-state population shows strong (Rabi) oscillations for the undamped case (dashed curve), and is close to zero in the presence of spontaneous emission (solid curve).

These results can be explained in an adiabatic dressed state picture [11-13]. For $v_z \approx 0$ the Hamiltonian matrix in Eq. (4) has a adiabatic dressed state eigenvalue $E = 0$. The associated eigenvector is

$$|E=0\rangle = \frac{\Omega_-}{(\Omega_+^2+\Omega_-^2)^{1/2}}|g_-,p_z-\hbar k\rangle - \frac{\Omega_+}{(\Omega_+^2+\Omega_-^2)^{1/2}}|g_+,p_z+\hbar k\rangle\,, \tag{6}$$

which is not contaminated by admixtures from the (decaying) excited state, and is independent of both Δ and κ. Furthermore it follows from Eq. (6) that

$$|E=0\rangle \rightarrow \begin{cases} |g_-,p_z-\hbar k\rangle, & \text{for } \Omega_+(t)/\Omega_-(t)\rightarrow 0 \quad (7)\\ -|g_+,p_z+\hbar k\rangle, & \text{for } \Omega_-(t)/\Omega_+(t)\rightarrow 0. \quad (8)\end{cases}$$

Thus for a pulse sequence where the $\Omega_+(t)$ pulse is time delayed with respect to $\Omega_-(t)$, (6) will satisfy condition (7) at the beginning and (8) at the end of the interaction.

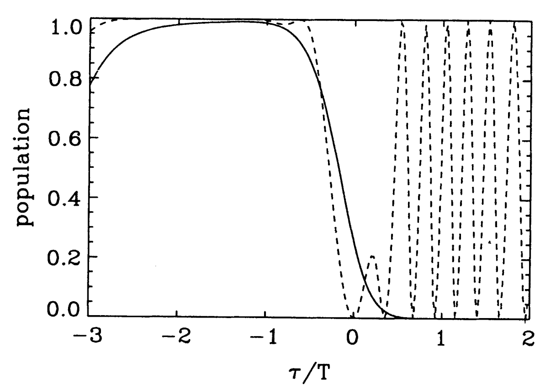

FIG. 2. Population of the final $|g_{+1}\rangle$ state as a function of the time delay τ between the Gaussian laser pulses. The Rabi frequencies are $\Omega_\pm T = 50$. The solid curve corresponds to a spontaneous decay rate $\kappa T = 5$, while the dashed curve is for $\kappa T = 0$. T is the pulse duration of both laser pulses.

Therefore the pulses transform $|g_-,p_z-\hbar k\rangle$ adiabatically into $|g_+,p_z+\hbar k\rangle$. For $\Delta=0$ the other two dressed state eigenvalues are $E_\pm=\pm\frac{1}{2}(\Omega_+^2+\Omega_-^2)^{1/2}$. The condition for the initial state to follow $|E=0\rangle$ adiabatically is [11-13]

$$\Omega_\pm T \gg 1, \;\; \Omega_\pm \gg \kappa . \tag{9}$$

The pulses must have significant overlap in time so that the dressed energy level $|E=0\rangle$ is well separated from the states $|E_\pm\rangle$ for all times during the interaction. For transverse Doppler shifts the above argument is still valid if $\Omega_\pm \gg kv_z$; typically, we expect v_z to be of the order of a few recoil velocities $v_r=\hbar k/M$ so that this assumption is extremely well fulfilled (Raman-Nath approximation).

The attraction of the above scheme for atomic interferometry is based on the following features: We have close to 100% scattering to a *single* final state corresponding to a fixed momentum transfer $2\hbar k$. Thus there is no splitting of the incident wave packet into a superposition corresponding to many momentum peaks $\pm\hbar k$, $\pm 2\hbar k$, etc., as for scattering of a two-level system from a standing light wave. Since the process is adiabatic, it is, within the validity of the adiabaticity condition equation (9), quite robust against changes in laser parameters (Rabi frequencies and detunings), interaction time (atomic velocity), etc.; this is in contrast to transfer by π pulses with Rabi oscillations which is very sensitive to the exact value of the pulse area. Furthermore, although the interaction is resonant, the scattering process is immune to spontaneous decay, as the excited state is never populated during the transition. This avoids the common approach of eliminating spontaneous decay by detuning the laser far off resonance to reduce the excited-state population, at the expense of introducing a weak atom-laser coupling. Finally, the Λ configuration of Fig. 1 has the advantage that it requires only a single laser since the σ^- wave can be derived by reflecting the σ^+ laser light.

The scheme outlined above for three-level systems is readily generalized to a chain of Raman transitions. This leads to an increased deflection angle due to the momentum transfer associated with the multiphoton transition. As an example consider the five-level system shown in Fig. 3 corresponding to a $J_g=2$ to $J_e=2$ transition with a transfer of $4\hbar k$. We assume again that the atoms are initially prepared in the $|g_{-2}\rangle$ state. Two counterpropagating circularly polarized waves of frequency ω couple a chain of two Λ transitions from $|g_{-2}\rangle$ to $|g_{+2}\rangle$. We can show again that there is an adiabatic dressed state eigenvalue $E=0$ with eigenvector

$$|E=0\rangle = N(\Omega_-^2\Omega_-^{2'}|g_{-2},p_z-2\hbar k\rangle - \Omega_+^2\Omega_-^2|g_0,p_z\rangle + \Omega_+^2\Omega_+^{2'}|g_{+2},p_z+2\hbar k\rangle) , \tag{10}$$

with N a normalization constant, and $\Omega_\pm$, $\Omega'_\pm$ Rabi frequencies related by appropriate Clebsch-Gordan coefficients (compare Fig. 3). From (10) it follows again that a delay of the σ^+ pulse with respect to the σ^- wave gives complete adiabatic population transfer from $|g_{-2}\rangle$ to the final state $|g_{+2}\rangle$ with the atom absorbing a momentum $4\hbar k$. In particular, we emphasize that there is no population left in the middle ground state $|g_0\rangle$ after the interaction. We have again no admixture from the excited states during the process and find the transfer to be insensitive to variations of the Rabi frequencies and detunings within the validity of the adiabaticity condition. This adiabatic four-photon process is illustrated in Fig. 4, which shows the time evolution of the atomic populations (the modulus squared of the vacuum amplitudes) during the pulse, obtained by numerical integration of the Schrödinger equation for two time-delayed Gaussian pulses. The parameters are $\Omega T=50$ and $\kappa T=5$ with time delay $\tau=-1.2T$ and $T=T_1=T_2$ the pulse duration. In both analytical and numerical work we have found that the adiabatic population transfer works extremely well even for very-high-order Raman transitions achieving high-momentum transfer in a single laser interaction zone.

Another attractive feature of adiabatic passage is the possibility to achieve large momentum transfer by deflection of atoms in several successive interaction zones. We discuss the idea for the case of a Λ transition interacting with $\sigma^\pm$ light waves (discussed in the context of Figs. 1 and 2). Again we consider an incident atomic wave packet prepared in a $|g_-,p_z\rangle$ state. The wave packet is scattered in a first interaction zone into the $|g_+,p_z+2\hbar k\rangle$ state with the $\sigma^\pm$ waves propagating in the $\pm z$ direc-

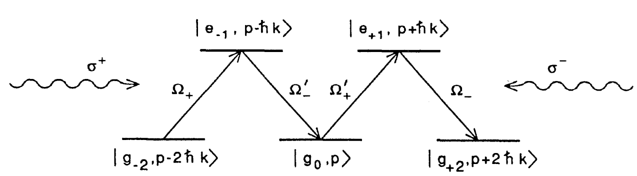

FIG. 3. Five-level system with Zeeman ground states $|g_{\pm 2,0}\rangle$ coupled by $\sigma^\pm$-polarized laser light to the upper states $|e_{\pm 1}\rangle$.

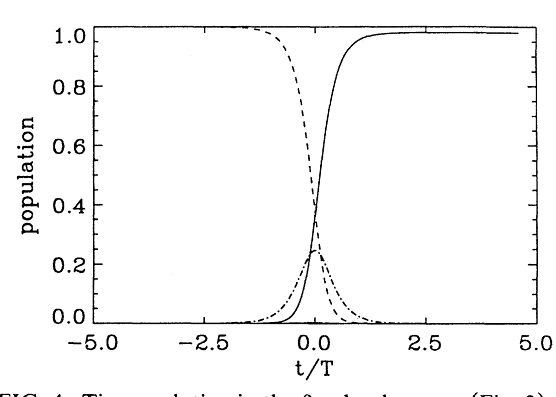

FIG. 4. Time evolution in the five-level system (Fig. 3) corresponding to a $J_g=2$ to $J_e=2$ transition. According to the Clebsch-Gordan coefficients the Rabi frequencies are related by $\Omega'_\pm(t)=\sqrt{\frac{3}{2}}\,\Omega_\pm(t)$. The population of the initial state $|g_{-2}\rangle$ (solid line), the middle state $|g_0\rangle$, and the final state $|g_{+2}\rangle$ is plotted as a function of the interaction time for Gaussian laser pulses with duration $T=T_1=T_2$. The parameters are $\Omega_\pm T=50$, $\tau=-1.2T$, $\Delta T=0$, $\kappa T=5$.

tions, respectively. In a second interaction zone with the lasers propagating in opposite directions the electron is transferred back to the $m=-1$ state, corresponding to $|g_-,p_z+4\hbar k\rangle$. Thus we have the sequence

$$|g_-,p_z\rangle \underset{\text{zone 1}}{\longrightarrow} -|g_+,p_z+2\hbar k\rangle \underset{\text{zone 2}}{\longrightarrow} |g_-,p_z+4\hbar k\rangle \rightarrow \cdots . \tag{11}$$

Note that after two laser zones the atom is in the same internal atomic state $|g_-\rangle$ that it occupied initially, but it has received a momentum kick of $4\hbar k$. An advantage of adiabatic passage is that its efficiency factor for the transfer is close to unity, which allows us to combine a large number of these interaction zones. Finally, replacing the Λ transitions (Fig. 1) by a chain of Raman transitions in a single zone (Fig. 3) promises very-high-order accumulated momentum transfer into a single final momentum state of the atom.

So far our discussion has concentrated on coherent beam deflection corresponding to an atomic mirror. There are several possibilities to realize a beam splitter. Ideally, one expects an atomic-beam splitter to produce a superposition state of two wave packets which differ by a large center-of-mass momentum. As a first possibility, Eq. (6) suggests that in a three-level Λ system (Fig. 1) a coherent superposition can be formed by pulse shapes $\Omega_+(t)/\Omega_-(t)\rightarrow 1$,

$$|E=0\rangle \rightarrow \frac{1}{\sqrt{2}}(|g_-,p_z-\hbar k\rangle+|g_+,p_z+\hbar k\rangle) . \tag{12}$$

A second possibility is to create a coherent superposition of two atomic ground states prior to the interaction with the laser light using a radio frequency (rf) field. Then *one* component of this superposition state is deflected *selectively* by adiabatic passage through a sequence of laser interaction zones. Consider again the Λ system $\{|g_-\rangle,|e_0\rangle,|g_+\rangle\}$ described in Fig. 1. Typically, $|g_m\rangle$ will be hyperfine structure components for a particular F value. Consider now a situation where in addition to F we have a second hyperfine structure state F' with states $|g'_m\rangle$ separated from the F states by a rf transition. Then a possible sequence of transitions to generate a superposition state is

$$\begin{aligned}|g_-,p_z\rangle &\underset{\text{rf}}{\longrightarrow} \cos\theta|g_-,p_z\rangle+\sin\theta|g'_-,p_z\rangle\\ &\underset{\text{zone 1}}{\longrightarrow} -\cos\theta|g_+,p_z+2\hbar k\rangle+\sin\theta|g'_-,p\rangle\\ &\underset{\text{zone 2}}{\longrightarrow} \cos\theta|g_-,p_z+4\hbar k\rangle+\sin\theta|g'_0,p_z\rangle\rightarrow \cdots .\end{aligned} \tag{13}$$

Here θ is the mixing angle due to the rf field. $|g'_-\rangle$ denotes a ground state which is assumed nonresonant with the laser light and, thus, remains undeflected. The last line in Eq. (13) corresponds to a macroscopic superposition state of two wave packets differing by a center-of-mass momentum $4\hbar k$.

To summarize, we have shown that adiabatic passage in multilevel systems leads to atomic-beam deflection with the possibility of achieving high-momentum transfer from the laser to the atomic wave packet to prepare a single momentum final state with high efficiency, while at the same time avoiding the momentum diffusion associated with spontaneous decay. Details of the theoretical derivations and further results will be published elsewhere. An analysis of an atomic interferometer in a triple Laue configuration with interaction zones using counterpropagating $\sigma^\pm$ light to couple Zeeman levels by Raman transitions has been given in Ref. [2(a)].

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*Quantum Physics Division, National Institute of Standards and Technology, Boulder, CO 80309-0440.

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