Storing and releasing light in a gas of moving atoms

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We propose a scheme of storing and releasing pulses or cw beams of light in a moving atomic medium illuminated by two stationary and spatially separated control lasers. The method is based on electromagnetically induced transparency but in contrast to previous schemes, storage and retrieval of the probe pulse can be achieved at different locations and without switching off the control laser.

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Over the past couple of years, there has been a great deal of interest in storing and releasing light pulses in atomic gases exhibiting electromagnetically induced transparency (EIT) [1–7]. EIT is a phenomenon in which a weak probe pulse travels slowly and almost without dissipation in a resonant medium controlled by another laser [8–14]. Following the proposal of Ref. [1], storage and release of a probe pulse has been demonstrated [2,3,15] by dynamically changing the intensity of the control laser. In the read-out process it is possible to alter certain properties of the stored probe pulse, such as its carrier frequency, propagation direction, or pulse shape [5,7].

Recently, an application of the light-storage scheme to generate continuous beams of *atoms* in nonclassical quantum states was proposed [16]. In that scheme the probe field propagates coparallel to a beam of atoms in a spatially varying control field. This corresponds in the rest frame of the atoms to an explicitly time dependent control laser and thus allows for a complete and loss-free adiabatic transfer of probe-field excitations to the matter wave [6]. We consider here another situation in which both the control and the incoming probe beams are perpendicular to the atomic motion. Such a scheme is better suited to eliminate effects of Doppler broadening and thus to alleviate limitations to the allowed velocity spread of the atomic beam. Furthermore, in contrast to previous setups [1–7], storing and releasing of a probe beam can now be achieved with a pair of stationary control lasers, i.e., there is no need to switch "off" and "on" a control laser at precise times. This is advantageous if one does not know the exact time of arrival of the probe pulse. Thus, the present setup could be used as an example to store entangled photons generated by a continuous optical parametric oscillator (OPO) below threshold and frequency selected by a cavity [17,18]. Finally, storage and release of the probe field are spatially separated in the present scheme.

Consider a stream of atoms moving along the x axis as shown in Fig. 1 (bottom part). The atoms are characterized by two hyperfine ground levels g and q, as well as an electronic excited level e (top part of Fig. 1). The atoms in different internal states are described by the field operators $\Psi_g \equiv \Psi_g(\mathbf{r},t)$, $\Psi_q \equiv \Psi_q(\mathbf{r},t)$ and $\Psi_e \equiv \Psi_e(\mathbf{r},t)$ obeying Bose-Einstein or Fermi-Dirac commutation relations (depending on the type of atoms). The matter fields interact with several light fields propagating along the z axis. A strong classical control laser centered at $x = x_1$ drives the transition $|e\rangle \rightarrow |q\rangle$, whereas a weaker quantum probe field (also entering the medium at $x = x_1$) is coupled with the transition $|g\rangle \rightarrow |e\rangle$. Finally there is a second spatially separated control laser (centered at $x = x_2$) that is used to release the stored probe pulse.

The first control laser has a frequency ω_c, a wave vector $\mathbf{k}_c = \hat{\mathbf{z}} k_c$, and a Rabi frequency $\Omega(\mathbf{r},t) = \Omega_1(x) e^{-i(\omega_c t - k_c z)}$, where the y dependence of the amplitude $\Omega_1(x) \equiv \Omega_1$ will be kept implicit. The same applies to other parameters, such as the mean atomic number density $n \equiv n(z)$.

The probe beam is described by the electric field operator: $\mathbf{E}(\mathbf{r},t) = \hat{\mathbf{e}}(\hbar c k / 2\varepsilon_0)^{1/2} \mathcal{E}(\mathbf{r},t) e^{-i(\omega t - kz)} + \text{H.c.}$, where $\mathcal{E}(\mathbf{r},t) \equiv \mathcal{E}$ is the slowly varying amplitude, $\omega = ck$ is the central frequency of probe photons, $\mathbf{k} = \hat{\mathbf{z}} k$ is the wave vector, and $\hat{\mathbf{e}} \perp \hat{\mathbf{z}}$ is the unit polarization vector. The dimension of $\mathcal{E}$ is such that the operator $\mathcal{E}^\dagger \mathcal{E}$ represents the number density of probe photons.

The atoms (initially in the ground level g) are moving along the x axis with an average speed $\mathbf{v}_0 = \hbar \mathbf{k}_0 / m$ and the kinetic energy $\hbar \omega_{at} = \hbar^2 k_0^2 / 2m$. The atomic velocities are assumed to be spread over a narrow range around v_0 so that we can introduce slowly varying atomic amplitudes as: $\Phi_g = \Psi_g e^{i(\omega_{at} t - \mathbf{k}_0 \cdot \mathbf{r})}$, $\Phi_e = \Psi_e e^{i(\omega_{at} + \omega)t - i(\mathbf{k}_0 + \mathbf{k}) \cdot \mathbf{r}}$, and $\Phi_q = \Psi_q e^{i(\omega_{at} + \omega - \omega_c)t - i(\mathbf{k}_0 + \mathbf{k} - \mathbf{k}_c) \cdot \mathbf{r}}$.

Consider first storing of the probe by the first control laser. The following equations hold for the slowly varying electromagnetic and matter field operators:

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z} \right) \mathcal{E} = i g \Phi_g^\dagger \Phi_e , \tag{1}$$

$$\left(\frac{\partial}{\partial t} + v_0 \frac{\partial}{\partial x} \right) \Phi_g = i g \mathcal{E}^\dagger \Phi_e , \tag{2}$$

$$\left(\frac{\partial}{\partial t} + i\Delta + v_0 \frac{\partial}{\partial x} + v_r \frac{\partial}{\partial z} \right) \Phi_e = i \Omega_1 \Phi_q + i g \mathcal{E} \Phi_g , \tag{3}$$

$$\left(\frac{\partial}{\partial t} + i\delta + v_0 \frac{\partial}{\partial x} + \Delta v_r \frac{\partial}{\partial z} \right) \Phi_q = i \Omega_1 \Phi_e , \tag{4}$$

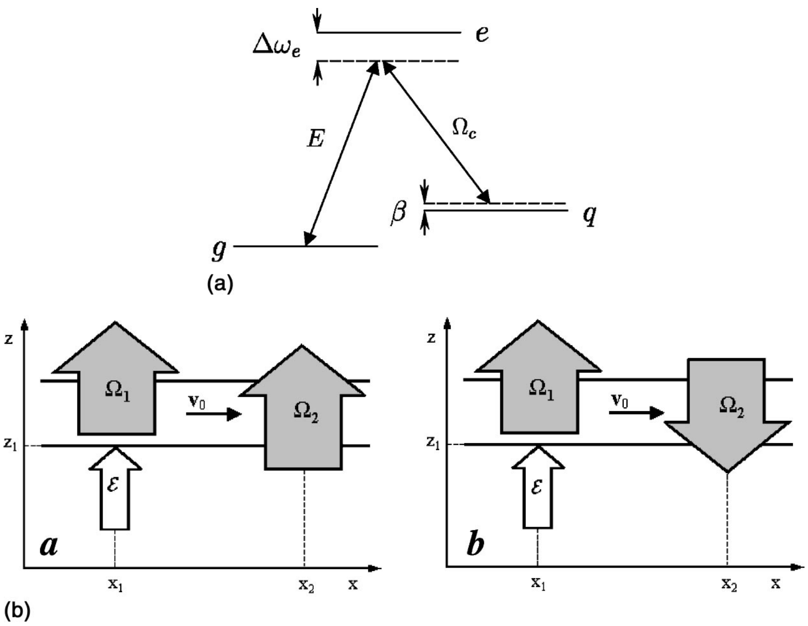

FIG. 1. Top: Atomic level scheme. Bottom: a stream of atoms moving along the x axis with an average speed $\mathbf{v}_0$. Two strong classical control lasers characterized by the Rabi frequencies Ω_1 and Ω_2 are centered at $x=x_1$ and $x=x_2$, respectively. The first control beam propagates along the z-axis. The second beam propagates either parallel (a) or antiparallel (b) to the z axis. A quantum probe field ($\mathcal{E}$) enters the medium at $x=x_1$.

where $g=\mu\sqrt{ck/2\varepsilon_0\hbar}$ characterizes the strength of the radiation-matter coupling, and $\delta=\omega_q+\omega_c-\omega-\omega_g+\Delta\omega_r$ and $\Delta=\omega_e-\omega-\omega_g+\omega_r$ are the two- and single-photon detunings, $v_r=\hbar k/m$ is the atomic recoil velocity due to absorption of a probe photon, $\omega_r=\hbar k^2/2m$ is the recoil frequency, and $\Delta\omega_r=\hbar(k-k_c)^2/2m$. The probe and control beams are assumed to be copropagating. In this case, the overall recoil velocity $\Delta v_r=\hbar(k-k_c)/m$ is small and can be neglected in Eq. (4). Dissipation of the excited state $|e\rangle$ can be included in Eq. (3) replacing Δ by $\Delta-i\gamma$ and adding the appropriate noise operator.

Assume that the probe field $\mathcal{E}$ is sufficiently weak so that one can disregard the depletion of the gound level $|g\rangle$. Neglecting the last term in Eq. (2), one has

$$\frac{\partial}{\partial t}\Phi_g(\mathbf{r},t)=-v_0\frac{\partial}{\partial x}\Phi_g(\mathbf{r},t). \tag{5}$$

It is convenient to introduce the field operators describing annihilation of atomic and spin excitations (excitons)

$$\psi_u\equiv\psi_u(\mathbf{r},t)=n^{-1/2}\Phi_g^\dagger\Phi_u, \tag{6}$$

with $u=e,q$, where $n\equiv n(z)=\langle\Phi_g^\dagger\Phi_g\rangle$ describes the spatial profile of the number density of ground-state atoms. Exploiting relation (5), the equations of motion for the new operators have the form of Eqs. (1), (3), and (4) subject to the replacement $\Phi_e\rightarrow\psi_e$, $\Phi_q\rightarrow\psi_q$ and $\Phi_g\rightarrow n^{1/2}$. Note that the exciton operators ψ_e and ψ_q obey approximately Bose-commutation relations even through the constituent atoms may be fermions.

Let us consider the case of exact two-photon resonance ($\delta=0$). This is well justified as the two-photon Doppler shift due to longitudinal and transversal velocity spreads of the atoms is strongly diminished due to the chosen geometry. Neglecting terms containing ψ_e and $\dot{\psi}_e$ in Eq. (3), one arrives at the adiabatic approximation relating ψ_q to the electric-field amplitude $\mathcal{E}(\mathbf{r},t)$ as

$$\psi_q(\mathbf{r},t)=-gn^{1/2}\tilde{\mathcal{E}}(\mathbf{r},t), \tag{7}$$

with $\tilde{\mathcal{E}}(\mathbf{r},t)=\mathcal{E}(\mathbf{r},t)/\Omega_1(x)=-\psi_q(\mathbf{r},t)/gn^{1/2}$ being an auxiliary field. Assuming a stationary flow of atoms in the x direction, the atomic density $n\equiv n(z)$ does not depend on x, and thus Eqs. (4) and (7) yield

$$\psi_e(\mathbf{r},t)=i\frac{gn^{1/2}}{\Omega_1}\left(v_0\frac{\partial}{\partial x}+\frac{\partial}{\partial t}\right)\tilde{\mathcal{E}}(\mathbf{r},t). \tag{8}$$

On the other hand, the Rabi frequency $\Omega_1\equiv\Omega_1(x)$ is z independent, so Eqs. (1) and (8) lead to the following equation for $\tilde{\mathcal{E}}(\mathbf{r},t)$:

$$\left(\frac{\partial}{\partial t}+\frac{c}{1+n_g}\frac{\partial}{\partial z}+\frac{v_0 n_g}{1+n_g}\frac{\partial}{\partial x}\right)\tilde{\mathcal{E}}(\mathbf{r},t)=0, \tag{9}$$

with $n_g(x,z)=g^2n(z)/\Omega_1^2(x)$ being the group index. Assuming slow propagation, the group index is much larger than unity, so that Eq. (9) simplifies to

$$\left(\frac{\partial}{\partial t}+\tilde{v}_g(z)a(x)\frac{\partial}{\partial z}+v_0\frac{\partial}{\partial x}\right)\tilde{\mathcal{E}}(\mathbf{r},t)=0, \tag{10}$$

where $v_g\equiv\tilde{v}_g(z)a(x)=c/n_g(x,z)$ is the group velocity ($v_g \ll c$). The dimensionless quantity $a(x)\equiv a_1(x) =[\Omega_1(x)/\Omega_1(x_1)]^2$ characterizes the spatial shape of the first control laser centered at $x=x_1=0$. The z dependence of $\tilde{v}_g(z)\equiv v_g(z,x_1)$ emerges through the atomic density $n \equiv n(z)$. It is noteworthy that $\mathcal{E}(\mathbf{r},t)=-\psi_q(\mathbf{r},t)\sqrt{v_g/c}$, so the ratio between the number density of photons and the spin excitations is given by the relative group velocity v_g/c. In other words, the slowly propagating probe beam is made of the EIT (dark state) polaritons [1,5,6] comprising predominantly of the spin excitations.

Assume that the spatial width of the first control beam Δx_{c1} is much larger than that of the incomming probe beam Δx_p centered at $x=x_1=0$. Under this condition, the solution of Eq. (10) can be expressed in terms of the incoming electric field at an entry point $z=z_1$ as

$$\tilde{\mathcal{E}}(\mathbf{r},t)=\frac{\mathcal{E}[\xi(x,z),z_1,\tau(t,x,z)]}{\Omega_1(x_1)}, \tag{11}$$

where $\xi(x,z)=\int_0^x a(x')dx'-v_0\int_{z_1}^z[\tilde{v}_g(z')]^{-1}dz'$, and $\tau(t,x,z)=t+\xi(x,z)/v_0-x/v_0$ obey the proper boundary conditions: $\xi(x,z=z_1)\approx x$ and $\tau(t,x,z=z_1)\approx t$ for the incoming probe field.

Using the definition of $\tilde{\mathcal{E}}$ and Eq. (11), one finds the temporal and spatial behavior of the operators for electric field and spin excitations:

$$\mathcal{E}(\mathbf{r},t)=\sqrt{|a(x)|}\mathcal{E}[\xi(x,z),z_1,\tau(t,x,z)], \tag{12}$$

$$\psi_q(\mathbf{r},t)=-\frac{gn^{1/2}}{\Omega_1(x_1)}\mathcal{E}[\xi(x,z),z_1,\tau(t,x,z)]. \tag{13}$$

Equations (12) and (13) define the electric and spin components of the EIT polariton. A set of trajectories for such a polariton in the x-z plane is given by $\xi(x,z)=\xi_0$, with $\xi_0=0$ corresponding to the central trajectory. Due to the motion of the atoms parallel to the x axis, the polariton is dragged in that direction. Following the spatial profile $a(x)$ of the coupling beam, its velocity component in the z direction, $\tilde{v}_g(z)a(x)$, is further reduced. As soon as $\tilde{v}_g(z)a(x)$ becomes less than v_0, the flow of excitation is predominantly determined by the velocity of the atoms v_0 in the x direction. If the atomic beam is optically thick in the z direction $[\xi(\infty,z_{max})<0]$, the propagation direction of the polariton may be completely converted from $\hat{\mathbf{z}}$ to $\hat{\mathbf{x}}$. The probe field is then stored in the form of a spin excitation that moves with the atoms parallel to the x axis and is centered at $z=z_\infty$, which is a solution of $\xi(x\to\infty,z_\infty)=0$. This is illustrated in Figs. 2 and 3, where we have shown the amplitude of the electric field at different instances of time. Also shown is the effect of the second regeneration laser, which will be discussed later on.

As one can see from the figures, the dragging of the polariton along with the moving atoms leads to a deformation of the excitation. For large values of x, such that $a(x)\to 0$ one has in the vicinity of $z=z_\infty$,

$$\xi(x\to\infty,z)\approx-(z-z_\infty)\frac{v_0}{\tilde{v}_g(z_\infty)}, \tag{14}$$

$$\tau(t,x\to\infty,z)\approx-\frac{z-z_\infty}{\tilde{v}_g(z_\infty)}-\frac{x-x_0}{v_0}, \tag{15}$$

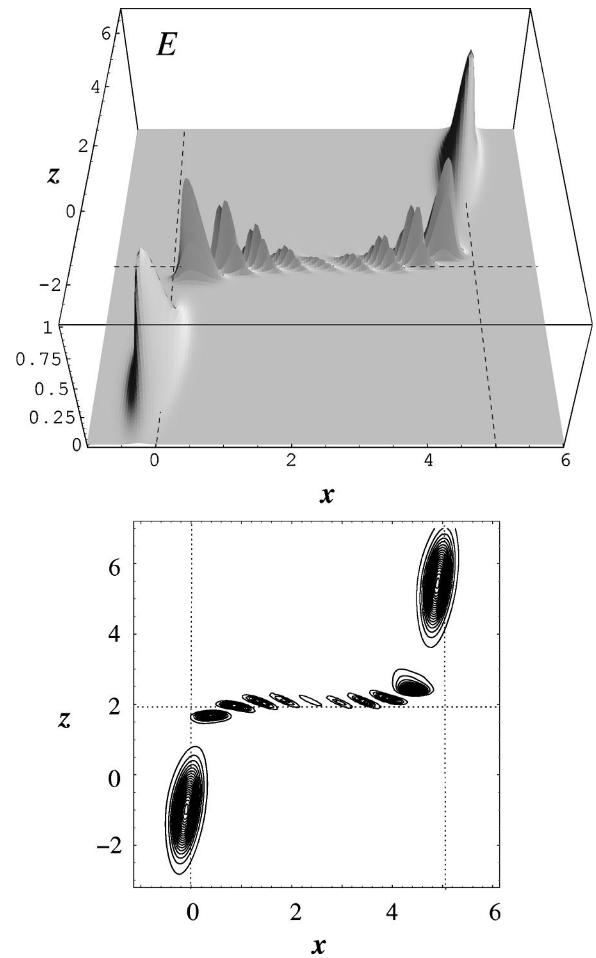

FIG. 2. Propagation of a probe pulse in a moving EIT medium at times $t=0$, 5, 10, 15, 20, 25, 30, 35, 40, 45, 50 illuminated by a pair of control lasers propagating in the $+z$ direction at $x_1=0$ and $x_2=5$. The control lasers have equal amplitudes and Gaussian profiles with unity width. The flow and group velocities are $v_0=0.1$, and $\tilde{v}_g(z)=(1-0.95\exp[-(z-2)^2])$.

where $x_0\equiv x_0(t)=v_0t$ is the "center of mass" of a stored polariton moving in the x direction. Thus, if Δx_p and $2\Delta\tau_p$ denote the half width (in x direction) and the duration of the input probe pulse, the dimensions of the spin wave are in the limit $x\to\infty$,

$$\Delta x_s\approx v_0\Delta\tau_p, \quad \Delta z_s\approx\Delta x_p\frac{\tilde{v}_g(z_\infty)}{v_0}, \tag{16}$$

where we have assumed that $v_0\ll\Delta x_p/\Delta\tau_p$. Since the spatial profile of the polariton in the y direction is not changed, the transfer from a electromagnetic to a matter-wave pulse is associated with a spatial compression of the excitation volume by a factor $\tilde{v}_g(z_\infty)/c$. From Eq. (16), one can also easily obtain a necessary condition for a complete transfer of the electromagnetic excitation to the atomic beam: Since Δz_s should be less than the half width of atoms beam Δz_{atom}, one finds a minimum ratio of v_0 to the group velocity at the center

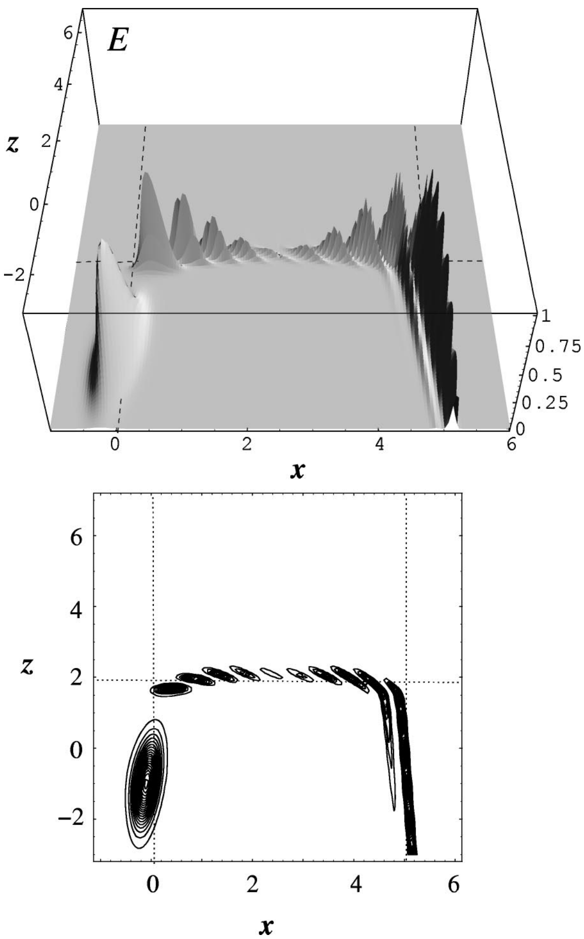

FIG. 3. Propagation of a probe pulse in a moving EIT medium for conditions of Fig. 2 but with a counterpropagating (i.e., in $-z$ direction) second control field.

$$\frac{v_0}{\tilde{v}_g(z_\infty)} \geqslant \frac{\Delta x_p}{\Delta z_{\text{atom}}}. \tag{17}$$

If this condition is fulfilled, a pulse is completely stored within the atom beam and does not exit from the back side.

Consider now the regeneration of the probe beam by a spatially separated, i.e., nonoverlapping, second control laser centered at $x=x_2$, and characterized by a Rabi frequency $\Omega_2(x)$. In this case, the previous Eqs. (10)–(14) describe the propagation of a probe beam within the entire system, subject to the following replacement: $a(x)\rightarrow a(x)=a_1(x) \pm a_2(x)$, where $a_2(x)=[\Omega_2(x)/\Omega_2(x_1)]^2$ characterizes the shape of the second control laser. The upper (lower) sign corresponds to the case where the second control laser propagates in the same (opposite) direction, as compared to the first one [19]. This is because the radiative group velocity of the regenerated probe beam changes sign in the latter case. It is noteworthy that the reversed probe beam experiences a slight shift out of the EIT resonance [5]. Yet such a shift can be neglected in the case where the initial control and probe beams are copropagating. Note also that the previously considered regeneration of a probe beam is due to the temporal switching of the control laser [1–7], whereas now the regeneration is induced by a second spatially separated control laser that can be stationary.

Finally, let us analyze the output field at a spatial point where $z=z_2$. The maximum electric and spin fields are then concentrated at $x=x_2$, for which $\xi(x_2,z_2)=0$. In the vicinity of this point, one has $\xi(x,z_2)\approx \pm a_2(x_2)(x-x_2)$, i.e., the output field has a width $\Delta x_{p2}=\Delta x_{p1}/a_2(x_2)$.

In summary, we have investigated a different scheme of storing and releasing a beam of probe light in a moving atomic medium illuminated by two spatially separated control lasers depicted in Fig. 1. Beyond the area illuminated by the first control laser, the probe beam transforms into a beam of pure spin excitations moving along the x axis, as one can see from Figs. 2 and 3. The regeneration of the probe beam is accomplished by applying the second continuous control laser. Depending n the direction of the latter, the restored probe beam moves either parallel (Fig. 2), or antiparallel (Fig. 3) to the initial probe beam.

In contrast to the previous schemes [1–7], storing and releasing of a probe beam can be now accomplished without switching off and on of a control laser. This is advantageous if one does not know the exact time of arrival of the probe photons, e.g., if the latter are created via spontaneous processes.

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