Storage and retrieval of light pulses in atomic media with "slow" and "fast" light

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We report experimental evidence that light storage, understood as the controlled release of a light pulse by an atomic sample dependent on the past presence of a writing pulse, is not restricted to small-group-velocity media but can also occur in a negative-group-velocity medium. We present a numerical modeling in close agreement with our observations and a simple physical picture applicable to light storage experiments in both "slow" and "fast" light media.

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I. INTRODUCTION

All-optical information processing requires the use of photons as fast and reliable carriers of information. Photons are quantum objects and the search for media in which the quantum state of photons can be preserved and processed is of great significance. Recently, broad attention has focused on the possibility of "light storage" (LS) which is the preservation of the information carried by a light pulse for controllable later release. Such a possibility was suggested in theory [1] and subsequent experimental results were presented in support of this suggestion [2–5]. In all of these experiments, a weak light pulse was "written" into an atomic medium driven by a stronger field and, after a dark interval, retrieved from the medium by turning on the strong (drive) field.

All observations of LS were achieved under conditions of electromagnetically induced transparency (EIT) since, according to [1], EIT and "slow light" (small group velocity associated with EIT) play a key role. The storage effect is seen as a consequence of the slowing and compression of the light pulse in the atomic medium, the propagation of a mixed light-matter excitation (dark-state polariton), the transformation of the dark-state polariton in the absence of light into a pure spin excitation, and finally the release of a light pulse once the drive field is turned on [6].

The purpose of this paper is to present experimental results and theoretical considerations which broaden the scope of the subject by demonstrating that a LS effect, analogous to that previously observed [2–5], can take place in media where EIT does not occur and where the probe pulse group velocity is negative ("fast light") as a result of large anomalous dispersion.

Steep anomalous dispersion exists in driven atomic media in connection with electromagnetically induced absorption (EIA) [7]. EIA occurs when resonant light interacts with a two-level atomic transition in which the Zeeman degeneracy of the excited level is higher than that of the lower level; namely, $F_e > F_g > 0$ where F_e and F_g are the total angular momenta of the excited level and ground level, respectively. A resonant increase in the probe absorption occurs under the condition of a Raman resonance with ground-state Zeeman sublevels [8]. In particular, at zero magnetic field the EIA resonance condition is achieved for the two orthogonal polarization components of a single monochromatic optical field. Superluminar pulse propagation in an atomic vapor under the conditions of EIA has been demonstrated in [9].

II. EXPERIMENT

The experimental scheme used is similar to the one presented in Ref. [2] (Fig. 1). LS was studied in a 5-cm-long vapor cell containing a natural isotopic mixture of rubidium. The cell was placed at the center of a cylindrical μ-metal shield. A solenoid inside the magnetic shield allows tuning of the longitudinal magnetic field B. Two extended cavity diode
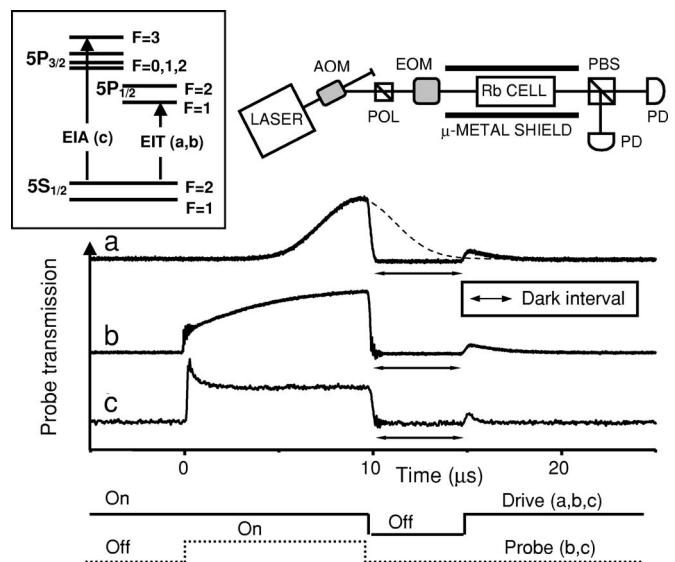

FIG. 1. Upper: Level scheme and experimental setup. Center: observed signals for the probe field intensity transmission. Trace *a* transition $5S_{1/2}(F=2) \rightarrow 5P_{1/2}(F'=1)$ of [87]Rb with Gaussian probe pulse; dashed, Gaussian envelope of transmitted probe pulse without dark interval. Trace *b*: Same transition; square probe pulse. Trace *c*: Transition $5S_{1/2}(F=2) \rightarrow 5P_{3/2}(F'=3)$ of [87]Rb; square probe pulse. Bottom: timing sequence for the drive and probe fields.

lasers were alternatively used for the D_1 and D_2 transitions. Fast switching on and off of the laser light was achieved with an acousto-optic modulator (AOM). After the AOM, a polarizer fixes the light polarization. An electro-optic modulator (EOM) was used to produce a temporal rotation of the light polarization and consequently to generate a probe pulse with orthogonal polarization relative to the drive field. Care was taken to use the EOM in the $\lambda/2$ configuration to produce in-phase probe and drive fields. The probe pulse intensity was five times weaker than the drive component. The light beam was expanded to a 1 cm diameter before the cell. The maximum laser intensities at the cell were 0.6 and 2 mW/cm^2 for the D_1 and D_2 lines, respectively. The cell was heated (~70 °C) to produce almost 100% linear absorption and 50–80 % absorption at maximum light intensity. After the atomic vapor, the drive and probe polarization components of the light were separated by a polarizing beam splitter and collected by fast photodiodes. The electronic control and detection response times were shorter than 1 μs.

Our setup allows the use of linear perpendicular or opposite circular drive and probe field polarization combinations. Qualitatively similar LS effects were observed for both choices. We describe here the signals obtained with perpendicular linear drive and probe field polarizations. LS under conditions of EIT was observed on both the transitions $5S_{1/2}(F=1)\rightarrow5P_{3/2}$ (D_2 line) and $5S_{1/2}(F=2)\rightarrow5P_{1/2}(F'=1)$ (D_1 line) of ^{87}Rb. We present results obtained for the latter transition which was also used in previous experiments [2,4,5]. The width of the EIT resonance, measured at maximum light intensity by scanning the magnetic field, was $\Delta B=46$ mG (corresponding to 65 kHz). Trace *a* of Fig. 1, obtained with a Gaussian-shape probe pulse, reproduces the essential features of the previous LS experiments [2–5]. Since the inverse of the pulse duration is small compared to the EIT resonance width, the Gaussian pulse shape is well preserved during propagation. A pulse delay (relative to vacuum propagation) of approximately 1 μs is observed corresponding to a group velocity $v_g\sim5\times10^4$ m/s. We subsequently used square probe pulses to simplify further comparison with numerical modeling. Trace *b* was obtained for the same transition as for trace *a* with a square probe pulse. The transmitted probe pulse exhibits a characteristic distortion, indicative of a slow light medium [9]. In both traces the retrieved pulse has an exponential decay with approximately the same decay time (~2 μs). An exponential decay time of 8 ± 2 μs was measured for the retrieved pulse amplitude (or area) as a function of the dark interval, in agreement with the transverse atomic time-of-flight estimate.

We now turn our attention to a fast light medium which can be achieved for the D_2 line in connection with EIA. Trace *c* of Fig. 1 was obtained with the laser tuned to the $5S_{1/2}(F=2)\rightarrow5P_{3/2}(F'=3)$ transition of the D_2 line. We notice a characteristic probe pulse distortion with a leading edge less absorbed than the rest of the pulse. Such distortion is indicative of the pulse advancement expected in a fast light medium [9]. We observe that in the fast light medium the retrieved pulse has similar characteristics to the

pulse retrieved from the slow light medium. In either case, in the absence of a magnetic field, the retrieved pulse has an exponentially decaying slope with a time constant dependent on the drive field intensity. In the presence of a nonzero magnetic field similar damped oscillations are seen on the retrieved pulse envelope for the two types of transitions. The same dependence of the retrieved pulse amplitude (and area) on the dark interval duration is observed in the two cases.

III. NUMERICAL MODELING

The observations described above can be theoretically modeled with the help of the Bloch equations for the atom-field interaction [10]. Following the notation in [11,12], we consider atoms at rest with a ground level g and an excited level e and angular momenta F_g and F_e, respectively, and energy separation $\hbar\omega_0$. Spontaneous emission from e to g occurs at a rate Γ. The finite interaction time is accounted for by the ground-level relaxation rate γ ($\gamma\ll\Gamma$). The atoms are subjected to a magnetic field B and a classical monochromatic electromagnetic field $\vec{E}_I(t)=(E_1\hat{e}_1+E_2\hat{e}_2)\exp(i\omega t)$, where $\hat{e}_1$ and $\hat{e}_2$ are two orthogonal unit polarization vectors (generally complex) and E_1 and E_2 are the corresponding field amplitudes.

Introducing the slowly varying matrix $\sigma=P_g\rho P_g+P_e\rho P_e+P_g\rho P_e\exp(-i\omega t)+P_e\rho P_g\exp(i\omega t)$ (where ρ is the density matrix in the Schrödinger representation and P_g and P_e are projectors on the ground and excited subspaces, respectively), the time evolution of the system (in the rotating-wave approximation) may be expressed as

$$\frac{d\sigma}{dt}=-\frac{i}{\hbar}[H_Z+\hbar\,\Delta P_e+V,\sigma]-\frac{\Gamma}{2}\{P_e,\sigma\}$$
$$+b\Gamma(2F_e+1)\sum_{q=-1,0,1}Q_{ge}^q\sigma Q_{eg}^q-\gamma(\sigma-\sigma_0)\quad(1)$$

where $H_Z=(\beta_g P_g+\beta_e P_e)F_z B$ is the Zeeman Hamiltonian (β_g and β_e are the ground- and excited-state gyromagnetic factors and F_z is the total angular momentum projection operator along the magnetic field); $Q_{ge}^q=Q_{eg}^{q\dagger}(q=-1,0,1)$ are the standard components of the vectorial operator defined by $\vec{Q}_{ge}=\vec{D}_{ge}\langle g\|\vec{D}\|e\rangle^{-1}$ where $\vec{D}_{ge}\equiv P_g\vec{D}P_e$ and $\langle g\|\vec{D}\|e\rangle$ is the reduced matrix element of the dipole operator between g and e; and $\Delta\equiv\omega_0-\omega$ is the optical field detuning. V is the atom-field coupling Hamiltonian: $V=\frac{1}{2}(\Omega_1\hat{e}_1+\Omega_2\hat{e}_2)\cdot\vec{Q}_{ge}+$H.c., where Ω_j ($j=1,2$) is the reduced Rabi frequency: $\Omega_j=E_j\langle g\|\vec{D}\|e\rangle\hbar^{-1}$. The coefficient b ($0\leq b\leq1$) determines the branching ratio for spontaneous transitions from the excited to the ground level ($b=1$ for a closed transition). The term $\gamma\sigma_0$ represents a constant pumping rate of fresh atoms in the isotropic state $\sigma_0=P_g/(2F_g+1)$. For a given solution of Eq. (1), the transmitted field with polarization $\hat{e}_j$ is given to the lowest order in the atomic density by [13]
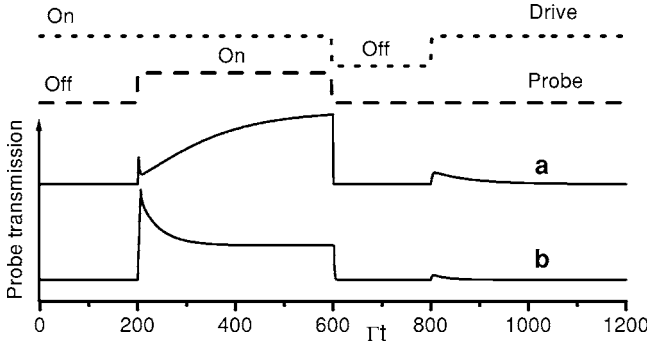
FIG. 2. Dotted (dashed): drive (probe) time sequence. Trace a: Calculated transmitted probe field intensity for the transition $5S_{1/2}(F=2) \rightarrow 5P_{1/2}(F'=1)$ of ^{87}Rb. Trace b: Same for the transition $5S_{1/2}(F=2) \rightarrow 5P_{3/2}(F'=3)$. Parameters (see [14]): $\Omega_D/\Gamma=0.3$ (trace a), $\Omega_D/\Gamma=0.6$ (trace b), $\Omega_D/\Omega_P=2$, $\gamma/\Gamma=10^{-3}$, $\alpha=8$, $B=0$ (Ω_D and Ω_P are the drive and probe field Rabi frequencies).

$$E_{Tj}(t) = E_{Ij} + i \frac{\alpha \hbar \Gamma}{\langle g \| \vec{D} \| e \rangle} \mathrm{Tr}[\sigma_{ge}(\hat{e}_j^* \cdot \vec{Q}_{eg})] \qquad (2)$$

where $\sigma_{eg} = \sigma_{ge}^\dagger = P_e \sigma(t) P_g$ and α is a real positive dimensionless constant (proportional to the atomic density) that is adjusted to fit the observed field attenuation.

For simplicity, we have chosen to simulate the atomic response for excitation sequences where the two field components along $\hat{e}_1$ and $\hat{e}_2$ have stepwise variations. In this case, the time can be divided into intervals during which the amplitude and the polarization of the total incident field $\vec{E}_I(t)$ are constant. Equation (1) then represents a system of coupled first-order linear differential equations with constant coefficients for σ. Using the Liouville method the matrix elements of σ can be organized into a vector $\mathbf{y}$ and Eq. (1) rewritten in the form $d\mathbf{y}/dt = M\mathbf{y} + \mathbf{p}_0$ where $M(e, \Delta, \Omega, \Gamma, \gamma, B)$ is a matrix and $\mathbf{p}_0$ a constant vector corresponding to the pumping term $\gamma \sigma_0$.

The solution for $\mathbf{y}(t)$ is found following a standard procedure. First the eigenvalues Λ_n and eigenfunctions $\mathbf{v}_n$ of M are numerically calculated and then $\mathbf{y}(t)$ is found as $\mathbf{y}(t) = \mathbf{y}_0 + \Sigma a_n \mathbf{v}_n e^{\Lambda_n t}$ where a_n are constants adjusted to the initial conditions and $\mathbf{y}_0$ is the steady-state solution. We note that for $B=0$ and $\Delta=0$ all the Λ_n are real and negative which implies exponential decay to the steady state dominated by the Λ_n with the smallest absolute value. Once $\mathbf{y}(t)$ has been determined the transmitted field for a given polarization component is calculated using Eq. (2).

The pulse sequence used for the calculation is shown in the upper traces of Fig. 2. Trace a corresponds to the calculated transmission (at the probe field polarization) for parameters corresponding to the transition $5S_{1/2}(F=2) \rightarrow 5P_{1/2}(F'=1)$ of ^{87}Rb. Trace b was calculated for the transition $5S_{1/2}(F=2) \rightarrow 5P_{3/2}(F'=3)$ of ^{87}Rb. The pulse distortions observed in the experiment are well reproduced and a retrieved pulse is obtained for both transitions after the dark interval. The amplitude of the retrieved pulse depends on the probe pulse amplitude and duration and decays exponentially with a decay time depending on the drive field intensity alone.
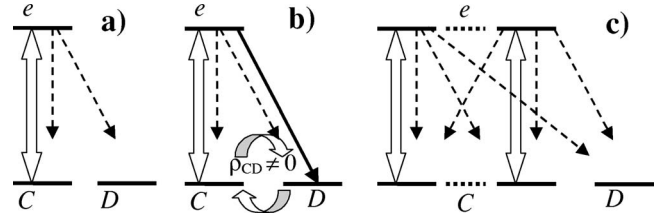

FIG. 3. Basic level schemes for the light-atom interaction. White arrow, applied field; dashed arrow, spontaneous emission; solid arrow, stimulated Raman emission; curved arrow, ground-state coherence. (a) Incoherent optical pumping, (b) optical pumping with coherence, and (c) generalization of (a) for two levels with Zeeman degeneracy.

Similar results to those presented in Fig. 2 are obtained if orthogonal circular polarizations are considered for both fields.

IV. PHYSICAL PICTURE

In the previous sections we have presented experimental data illustrating the storage and retrieval of light pulses in both slow and fast light media. As shown in the previous section, these observations, which are in many respects similar to previously reported light storage experimental results, are fully consistent with a simple theoretical model based on the standard Bloch equations for an optically thin medium. The aim of this section is to present a simple picture that provides a physical insight into the underlying common mechanism giving rise to light storage in both types of media.

Consider the basic light-atom interaction process presented in Fig. 3(a) where an optical field couples the atomic ground state C (coupled state) to the excited state e while the second ground state D (dark state) is unaffected by the field. After excitation in state e the atom can decay spontaneously to both lower states. This is an optical pumping process [15]. If the atoms interact with the light for a sufficiently long time and if states C and D are stationary, the atoms will eventually end in state D and the medium will become transparent. During this process, transitions from state e to state D occur spontaneously but may also take place through stimulated Raman emission. The condition for stimulated Raman emission is the existence of nonzero coherence between states C and D, i.e., $\rho_{DC} \neq 0$, where ρ_{DC} is the element of the density matrix between states D and C. Under such a condition the interaction with the applied field results in optical coherence between states e and D ($\rho_{eD} \neq 0$) which in turn radiates a field for this transition coherent with the initial field. As a general rule, a light field interacting with a coherently prepared atom in the ground state will induce coherent emission of light in all allowed optical transitions as long as the (ground-state) atomic coherence is preserved [Fig. 3(b)].

Direct generalization of the situation represented in Fig. 3(a) to the more complex case of a two-level system with arbitrary Zeeman degeneracy interacting with an optical field is possible. In general, given the field polarization, after a suitable basis transformation, one can identify several (C)

states in the ground level, each coupled to a corresponding excited state sublevel, and two, one or zero dark (*D*) states [Fig. 3(c)]. The last case occurs when the Zeeman degeneracy of the excited level is higher than that of the ground level [16,17]. In all cases the general rule stated above applies.

Two types of systems have commonly been considered for the experimental study of coherent light-atom interaction dynamics. In the first class one has the Hanle-type experiments [18–22] where a single optical field with well-defined polarization interacts with two atomic levels with Zeeman sublevels whose energy is tuned with a static magnetic field. If the light polarization is linear it is straightforward that a scheme similar to that of Fig. 3(c) results by choosing the quantization axis along the direction of the light polarization. Even in the general case of arbitrary elliptical light polarization the scheme of Fig. 3(c) can be applied after a suitable transformation of the state basis that depends on the incident light polarization [17]. An immediate consequence of this description is that in the presence of ground-state coherence the system will respond by the coherent emission of light with a polarization orthogonal to that of the incident field. Such a process was recently discussed in terms of the polarization change that the light experiences while propagating in an atomic medium with light-induced anisotropy and an alternative explanation of the LS effect in such terms was proposed [23].

A second class of experiments for which the picture presented in Figs. 3(a) and 3(b) can be conveniently applied concerns a system of an excited level and two ground levels with energy separation Δ driven by two optical fields of frequencies ω_1 and ω_2. This system (hereafter designated a Λ system), has been extensively studied theoretically and experimentally in connection with EIT. Since in a Λ system each field interacts with a different transition, it is possible after a time-dependent unitary transformation to describe the atomic dynamics by a time-independent Hamiltonian and to identify a dark state *D* and a coupled state *C* [24]. This "coupled-uncoupled" description exactly corresponds to the situation of Figs. 3(a) and 3(b). Both levels *C* and *D* are linear combinations of the two lower levels that explicitly depend on the amplitudes and phases of the applied fields. Here again, if coherence is present between levels *C* and *D*, the atomic medium will react to the optical excitation by the coherent emission of fields at frequencies ω_1 and ω_2. These fields are "orthogonal" to the applied fields in the sense that they *together* only couple to state *D* (while the incident fields only couple to *C*). As an example we consider the situation corresponding to the retrieval process in the LS experiment reported in [3]. After the dark interval, the drive field (ω_1) alone excites the Λ system. Since (previously created) coherence is present between the ground levels (that coincide with *C* and *D* in this case), a field of frequency $\omega_2 = \omega_1 - \Delta$ is emitted by the medium. In the general case, in the presence of coherence between *C* and *D*, two fields will be emitted with frequencies ω_1 and ω_2 that will interfere with the incident fields.

The previous discussion leads to a very general qualitative understanding of the transient behavior of coherently driven atomic systems. For given excitation conditions, after a long enough time, the system will reach a steady state. If one or more dark states exists the steady state of the system will be the dark state(s) *D*. If no dark state exists the steady state will generally be a statistical mixture (diagonal density matrix) of states *C*. A rapid modification of the excitation conditions, for instance a light polarization change in a Hanle experiment or a change in the amplitude or phase of the fields in a Λ scheme, will determine a new set (C', D') of coupled and dark states. The change in the excitation conditions needs to be rapid with respect to the optical pumping time but may otherwise be slow with respect to other characteristic times. Since the previous state of the system is generally a linear combination of the new C', D' states, coherence among the latter states exists and results in the transient emission of an orthogonal field. Such emission will last as long as the coherence between C' and D' survives. The decay of the transient emission will be purely exponential with a time constant given by the optical pumping time between states C' and D' if those states are stationary. If they are not stationary (nonzero magnetic field for a Hanle experiment or nonzero Raman detuning in a Λ system), the decay will show damped oscillations [14]. If the two different excitation conditions are separated in time by a dark interval then the transient emission of an orthogonal field will occur after the dark interval provided that this interval is not long compared to the ground-state coherence lifetime. This physical description of the light storage process is fully compatible with the numerical modeling presented in Sec. III. In such a calculation, the atoms are submitted to constant excitation conditions during finite time intervals. During each interval, the physical description above directly applies. The physical picture presented here can be applied to the LS experiments previously reported [2–5].

V. DISCUSSION

The experimental results and the theoretical considerations presented above show that EIT and slow light propagation are not essential requirements for the retrieval of a light pulse from a previously prepared coherent atomic medium. In fact, we have shown that the retrieval of a light pulse from a previously prepared atomic medium will occur, under very general conditions (including EIT or EIA), as a consequence of the transient response of the coherently prepared atomic system to new optical excitation conditions. Such a storage-retrieval process does not require the probe pulse to be spatially contained during propagation inside the atomic medium. As shown, the light storage process can occur in an optically thin medium where only a fraction of the writing pulse is present at a given time. Under such circumstances, the light retrieval transient appears as a consequence of the irreversible relaxation of the system toward a new steady state. In this context, only exponential decaying pulses can be retrieved. This prevents the recovery of full information about the state of the writing probe pulse. For instance, information on the shape of the probe pulse envelope is lost. Nevertheless, such simple light pulse storage under EIT or EIA conditions may result in interesting applications for optical delaying and buffering [25].

In view of the previously reported light storage experiments [2–5] and the results presented here, we must conclude that further theoretical and experimental work is needed for the practical realization of *information-preserving* storage in an atomic medium. Such reversible information storage would, in principle, be only possible under conditions of EIT since dissipation (absorption) is unavoidable for EIA. In addition, a necessary practical requirement is the realization of an atomic medium in which more than one probe pulse (more than one bit of information) can be spatially contained in the atomic sample and propagate without significant distortion [25,26].

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