# Storage and retrieval of light pulses at moderate powers

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(Received 17 September 2002; published 26 March 2003)

We investigate whether it is possible to store and retrieve an intense probe pulse using a medium that can be modeled as a set of atoms with the relevant energy levels in $\Lambda$ configuration. We demonstrate that it is indeed possible to store and retrieve the probe pulses that are not necessarily weak. We find that the retrieved pulse remains a replica of the original pulse, although there is overall broadening and loss of the intensity. The loss of intensity can be understood in terms of the dependence of absorption on the intensity of the probe. Our calculations include the dynamics of the control field, which becomes especially important as the intensity of the probe pulse increases. We use the adiabatic theory of Grobe *et al.* [Phys. Rev. Lett. **73**, 3183 (1994)] to understand our numerical results on the storage and retrieval of light pulses at moderate powers.

DOI: 10.1103/PhysRevA.67.033813 PACS number(s): 42.50.Gy, 32.80.Qk

## I. INTRODUCTION

Electromagnetically induced transparency (EIT) [1] continues to lead to newer applications in optical physics. It was shown how EIT can lead to ultraslow light [2–6]. There are proposals where EIT and other ideas are used to produce light with zero group velocity [7–15]. Fleischhauer and Lukin [11] introduced the idea of a dark state polariton which is a combination of atomic coherence and the probe pulse. The relative weightage of the dark state polariton depends on the temporal shape of the control field. On switching off the control field, the polariton prominently has atomic character. On switching on the control field, the polariton starts acquiring the field character that eventually dominates. The storage and retrieval of weak pulses thus become possible. Matsko *et al.* [14] extended this analysis to consider the storage and retrieval for nonadiabatic switching of the control field. Further generalizations to include the detunings of the probe and pump fields have appeared [15]. Previous works [9,11–15] on this problem have been carried out in the linear regime, i.e., where the probe pulse is much weaker than the control field. The natural question then is: whether it is still possible to store and retrieve an intense probe pulse. Here, the nonlinearity of the medium with respect to the probe pulse amplitude becomes quite significant and this might lead to the distortion in the pulses. In this paper, we address this question by characterizing the probe pulse propagation for different types of pulses in a medium that can be modeled as an ensemble of atoms with the relevant energy levels in $\Lambda$ configuration. The paper is organized as follows: In Sec. II, we present the Maxwell-Bloch equations governing the propagation dynamics of optical pulses at moderate power. In Sec. III, we present the results obtained by numerical solutions of the relevant Maxwell-Bloch equations. We include the radiative decay of the atoms as well as the dynamical evolution of the control field. In Sec. IV, we describe how the adiabatic theory of Grobe *et al.* [17] can be used for understanding the storage and retrieval of probe pulses at moderate power.

## II. DYNAMICS OF PULSE PROPAGATION

Consider the propagation of a pulse (called probe) defined by the electric field

$$\vec{E}_p(z,t)=\vec{\mathcal{E}}_p(z,t)\ e^{-i(\omega_1 t-k_1 z)}\ +\ \text{c.c.} \tag{1}$$

Here $\vec{\mathcal{E}}_p$ is the slowly varying envelope of the probe field. We assume that a pump pulse defined by

$$\vec{E}_c(z,t)=\vec{\mathcal{E}}_c(z,t)\ e^{-i(\omega_2 t-k_2 z)}\ +\ \text{c.c.}, \tag{2}$$

will act as a control on the propagation of the probe pulse. We assume a medium whose atomic transitions are as shown in Fig. 1. The probe pulse is tuned close to the transition $|1\rangle\leftrightarrow|3\rangle$. The control field is tuned to the transition $|1\rangle\leftrightarrow|2\rangle$. We further assume that the states $|2\rangle$ and $|3\rangle$ are metastable states. We work with the density-matrix equations for the medium as we fully take into account the radiative decay of the excited state. For the case when the fields are in resonance with their respective transitions, the density-matrix equations are

$$\begin{aligned}
\dot{\rho}_{11}&=-2(\gamma_1+\gamma_2)\rho_{11}+iG\rho_{21}+ig\rho_{31}-iG^*\rho_{12}-ig^*\rho_{13},\\
\dot{\rho}_{22}&=2\gamma_2\rho_{11}+iG^*\rho_{12}-iG\rho_{21},\\
\dot{\rho}_{12}&=-[\gamma_1+\gamma_2]\rho_{12}+iG\rho_{22}+ig\rho_{32}-iG\rho_{11},\\
\dot{\rho}_{13}&=-[\gamma_1+\gamma_2]\rho_{13}+iG\rho_{23}+ig\rho_{33}-ig\rho_{11},\\
\dot{\rho}_{23}&=iG^*\rho_{13}-ig\rho_{21}.
\end{aligned} \tag{3}$$

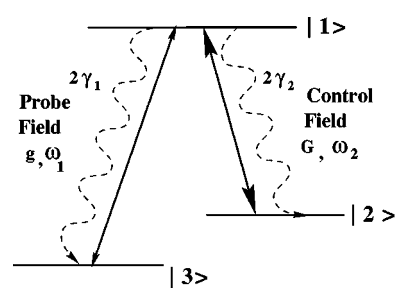


FIG. 1. Three-level $\Lambda$-type medium resonantly coupled to a control field with Rabi frequency $2G$ and probe field $2g$.

 

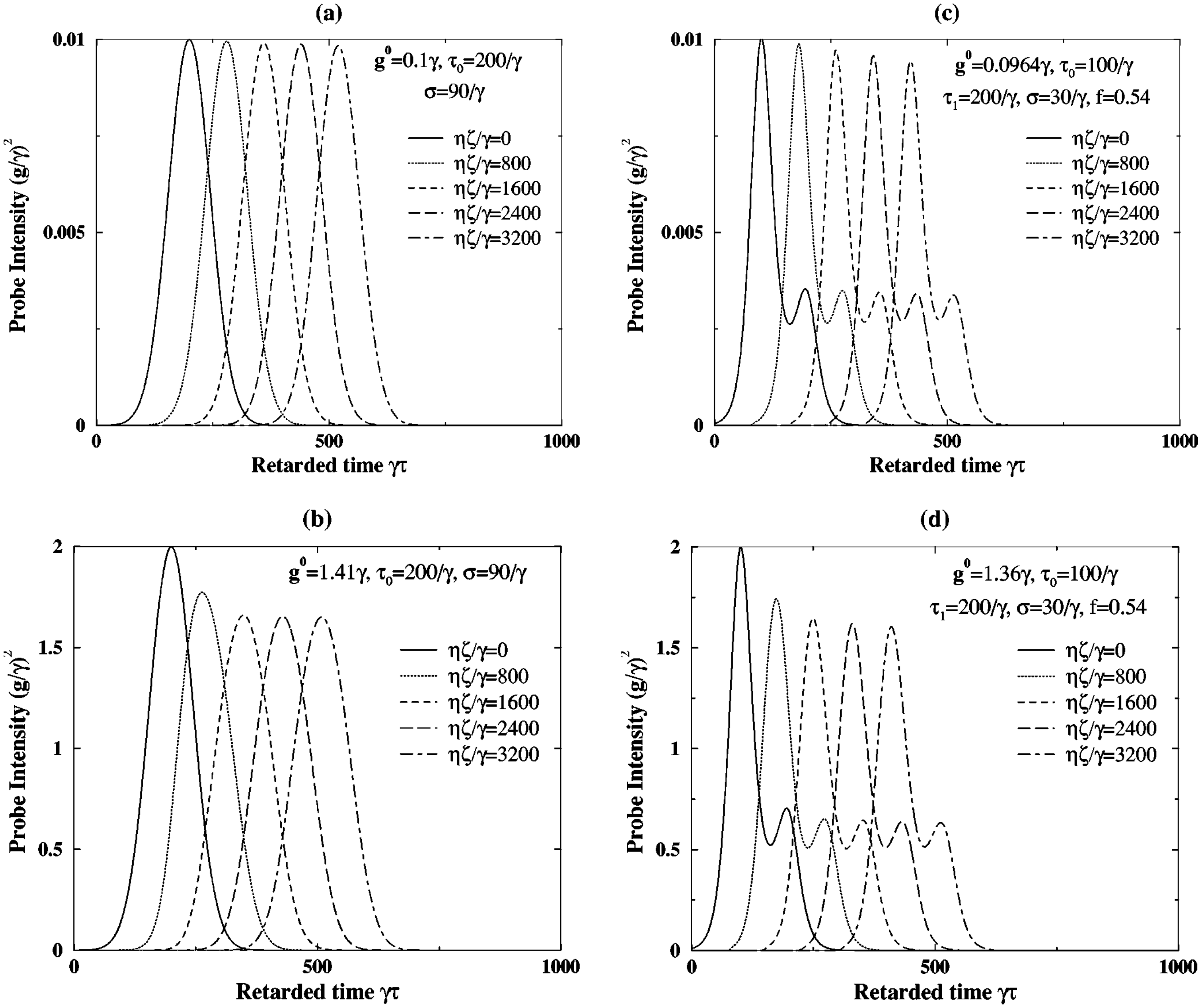


FIG. 2. The probe field intensity in the medium is plotted against retarded time at different propagation distances within the medium. (a) and (c) show the probe pulse propagation with nondiminishing amplitude, for small intensities. (b) and (d) depict the broadening and loss of intensity in the case of an intense probe pulse. In all the cases, the control field is taken as a cw with $G=3.16\gamma$.

Here $2\gamma_1$ ($2\gamma_2$) is the rate of decay of the state $|1\rangle$ to $|3\rangle(|2\rangle)$. The Rabi frequencies $2g$ and $2G$ are defined by

$$2g=\frac{2\vec{d}_{13}\cdot\vec{\mathcal{E}}_{\rm p}}{\hbar}, \tag{4}$$

$$2G=\frac{2\vec{d}_{12}\cdot\vec{\mathcal{E}}_{\rm c}}{\hbar}, \tag{5}$$

where $\vec{d}_{ij}$ is the dipole-matrix element. The induced polarization, say at the frequency $\omega_1$, is given in terms of the density-matrix element $\rho_{13}$:

$$\vec{P}=N\vec{d}_{13}\rho_{13}, \tag{6}$$

where $N$ is the density of the medium. It should be borne in mind that $g$ and $G$ are dependent on both space and time coordinates. Therefore, the polarization $\vec{P}$ is also a slowly varying function of both space and time coordinates. We can now obtain the equations for the propagation of both probe and control pulses. We work in the slowly varying envelope approximation to obtain the temporal and spatial evolution of the pulses. On converting the equations for the Rabi frequencies, we obtain

$$\frac{\partial g}{\partial z}+\frac{\partial g}{\partial ct}=i\eta\rho_{13},$$

$$\frac{\partial G}{\partial z}+\frac{\partial G}{\partial ct}=i\eta\rho_{12}, \tag{7}$$

where $\eta$ is the coupling constant

$$\eta=3\lambda^2N\gamma/8\pi, \tag{8}$$

and the wavelength of the atomic transitions $\lambda_{13}\simeq\lambda_{12}\simeq\lambda$.

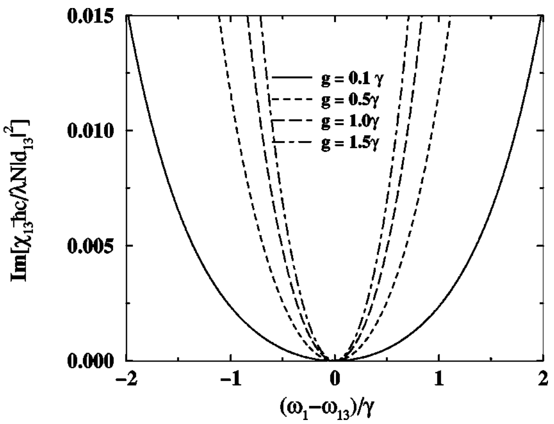


FIG. 3. Imaginary part of susceptibility $[\chi_{13}\hbar c/N\lambda|d_{13}|^2]$ as a function of probe frequency $\omega_1$ in the presence of a cw control field with $G=3.16\gamma$. The width of the transparency window decreases with increase in the intensity of the probe field.

For simplicity we will assume the same coupling constant for the two transitions. The solutions of Eqs. (3) and (7) give the complete evolution of the atom-field system. The analytical solutions of the Maxwell-Bloch equations are not known except under very special conditions [16]. Therefore, we study the pulse-propagation problem numerically.

## III. NUMERICAL SIMULATIONS

### A. Pulse propagation: Effect of nonlinearities

We solve the propagation problem numerically for a homogeneously broadened gas of cold atoms. We will use the traveling coordinates: $\tau=t-z/c$ and $\zeta=z$. We consider propagation of the probe pulse with two different shapes, namely, a Gaussian pulse

$$g(0,\tau)=g^0 e^{-[(\tau-\tau_0)/\sigma]^2} \tag{9}$$

and a combination of Sech pulses

$$g^0\left[\operatorname{sech}\left(\frac{\tau-\tau_0}{\sigma}\right)+f\operatorname{sech}\left(\frac{\tau-\tau_1}{\sigma}\right)\right]. \tag{10}$$

Here, $g^0$ is a real constant characterizing the peak amplitude of the Rabi frequency before the pulse enters the homogeneous atomic medium, $\sigma$ is the temporal width of the input pulse and $\tau_i(i=0,1)$ gives the location of the peaks. All the atoms are initially in the state $|3\rangle$ and thus $\rho_{33}(\zeta,0)=1$ with all other density-matrix elements equal to zero.

In order to appreciate the effect of nonlinearities due to the moderate power of the probe pulse, we first consider the control field as a continuous wave (cw): $G(0,\gamma\tau)=\text{constant}[(G/\gamma)^2=10]$. We work under the condition of EIT, i.e., $\omega_1-\omega_2=\omega_{13}-\omega_{12}$. Figures 2(a) and 2(b) [2(c) and 2(d)] display the propagation of a Gaussian pulse (Sech combination pulse) through the medium. From Figs. 2(a) and 2(c), we see that the weak probe pulse propagates without any significant absorption and broadening inside the medium. Figures 2(b) and 2(d) show that the intense probe pulse suffers absorption and broadening. This is one of our key results. The shape of the pulse remains almost identical to the input pulse. This behavior of the intense probe pulse can be explained using the form of steady-state probe absorption spectra. In Fig. 3, we show the behavior of the probe absorption as a function of the probe detuning when the control field is on resonance. It is clear from Fig. 3 that an increase of the probe field intensity results in the increased absorption of the probe for a given frequency in the neighborhood of the frequency satisfying the two-photon resonance condition. Note that the width of the transparency window depends on the intensities of the control and probe fields. For a fixed intensity of the control field, the width of the transparency window becomes smaller when the probe field intensity is increased. Note that the condition for distortionless pulse propagation is that the spectral width of the probe pulse should be contained within the transparency window of the medium. If the pulse becomes too short, or its spectrum too broad relative to the transparency window of the medium, absorption and also the higher order dispersion need to be taken into account. The dispersion of the medium also causes the distortion of the probe pulse.

### B. Storage and retrieval of electromagnetic fields at moderate powers

Fleischhauer and Lukin [11] showed that it is possible to store and retrieve weak pulses of the electromagnetic radiation by using atomic coherences. They demonstrated how a control pulse and a probe pulse create atomic coherence and that by slowly switching off the control pulse, the probe pulse disappears and gets stored in the form of atomic coherence. Switching on the control field can retrieve the stored probe pulse. The smooth switch off and on of the control field is made possible by gradually varying the intensity of the control field with respect to time. Therefore, the switching off and on of the control field can be modeled by a super-Gaussian shape given by

$$G(0,\tau)=G^0[1-e^{-[(\tau-\tau_2)/\sigma']^{\alpha}}], \tag{11}$$

where the parameter $\alpha$ determines how the pulse is switched on. For $\alpha=4$ (100), we will have adiabatic (nonadiabatic) switching. Figure 4(a) shows the adiabatic switching of the control field. Switching off of the control field can give rise to the absorption of the probe pulse when the entire probe pulse is inside the medium. The group velocity of the probe pulse is reduced to zero and its propagation is stopped by switching off the control field. The stored probe pulse can be retrieved by switching on the control field. The time difference between switching off and on is dependent on the lifetime of the atomic coherence between the states $|2\rangle$ and $|3\rangle$. As seen from Figs. 4(b) and 4(d), for weaker probe pulse, the shape of the retrieved pulse is the same as the original one because the width of the probe pulse spectrum is very much less than the width of the EIT window. Therefore, almost perfect storage and retrieval of light is possible by adiabatic switching of the control field as pointed out by Fleischhauer

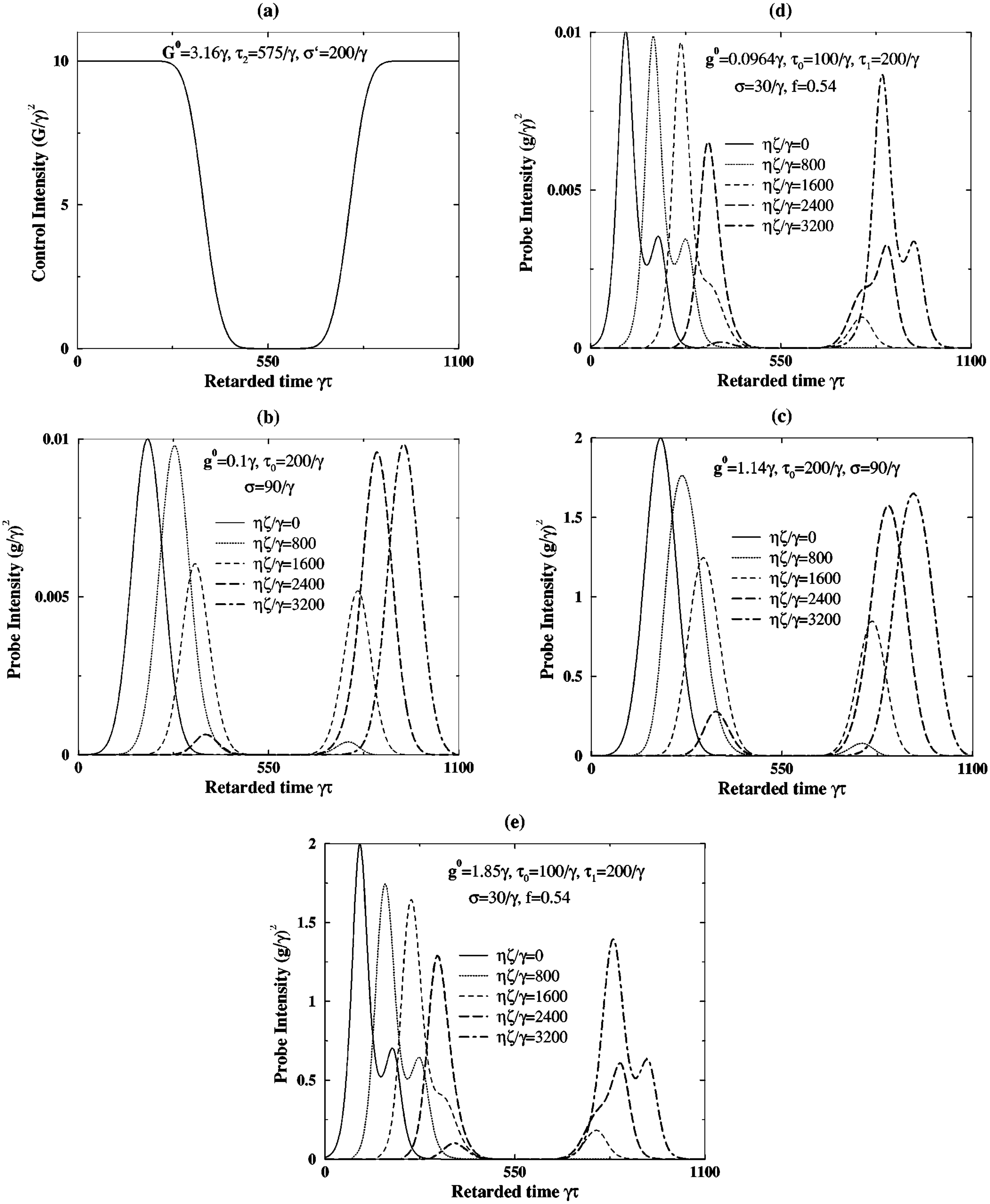


FIG. 4. (a) Shows the intensity of the control field as a function of retarded time at the entry surface of the medium at $\zeta = 0$. Switching mode of the super-Gaussian control field is adiabatic. The frames (b) and (d) show the time evolution of the weaker probe pulse at different propagation distances; and the frames (c) and (e) depict the temporal profile of the intense probe pulse at different propagation distances.

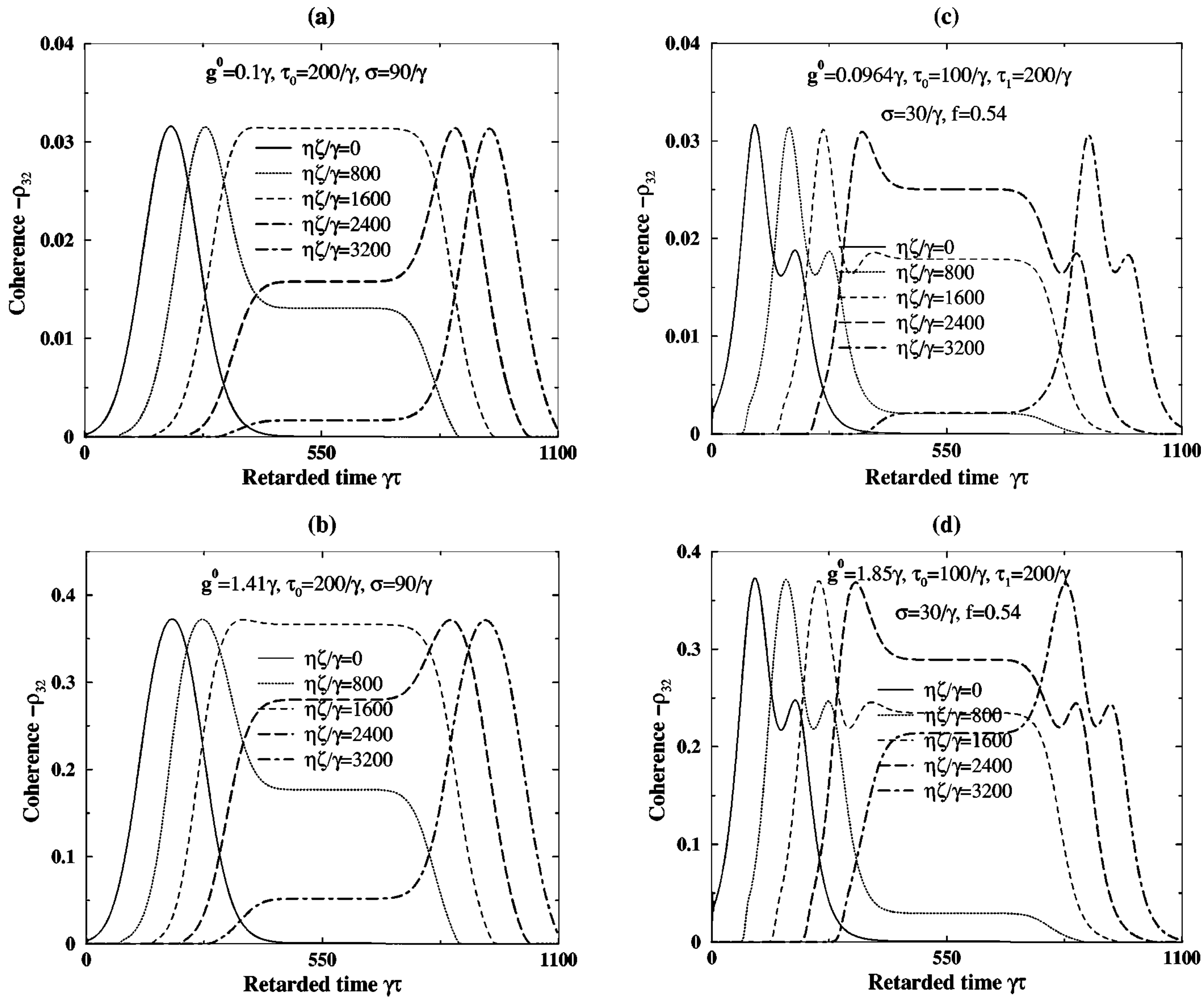


FIG. 5. (a)–(d) show the temporal profile of atomic coherence $-\rho_{32}$ against retarded time at different propagation distances.

*et al.* [11]. When the probe field intensity is large, we observe from Figs. 4(c) and 4(e) that the retrieved probe pulse suffers absorption as well as broadening because of a narrowing of the width of the EIT window. Remarkably, the probe pulse can be retrieved even for a probe that is not necessarily weak. Figure 5 shows the behavior of the atomic coherence $\rho_{32}$ as a function of retarded time at different distances. The switching off of the control field gives rise to a probe pulse that is stored inside the medium in the form of atomic coherence $\rho_{32}$. The atomic coherence retains the value that it attains before switching off. The atomic coherence starts generating the replica of the probe pulse at the moment the control field is switched on. Therefore, the atomic coherence is responsible for storage and retrieval of the probe pulse. In the presence of the control field, the temporal shape of the atomic coherence $\rho_{32}$ is same as the shape of the input probe pulse. A comparison of Figs. 4(b) and 5(a) shows very close connection between the probe pulse at different points in the medium and atomic coherence. For example, for $\eta\zeta/\gamma=3200$ and $\gamma\tau\leqslant 550$, the negligible amount of coherence leads to very little pulse power. Similar observations apply to the storage and retrieval of pulses at higher powers. On comparison of Figs. 5(a) [Fig. 5(c)] with 5(b) [Fig. 5(d)], we find that the generated atomic coherence $\rho_{32}$ is much more significant for larger values of the intensity of the probe.

### C. Nonadiabatic results

Matsko *et al.* [14] have shown that for any switching time of the control field, an almost perfect storage and retrieval of the weak probe pulse is possible. We have found that their results can be extended to probe pulses with moderate powers. We show the results in Fig. 6 for an intense probe pulse and nonadiabatic switching of the control field. For both adiabatic and nonadiabatic switching, the retrieved intense probe pulse is the same as the original one. However, there is a overall broadening and loss in the intensity of the retrieved

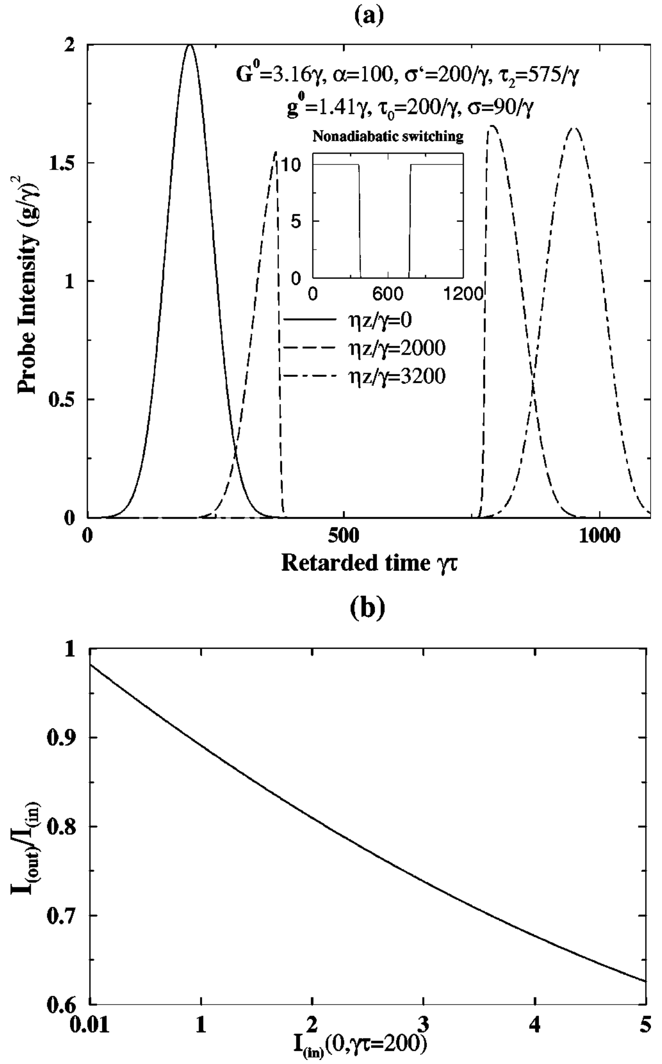


FIG. 6. (a) Shows storage and retrieval of an intense probe pulse even for nonadiabatic switching of the control field. Nonadiabatic switching is shown in the inset. (b) Drop in the intensity ratio of the probe retrieved to the input pulse as a function of the input probe intensity for the case of nonadiabatic switching of the control field; $I_{out}$ is measured at $\eta\zeta/\gamma = 3200$ and $\gamma\tau = 1000$.

probe pulse. We show in Fig. 6(b) the drop in the intensity of the output pulse as the intensity of the input pulse increases.

### D. Dynamical evolution of the control field

The dynamical evolution of the control field becomes important when the intensities of the control and probe fields are comparable. Figure 7 depicts time evolution of the control field at different distances. It is evident from this figure that a dip and a bump develop in the amplitude of the control field as it propagates through the medium. The shape of the bump and dip in the control field depends on the initial shape of the probe pulse at the entry of the medium. The changes in the control field are quite significant. Even for the cw control field, the output control field is a combination of both cw and a pulse.

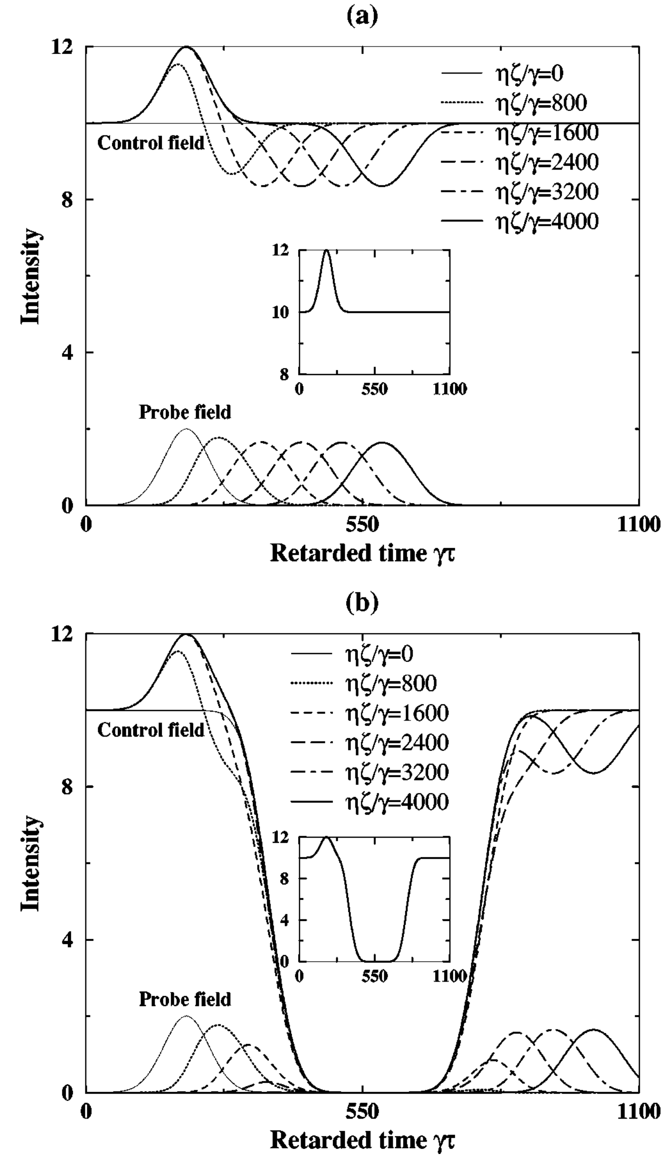


FIG. 7. (a), (b) Show temporal profiles of the control $(G/\gamma)^2$ and probe field $(g/\gamma)^2$ at different propagation distances within the medium. Temporal shape of $(V/\gamma)^2$ as shown in the inset does not depend on $\zeta$. In (a) the input control field is a cw. In (b) the input control field is a super-Gaussian shape with parameters $\tau_2 = 575/\gamma$ and $\sigma' = 200/\gamma$. The common parameters of the above two graphs are chosen as $G^0 = 3.16\gamma$, $g^0 = 1.414\gamma$, $\tau_0 = 200/\gamma$, and $\sigma = 90/\gamma$. The results of simulations using Maxwell-Bloch equations are indistinguishable from the results based on the adiabatic theory.

## IV. ADIABATIC THEORY OF GROBE, HIOE, AND EBERLY AND ITS RELATION TO LIGHT STORAGE

In a remarkable paper Grobe *et al.* [17] discovered pulse pair solutions that they called as adiabatons. These are the pulse pairs that are generated in a $\Lambda$ system under conditions of adiabaticity. We show the deep connection of the problem of storage and retrieval of pulses to the adiabatic theory. The control field is switched on before the probe field. This is to keep the system in the dark state, which is an essential condition for the formation of the adiabatic pulse pair. Under conditions of negligible damping, Grobe *et al.* [17] find that the response of the medium can be very well approximated by the solutions

$$\rho_{13} \approx \frac{i}{V}\frac{\partial}{\partial\tau}\left(\frac{g}{V}\right),$$

$$\rho_{12} \approx \frac{i}{V}\frac{\partial}{\partial\tau}\left(\frac{G}{V}\right), \tag{12}$$

$$\rho_{32} \approx -\frac{gG}{V^2},$$

where $V^2=(G^2+g^2)$. The approximate solution (12) holds provided the following adiabaticity condition is satisfied by the two fields:

$$G\frac{\partial g}{\partial\tau} - g\frac{\partial G}{\partial\tau} \ll V^3. \tag{13}$$

By inserting solution (12) into the Maxwell equations (7), we obtain a pair of coupled nonlinear wave equations

$$\frac{\partial g}{\partial\zeta} = -\frac{\eta}{V}\frac{\partial}{\partial\tau}\left(\frac{g}{V}\right),$$

$$\frac{\partial G}{\partial\zeta} = -\frac{\eta}{V}\frac{\partial}{\partial\tau}\left(\frac{G}{V}\right). \tag{14}$$

This pair of one-dimensional partial differential equations are nonlinearly coupled through the variable $V$. With the help of Eqs. (14), one can easily show that $V$ does not depend on the space variable, i.e., $\zeta$, during the propagation, $V$ satisfies the relation

$$V\left(\frac{\eta\zeta}{\gamma},\gamma\tau\right) = V(0,\gamma\tau). \tag{15}$$

Thus the conservation law would imply that any change in the probe field is compensated by a corresponding change in the control field for $V$ to remain independent of the spatial coordinates. The input fields determine the temporal shape of $V$. Analytical solution of Eqs. (14) can be obtained by changing the variable $\tau$ to $z(\gamma\tau) \equiv (1/\gamma^2)\int_{-\infty}^{\gamma\tau} V^2(0,\gamma\tau)d(\gamma\tau)$:

$$g\left(\frac{\eta\zeta}{\gamma},\gamma\tau\right) = V(0,\gamma\tau)F_g\left[z(\gamma\tau)-\frac{\eta\zeta}{\gamma}\right],$$

$$G\left(\frac{\eta\zeta}{\gamma},\gamma\tau\right) = V(0,\gamma\tau)F_G\left[z(\gamma\tau)-\frac{\eta\zeta}{\gamma}\right], \tag{16}$$

where $F_g[x]=g(0,z^{-1}(x))/V(0,z^{-1}(x))$ and $z^{-1}(x)$ denotes the inverse function of $z$. We have chosen the initial fields strong enough to ensure the formation of an adiabatic pulse pair. The input fields $g$ and $G$ are chosen such that $V$ is constant after a certain time $T$. Therefore, for $\tau \geqslant T$, the integral $z(\gamma\tau)$ can be analytically performed. For a cw control field and a Gaussian probe pulse, we find the explicit results for the probe and control fields:

$$g\left(\frac{\eta\zeta}{\gamma},\gamma\tau\right) = \frac{\sqrt{[g^{0^2}e^{-2(\gamma\tau-\gamma\tau_0)^2/(\gamma\sigma)^2}+G^{0^2}]}}{\sqrt{[g^{0^2}e^{-2(\gamma\tau-\gamma\tau_0-\gamma\eta\zeta/G^{0^2})^2/(\gamma\sigma)^2}+G^{0^2}]}} \times g^0 e^{-(\gamma\tau-\gamma\tau_0)^2/(\gamma\sigma)^2},$$

$$G\left(\frac{\eta\zeta}{\gamma},\gamma\tau\right) = \frac{\sqrt{[g^{0^2}e^{-2(\gamma\tau-\gamma\tau_0)^2/(\gamma\sigma)^2}+G^{0^2}]}}{\sqrt{[g^{0^2}e^{-2(\gamma\tau-\gamma\tau_0-\gamma\eta\zeta/G^{0^2})^2/(\gamma\sigma)^2}+G^{0^2}]}}G^0; \quad t \geqslant T. \tag{17}$$

For the case when the control field is taken as a super-Gaussian pulse and the probe field is taken as a Gaussian pulse, it is not possible to evaluate the function $z$ analytically. In Fig. 7, the solution of Eqs. (14) for both these cases is superimposed on the numerical results obtained from the density-matrix equations. It is remarkable that the solution of Eqs. (14) obtained under the adiabatic approximation matches extremely well with the numerical solution of the complete set of density-matrix equations. As shown in Fig. 5, the adiabatic approximation for the atomic coherence $\rho_{32}$ of the Eq. (12) is also indistinguishable from the result obtained from the density-matrix equations [18]. It is evident from the temporal profiles of the control and probe fields at different propagation distances that a dip and a bump develop in the control field intensity as it propagates through the medium. Figure 7 confirms unambiguously that the adiabatic pair pulse (consisting of the dip in the pump and the broadened probe) travels loss-free distances that exceed the weak probe absorption length (here typical value of $\eta\zeta/\gamma=2400$) by several orders of magnitude with an unaltered shape. In principle, $V$ could have both space and time dependences. Within the adiabatic approximation, $V$ does not depend on the space coordinate as shown in the inset of Fig. 7. From Fig. 7, it is very much clear that the temporal shape of $V$ depends on the input shape of the control and probe fields and it propagates inside the medium with unaltered shape. To keep $V$ constant in the space domain, any change in the temporal shape of the control field is compensated by a change in the temporal shape of the probe field. When $V^2$ and the control field are zero, then the probe field is also zero that suggests that the probe field gets stored inside the medium. The retrieved probe pulse that is a replica of the input probe pulse is a part of the adiabatic-pulse pair. The numerical results on the storage and retrieval of light obtained from density-matrix formalism match extremely well with those obtained from the adiabatic approximation. Clearly, the propagation of the pulse pair in adiabatic approximation is quite important for understanding the storage and retrieval of light.

## V. CONCLUSIONS

We have investigated the possibility of storage and retrieval of moderately intense probe pulses in a system with relevant atomic transition in a $\Lambda$ configuration. We integrate numerically the full set of the density-matrix equations and the Maxwell equations for both control and probe fields. The numerical results show that the storage and retrieval of probe pulses with moderate powers are possible. The dynamical evolution of the control field is important. It may be worth noting that a cw control field becomes pulsed due to its coupling to the probe pulse via the atomic polarization. We find that even though the storage and retrieval at larger powers are possible, the probe field gets absorbed and broadened. However the absorption and broadening are not very significant. This behavior is explained in terms of the narrowing of the EIT window as the power of the probe increases. We further show how the theory of Grobe *et al.* [17] enables us to understand the storage and retrieval of pulses.


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