From Storage and Retrieval of Pulses to Adiabatons

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Abstract

We investigate whether it is possible to store and retrieve the intense probe pulse from a Λ-type homogeneous medium of cold atoms. Through numerical simulations we show that it is possible to store and retrieve the probe pulse which are not necessarily weak. As the intensity of the probe pulse increases, the retrieved pulse remains a replica of the original pulse, however there is overall broadening and loss of the intensity. These effects can be understood in terms of the dependence of absorption on the intensity of the probe. We include the dynamics of the control field, which becomes especially important as the intensity of the probe pulse increases. We use the theory of adiabatons [Grobe et al. Phys. Rev. Lett. 73, 3183 (1994)] to understand the storage and retrieval of light pulses at moderate powers.
I. INTRODUCTION

Photons are ideal carriers of information. However it is extremely difficult to store photons for a long time and to retrieve them from a medium. The basic principle of photon storage is based on ultraslow group velocity of the light [1–5]. The later is made possible by electromagnetically induced transparency (EIT) [6] where an external control field is used to make an atomic medium transparent even for a frequency near atomic resonance. Under such conditions, a probe pulse at a particular frequency and polarization can propagate with a substantially reduced group velocity. Experimental and theoretical studies [7–9] show that the ultraslow group velocity of the pulse causes spatial compression of a probe pulse that has a spatial length of several kilometers in free space. This spatial compression leads to localization of the pulses within the atomic medium. Within the spatially localized pulse region, atoms are in a state of superposition characterized by the amplitudes and phases of the control and probe fields. Therefore, the induced transparency is associated with a substantial reduction of the group velocity of the probe pulse which is formed due to a coupled field-spin excitation, called dark state polaritons [9]. Switching off of the external control field converts the dark state polaritons into purely atomic coherence which is confined within the medium. By switching on the external control field at a later time, the atomic coherence can be transferred back into the radiation field. The regenerated radiation field can be an exact replica of the original one [10,11]. Therefore, the signal pulse information can be stored and retrieved by switching off and on the external optical field [7,9–12]. The mode of switching can be adiabatic [9] as well as nonadiabatic [12]. The past studies [7,9–12] on this problem have been carried out in the linear regime, where the probe pulse is much weaker than control field. Now the question is, whether it is possible to store and retrieve an intense probe pulse, in which case, the non-linearity of the medium with respect to the probe pulse amplitude becomes important? In this paper we address this question. We characterize the probe pulse propagation for different types of pulses in a homogeneously broadened Λ-type medium. The paper is organized as follows: In Sec. II, we obtain the Maxwell-Bloch equations that governs the propagation dynamics of the optical pulse at moderate power inside the Λ-type homogeneous medium. We include the radiative decay in simulations. The results of our numerical simulations which also includes the dynamical evolution of the control field are presented in Sec. III. In Sec. IV, we show how the results of numerical simulations can be understood in terms of the adiabaton theory of Grobe et al. [16].

II. DYNAMICAL EQUATIONS FOR PULSES PROPAGATION AT MODERATE POWERS

We consider an atomic system where relevant atomic transitions are taken in the configuration as shown in the Fig. 1. wherein, one of the transition is coupled to a external control field $E_c$ and the other transition is coupled to a probe field $E_p$. We define all fields as

$$\tilde{E}(z, t) = \tilde{E}(z, t)e^{-i(\omega t - kz)} + \text{c.c.},$$  \hspace{1cm} (1)
where $\vec{E}$ is a slowly varying envelop, $\omega$ is the carrier frequency and $k$ is the wave number of the field. Here we assume the frequencies of the carrier waves coincide with the frequencies of the corresponding atomic transitions $\omega_1 = \omega_{13}$ and $\omega_2 = \omega_{12}$.

We are using the Maxwell-Bloch Equation to describe the propagation dynamics of light pulses through the atomic vapour. In the slowly varying envelope approximation the temporal and spatial evolution of the field envelops is governed by

$$\frac{\partial g}{\partial z} + \frac{\partial g}{\partial ct} = i\eta \rho_{13},$$
$$\frac{\partial G}{\partial z} + \frac{\partial G}{\partial ct} = i\eta \rho_{12},$$

(2)

where $2g = 2d_{13} \cdot \vec{E}_p/\hbar$ and $2G = 2d_{12} \cdot \vec{E}_c/\hbar$ are the Rabi frequencies of the probe and drive fields, $d_{13}$ and $d_{12}$ are the dipole moments, $\rho_{13}$ and $\rho_{12}$ are the density matrix elements of the corresponding atomic transitions, and $c$ is the velocity of light in vacuum. The propagation constant is taken to be equal for all transitions and is given by $\eta = 3\lambda^2 N \gamma/8\pi$. It depends on the density of the atomic vapor $N$ and the atomic transition wavelength $\lambda$.

Within the rotating wave approximation the interaction between the atoms and the fields can be describe by the Bloch equations in a rotating frame

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

(3)

Here $\gamma$'s govern the radiative decay of the state $|1\rangle$; we further assume the equality $\gamma_1 = \gamma_2 = \gamma/2$. These Bloch equations are to be supplemented by the population conservation law

$$\rho_{11} + \rho_{22} + \rho_{33} = 1$$

(4)

The solution of Eqs.(2)-(3) gives the complete evolution of the atom-field system at any instant of time. The analytical solution of the Maxwell-Bloch equations not known though under special conditions some solutions are known [15]. Therefore we study the pulse-propagation problem only numerically.

III. NUMERICAL SIMULATIONS

A. Pulse Propagation: Effect of nonlinearities

We solve the propagation problem numerically for a homogeneously broadened gas of cold A-atoms in the travelling window frame of reference: $\tau = t - z/c, \zeta = z$. We consider initial probe pulses of two different shapes given by
\[ g(0, \tau) = \begin{cases} g_0 e^{-\left(\frac{\tau}{\sigma}\right)^2} & \text{Gaussian pulse} \\ g_0 \left[ \text{sech}\left(\frac{\tau - \tau_0}{\sigma}\right) + f \times \text{sech}\left(\frac{\tau - \tau_1}{\sigma}\right) \right] & \text{Sech pulse.} \end{cases} \]

Here, \( g_0 \) is the real constant characterizing the peak amplitude of the Rabi frequency before the pulse enters the homogeneous medium, \( \sigma \) is the bandwidth of the input pulse, and \( \tau_i \) determines the number of peaks as well as their position. The initial condition for the atomic system is to be taken as \( \rho_{33}(\eta/\gamma, 0) = 1 \), with all other density matrix elements equal to zero. In order to appreciate the effect of nonlinearities, we first consider the control field as a continuous wave \( G(0, \gamma \tau) \equiv \text{constant} \left[ (G/\gamma)^2 = 10 \right] \). We work under the condition of electromagnetically induced transparency, Figures 2(a), 2(b), 2(c) and 2(d) display the propagation of two types of probe pulses inside the medium. From Figs. 2(a) and 2(c), we see that the weaker probe pulse would propagate without any significant absorption and broadening inside the medium. The Figs. 2(b) and 2(d) show that the intense probe pulse suffers absorption and broadening. The shape of the pulse remains almost identical to the input pulse. This behavior of the intense probe pulse can be explained with help 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 the Fig. 3 that, 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 two photon resonance condition. Note that the transparency window that appears in the absorption spectrum has a finite bandwidth which depends on the intensities of the control and probe fields. The width of the transparency window becomes narrow when the probe field intensity is increased while keeping the control field intensity constant. Therefore, the condition for distortionless pulse propagation is that the bandwidth of the probe pulse should be contained within the transparency window of the EIT medium. If the pulse becomes too short, or its spectrum too broad relative to the transparency window of the EIT medium, absorption and higher order dispersion need to be taken into account. The group velocity dispersion of the medium causes the broadening of the probe pulse.

**B. Storage and Retrieval at Moderate powers**

The light pulse propagating in a homogeneously broadened \( \Lambda \)-type medium, suitably driven by another control field, can be stooped and later released in a controlled way [7,9,13,14]. The ultra slow group velocity of the light is the main key issue of the “light-storage” technique. In particular, changing the control field intensity results in changing the group velocity of the light 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 modelled by a super-Gaussian shape given by

\[ G(0, \tau) = G_0 \left[ 1 - e^{-\left(\frac{\tau}{\sigma}\right)^\alpha} \right] \]

\[ \alpha = 4 \quad \text{adiabatic switching} \]

\[ \alpha = 100 \quad \text{nonadiabatic switching}. \]

In Fig. 4(a), we consider the adiabatic switching of the control field. Switching off of the control field 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 state $|2\rangle$ and $|3\rangle$. As seen from Figs. 4(b) and 4(d), for weaker probe pulse, the shape of the retrieved pulse is same as the original one because the width of the probe pulse spectrum is very much less than 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 et al. [9]. 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 narrowing of the width of EIT window. Remarkably, the probe pulse can be retrieved even for probe that is not necessarily weak. We also present Figure 5 to show the behavior of the atomic coherence $\rho_{32}$ as a function of retarded time, which depends on the behavior of the input probe pulse. In presence of control field, the temporal shape of the atomic coherence $\rho_{32}$ is same as the shape of the input probe pulse. Remarkably enough the $\Lambda$ system leads to storage and retrieval even at moderate powers of the probe field.

C. Nonadiabatic Results

Scully et al. [12] have shown that for any switching time of the control field, an almost perfect storage and retrieval of weak probe pulse is possible. We extend storage and retrieval of the probe pulse in the nonlinear regime. The results have been shown 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 same as original one. However, there is overall broadening and loss in intensity of the retrieved probe pulse. The Fig. 6(b) shows that the drop in intensity ratio of the output to input probe pulse as function of the input pulse intensity.

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 of comparable strength. Figure 7 depicts time evolutions of the control field at different distances. It is evident from this figure that a dip and a bump develops in the amplitude of the control field as its propagates through the medium. The shape of the bump and dip in the control field depends on the initial shape of the input probe pulse at the entry surface of the medium. The propagation dynamics of the dip of the control field and broadened probe field together can be understood in terms of adiabaton pair [16].

IV. ADIABATON THEORY AND ITS RELATION TO LIGHT STORAGE

In a remarkable paper Grobe et al. [16] discovered what they called as adiabatons. These are the pulse pairs which 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 theory of
adiabatons. The control field is switched on before the probe field, to keep the system in the
dark state and which is an essential condition for the adiabaton formation. Under conditions
of negligible damping, the response of the medium then can be very well approximated by
the solutions

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

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

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

provided the following adiabacity relation is satisfied by the two field:

\[
G \frac{\partial g}{\partial \tau} - g \frac{\partial G}{\partial \tau} \ll V^3,
\]

where \(V^2 = (G^2 + g^2)\). By inserting solution (5) into the Maxwell equation (2), we obtain a
pair of coupled nonlinear wave equations

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

\[
\frac{\partial G}{\partial \zeta} = -\frac{\eta}{V} \frac{\partial}{\partial \tau} \left( \frac{G}{V} \right).
\]

Note that the radiative decay constant does not play any role in the above two equations.
These two, one dimensional PDE are nonlinearly coupled through the variable \(V\). With the
help of equation (7), one can easily show that \(V\) does not depend on \(\zeta\), implying that \(V\) is
remains constant during the propagation.

\[
V \left( \frac{\eta\zeta}{\gamma}, \gamma \tau \right) = V (0, \gamma \tau)
\]

Thus the conservation law would imply that \(V\) any change in probe field is compensated by
a corresponding in change the control field. Analytical solutions of equation (7) can also be
obtained by changing the variable \(\tau\) to \(z(\gamma \tau) \equiv \frac{1}{\gamma} \int_{-\infty}^{\tau} V^2(0, \gamma \tau) d(\gamma \tau)\):

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

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

\(F_g[x] = g[0, z^{-1}(x)]/V[0, z^{-1}(x)]\) and \(z^{-1}(x)\) denotes the inverse function of \(z\). We have
chosen the initial fields strong enough to ensure the formation of an adiabaton pair. The
input fields \(g\) and \(G\) are chosen such that \(V\) is constant after a certain time \(T\); hence, for
\(\tau \geq T\) the integral \(z(\gamma \tau)\) can be analytically performed. First we take the control field of
a constant amplitude \((G/\gamma = 3.16)\) and the probe field as Gaussian pulse, we then obtain
(for \(\tau \geq T\)
\begin{align*}
g \left( \frac{\eta \zeta}{\gamma}, \gamma \tau \right) &= \frac{\sqrt{[g^2 e^{-\frac{2(\gamma \tau - \gamma \tau_0)^2}{\sigma^2}} + G^0]^2}}{\sqrt{[g^2 e^{-\frac{2(\gamma \tau - \gamma \tau_0 - \frac{\eta \zeta \gamma G^0}{g^2})^2}{\sigma^2}} + G^0]^2}} g^0 e^{-\frac{(\gamma \tau - \theta_0)^2}{\sigma^2}} \\
G \left( \frac{\eta \zeta}{\gamma}, \gamma \tau \right) &= \frac{\sqrt{[g^2 e^{-\frac{2(\gamma \tau - \gamma \tau_0)^2}{\sigma^2}} + G^0]^2}}{\sqrt{[g^2 e^{-\frac{2(\gamma \tau - \gamma \tau_0 - \frac{\eta \zeta \gamma G^0}{g^2})^2}{\sigma^2}} + G^0]^2}} G^0.
\end{align*}

Next the control filed is taken as a super Gaussian shaped pulse and the probe field as Gaussian pulse as before. The solutions of equation (7) for both these cases, are superimposed on the numerical results obtain from density matrix formalism in Figure 7. It is remarkable that, the solutions of equations (7) obtain under adiabatic approximation, matches extremely well with the numerically solutions of the density matrix formalism containing seven nonlinear equations. The adiabatic calculation of atomic coherence \( \rho_{32} \) of the equation(5) is indistinguishable from the density matrix formalism as shown in Fig. 5 [17]. It is evident from the temporal profiles of the control and probe at different propagation distances that a dip and a bump develops in the control field intensity as it propagates through the medium. Figure 7 unambiguously confirms that the adiabaton pair (consisting of the dip in the pump and broadened probe) travels loss-free distances which exceed the weak probe absorption length(there typical value of \( \eta \zeta / \gamma = 2400 \))by several orders of magnitude with an unaltered shape. A gradual decrease of the control field intensity \((G/\gamma)^2\) results in a corresponding decrease of the intensity of the probe pulse \((g/\gamma)^2\) maintaing the constancy of \(V^2\), as shown in Fig. 7(b). Therefore, it is clear that when the control field intensity becomes zero the probe pulse gets stored inside the medium. The reverse phenomena of the retrieval of the probe pulse is achieved by gradual increase of the control field intensity. The fact that the numerical results from density matrix formalism matches extremely well with the adiabatic approximation clearly indicates that the storage and retrieval of light can be understood in terms of the adiabaton pair propagation.

V. CONCLUSIONS

We have investigated and answered in affirmative, the possibility of storage and retrieval of moderately intense probe pulses in a system with \( \Lambda \) configuration. The propagation of two type probe pulses is first analyzed numerically, using Maxwell Bloch equation. It was found that the intense probe pulses are absorbed and broadened due to the non-linear dependence of susceptibility in the medium. In addition, for the constant control field intensity, the width of EIT window becomes narrower with a simultaneous increase of the probe field intensity. Therefore, when the intensity of the probe pulse is large, the non-linear effect of the medium becomes more important. For larger intensity probe pulses, storage and retrieval is possible for both adiabatic as well as nonadiabatic switching of the control field. We studied the dynamics of the control field and found that the dip of the control field and broadened probe pulse propagate together inside the medium as an adiabaton pair. We show that retrieval probe pulse is the part of propagating adiabaton pair. Hence, the storage and retrieval of light can be clearly understood in terms of adiabaton pair propagation. The same was also studied by an adiabatic approximation. Remarkably, the result of the both cases are indistinguishable from each other in the wide parameters range studied here.
APPENDIX: NUMERICAL INTEGRATION PROCEDURE

The integration of the equation is performed by the fourth order Runge-Kutta method [18]. In our numerical calculation, sampling point along $\tau$ and $\zeta$ direction are $4 \times 10^6$ and $2 \times 10^6$ respectively. We prefer the parallelization of the sequential code because of the large number of sampling points which requires large execution time. The sequential code can be easily parallelized if each iteration is independent of the other, that is no variables that are written in some iteration will be read and/or written in another iteration. But in our code, each iteration is dependent on the previous iteration. Thus for our flow dependence case, the parallelization of the code is difficult and we do the parallelization manually by using OPEN MP directives in RS6000 in IBM machine.
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FIG. 1. Three-level Λ-type medium resonantly coupled to a control field with Rabi frequency $G$ and probe field $g$. 
FIG. 2. The probe field intensity in the medium is plotted against retarded time at different propagation distances within the medium. Fig (a) and (c) show the probe pulse propagation with non-diminishing amplitude, for small intensities. Fig (b) and (d) depict the broadening and loss of intensity in case of an intense probe pulse case. In all the cases the control field is taken as CW($G = 3.16\gamma$).
FIG. 3. Imaginary parts of susceptibility $\chi_{13}\hbar\nu/N|d_{13}|^2$ at a probe frequency $\omega_1$ in the presence of control field as a CW $G = 3.16\gamma$. The width of the transparency window decreases with increase in the intensity of the probe field.
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 time profile of the intense probe pulse at different propagation distances.
FIG. 5. Fig. (a)-(d) shows the temporal profile of atomic coherence $-\rho_{32}$ against retarded time at different propagation distances.
FIG. 6. (a) shows storage and retrieval of intense probe pulse even for nonadiabatic switching of the control field. (b) Drop in intensity ratio of the probe retrieved to the input pulse as a function of input probe intensity for the case of nonadiabatic switching of the control field; the \(I_{\text{out}}\) is measured at \(\eta \zeta / \gamma = 3200\) and \(\gamma \tau = 1000\).

FIG. 7. (a) and (b) shows temporal profiles of the control \((G/\gamma)^2\) and probe field \((g/\gamma)^2\) at different propagation distances within the medium. In Fig. (a) the input control field is a CW. In Fig (b) the input control field is a Super-Gaussian shape with parameter \(\tau_2 = 575/\gamma, \sigma' = 200/\gamma\). The common parameter of the above two graphs are chosen as: \(G^0 = 3.16\gamma, g^0 = 1.14\gamma, \tau_0 = 200/\gamma, \sigma = 90/\gamma\). The results of simulations using Maxwell-Bloch equations are indistinguishable from the results based on adiabaton theory.