Manipulating the retrieval of stored light pulses

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We consider storage and retrieval of light based on electromagnetically induced transparency in an atomic medium using nonadiabatic switching of a control field. We derive various conditions for writing (and reading) the light information to (and from) the atomic coherence. We obtain an analytical solution for the retrieval of the stored pulse that is in excellent agreement with the full numerical results. We identify the origin of distortion at the output and derive a condition to correct the distortions by manipulating the retrieval process.

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A weak signal pulse when passes through a three-level Λ system, along with a strong control field, generates a coherence between the ground states of the system that follows the field evolution under certain conditions. Therefore, if the medium is lossless for the signal, due to electromagnetically induced transparency (EIT) [1], the ground-state coherence along the propagation direction of the probe field is a spatial map of the time evolution of the signal pulse. It has been demonstrated that this map can be preserved to store the information of the signal by dynamically turning off the control field [2–5]. Later, the control field is turned on to beat with the coherence of the medium, which maps back the information written in the medium into an output pulse [6] that resembles the signal pulse shape [2]. It has been demonstrated that the light storage and retrieval is possible by switching the control field both adiabatically [7] and nonadiabatically [8,9]. EIT based storage has attracted so much attention due to its potential application to store a quantum state [6–8], unlike the earlier light storage proposals based on photon echo [10].

In this Brief Report, we consider the light storage via nonadiabatic switching of the control field [8] to manipulate the retrieval process. We obtain an analytical expression for the retrieved pulse that gives deep insight into the retrieval process and determines the crucial parameters of the field and medium that cause distortion in the output pulse shape. We derive a condition which shows that the best retrieval is possible if a ramp field is used as the reading control field, see Fig. 1(b)}.

We consider propagation of a weak probe (signal) pulse $E_p(z,t)$, in the presence of a strong control field $E_c(z,t)$ inside a medium consisting of atoms having Λ configuration, shown in Fig. 1(a)}. Here $E_\alpha(z,t)=\mathcal{E}_\alpha(z,t)e^{ik_\alpha z-i\omega_\alpha t}+\text{c.c.}$, with $\mathcal{E}_\alpha(z,t)$ being the slowly varying amplitude of the pulse envelop and k_α the propagation constant with central frequency ω_α; $\alpha=p,c$ refers to the probe and control field, respectively. Both the fields are on resonance with their respective transitions, see Fig. 1(a)}. Assuming the amplitude and phase of the fields and that of the medium polarization are slowly varying, and the fields do not have any transverse dependence, the Maxwell equations for the Rabi frequencies $\Omega_p=d_{13}\mathcal{E}_p/\hbar$ and $\Omega_c=d_{12}\mathcal{E}_c/\hbar$ of the probe and control fields, respectively, are

$$c\partial_z\Omega_p+\partial_t\Omega_p=ic\,\eta_p\rho_{13}, \tag{1}$$

$$c\partial_z\Omega_c+\partial_t\Omega_c=ic\,\eta_c\rho_{12}. \tag{2}$$

Here, $\partial_\xi\equiv\partial/\partial\xi$ and, $\eta_p=k_pNd_{13}^2/(2\hbar\epsilon_0)$ and $\eta_c=k_cNd_{12}^2/(2\hbar\epsilon_0)$ are the coupling constants that depend on the atomic density N, the dipole matrix element d_{ij}. The induced polarizations of the medium, determined by the coherences ρ_{ij}, are obtained from the equations [8]

$$\partial_t\rho_{12}=-\gamma\rho_{12}-i\Omega_c(\rho_{11}-\rho_{22})+i\Omega_p\rho_{32}, \tag{3a}$$

$$\partial_t\rho_{13}=-\gamma\rho_{13}+i\Omega_c\rho_{23}-i\Omega_p(\rho_{11}-\rho_{33}), \tag{3b}$$

$$\partial_t\rho_{23}=i\Omega_c^*\rho_{13}-i\Omega_p\rho_{21}. \tag{3c}$$

Here we have the spontaneous decays $\gamma_1=\gamma_2=\gamma/2$, equality taken for simplicity; ρ_{ii} is the population of $|i\rangle$.

To determine the evolution of the probe field, we make the following assumptions: $|\Omega_p|\ll\gamma$ and $\rho_{33}\approx 1$. Thus the coherence $\rho_{12}\approx 0$, and hence the control field propagation is unaffected by the medium state. In the presence of the control field $\Omega_c\neq 0$, from Eq. (3c), the coherence ρ_{13} has a for-

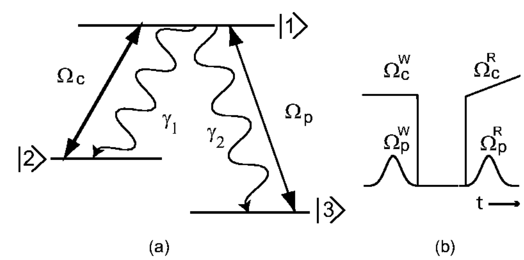

FIG. 1. (a) The level diagram of the three level Λ scheme under consideration. Here γ_i's are the spontaneous decay rates and Ω_α represents the Rabi frequency of the control (probe) field with $\alpha\to c\ (p)$. (b) The amplitude of the writing (reading) control fields $\Omega_c^W(\Omega_c^R)$, and the amplitude of input probe pulse Ω_p^W and the retrieved signal Ω_p^R. Note that $\Omega_c^R\neq\Omega_c^W$ for the exact retrieval of the signal.

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mal solution $\rho_{13}=-(i/\Omega_c^*)\partial_t\rho_{23}$, which governs the probe evolution. Clearly, from Eq. (1), the ground-state coherence ρ_{23} plays a vital role in the evolution process. Substituting the above ρ_{13} into Eq. (3b) and assuming that $\partial_t\Omega_c$ is negligible, we get

$$\partial_t^2\rho_{23} + \gamma\partial_t\rho_{23} + |\Omega_c|^2\rho_{23} + \Omega_p\Omega_c^* = 0. \tag{4}$$

Next we describe the light storage and retrieval process using the above basic equations. The three phases involved in this process are writing, storage, and reading of the signal—denoted as the *W*, *S*, and *R* phase, respectively. We consider that the control field is abruptly turned off after the *W* phase and turned on after the *S* phase. The probe and control fields in *W* and *R* phases are distinguished by using the notation $\Omega_p\to\Omega_p^\mu$ and $\Omega_c\to\Omega_c^\mu$, respectively, $\mu\to W,R$. We consider that the control field has different values of constant or extremely slowly varying amplitudes in each phase of the process.

W phase: For the probe field to be written to the atomic coherence ρ_{23}, the first two terms in the left-hand side of Eq. (4) should be small, so that $\rho_{23}\propto\Omega_p^W$. The required conditions are

$$T_p \gg \gamma/|\Omega_c^W|^2, \quad T_p^2 \gg 1/|\Omega_c^W|^2, \tag{5}$$

where, T_p is the characteristic width of the probe pulse. Thus, clearly the ground-state coherence, that becomes $\rho_{23}(t)=-\Omega_p^W(t)/\Omega_c^W$, can adiabatically follow the evolution of the probe field, if the conditions in Eq. (5) are satisfied. Under EIT condition, however, the second condition is redundant, because $|\Omega_c^W|\lesssim\gamma$ and broad pulses ($T_p\gg\gamma^{-1}$) are used to limit the spectral width of the probe field within the transparency window.

Using the above ρ_{23}, we obtain the evolution equation for the probe in the *W* phase as $v_g\partial_z\Omega_p^W+\partial_t\Omega_p^W=0$. The solution can be easily evaluated as $\Omega_p^W(z,t)=\Omega_p^W(0,t-z/v_g)$—that is, the probe pulse inside the medium having group velocity $v_g=c/(1+c\eta_p/|\Omega_c^W|^2)$. The group velocity is drastically reduced inside the medium. The slow group velocity and the resulting spatial compression of the probe pulse [11] are critically important for a complete mapping of the signal pulse onto the coherence. The coherence $\rho_{23}(z,t)$ in the *W* phase is calculated using the above solution of $\Omega_p^W(z,t)$.

S phase: In this phase of the process, the mapped coherence in the *W* phase is frozen by abruptly turning off Ω_c^W. Obviously, there is no adiabatic following at the moment of turning off. Once Ω_c^W is turned off, we numerically observe (not presented here) that the probe pulse is absorbed by the medium, transferring its residual energy back and forth between the probe and control transitions. However, as the switching time is too short and probe field is weak, the medium state is almost unchanged and thus the coherence ρ_{23} is frozen at its previous value (i.e., at the value just before Ω_c^W was turned off at $t=t_w$)—say $\tilde{\rho}_{23}$. We emphasize that such an approximation is valid only if the probe field is weak ($\Omega_p^W\ll\gamma$), unlike in the adiabatic turning off of the control field, where even a relatively strong probe field can also be used [12]. The typical storage timescale is limited by the ground-state dephasing, which we have neglected in our calculation.

R phase: The signal pulse is retrieved in this phase by beating the frozen coherence $\tilde{\rho}_{23}$, with a reading control field Ω_c^R. Thus the initial conditions in the *R* phase are $\rho_{23}=\tilde{\rho}_{23}$ and $\Omega_p^R=0$ at $t'=0$; where $t'=t-(t_w+t_s)$. The equations that govern the retrieval of the probe pulse are obtained as follows: from Eq. (4), neglecting $\partial_t^2\rho_{23}$ in the *R* phase we get

$$\partial_{t'}\rho_{23} = -(\Omega_c^{R*}/\gamma)(\Omega_c^R\rho_{23} + \Omega_p^R). \tag{6}$$

Using the above, the evolution equation (1) becomes

$$c\partial_z\Omega_p^R + \partial_{t'}\Omega_p^R = -(c\eta_p/\gamma)(\Omega_c^R\rho_{23} + \Omega_p^R). \tag{7}$$

The above two coupled equations determine the retrieval of the stored pulse. To solve, we follow Matsko *et al.* [8] to Fourier decompose the spatial components of Ω_p^R and ρ_{23} using the definition $\alpha(z,t')=\int_{-\infty}^{\infty}\alpha_k(t')e^{ikz}dk$. Here $\alpha_k\equiv\Omega_{k,p}^R$ (and $\rho_{k,23}$) correspond to the Fourier transform of Ω_p^R (and ρ_{23}), with k being the Fourier conjugate variable of z. From the above definition, Eqs. (6) and (7) for a single component k reduce to

$$\partial_{t'}X(t') = MX(t'), \tag{8}$$

where $X(t')=[\Omega_{k,p}^R(t'),\rho_{k,23}(t')]^T, M=\begin{pmatrix} a_{11} & a_{12} \\ a_{21} & a_{22}\end{pmatrix}$, with $a_{11}=-(c\eta_p+ikc\gamma)/\gamma, a_{12}=-c\eta_p\Omega_c^R/\gamma, a_{21}=-\Omega_c^{R*}/\gamma$, and $a_{22}=-|\Omega_c^R|^2/\gamma$. The general solution of the above equation is $X(t')=C_+e^{\lambda_+t'}\psi_+ + C_-e^{\lambda_-t'}\psi_-$. Here $C_\pm$ are arbitrary constants, $\lambda_\pm$ are the eigenvalues of the matrix M. Assuming k to be small such that $kc\gamma\ll c\eta_p$. Physically this condition means that the pulse length k^{-1} is much larger than the linear absorption length $(\eta_p/\gamma)^{-1}$. Taylor expanding $\lambda_\pm$ around the small parameter $a=ikc\gamma/(\eta_p c)\approx 0$ and keeping up to second order in a, the eigenvalues are evaluated as

$$\lambda_+ = -ikc\frac{|\Omega_c^R|^2}{|\Omega_c^R|^2 + c\eta_p} - 2k^2c^2\gamma\frac{|\Omega_c^R|^2 c\eta_p}{(|\Omega_c^R|^2 + c\eta_p)^3}, \tag{9}$$

and $\lambda_-=-(|\Omega_c^R|^2+c\eta_p)/\gamma+O(a)$. Here λ_+ is the slow component at the output having group velocity $v_g'=c|\Omega_c^R|^2/(c\eta_p+|\Omega_c^R|^2)$. Substituting the above λ_- in the solution $X(t')$, it is clear that the λ_- component suffers a large absorption inside the medium, contributes little to the output field, and hence is dropped hereafter. Note that, by truncating the series at a^2, we have limited the spatial variation of the field inside the medium upto ∂_z^2. However, our calculation takes into account all orders of variation in time, subject to the slowly varying approximation. The eigenfunctions corresponding to $\lambda_\pm$ are $\psi_\pm=(\phi_\pm,1)^T$, where $\phi_\pm=-[\lambda_\pm\gamma+|\Omega_c^R|^2]/\Omega_c^{R*}$. The constants $C_\pm$ in $X(t')$ are obtained using the initial conditions $\Omega_{k,p}^R=0$ and $\rho_{k,23}=\tilde{\rho}_{k,23}$ at $t'=0$. Here $\tilde{\rho}_{k,23}$ is the spatial Fourier transform of $\tilde{\rho}_{23}$ in the *S* phase, obtained as $\tilde{\rho}_{k,23}=-\tilde{\Omega}_{k,p}^W/\Omega_c^W$. Here

$$\tilde{\Omega}_{k,p}^W = \frac{1}{2\pi}\int_0^L \Omega_p^W(0,t_w - z/v_g)e^{-ikz}dz \tag{10}$$

is the value of $\Omega_{k,p}^W$ at $t=t_w$. Note that the above integration limits are changed to $[0,L]$ using an important condition in the *W* phase: when Ω_c^W is turned off at $t=t_w$, the whole pulse should be present inside the medium. Here L is the length of the medium.

With these initial conditions, the generated output field is calculated from $X(t')$ as

$$\Omega^R_{k,p}(t') = \frac{\tilde{\Omega}^W_{k,p}}{\Omega^W_c}\exp[\lambda_+ t' + \ln A], \tag{11}$$

where $A=(\phi_+\phi_-)/(\phi_+-\phi_-)$. To the second order in a, the amplitude part A is calculated as

$$\ln A \equiv \ln\left(\frac{\Omega^R_c c\eta_p}{c\eta_p + |\Omega^R_c|^2}\right) - \frac{ikc\gamma(c\eta_p - |\Omega^R_c|^2)}{(c\eta_p + |\Omega^R_c|^2)^2} - \frac{k^2c^2\gamma^2}{(c\eta_p + |\Omega^R_c|^2)^2}\left[2\frac{(c\eta_p - |\Omega^R_c|^2)^2}{(c\eta_p + |\Omega^R_c|^2)^2} - 1\right]. \tag{12}$$

Equation (11), along with Eqs. (9) and (12), determines the retrieved field for a single Fourier component k. The evolution of the retrieved field in the real space can be calculated by assuming a specific profile for the input signal pulse Ω^W_p. Equation (11) will reduce to the result of Matsko *et al.* [8] when λ_+ and $\ln A$ are considered to their lowest order in k.

Next, we consider a definite input pulse shape, a Gaussian as input probe, given by $\Omega^W_p(t)=\Omega^W_{p0}e^{-(t-t_0)^2/T_p^2}$, having its peak Ω^W_{p0} at $t=t_0$ and pulse width $2T_p$. The probe field inside the medium in the W phase is $\Omega^W_p(z,t)=\Omega^W_{p0}e^{-(t-t_0-z/v_g)^2/T_p^2}$. We substitute the above into Eq. (10) and evaluate the Gaussian integral, which on simplification, we get

$$\tilde{\Omega}^W_{k,p} = \frac{v_gT_p\Omega^W_{p0}}{2\sqrt{\pi}}\exp\left[-\frac{k^2v_g^2T_p^2}{4} - ikv_g\tau_w\right]. \tag{13}$$

The above simplification is obtained by imposing the following conditions: $L \gg v_g\tau_w, \tau_w \gg T_p, k < 2/(v_gT_p)$; here $\tau_w = t_w - t_0$. Here the first two conditions ensure that the whole pulse is there inside the medium when Ω^W_c is switched off. The third condition determines the limit of the maximum value of Fourier variable k, that validates the assumption of small a. It is important to note that, the coefficient of k^2 in Eq. (13) is associated with the pulse width T_p, and hence, it carries all the information of the pulse shape modification during the writing and reading process.

Substituting Eqs. (9), (12), and (13) into Eq. (11) and transforming back to z space, the retrieved probe pulse is obtained as

$$\Omega^R_p(z,t') = \frac{\Omega^R_c c\eta_p(\Omega^W_{p0}/p(t'))}{\Omega^W_c(c\eta_p + |\Omega^R_c|^2)}\exp\left[\frac{-(t'-q)^2}{\{T_pp(t')v_g/v'_g\}^2}\right], \tag{14}$$

where

$$p(t')^2 = 1 + \frac{8\gamma\eta_p v_g'^3 t'}{|\Omega^R_c|^4v_g^2T_p^2} + \frac{2\gamma^2v_g'^2}{|\Omega^R_c|^4v_g^2T_p^2}\left[\frac{(c\eta_p - |\Omega^R_c|^2)^2 - 4c\eta_p|\Omega^R_c|^2}{(c\eta_p + |\Omega^R_c|^2)^2}\right],$$

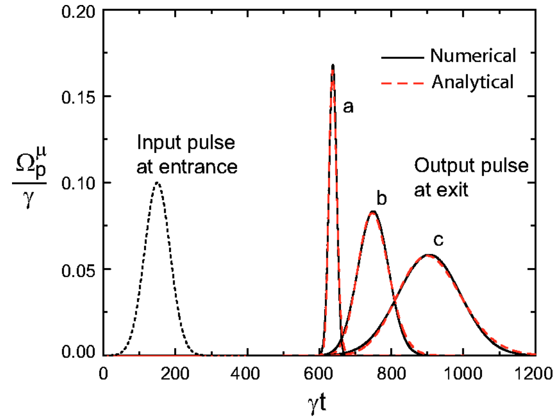

FIG. 2. The storage and retrieval of a Gaussian field. Both numerical and analytical results are shown for $\Omega^R_c = 2\Omega^W_c$ (a), Ω^W_c (b), and $0.7\Omega^W_c$ (c). Here $z\eta_p/c=320$, $T_p=50/\gamma$, $\Omega^W_c=\gamma$, $t_w=300/\gamma$, and $t_w+t_s=600/\gamma$. The numerical value of switching time of the control field is $0.12/\gamma$.

$$q = \frac{z}{v'_g} - \frac{v_g}{v'_g}\tau_w - \frac{\gamma c}{v'_g}\frac{c\eta_p - |\Omega^R_c|^2}{(c\eta_p + |\Omega^R_c|^2)^2}. \tag{15}$$

Assuming $c\eta_p \gg |\Omega^\mu_c|^2$, we have $v_g/v'_g \approx |\Omega^W_c|^2/|\Omega^R_c|^2$ and hence

$$\Omega^R_p(z,t') = \frac{\Omega^R_c\Omega^W_{p0}}{p(t')\Omega^W_c}\exp\left[\frac{-(t'-q)^2}{\{T_pp(t')|\Omega^W_c|^2/|\Omega^R_c|^2\}^2}\right], \tag{16}$$

with $p(t')^2 = 1 + \{8\gamma|\Omega^R_c|^2/(|\Omega^W_c|^4T_p^2)\}\{t' + \gamma/(4|\Omega^R_c|^2)\}$, and $q = z/v'_g - \gamma/|\Omega^R_c|^2 - (|\Omega^W_c|^2/|\Omega^R_c|^2)\tau_w$. As seen clearly from Eq. (16), the amplitude of the generated pulse is increased (decreased) and pulse width is compressed (broadened) when $|\Omega^R_c| > |\Omega^W_c|$ ($|\Omega^R_c| < |\Omega^W_c|$), which was reported in the experiment of Liu *et al.* [2]. In Fig. 2, we have plotted both numerical and analytical results of Ω^R_p for different Ω^R_c values—that shows an excellent agreement between the two. All numerical calculations are performed using Eqs. (1)–(3), and $\eta_c = \eta_p$ is used for simplicity.

If we drop the second order term containing k^2 in Eq. (11), and with $\Omega^R_c = \Omega^W_c = \Omega_c$, we get $p=1$ and $q = z/v_g - \tau_w$. Thus it would appear as if input signal field is retrieved exactly at the output [8], which is true when $T_p \to \infty$. However, from Eq. (16), the probe pulse would experience absorption and broadening, if $p \neq 1$—that can happen when $1/T_p^2 \neq 0$. Thus finite spectral width of the input probe pulse is the origin of pulseshape distortion in the retrieved output. Such distortions could be analytically described almost exactly, by taking into account the second order k^2 correction, as depicted in Eq. (16) and Fig. 2(b). From Eq. (16), the generated pulse Ω^R_p is broadened by the factor $p|\Omega^W_c|^2/|\Omega^R_c|^2$ and has an elongated tail due to the time dependence of $p(t')$, though numerically the elongation is insignificant. We derive a condition for its correction as

$$|\Omega_c^R|^2 = \left(|\Omega_c^W|^4 + \frac{2\gamma^2}{T_p^2} + \frac{16\gamma^2 t'^2}{T_p^4} \right)^{1/2} + \frac{4\gamma t'}{T_p^2} \qquad (17)$$

by equating $p=|\Omega_c^R|^2/|\Omega_c^W|^2$. For moderately long pulses used in the experiments [2–5], for $T_p=x/\gamma$, $x \gg 1$ (typically $x \sim 10$ to 100). Thus with $|\Omega_c^W| \sim \gamma$, $(2\gamma^2/T_p^2 + 16\gamma^2 t'^2/T_p^4) < |\Omega_c^W|^4$. Tayler expanding the first term in the right-hand side and keeping terms upto first order,

$$|\Omega_c^R|^2 \approx |\Omega_c^W|^2 + \frac{\gamma^2}{|\Omega_c^W|^2 T_p^2} + \frac{8\gamma^2 t'^2}{|\Omega_c^W|^2 T_p^4} + \frac{4\gamma t'}{T_p^2}. \qquad (18)$$

However, since $\max(t') \sim O(T_p)$, considering the terms only upto $1/x$ in the expansion, Eq. (17) can be reduced to $|\Omega_c^R|^2 \approx |\Omega_c^W|^2 + \beta t'$, i.e., a ramp field, with $\beta = 4\gamma/T_p^2$. We demonstrate in the Fig. 3 (from full numerical calculation) that such a ramp as the reading control field could recover a near replica of the input pulseshape of a probe having width $T_p=50/\gamma$. Note that, since a part of the signal filed is irrecoverably lost due to absorption in the medium, the total energy (area under the curve of $|\Omega_p^\mu|^2$) of the retrieved pulse is less than that of the original signal pulse. Hence what we show in Fig. 3(b) is the best possible output one can retrieve using the ramp field. As the tail elongation is insignificant, one can also use a constant reading field $\Omega_c^R > \Omega_c^W$ to approximately correct the shape of the output pulse. The value of the constant Ω_c^R could be choosen, for example, as the value of the ramp field at which the peak of the retrieved pulse appears. For very long input pulse, $\beta \approx 0$, i.e., the retrieved pulseshape can be preserved even for $|\Omega_c^W| = |\Omega_c^R|$. For too short pulse ($T_p < 10/\gamma$), k becomes large and hence this calculation fails. This calculation can also be generalized for a two- or multi-peak structure to obtain a condition for the best retrieval of the input pulse shape.

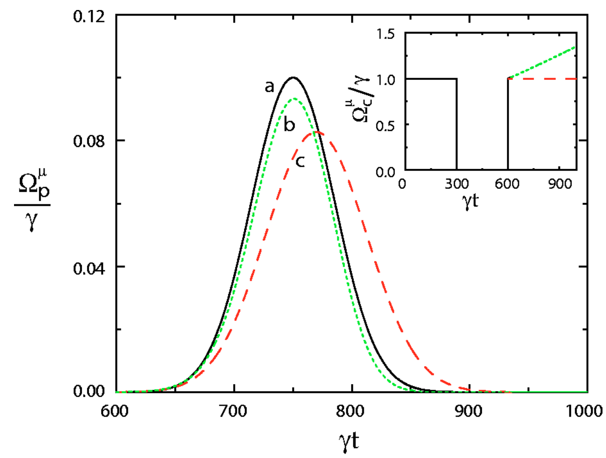

FIG. 3. The correction of distortion by a ramp field (see inset) given by the condition (17). (a) The input pulse shape at the entrance of the medium, shifted by $\gamma t=600$; (b) the retrieved field using a ramp field shown as dotted line in the inset; (c) the retrieved field using the constant reading field $\Omega_c^R=\Omega_c^W$. Here $z\eta_p/\gamma=300$ and all other parameters are same as in Fig. 2. Here the value of $\beta \approx 0.0016\gamma^2$.

In summary, we have discussed light pulse storage and retrieval in an EIT based storage medium via nonadibatic switching of the control field. We have obtained an analytical solution of the retrieved pulse, identified the source of the distortion in the retrieved pulse and obtained a condition for its remedy. We have substantiated the analytical results with the full numerical calculation—that shows excellent agreement. We have demonstrated analytically that a ramp field as the reading control field can retrieve an output pulse with a shape that is a near replica of the input pulse shape. We hope such a correction of the distortion can make nonadiabatic switching based light storage an efficient method and may find useful applications to extend the light storage to other potential systems, such as in thin optical fiber [13], etc.

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