Propagation of ultrashort pulses in multilevel systems under electromagnetically induced transparency

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In this paper, I numerically investigate the propagation of ultrashort pulses through an extended collection of multilevel systems under the condition of electromagnetically induced transparency. The transparency of a weak probe pulse is induced by a much stronger coupling pulse. In the limit of a weak probe excitation and a large number of excited states, I show that the free-induction decay signal that would naturally follow the excitation pulse is strongly suppressed, and the envelope of the probe pulse suffers a distortion in its temporal profile as the pulse propagates into the medium. This distortion is characterized by a stronger absorption of the leading edge of the pulse than that of its trailing edge. The temporal phase of the pulse, however, remains constant throughout propagation.

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I. INTRODUCTION

Absorption and dispersion in an optically thick medium can change the characteristics of a resonant laser pulse propagating through that medium. For example, a weak ultrashort pulse (whose temporal width is much shorter than the polarization decay time T_2) is converted into an oscillatory function of time as it propagates through an extended collection of two-level atoms [1–4]. In a multilevel medium, an ultrashort pulse can create a nonstationary superposition of states, or wave packet, which emits secondary pulses as it oscillates modifying the shape of the input driving pulse [5–8].

Electromagnetically induced transparency (EIT) [9], the phenomenon by which a strong coupling field turns an optically thick medium transparent to a resonant weak field, can strongly modify the optical properties of a medium and significantly alter the propagation of a laser pulse. The propagation of optical pulses in a medium of three-level Λ-type atoms under EIT has been investigated extensively [10–16]. In general, these studies considered pulses whose widths were long compared to the relaxation time of the intermediate excited state [10–14]. The case of pulses whose widths were shorter than the excited state's lifetime was investigated by Arkhipkin and Timofeev [15] and Kozlov and Eberly [16]. Arkhipkin and Timofeev studied the spatial and temporal evolution of two short laser pulses of identical shapes, but different durations. The probe pulse, which was shorter than the coupling pulse, was shown to propagate initially without major changes in its envelope, but it was gradually depleted, while the coupling pulse got stronger. Kozlov and Eberly studied the propagation of ultrashort pulses in an inhomogenously broadened phaseonium medium. They investigated the interplay of self-induced transparency and EIT for a number of starting conditions, such as pulse shapes and population distribution. Neither study addressed the issue of how the dynamics of propagation is affected if the pulses carry a temporal phase. A common feature to these studies is the three-level Λ-type medium with one intermediate excited state. But an ultrashort pulse may have a bandwidth large enough to be simultaneously on resonance with two or more excited states of the atom. These "extra" excited states may change the characteristics of the induced transparency and pulse propagation. EIT of an ultrashort pulse in a "multi-Λ" medium, whose intermediate excited state is replaced by a manifold of states, was studied recently [17].

In this paper, I study the propagation of an ultrashort probe pulse through an extended collection of multilevel atomic systems under the condition of electromagnetically induced transparency. The two pulses (probe and coupling) are initially matched and have, at the entrance to the medium, the same envelope and phase. I investigate how the temporal profile of the probe pulse is modified by propagation in the presence of a much stronger coupling pulse.

Figure 1 shows the model atomic system considered here. It consists of two lower states and a manifold of equally spaced excited states. In this multi-Λ medium, a weak probe

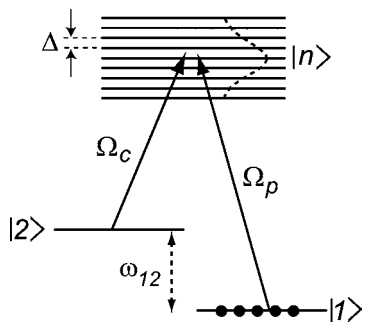

FIG. 1. A "multi-Λ" system excited by a pair of coupling and probe pulses, both resonant with the same excited state and with Rabi frequencies Ω_c and Ω_p, respectively. Both pulses are initially matched and have a spectrum (dashed line) large enough to overlap many states in the excited manifold; at the same time, their spectrum is much smaller than the energy difference ω_{12} of the lower levels. Energy separations are not drawn to scale; solid circles represent the initial population distribution, and Δ is the level spacing.

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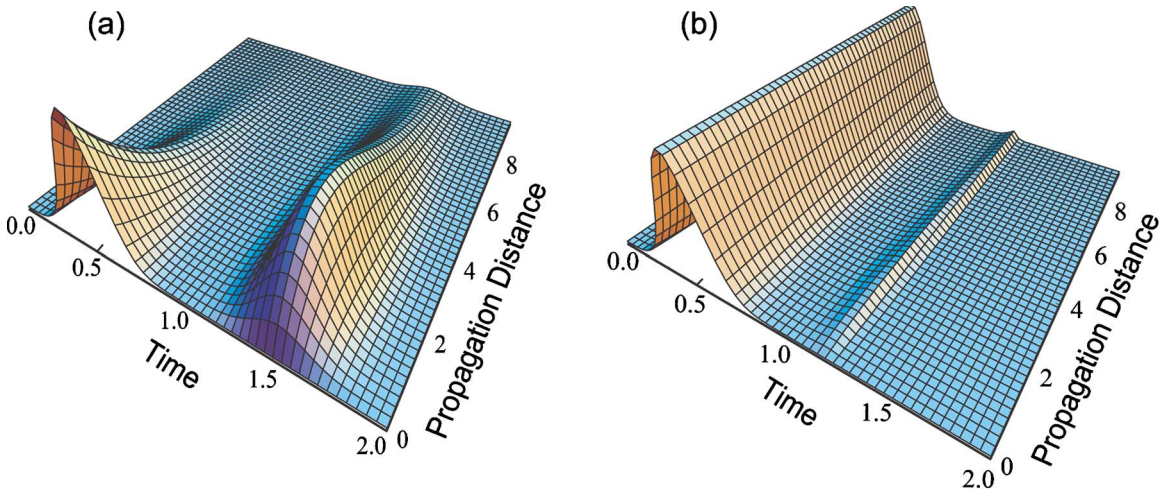

FIG. 2. (Color online) Propagation of a weak ultrashort Gaussian pulse in an extended collection of multi-Λ media in (a) the absence of a coupling pulse and (b) the presence of a matched coupling pulse with $\Omega_c = 100\Omega_p$. Time is measured in units of the excited states' beat period T, and propagation distance is measured in units of the inverse absorption coefficient α^{-1}. These results were obtained by numerically integrating the Maxwell-Bloch equations (1) and (2) for the input probe pulse given by Eq. (8) with phase $\varphi(\tau)=0$.

pulse connects state $|1\rangle$ to the excited manifold of states $|n\rangle$, while a stronger coupling pulse connects state $|2\rangle$ to the same excited states. The input pulses to be considered here are many orders of magnitude shorter than any relaxation time of the medium. Their bandwidths are large enough to overlap many of the excited states, but at the same time, they are much smaller than the energy splitting ω_{12} of the lower levels. Such an excitation scheme could be implemented, for example, in a Cs atom: starting from the $7s$ electronic state, a 150-fs ultrashort pulse with a central wavelength of 794 nm can excite a wave packet centered at $\bar{n}=44p$. (Nanosecond dye lasers can be used to populate the initial $7s$ state.) A coupling ultrashort pulse with a central wavelength of 1.04 μm would then connect the $8s$ state to the upper manifold of states, which is approximately harmonic. Although both pulses have large bandwidths, the $7s$ and $8s$ states are spaced far enough apart that neither pulse can alone connect both energy levels. The hyperfine structure of the Cs atom (or any other alkali-metal atom for that matter) can be ignored and plays no role since this structure is many (four or five) orders of magnitude smaller than the ultrashort pulses' bandwidth. (Possible losses induced by ionization from the intermediate states will be neglected. The typical pulse intensities to be considered here are on the order of 10^9 W/cm^2, and thus five orders of magnitude lower than the intensities at which ionization becomes significant [19].)

In the absence of the coupling pulse, when an ultrashort probe pulse is applied to the ground-state atom, many of the states within the excited manifold are coherently excited, and a wave packet is created. This wave packet oscillates in the excited manifold with a characteristic time $T=2\pi/\Delta$—where Δ is the level spacing—corresponding to the quantum beat period between the excited states. The motion of the wave packet reshapes the total electric field in the medium through the emission of secondary pulses that appear in the free-induction decay that follows the input probe pulse [5–8]. These secondary pulses, which are emitted by the induced coherent polarization and correspond to wave-packet recurrences, are separated in time by the beat period T. In the limit of a large number of excited states, the secondary pulses have the same temporal profile as the input pulse. As these impulses propagate through an extended collection of such multi-Λ atoms, their temporal profile (envelope and phase) remains unchanged [8]. The amplitude of the pulses is, however, modified due to propagation. Figure 2(a) shows the propagation of a weak input Gaussian pulse along with the first secondary pulse emitted by the atoms. While the input probe pulse experiences an exponential decay, the secondary pulse exhibit a more complex dependence on propagation distance, first showing an increase in amplitude, followed by decay [8]. Inside the medium, the secondary pulse grows to a peak amplitude comparable to that of the input pulse. The propagation of both impulses is characterized by an absorption coefficient α that is determined by the beat period T [8].

However, when a much stronger coupling pulse is present, the propagation dynamics is modified significantly, as shown in Fig. 2(b). In the presence of the coupling pulse, the atom-field interaction can be described in terms of "bright" and "dark" superpositions of states $|1\rangle$ and $|2\rangle$. Since the coupling pulse is much stronger than the probe pulse, the atom is, initially, mostly in the dark state; electromagnetically induced transparency of the probe pulse takes place, greatly suppressing the excitation of a wave packet in the excited manifold [17]. The very small population initially in the bright state is transferred to the excited manifold, creating a weak polarization in the medium, and subsequently, very weak secondary pulses. The weak secondary pulses are amplified during propagation, but to a significantly smaller peak amplitude than that reached in the absence of the coupling pulse. Also, up to the propagation distance shown in the figure, the probe pulse propagates undamped.

The paper is organized as follows: In Sec. II, the Maxwell-Bloch equations describing the propagation of the ultrashort pulses in these multi-Λ media are derived. In Sec. III, the propagation of the probe pulse is then investigated for the case when the coupling pulse is absent. And then in Sec.

IV, propagation of the probe pulse under EIT is considered. Section V concludes the paper.

II. THE MAXWELL-BLOCH PROPAGATION EQUATIONS

To derive the Maxwell-Bloch propagation equations, I start by writing the electric field of the probe and coupling pulses as $E_p(z,t)=E_{p,0}f(z,t)\exp[-i(kz-\omega_p t)]+\text{c.c.}$ and $E_c(t)=E_{c,0}g(z,t)\exp[-i(kz-\omega_p t)]+\text{c.c.}$, respectively. Both dimensionless functions $f(z,t)$ and $g(z,t)$ are complex and slowly varying (envelope as well as phase); $\omega_{c,p}$ are their carrier frequencies. At the entrance to the medium, the two pulses are matched: $f(0,\tau)=g(0,\tau)$. Both input pulses have temporal widths that are shorter than the quantum beat period T of the excited states and many orders of magnitudes shorter than the polarization decay time T_2 or any other relaxation time of the system.

If a, c, and b_n are the probability amplitudes of the bare $|1\rangle$, $|2\rangle$, and $|n\rangle$ states, respectively, then, under the rotating-wave approximation, Schrödinger's equations for the amplitudes are [8]

$$\dot{a}(z,t)=0.5i\Omega_p f^*(z,t)\sum_n b_n(z,t)e^{-i\delta_n t}, \tag{1a}$$

$$\dot{c}(z,t)=0.5i\Omega_c g^*(z,t)\sum_n b_n(z,t)e^{-i\delta_n t}, \tag{1b}$$

$$\dot{b}_n(z,t)=0.5i[\Omega_p f(z,t)a(z,t)+\Omega_c g(z,t)c(z,t)]e^{i\delta_n t}, \tag{1c}$$

where $\Omega_p\equiv 2d_1E_{p,0}/\hbar$ and $\Omega_c\equiv 2d_2E_{c,0}/\hbar$ are real Rabi frequencies; $\delta_n=\omega_n-\omega_p$ is the detuning between the probe's carrier frequency and level n; the eigenfrequency of level n is ω_n; and the summations are carried out over the number of levels in the excited manifold of states. The electric dipole moments $d_{1,2}$ for transitions between the lower and excited states were assumed to be constant over the excited states accessed by the pulses' spectra. The atom is initially in its ground state: $a(0,0)=1$.

In an optically thick medium, the spatial evolution of the probe and coupling fields is, under the slowly-varying-envelope approximation, dictated by the induced polarizations in the medium. Propagation of these fields is described by the reduced wave equations [6,8]

$$\frac{\partial}{\partial z}[\Omega_p f(z,\tau)]=2i\mu_1\sum_n P_{1,n}(z,\tau), \tag{2a}$$

$$\frac{\partial}{\partial z}[\Omega_c g(z,\tau)]=2i\mu_2\sum_n P_{2,n}(z,\tau), \tag{2b}$$

where $P_{1,n}(z,\tau)=[b_n(z,\tau)\exp(-i\delta_n\tau)]a^*(z,\tau)$ and $P_{2,n}(z,\tau)=[b_n(z,\tau)\exp(-i\delta_n\tau)]c^*(z,\tau)$ are the polarizations induced in the medium by the probe and coupling fields, respectively; $\mu_{1,2}=\omega_L\mathcal{N}d_{1,2}^2/\epsilon_0\hbar c$ is the coupling constant between the field and the polarization; $\mathcal{N}$ is the number density; ϵ_0 is the permittivity of the multilevel medium for nonresonant transitions; and $\tau=t-z/v_g$ is the local time, with v_g being the group velocity of the pulse at the center of the spectrum.

Equations (1) and (2) are the Maxwell-Bloch equations, and they are to be solved together [with t replaced by the local time τ in Eqs. (1)]. In the case of a large number of states in the excited manifold, these equations can be simplified by eliminating the excited states. Integrating Eq. (1c) gives as the total probe-induced polarization

$$\sum_n b_n(z,\tau)a^*(z,\tau)e^{-i\delta_n\tau}=0.5ia^*(z,\tau)\int_{-\infty}^{\tau}ds[\Omega_p f(z,s)a(z,s)+\Omega_c g(z,\tau)c(z,s)]\sum_n e^{-i\delta_n(\tau-s)}. \tag{3}$$

Following the procedure described in Ref. [8], the Poisson sum formula [20] can then be used to show that

$$\sum_n e^{-i\delta_n(\tau-s)}\approx T\sum_{m=-\infty}^{\infty}\delta(\tau-s-mT), \tag{4}$$

where it was assumed that the probe's center frequency is tuned to a resonance. In deriving the above equation, it was assumed that the number of levels in the excited manifold could be extended to an infinitely large number since the contributions of levels with large detunings δ_n are cut off by the exponentials. Equation (3) can now be easily integrated. Substituting the resulting polarization into Eq. (2a), gives, for times $0<\tau\leqslant T$,

$$\frac{\partial}{\partial z}f(z,\tau)=-\frac{\alpha}{2}[f(z,\tau)a(z,\tau)+(\Omega_c/\Omega_p)g(z,\tau)c(z,\tau)]a^*(z,\tau), \tag{5}$$

where $\alpha=\omega_L\mathcal{N}d_1^2T/\epsilon_0\hbar c$ is the absorption coefficient of the medium. Unlike the conventional two-level absorption coefficient, which is determined by the polarization decay time T_2, α is here determined by the beat period T of the excited-state manifold [8]. [The assumption of an infinitely large number of levels is not necessary to arrive at Eq. (5). For times $0<\tau\leqslant T$, the sum of exponentials in the right-hand side of Eq. (3) will behave very much like Dirac's δ function, even for anharmonic systems [19].]

The excited states were thus removed from the wave equation. By restricting the local time to the interval $0<\tau\leqslant T$, the resulting wave equation describes only the propagation of the input pulse; the secondary emitted pulses are left out, and their propagation is not described by Eq. (5).

III. PROPAGATION IN THE ABSENCE OF THE COUPLING PULSE

When the coupling pulse is absent ($\Omega_c=0$), the probe pulse evolves according to

$$\frac{\partial}{\partial z}f(z,\tau)=-\frac{\alpha}{2}f(z,\tau)|a(z,\tau)|^2, \tag{6}$$

with the dynamics of the ground-state probability amplitude being given by

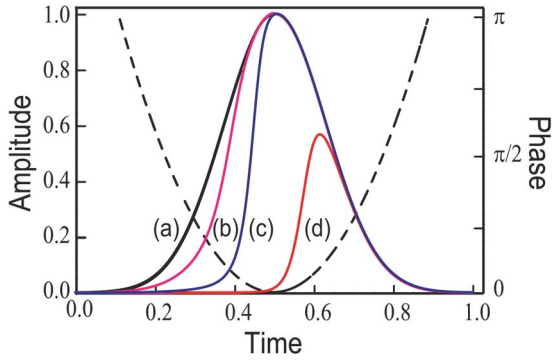

FIG. 3. (Color online) The probe pulse, with envelope function given by Eq. (8), at different positions inside the medium: αz= (a) 0, (b) 2, (c) 6.5, and (d) 19 in the nonperturbative regime of excitation. Solid lines represent envelope, and the dashed line represents phase. Here, $\Omega_c=0$. Time is measured in units of the excited states' beat period T.

$$\dot{a}(z,\tau) = -\frac{1}{8}(\Omega_p^2 T)|f(z,\tau)|^2 a(z,\tau). \tag{7}$$

Equation (7) is obtained by integrating Eq. (1c) with the help of Eq. (4) and substituting the result into Eq. (1a) for the ground-state probability amplitude.

In the perturbative regime of excitation, the ground state is weakly depleted, and $|a(z,\tau)|^2 \approx 1$ throughout excitation. It is then straightforward to show from Eq. (6) that the input probe pulse follows Beer's law and is attenuated exponentially during propagation, with its input temporal form (both envelope and phase) preserved [8].

In the nonperturbative regime, Eq. (7) can be integrated to show that the ground-state population decays as $|a(z,\tau)|^2 \approx \exp[-(\Omega_p^2 T/8)\int_{-\infty}^{\tau} ds\,|f(z,s)|^2]$. The rate of ground-state decay is thus independent of the pulse's phase and is proportional to its energy up to time τ. This exponential decay of the ground-state population has been discussed previously in the context of electronic-Rydberg and vibrational wave-packet excitations [18,19]. Since $|a(z,\tau)|^2$ is real, the population decay will modify the spatial evolution of the probe's amplitude, but will leave its phase unaltered.

I numerically solved Eqs. (6) and (7) for a Gaussian test probe pulse with a quadratic phase. The input test pulse was written as

$$f(0,\tau) = \exp[-(4\ln 2)(\tau-\tau_0)^2/a^2 + i\varphi(\tau-\tau_0)], \tag{8}$$

where $a=0.3T$ is the pulse's full width at half maximum, $\tau_0=0.5T$, and $\varphi(\tau)=0.6\pi(\tau/a)^2$. The amplitude $E_{p,0}$ was set so that the pulse area was $\Omega_p\int_{-\infty}^{\infty}|f(0,\tau)|d\tau=2\pi$. [With such a pulse width and Rabi frequency, a numerical solution of the complete Maxwell-Bloch equations (1) and (2) shows that the spectrum of the probe pulse overlaps about ten states in the excited manifold.] Under these conditions, the probe pulse transfers all the ground-state population to the excited states.

Figure 3 shows the probe pulse at different positions inside the medium. When the probe pulse has traveled a distance $\alpha z \approx 6.5$, the leading edge is almost completely removed from the pulse [curve c]. Only after this propagation distance does the trailing edge start to be absorbed, followed by a decrease in the pulse's peak amplitude. Unlike the envelope, however, the phase remains the same throughout propagation. The distortion in the temporal profile of the probe pulse can be explained as resulting from the nonlinear depletion of the ground-state population. In the perturbative excitation regime, every point on the profile of the pulse propagates in the same way, following a Beer's law exponential attenuation; there is no pulse distortion in this regime. However, in the limit of strong excitation, different points in the temporal profile will experience different attenuations, leading to distortions.

IV. PROPAGATION IN THE PRESENCE OF THE COUPLING PULSE

The spatial evolution of a weak probe pulse can be altered significantly by the presence of a coupling field. Substituting Eqs. (3) and (4) into Eq. (1b), yields, for times $0<\tau\leqslant T$,

$$\dot{c}(z,\tau) = -\frac{\Omega_c\Omega_p}{8}Tg^*(z,\tau)f(z,\tau) - \frac{\Omega_c^2}{8}T|g(z,\tau)|^2 c(z,\tau). \tag{9}$$

The spatial evolution of the probe pulse is then determined by solving the system of equations comprised of Eqs. (5) and (9), and an equation equivalent to Eq. (5) for the coupling field $g(z,\tau)$. However, when all the atomic population is initially in the ground state, as has been assumed, the coupling field does not suffer any significant change in its temporal profile as a result of propagation. This is because, in the quasi-impulsive excitation regime I am considering here, very little probability amplitude is transferred to state $|2\rangle$ [17]. Thus, a negligible polarization between this lower state and the excited states is created, even for $\Omega_c \lesssim \Omega_p$. To a good approximation $g(z,\tau)\approx f(0,\tau)$, since the pulses are matched at the input. Therefore, Eqs. (5) and (9) are sufficient to study the spatial evolution of the probe pulse. [These assumptions were confirmed by numerically solving the complete Maxwell-Bloch equations (1) and (2).]

Equations (5) and (9) were obtained under the assumption of a harmonic manifold, but these equations can be applied to even mildly anharmonic systems. This observation is true even though Eq. (4) is not valid for anharmonic systems for arbitrary times. Equations (5) and (9) hold for anharmonic systems because up to the first period T, the sum of exponentials in the right-hand side of Eq. (3) still behaves like a Dirac δ function for these systems [19]. Thus, the results to be derived here with these equations are applicable to systems such as Rydberg atoms, for example.

I numerically solved Eqs. (5) and (9) for the Gaussian test probe pulse of Eq. (8). The amplitude $E_{p,0}$ was set so that the pulse area was $\Omega_p\int_{-\infty}^{\infty}|f(0,\tau)|d\tau=0.1\pi$. (A numerical calculation shows that, for $\Omega_c=0$, the probe pulse transfers a small fraction, about 5%, of the ground-state population to the excited states, and its spectrum overlaps about seven states in the excited manifold.) The coupling pulse's envelope function was set to $g(z,\tau)=f(0,\tau)$, so both pulses are initially matched in envelope and phase.

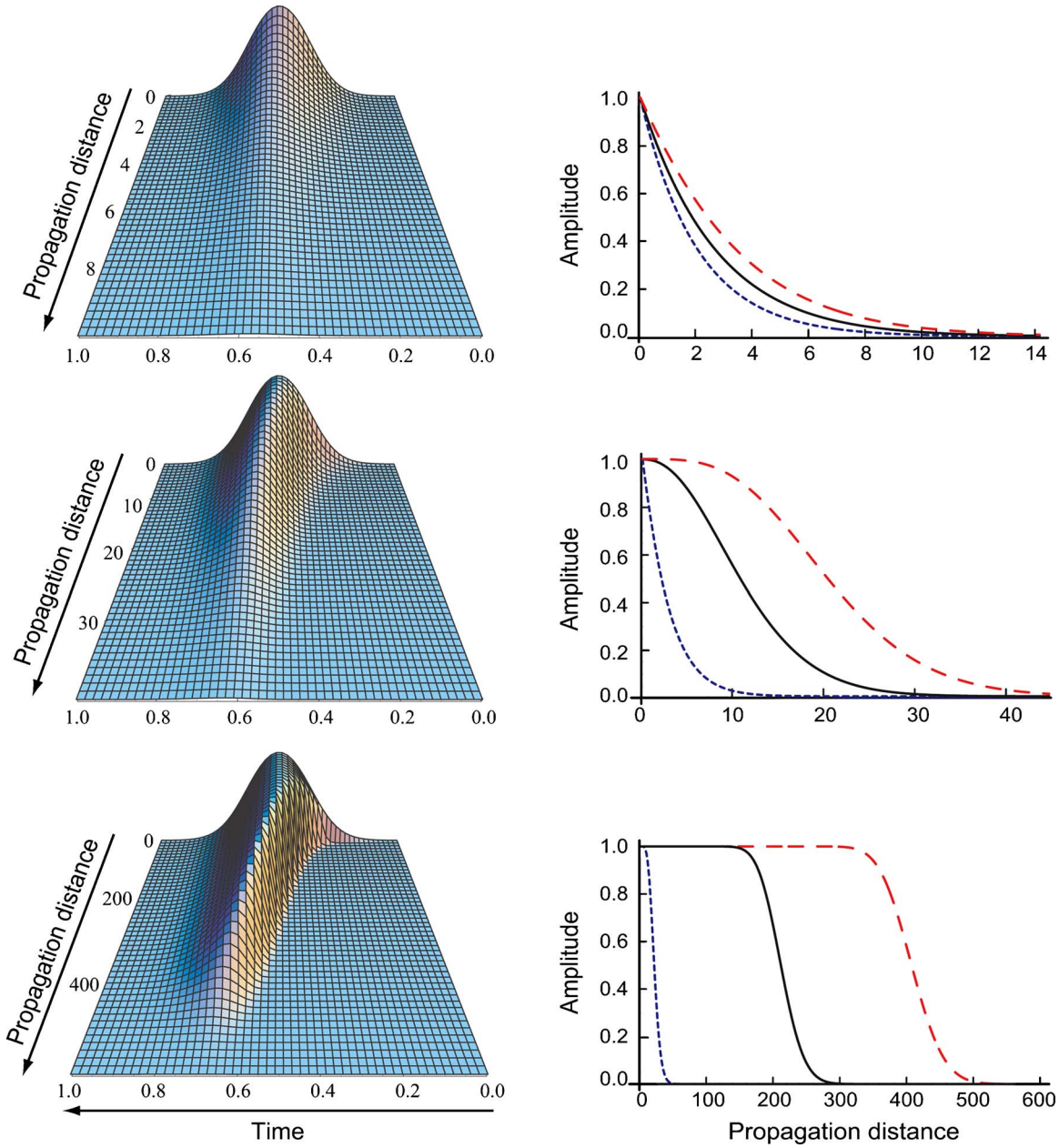

FIG. 4. (Color online) The left column shows the envelope $|f(z,\tau)|$ of a weak probe pulse, with envelope function given by Eq. (8), as a function of propagation distance for $\Omega_c=5\Omega_p$ (top row), $20\Omega_p$ (middle row), and $100\Omega_p$ (bottom row). Time is measured in units of the excited states' beat period T, and propagation distance is measured in units of the inverse absorption coefficient α^{-1}. The right column shows the propagation of three distinct points in the temporal profile of the probe pulse: $\tau-\tau_0=-a$ [dotted (blue) line], 0 [solid (black) line], and $+a$ [dashed (red) line]. The input value of each of these points was normalized to unity for a better comparison.

The left column of Fig. 4 shows the temporal shape of the probe pulse as a function of propagation distance for three values of the coupling Rabi frequency $\Omega_c=5\Omega_p$, $20\Omega_p$, and $100\Omega_p$. As Ω_c increases, the pulse propagates a longer distance into the medium, due to the induced transparency. But the transparency is not complete, and the probe pulse is still absorbed by the medium, eventually disappearing. Unlike the results of Arkhipkin and Timofeev [15], there is no transfer of energy from the probe to the coupling pulse; $g(z,\tau)\approx g(0,\tau)$ throughout propagation. The energy removed from the probe pulse is used to transfer population to the excited states as well as to state $|2\rangle$. (A coherence is created between the two lower levels.) Half of the very small population that is removed from the ground state is placed in the excited manifold and the other half in state $|2\rangle$ through the coupling pulse. Furthermore, a distortion of the temporal shape of the probe pulse also takes place, even though the probe pulse is weak. In the absence of the coupling pulse, every point on the temporal profile follows a Beer's law exponential attenuation. However, when the coupling pulse is present, different points in the temporal profile will experience different attenuations, leading to distortions. The right column of Fig. 4 shows the spatial evolution of three distinct points (the center of the pulse and the two points of half maximum, as defined

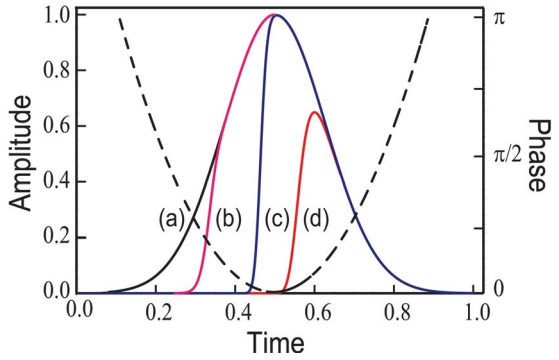

FIG. 5. (Color online) A weak probe pulse, with envelope function given by Eq. (8), at different positions inside the medium: αz = (a) 0, (b) 15, (c) 150, and (d) 330. Solid lines represent envelope, and dashed line represents phase. Here, $\Omega_c = 100\Omega_p$. Time is measured in units of the excited states' beat period T.

for the pulse at $z=0$) on the temporal profile of the probe pulse for the same three values of Ω_c. In general, in the presence of the coupling field, every point in the profile travels a longer distance than that traveled when $\Omega_c=0$ due to the induced transparency. But, even for a relatively low Ω_c, the leading edge of the pulse tends to be absorbed faster than the trailing edge. The higher Ω_c, the more pronounced the difference in absorption between the trailing and leading edges.

The dynamics of the pulse distortion can be seen more clearly in Fig. 5, which shows the probe pulse at different positions inside the medium for the $\Omega_c=100\Omega_p$ case. When the probe pulse has traveled a distance $\alpha z \approx 150$, the leading edge is completely removed from the pulse [curve c]. Further propagation into the medium results in absorption of the trailing edge, followed by a decrease in the pulse's peak amplitude. But the phase remains the same as the probe pulse moves into the medium. The same results were observed for different phases [other than the Gaussian phase of Eq. (8)] of the input pulse, even for a probe phase with a discontinuous π phase change at the center of its envelope. That is, the probe's phase is not altered during propagation, nor is the absorption dynamics modified by the particular form of the phase. This behavior is very similar to that observed in the propagation of a strong probe pulse. Although the probe pulse by itself operates in the weak-excitation regime, the presence of the strong coupling pulse induces a nonlinear behavior in the absorption of the probe pulse. Figures 3 and 5 are very similar. When there is a coupling pulse present, and describing the interaction in terms of the bright superposition state, the population in this bright state decays exponentially, with its rate of decay being determined by the combined energy of probe and coupling pulses [17]. It is possible the probe pulse distortion can be attributed to the nonlinear decay of the bright superposition of states $|1\rangle$ and $|2\rangle$, although this assertion requires further checking. This pulse distortion behavior is the opposite of what happens in the propagation of a long pulse in a three-level medium. In that case, the pump pulse "protects" the probe pulse from changing as it propagates along the medium [11].

Even if the coupling pulse is not matched to the probe pulse, but delayed with respect to the probe, the dynamics of the total field inside the medium is altered. Figure 6 shows the propagation of a probe pulse in the presence of a coupling pulse that is delayed with respect to the probe by one beat period T. Both input pulses are given by Eq. (8), with $\varphi(\tau)=0$, $\Omega_p\int_{-\infty}^{\infty}|f(0,\tau)|d\tau=0.1\pi$, and $\Omega_c=100\Omega_p$. The coupling pulse strongly suppresses emission by the induced polarization, and no impulse is observed at $\tau=1.5T$. The coupling pulse does not simply remove population from the excited states since reemited impulses are observed at later times. The first reemited impulse occurs at $\tau=2.5T$. This impulse corresponds to the first impulse observed in the absence of the coupling pulse, and not to the second impulse. In other words, the impulses are shifted in time by one beat period by the coupling pulse. For greater probe-coupling delays, the reemited impulses that occur before the coupling pulse propagate unchanged; the other impulses are shifted in time by one beat period as well.

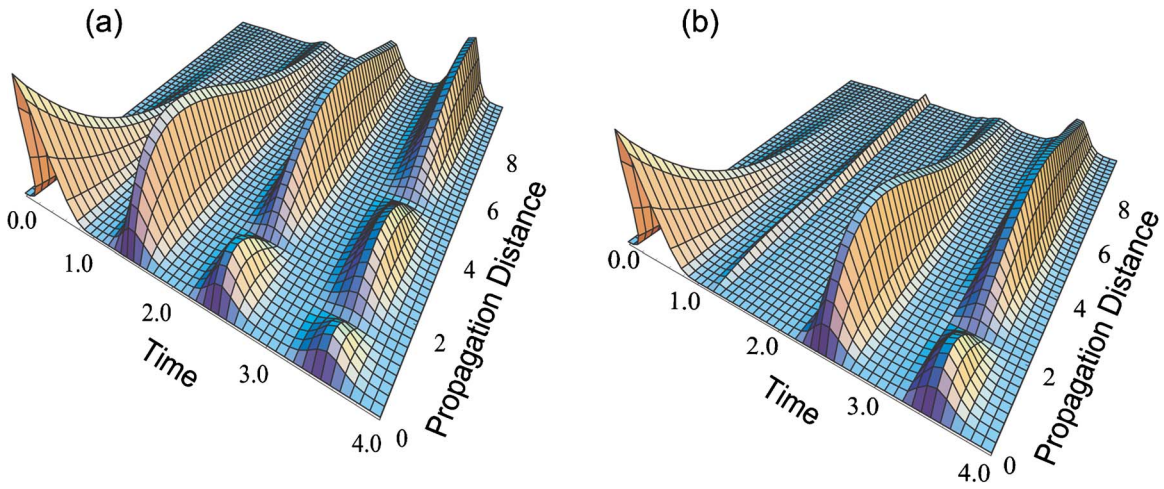

FIG. 6. (Color online) The envelope $|f(z,\tau)|$ of a weak ultrashort Gaussian pulse propagating in an extended collection of multi-Λ media in (a) the absence of a coupling pulse and (b) the presence of a coupling pulse delayed by one period T with respect to the input probe pulse. Time is measured in units of the excited states' beat period T, and propagation distance is measured in units of the inverse absorption coefficient α^{-1}.

V. CONCLUSIONS

I have here studied the propagation of a weak ultrashort probe pulse in a medium (dubbed a multi-Λ system) consisting of two lower states and a manifold of excited states. It is shown that the propagation of such pulses is greatly modified by the presence of a stronger coupling pulse. The coupling pulse—initially matched to the probe pulse in envelope and phase—distorts the probe's envelope by gradually removing its leading edge as the pulse propagates into the medium. However, the probe's phase does not change during propagation, nor does its particular form affect the absorption dynamics. A similar behavior is observed in the propagation of a strong probe pulse without the coupling pulse. Furthermore, the coupling pulse strongly suppresses the excitation, by the probe pulse, of a wave packet in the excited states throughout the medium. In this way, only extremely weak secondary pulses are generated by the induced polarization in the medium. However, if the coupling pulse is delayed from the probe pulse by an integer multiple of the beat period, polarization-induced emission can be suppressed at the time the coupling pulse enters the medium. The reemitted impulses that follow are then temporally shifted by one beat period.

The matched coupling pulse also transfers a very small probability amplitude from the ground state to the other lower state, thus creating a coherence between the two states. This coherence can possibly be exploited as a way to store and retrieve both the envelope and phase of an ultrashort pulse in an atomic medium. For example, reapplying the coupling pulse could create a polarization between the excited and ground states that would recreate the probe pulse.

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