Mixed electromagnetically and self-induced transparency

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Abstract: We show that application of self-induced transparency (SIT) solitons as a driving field in V-type electromagnetically induced transparency (EIT) leads to "mixed induced transparency" (MIT) that nicely combines the best features of both SIT and EIT.

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1 Introduction

Among the most important milestones concerning propagation of coherent optical pulses through multilevel absorbers are self-induced transparency in two-level systems by McCall and Hahn [1, 2], simultons by Konopnicki and Eberly [3], counterintuitive pulse sequences by Oreg, Hioe and Eberly [4, 5], matched pulses via electromagnetically induced transparency (EIT) by Harris [6, 7], and the dressed-state pulses by Eberly, Pons, and Haq [8]. The phenomena of matched and dressed-state pulses are linked to electromagnetically induced transparency (EIT), which is particularly interesting because it offers a wide variety of applications ranging from lasers without population inversion (for the earliest papers on EIT/LWI see [9, 10, 11, 12, 13], for reviews on EIT/LWI see [14, 15, 16, 17, 18, 19]) to new trends in nonlinear optics [20, 21, 22, 23, 24, 25].

The EIT of a weak probe pulse relies on a two-photon coherence which is induced by the joint action of the probe field and a strong driving field in a three-level system. In order to make an optically thick medium transparent, the driving field must preserve its intensity all along the path. This condition appears rather demanding, especially for V-type systems, because the driving field couples fully populated state and empty excited state of the system. This is the main drawback of V EIT. The strong drive field provides the transparency for the probe field, but itself remains subject to resonant absorption and dispersion. In order to fix this we appeal to SIT [1, 2, 26, 27, 28, 29, 30] and apply the effect to achieve transparency of the driving transition in V EIT experiments. Then, both transitions appear to be transparent, one — in the sense of SIT, the other — in the sense of EIT. Such "mixed-induced transparency" (MIT) constitutes the subject of our study.

Previous related studies are associated with lossless propagation of simultons in three-level systems [3, 31, 32, 33] and N-type systems [34], Raman amplification of ultrashort pulses [35] in V configuration, theoretical and experimental studies on transparency enhancement for an ultrashort weak-pulse propagation in an inhomogeneously-broadened V-type medium [36, 37], and propagation of ultrashort pulses in phaseonium [38].

Equivalent durations of involved pulses is common feature of the above studies. In contrast, our interest here will be in suppression of absorption for both weak and *long* probe pulse when a sequence of *short* 2π-pulses drives the adjacent transition of the V-type atom, as shown in Fig. 1. We will demonstrate that the application of a sequence of 2π-pulses as the drive field in V-type configuration results in suppression of population transfer produced by a weak long pulse, thereby creating conditions of transparency for the weak field. During propagation, the population transfer on the probe transition is continuously minimized by a coherent 2π-pulse-induced rearrangement of energy (reshaping) inside the probe pulse.

2 Motivation

Let us consider a V-type three-level system with the ground state $|b\rangle$ and excited states $|a\rangle$ and $|c\rangle$, as shown in Fig. 1. Transition $|b\rangle \leftrightarrow |c\rangle$ of frequency ω_{bc} is driven by a field $E_\Omega(z,t) = \mathcal{E}_\Omega(z,t)\exp(ik_\Omega z - i\nu_\Omega t)$. A probe field $E_\alpha(z,t) = \mathcal{E}_\alpha(z,t)\exp(ik_\alpha z - i\nu_\alpha t)$ is applied to the transition $|b\rangle \leftrightarrow |a\rangle$ of frequency ω_{ab}. Here $\mathcal{E}_\Omega(z,t)$ and $\mathcal{E}_\alpha(z,t)$ are the slowly varying envelopes of the electric field. They are related to the Rabi frequencies Ω and α according to $\mathcal{E}_\Omega = \hbar\Omega/\wp_{cb}$ and $\mathcal{E}_\alpha = \hbar\alpha/\wp_{ab}$, where $\wp_{cb}$ and $\wp_{ab}$ are dipole matrix elements of the transitions $|c\rangle \leftrightarrow |b\rangle$ and $|a\rangle \leftrightarrow |b\rangle$. The carrier waves have wave numbers k_Ω and k_α, and frequencies ν_Ω and ν_α.

We start from a simple example which discloses advantages arising from the combination of SIT with EIT: a *single* atom of V-type with a weak long probe pulse applied

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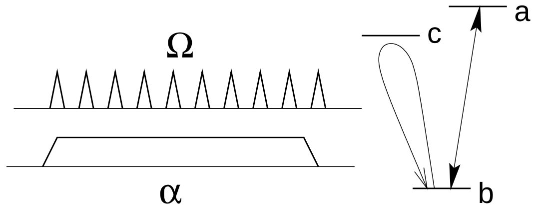

Fig. 1. Pictorial illustration of MIT via a V-type atom under the action of a sequence of 2π-pulses (Ω) and a continuous probe field (α). The overall effect of a 2π-pulse applied to $c \leftrightarrow b$ transition is in flipping the sign of the wave function of the ground state.

at one leg of the system and a sequence of strong sharp pulses driving the other transition, as shown in Fig. 1. The probability of the atomic decay is proportional to the population of the excited state. Hence, the more the atom is excited by the weak pulse to the upper state $|a\rangle$, the larger rate of spontaneous decay is expected. We shall show that while the adjacent transition $|c\rangle \leftrightarrow |b\rangle$ is driven by strong pulses, the excitation of $|a\rangle$ state by the weak pulse is greatly suppressed. Therefore, the absorption rate will be proportionally decreased.

The action of the probe pulse to promote atoms from $|b\rangle$ to $|a\rangle$, is described by the Hamiltonian

$$V = \hbar \frac{\alpha(t)}{2} \left(|a\rangle\langle b| e^{i\Delta t} + \text{adj.} \right), \tag{1}$$

with $\Delta = \omega_{ab} - \nu_\alpha$ as the offset of the field carrier frequency and the optical transition. Expanding state vector $|\Psi_t\rangle$ in eigenstates of the atomic Hamiltonian as $|\Psi_t\rangle = a(t)|a\rangle + b(t)|b\rangle + c(t)|c\rangle$ and substituting into the Schrödinger equation gives the standard set of coupled equations for the probability amplitudes

$$\dot{a} = -\frac{i}{2}\alpha\, \mathrm{e}^{i\Delta t} b\,, \tag{2}$$

$$\dot{b} = -\frac{i}{2}\alpha\, \mathrm{e}^{-i\Delta t} a\,, \tag{3}$$

$$\dot{c} = 0\,. \tag{4}$$

Only three relevant states $|a\rangle$, $|b\rangle$, and $|c\rangle$ are kept for our discussion.

It is convenient to represent the time evolution of the atomic state $|\Psi_t\rangle$ by unitary transformation $|\Psi_t\rangle = U_\alpha(t)|\Psi_0\rangle$ of initial state $|\Psi_0\rangle$ given by

$$U_\alpha(t) = \mathcal{A}|a\rangle\langle a| + \mathcal{B}|b\rangle\langle b| + |c\rangle\langle c| + (\mathcal{C}|a\rangle\langle b| - \text{adj.})\,. \tag{5}$$

Time-dependent coefficients $\mathcal{A}$, $\mathcal{B}$ and $\mathcal{C}$ are derived from (2)-(4):

$$\mathcal{A}(t) = \left(\cos\frac{\tilde{\alpha}}{2}t - i\frac{\Delta}{\tilde{\alpha}}\cos\frac{\tilde{\alpha}}{2}t\right)\exp\left[i\frac{\Delta}{2}t\right] \tag{6}$$

$$\mathcal{B}(t) = \left(\cos\frac{\tilde{\alpha}}{2}t + i\frac{\Delta}{\tilde{\alpha}}\cos\frac{\tilde{\alpha}}{2}t\right)\exp\left[-i\frac{\Delta}{2}t\right] \tag{7}$$

$$\mathcal{C}(t) = -i\frac{\alpha}{\tilde{\alpha}}\sin\frac{\tilde{\alpha}}{2}t\exp\left[i\frac{\Delta}{2}t\right] \tag{8}$$

with $\tilde{\alpha} = \sqrt{\alpha^2 + \Delta^2}$ as effective Rabi frequency.

For simplicity, we consider the case of exact resonance. Then, Eqs. (5)-(8) are reduced to

$$U_\alpha(t) = \cos\left[\frac{\alpha}{2}t\right](|a\rangle\langle a| + |b\rangle\langle b|) - i\sin\left[\frac{\alpha}{2}t\right]\sigma_x(a,b) + |c\rangle\langle c| \tag{9}$$

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with $\sigma_x(a,b) \equiv |a\rangle\langle b| + |b\rangle\langle a|$. It is instructive for further derivations to write Eq. (9) as

$$U_\alpha(t) = \exp\left[-i\frac{\alpha}{2}t\,\sigma_x\right] + |c\rangle\langle c| = \prod_{n=0}^{N} \exp\left[-i\frac{\alpha}{2}\tau\,\sigma_x\right] + |c\rangle\langle c|\,, \tag{10}$$

Bearing in mind further application of a sequence of N 2π-pulses, the time interval t in (10) is divided into N equal parts, each of duration τ. At time $t = N\tau = \pi/\alpha$, the atom starting in state $|b\rangle$ will be fully excited to the upper state $|a\rangle$. This is a direct indication of high rate of absorption of the probe field (if the latter would have been included in the model).

We now switch on a train of strong short pulses coupling $|b\rangle$ and $|c\rangle$ states. The pulses are separated by equal intervals τ and their durations τ_p are shorter than τ, see inset of Fig. 2. Since, the driving pulses are much shorter and stronger than the probe pulse, we may, to a reasonable approximation, ignore the dynamics on the probe transition during τ_p interval. Then, evolution operator during the excitation of the $|c\rangle \leftrightarrow |b\rangle$ transition reads as

$$U_\theta(t) = \cos\left[\frac{\theta(t)}{2}\right](|b\rangle\langle b| + |c\rangle\langle c|) + |a\rangle\langle a| - i\sin\left[\frac{\theta(t)}{2}\right](|b\rangle\langle c| + |c\rangle\langle b|)\ , \tag{11}$$

where $\theta(t) \equiv (\wp_{ac}/\hbar)\int \mathcal{E}_\Omega(t')dt'$. For specific pulses which span a total area of 2π, i.e. $\theta = 2\pi$ after integration over all pulse envelope, the evolution operator reduces to a simple form,

$$U_{2\pi}(t) = |a\rangle\langle a| - (|b\rangle\langle b| + |c\rangle\langle c|)\ . \tag{12}$$

The total evolution of the single atom can be written as a successive alternation of $|a\rangle \leftrightarrow |b\rangle$ and $|c\rangle \leftrightarrow |b\rangle$ excitations: first, caused solely by the weak pulse, and the second — solely by the 2π-pulse,

$$U(t) = U_{2\pi}^{(N)}U_\alpha(\tau)U_{2\pi}^{(N-1)}\ldots U_{2\pi}^{(2)}U_\alpha(\tau)U_{2\pi}^{(1)}U_\alpha(\tau)\,. \tag{13}$$

The unique feature of a 2π-pulse is in the coherent excitation and de-excitation of an atom in such a way, that the total population difference comes back to its original value after the pulse passed. The only change is in the π phase shift of the lower state. A successive application of two 2π-pulses returns the system accurately back to the original state. Then, it appears quite natural due to the symmetry of the problem that product of four evolution matrices $U_{2\pi}U_\alpha(\tau)U_{2\pi}U_\alpha(\tau)$ gives a characteristic evolution block. The product of the first three yields $U_{2\pi}U_\alpha(\tau)U_{2\pi} = \exp(i\frac{\alpha}{2}\sigma_x\tau)$, and after adding the fourth matrix one finally gets

$$U_{2\pi}U_\alpha(\tau)U_{2\pi}U_\alpha(\tau) = \exp\left(i\frac{\alpha}{2}\sigma_x\tau\right)\,\exp\left(-i\frac{\alpha}{2}\sigma_x\tau\right) = \mathbf{1}\,. \tag{14}$$

So, the atom comes back to the original state. Certainly, Eq. (13) gives the same result for any even N.

Concluding, in contrast to single weak pulse propagation where the analysis shows full excitation of $|a\rangle$ state in time $t = \pi/\alpha$, the SIT-assisted dynamics provides no population transfer from the lower $|b\rangle$ state during the same interval of time. Thus, for the latter configuration, the weak pulse remains insensitive to any decay process from the upper state. This is the keystone result which gives rise to new type of transparency — MIT.

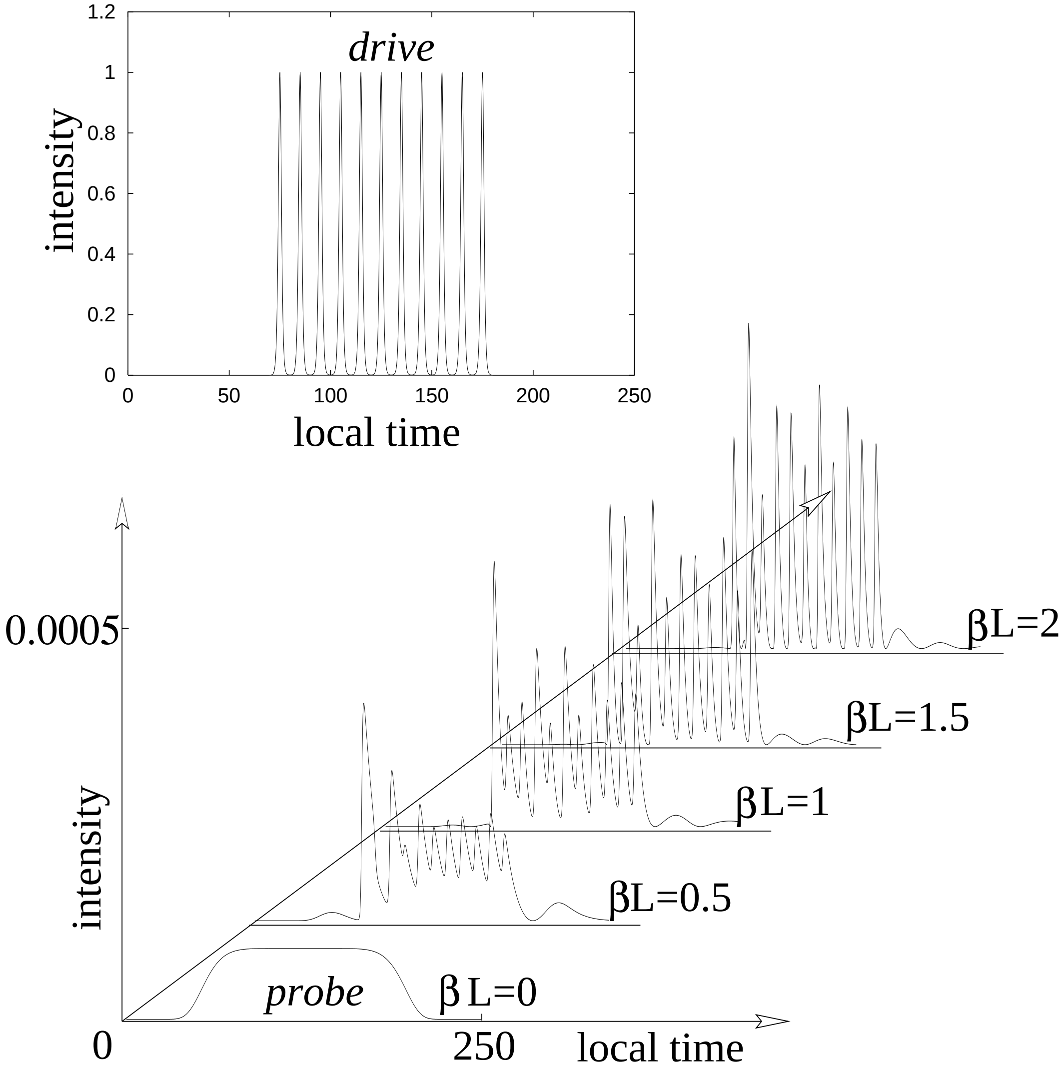

Fig. 2. Time dependent plots of intensity of the probe pulse showing the effect of propagation over 2 Beer's lengths in V-type system driven by a sequence of eleven 2π-pulses (shown in the inset). In all plots local time $(t - z/v)/\tau_p$ (with v as phase velocity of light in the medium) is measured in units of the strong pulse duration τ_p and distance in Beer's lengths, defined here as $\beta \equiv (\kappa_\Omega \tau_p)^{-1}$. For Rabi frequencies of the strong field, Ω, and the weak probe, α, we define intensity as $|\Omega\tau_p|^2$ and $|\alpha\tau_p|^2$, correspondingly. Initially, the weak pulse has a super-Gaussian shape $\alpha = 0.01 \exp[-(t/80\tau_p)^8]$, see the snapshot at $\beta L = 0$, and the strong pulses are identical 2π-solitons of self-induced transparency separated by 10 their own durations from each other: $\Omega = \sum_{n=-5}^{5} \mathrm{sech}[(t + 10n\tau_p)/\tau_p]$, see inset.

3 Self-enhancement of MIT during the propagation

In this section, we turn to a more accurate description of the transparency phenomenon by lifting major approximations assumed in the above analysis while keeping similarity with the original model as close as possible. Furthermore, generalizing the model from a single atom (thin layer) description to an optically thick V-type medium, we thereby introduce the important propagation effects into our consideration. The latter will allow us to follow evolution of the system, and finally conclude that the transparency is enhanced through the self-consistent dynamics of the matter and the fields.

For proper accounting for propagation effects, the matter equations driven by the optical fields should be supplemented with field equations driven by polarization of the matter. In the slowly varying envelope approximation, the temporal and spatial evolution of the field envelopes is governed by wave equations:

$$\left[\frac{\partial}{\partial z}+\frac{n_\Omega}{c}\frac{\partial}{\partial t}\right]\Omega(t,z)=i\kappa_\Omega\rho_{cb}\,, \tag{15}$$

$$\left[\frac{\partial}{\partial z}+\frac{n_\alpha}{c}\frac{\partial}{\partial t}\right]\alpha(t,z)=i\kappa_\alpha\rho_{ab}\,, \tag{16}$$

with propagation constants

$$\kappa_\Omega=k_\Omega\wp_{cb}^2N/\epsilon_0\hbar\,n_\Omega \quad \text{and} \quad \kappa_\alpha=k_\alpha\wp_{ab}^2N/\epsilon_0\hbar\,n_\alpha\,. \tag{17}$$

ρ_{ab} and ρ_{cb} are off-diagonal density matrix elements in the rotating frame, N is the density of resonant atoms, n_Ω and n_α are host refractive indices at frequencies of the pump and probe fields, respectively. Instead of Schrödinger description of the V-type atom in section II, we shall use a more general, density matrix approach [39]. The resultant coupled set of the field and matter equations is solved numerically.

For simplicity, coupling constants and nonresonant refractive indices on the two transitions are put equal: $\kappa_\Omega=\kappa_\alpha$ and $n_\Omega=n_\alpha$. We consider a weak pulse with a super-Gaussian shape shown in Fig. 2 at $\beta L=0$. Duration of the pulse is rather long in a sense that area under its envelope is only slightly less than π, but it is still shorter than any decay process in the system. In what following, we shall ignore all relaxation processes since they are not relevant to our discussion. The inset of Fig. 3 shows a significant population transfer produced by the pulse: almost all atoms appear to be excited. The absorbed energy is scattered very rapidly resulting in fast decay rate and total dispersion of the weak pulse, see curve (1) in Fig. 4(a). The pulse area also vanishes sharply, see curve (1) in Fig. 4(b).

Following analysis in the previous section we now apply a driving field in the form of a sequence of 2π-pulses on the adjacent transition. Bearing in mind further consideration of propagation effects, among all possible 2π-pulses we choose those which have a sech-shape envelope, see inset in Fig. 2. As has been shown by Matulic and Eberly in [28], they are the only shape-preserving 2π-solutions for an optically thick two-level absorber. Switching on the driving field dramatically decreases the amount of population transfer of atoms to $|a\rangle$ state, compare Fig. 3 at $\beta L=0$ and inset. Certainly, the suppression cannot be full as it is in the analytic model, since the 2π-pulses here are neither infinitely short nor infinitely strong. Nevertheless, the transparency effect finds its clear manifestation.

By taking into account propagation effects, we step into a domain which is beyond the analytical predictions. We find that the transparency effect is enhanced even more when the pulses are allowed to propagate deeper into the medium. Snapshots of population at $\beta L=0.5$, $\beta L=1$, $\beta L=1.5$, and $\beta L=2$ in Fig. 3 display orders of magnitude less excitations than it is at the entrance, i.e. at $\beta L=0$. Energy of the weak pulse experiences

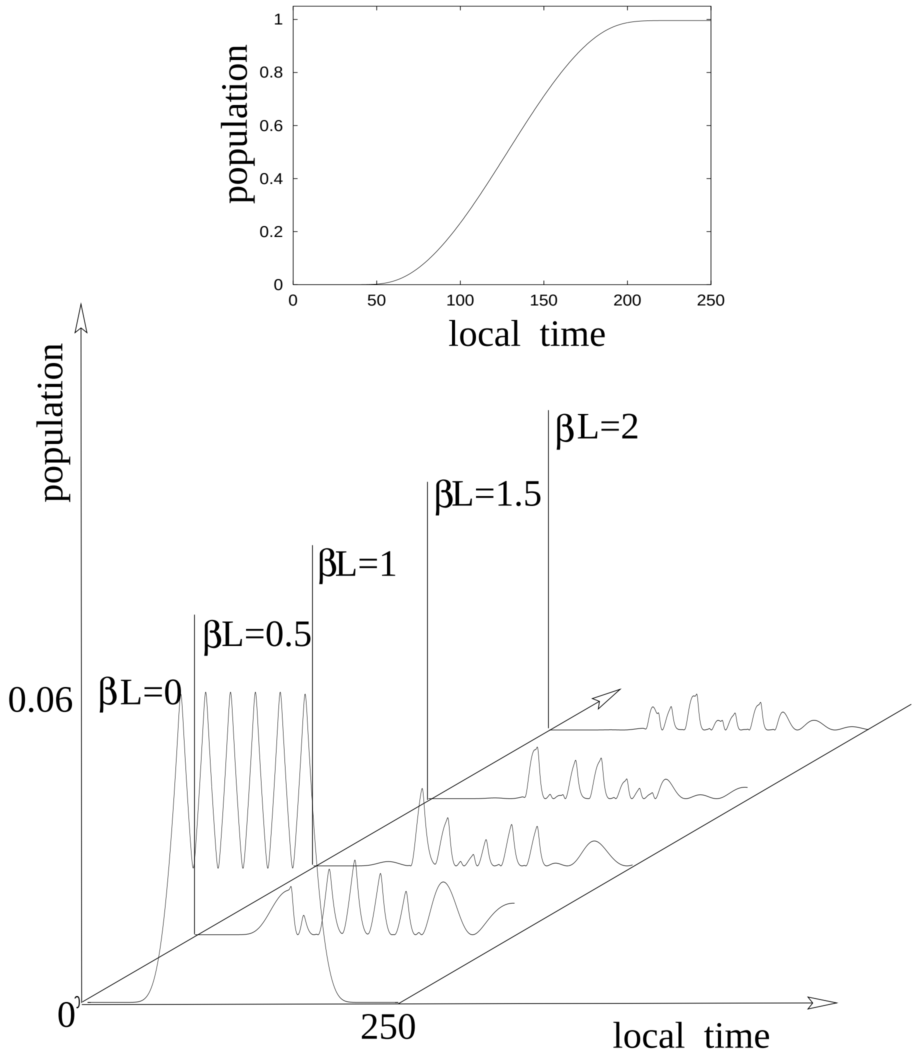

Fig. 3. Time dependent plots of population of the $|a\rangle$ state at the input and after propagating 0.5, 1, 1.5, and 2 Beer's lengths. Inset shows time evolution of ρ_{aa} at the entrance plane ($\beta L = 0$) in absence of driving pulses on the $|c\rangle \leftrightarrow |b\rangle$ transition.

extremely slow decay, see curve (2) in Fig. 4(a), and clearly the loss of energy (which is less than a half of initial value) cannot be the reason for such enormous decrease in population transfer.

The explanation of MIT starts with the observation of the ”pulse locking” effect, where a strong 2π-pulse and a weak pulse propagate together with the same group velocity, and without a substantial change of shape while propagating, [35]. Fig. 2 shows how the long weak pulse splits into a sequence of short more intense sub-pulses, the peaks of which are synchronized with locations of 2π-pulses. Energy of the probe pulse is continuously ”drawn” into the regions where intensity of the driving field reaches maximum, so that such master-slave combination of the two fields propagates through the medium with group velocity dictated by the 2π-pulses.

The effect of pulse locking assures no dispersive breakdown of a weak pulse and simultaneously its perfect spatio-temporal synchronization with a pump pulse. This creates ideal conditions for Raman amplification of ultrashort pulses in a V-type medium.

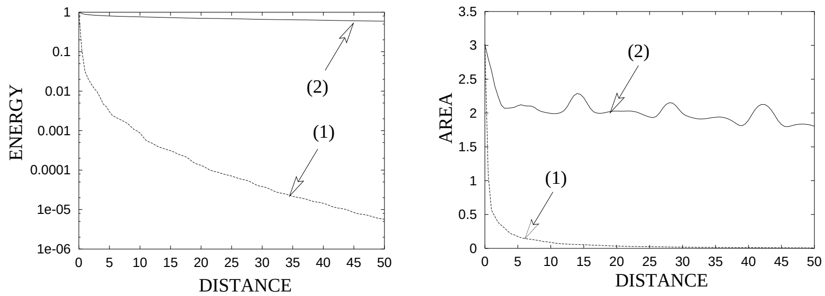

Fig. 4. Energy and Area versus distance for the weak pulse: (1) with no driving field; (2) with driving field in the form of 2π-pulses, see figure caption for Fig. 2. Energy is normalized to its initial value.

Ref. [35] reports about 100% efficiency of the amplification mechanism. Here, the pulse locking also constitutes a necessary condition for the enhanced transparency by preventing a dispersion of the weak pulse, however the main *source* of the transparency, that is *protection against spontaneous emission*, is of different nature, see also [40].

The less atoms are excited to the upper state, the less is the rate of the spontaneous emission, and therefore the better is the transparency for the field. That is, the true signature of the transparency is vanishing fraction of atoms excited to the upper state. With this criterion in hand, we compare the amount of population transfer shown in Fig. 3 before (at $\beta L = 0$) reshaping of the weak pulse and after (say at $\beta L = 2$). So, the population transfer appears to be greatly suppressed *solely due to redistribution of energy* inside the pulse. It is the key conclusion of the present Letter that the matter-field interaction *self-organizes* the shape of the weak field in such a way that the population of $|a\rangle$ state is *minimized*.

4 Discussion

We describe the new type of transparency, MIT, which is induced for a weak long pulse by a sequence of 2π-pulses applied on the adjacent transition of the medium of V-configuration. Based on calculations of the amount of population excited by the weak pulse to the upper state, we estimate the efficiency of the transparency effect. Analytical derivations for a single atom driven by train of infinitely sharp (δ-function-like) 2π-pulses predict perfect transparency for the weak field. Corresponding numerical computations for a non-ideal case confirms appearance of the transparency, though some residual excitation of $|a\rangle$ state arises due to finite intensity of the driving pulses. Then, taking propagation dynamics into account, the residual excitation is shown to be suppressed even further. The transparency enhancement is due to the coherent SIT-pulse-induced rearrangement of energy inside the weak pulse that ultimately results in *minimization* of the population transfer to the upper state $|a\rangle$.

The left plot in Fig. 4 displays a decay of energy of the probe pulse over distance of 50 Beer's lengths. The difference in 5 orders of magnitude between scenarios of free evolution and MIT-assisted propagation provides another direct proof of the transparency effect. Curve (1) on the right plot in Fig. 4 shows the evolution of area of the probe pulse towards 0π, which is characteristic ultimate value for a weak probe field propagating in a two-level medium. In contrast, in presence of 2π-pulses on the adjacent transition, see curve (2), the pulse area is stabilized near a constant nonzero value. This particular number is not unique and varies for different boundary conditions. Many numerical runs have not revealed a noticeable regularity in evolution of the pulse area like the famous Area Theorem for two-level systems, [1, 2, 41].

The pulse locking and the redistribution of energy inside the probe pulse are evidently of a coherent origin. The suppression of population transfer and associated decrease in energy loss of the probe field cannot be explained solely in terms of AC-Stark effect produced by a strong pulse driving the coupled transition. The transient coherent evolution in MIT will be discussed in detail elsewhere.

The idea of combination of SIT and EIT and the manifestation of its power with the above simple model, gives rises to further intriguing questions. What happens if the transitions are Doppler-broadened? How crucial is the condition of the two-photon resonance? What is the role of the two-photon Doppler effect? Such questions naturally follows from SIT physics which is known for its unique ability of making an inhomogeneously broadened medium transparent. On the other hand, the problem gives us a good chance to come back to that "golden oldie" SIT [26, 27, 28, 29, 30, 42], and to pick up some clues on further development of the subject.

We learn from Matulic and Eberly [28] that there are no steady-state single-pulse solutions to the absorber problem other than a 2π sech pulse. However, there are non-shape-preserving pulses described by Lamb in [29] which have total envelope area equal to 4π or 6π, for example. There are also 0π-pulses with non-stationary oscillating amplitude [42]. The infinite pulse-train solutions described by Eberly in [27] as early as in 1969, are especially relevant to our discussion. The train of pulses has an area of $2\pi\infty$, and the envelope shape is determined by certain elliptic functions. Physically such solutions correspond to a *continual* exchange of energy from a steady-state optical wave to the atoms and back. In contrast to our sequence of widely separated 2π-pulses, the elliptic functions provides a closer spacing between humps with no return to zero. Being applied to the driving transition instead of the 2π-pulse sequence the $2\pi\infty$-train may result in even better transparency for the weak field.

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