

## Electromagnetically Induced Transparency via Adiabatic Following of the Nonabsorbing State

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(Received 23 October 2000; published 25 September 2001)

It is shown that the adiabatic following of the dark, nonabsorbing state improves significantly the electromagnetically induced transparency performance and slows down the group velocity of the probe pulse. This concept can be used for fast selective gating of one pulse out of a pulse train.

DOI: 10.1103/PhysRevLett.87.153601

PACS numbers: 42.50.Md, 42.50.Gy

The first proposal [1] of electromagnetically induced transparency (EIT) dealt with the propagation of a train of short, single mode pulses resonant with two transitions forming a  $\Lambda$  scheme between one excited state 3 and two closely spaced initially populated ground state levels 1 and 2. This pulse train creates the coherence between the ground state levels that prevents absorption at the optical transitions. If the frequency of the pulse train is properly chosen with respect to the ground states splitting, the population is trapped in the dark, nonabsorbing, state [2,3]. In those papers, the two lower states are equally populated before the excitation, the fields have nearly equal amplitudes, and the optical coherences 1-3 and 2-3 decay fast. The problem is that to create the ground state slowly decaying coherence 1-2 uses an appreciable fraction of the input field energy dissipating via spontaneous emission.

Successful experimental testing of the EIT concept [4] employed a different scheme [the energy diagram is shown in the inset of Fig. 1(a)]. State 2 was an excited state, initially unpopulated. Therefore, the coupling field, having a large amplitude, drives the unpopulated states 2 and 3 and the sample is initially transparent for this field. Later the concepts of matched pulses [5], dressed-field pulses [6], and adiabats [7] (see also the review [8]) were introduced to treat the bichromatic pulse propagation in the medium with resonant absorbers being initially only in the ground state 1.

The matched pulses proposal considers the simultaneous excitation by the pump  $B_2(t)$  and probe  $B_1(t)$  pulses matched in time and shape. If the maximum amplitude of the  $B_2(t)$  pulse is much larger than that of the probe pulse  $B_1(t)$ , the absorption of the probe is much more reduced for the matched pulses than for the scheme proposed in [1]. However, for long propagation distances the amplitude of the probe pulse will be spatially and temporally modulated [5].

In this Letter we consider a way to follow adiabatically a dark state which coincides with state 1 before and after interaction with the pulse train. This significantly reduces the energy dissipation via luminescence from the state 3 during the excitation of the 1-2 coherence. One can define the adiabaticity parameter  $A$  via the probability amplitude of the state 3 during the excitation. If the following is adia-

batic ( $A \rightarrow 0$ ), the state 3 is almost empty throughout the excitation and the atom evolution is defined only by the dark state change. This gives a significant decrease of the atomic response to the probe field. Compared to the case of the matched pulses the adiabatic following provides an additional reduction factor of  $\xi/B_2$  for the maximum probability amplitude  $A$  of level 3, where  $\xi = \partial(\ln B_1)/\partial t$  characterizes the change rate of the probe pulse. The adiabatic following of the dark resonance was proposed for stimulated Raman adiabatic passage, transferring the atom from state 1 to state 2 without populating the excited state 3 [9]. In that context the counterintuitive Raman pulse sequence was proposed and experimentally tested. In contrast to the stimulated Raman adiabatic passage, we show that there are conditions where the atom is almost not excited and left in state 1 after the pulse train. To the best of our knowledge the application of the adiabatic following for EIT has not been analyzed yet, albeit proposed in [8] without emphasis on the difference between the maximum amplitudes of the coupling and probe pulses.

We consider the three-level atom interaction with the probe pulse  $\mathbf{E}_1 = \overline{\mathcal{E}}_1(z, t) \exp[i(k_1 z - \Omega_1 t)] + \text{c.c.}$ , resonant with the transition  $1 \leftrightarrow 3$ , and with the coupling pulse  $\mathbf{E}_2 = \overline{\mathcal{E}}_2(z, t) \exp[i(k_2 z - \Omega_2 t)] + \text{c.c.}$ , resonant with the transition  $2 \leftrightarrow 3$ . These pulses are copropagating along the  $z$  axis. The case of exact resonance for both pulses is taken for simplicity. In the rotating wave approximation the Hamiltonian of the three-level atom, excited by two resonant fields, can be made slowly varying by the unitary transformation to the interaction representation. The result is

$$\mathcal{H} = -\hbar B_1(t, z) e^{ik_1 z} \hat{P}_{31} - \hbar B_2(t, z) e^{ik_2 z} \hat{P}_{32} + \text{H.c.}, \quad (1)$$

where  $\hat{P}_{mn} = |m\rangle\langle n|$  is defined in the interaction representation. The Rabi frequencies  $B_1(t, z) = \mathbf{d}_{31} \cdot \overline{\mathcal{E}}_1(z, t)/\hbar$  and  $B_2(t, z) = \mathbf{d}_{32} \cdot \overline{\mathcal{E}}_2(z, t)/\hbar$  are defined in terms of the dipole transition matrix elements  $\mathbf{d}_{31}$  and  $\mathbf{d}_{32}$ , assumed to be real for simplicity. If we choose the new basis

$$|d\rangle = e^{ik_2 z} \cos\alpha |1\rangle - e^{ik_1 z} \sin\alpha |2\rangle, \quad (2)$$

$$|b\rangle = e^{-ik_1 z} \sin\alpha |1\rangle + e^{-ik_2 z} \cos\alpha |2\rangle, \quad (3)$$

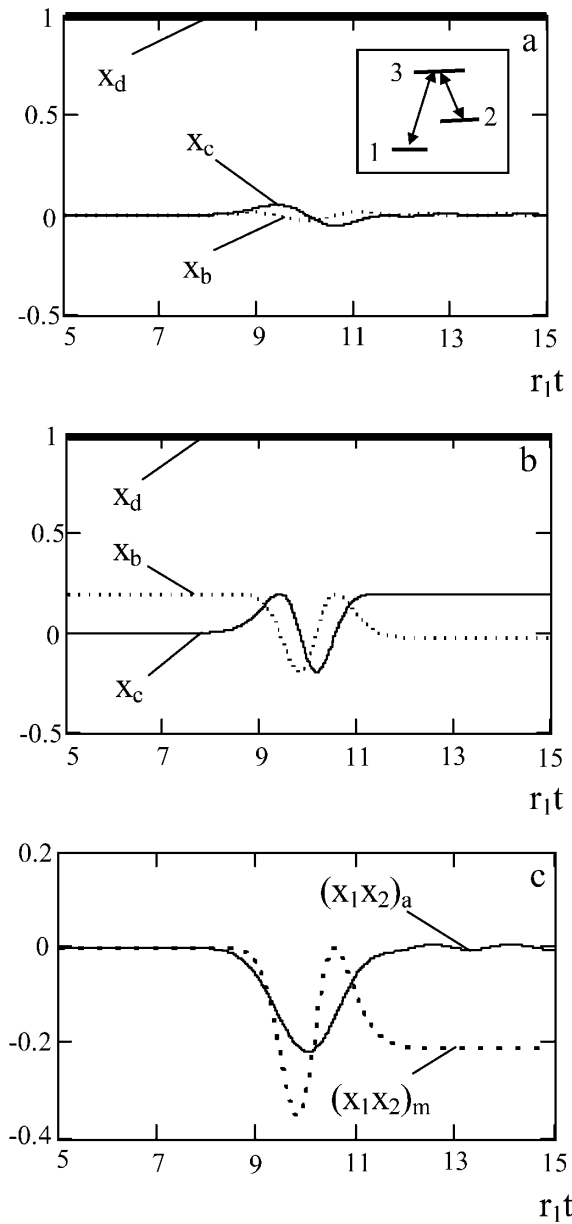


FIG. 1. The evolution of the probability amplitudes of the dark,  $x_d$ , bright,  $x_b$ , and common,  $x_c$ , states for the (a) adiabatic pulse train and (b) for the matched pulses. (c) The evolution of the coherence 1-2 for the matched pulses  $(x_1 x_2)_m$  and for the AFDS pulse sequence  $(x_1 x_2)_a$ .

and  $|c\rangle = |3\rangle$ , then the Hamiltonian (1) is transformed as

$$\mathcal{H}_{dbc} = S\mathcal{H}S^{-1} = -\hbar B(\hat{P}_{bc} + \hat{P}_{cb}), \quad (4)$$

where  $B = \sqrt{B_1^2 + B_2^2}$  and  $S$  is the transformation matrix for the basis change. We drop the variables  $t$  and  $z$  for simplicity. The time-varying mixing parameter  $\alpha$  is defined by  $\tan\alpha = B_1/B_2$ . The indices  $d$ ,  $b$ , and  $c$  denote the dark, bright, and common states. Only the dark state is an eigenstate of the atom-field interaction Hamiltonian (1). The bright state is chosen as a combination of states  $|1\rangle$  and  $|2\rangle$  which is orthogonal to the dark state. According to the Hamiltonian  $\mathcal{H}_{dbc}$  the bichromatic field  $\mathbf{E} = \mathbf{E}_1 + \mathbf{E}_2$

induces only the transitions  $b \leftrightarrow c$  and does not interact with the state  $d$ .

Initially states 2 and 3 are not populated and the sample is transparent for the coupling pulse. The time variation of the field amplitudes is chosen in such a way that for an atom with coordinate  $z$  the ground state 1 is the dark state, uncoupled to the bichromatic excitation at the beginning and at the end of the pulse train, i.e.,  $B_2(\pm\infty, z) \gg B_1(\pm\infty, z)$ . The coupling pulse is much longer than the probe and hence  $\alpha(\pm\infty, z) \approx 0$ . This is the first condition of the adiabatic following of the dark state specifying the time behavior of the pulses. When the probe pulse arrives, the dark state has an admixture of the state 2. The amplitude of this admixture can be made very small if the maximum amplitudes of the probe  $B_{1m} = B_1(t_m, z_0)$  and coupling  $B_{2m} = B_2(t_m, z_0)$  pulses, coinciding at time  $t_m$ , satisfy the condition  $B_{2m} \gg B_{1m}$ , where  $z_0$  is defined as a position of the front edge of the sample. If the relative time behavior of the pulses is chosen such that each atom adiabatically follows the changing dark state, then no population is transferred to the excited state 3 ( $A = 0$ ) and the atoms remain in the dark state, uncoupled to the pulses. Finally, the atoms are left in the ground state after the excitation. Adiabatic following has an essential advantage compared to the matched pulses with identical envelopes since for the latter, the bright state is initially populated and its population is completely transferred to the excited state 3 during the interaction, while in our case the population of the bright state is minimized by the adiabatic following.

The evolution of the state vector of the atom in the  $dbc$  basis,  $|\Phi_{dbc}(t)\rangle$ , is given by

$$|\Phi_{dbc}(t)\rangle = S(t)|\Phi(t)\rangle, \quad (5)$$

where  $|\Phi(t)\rangle$  is the state vector in the interaction representation. Both state vectors are defined for an atom localized in point  $z$ . Taking the time derivative of Eq. (5), one obtains the Schrödinger equation

$$\hbar \frac{d|\Phi_{dbc}\rangle}{dt} = -i\overline{\mathcal{H}}_{dbc}|\Phi_{dbc}\rangle, \quad (6)$$

with the modified Hamiltonian

$$\overline{\mathcal{H}}_{dbc} = \mathcal{H}_{dbc} + i\hbar\dot{S}S^{-1}, \quad (7)$$

$$i\dot{S}S^{-1} = -i\dot{\alpha}\hat{P}_{db}e^{-i(k_1+k_2)z} + \text{H.c.} \quad (8)$$

We have neglected relaxation processes, assuming that the time scale of the interaction with the probe field is short and the Rabi frequency of the coupling field is strong ( $t_p < T_2$  and  $B_{2m}T_2 > 1$ , where  $t_p$  is the probe pulse duration and  $T_2$  is the dephasing time of the induced polarization). The first part of the modified Hamiltonian,  $\mathcal{H}_{dbc}$ , induces transitions between  $|b\rangle$  and  $|c\rangle$  with the rate  $B$ . The second part,  $i\hbar\dot{S}S^{-1}$ , induces transitions between  $|d\rangle$  and  $|b\rangle$  with the rate  $\dot{\alpha} = (\dot{B}_1B_2 - B_1\dot{B}_2)/B^2$ . This rate is nonzero due to the relative variation of the field amplitudes. The state

vector  $|\Phi_{dbc}\rangle$  can be expressed as

$$|\Phi_{dbc}\rangle = x_d e^{-ik_2 z} |d\rangle + x_b e^{ik_1 z} |b\rangle + x_c i e^{ik_1 z} |c\rangle, \quad (9)$$

$$\dot{x}_d = -\dot{\alpha} x_b, \quad \dot{x}_b = \dot{\alpha} x_d - B x_c, \quad \dot{x}_c = B x_b. \quad (10)$$

For matched pulses  $\dot{\alpha}(t) = 0$  and  $x_b(-\infty) = B_{1m}/B_m$ , where  $B_m = \sqrt{B_{1m}^2 + B_{2m}^2}$ . Therefore, the matched pulses transfer the population of the bright state to the excited state  $c$ , i.e.,  $x_c(t) = x_b(-\infty) \sin[\theta(t)/2]$  and  $x_b(t) = x_b(-\infty) \cos[\theta(t)/2]$ , where  $\theta(t) = 2 \int_{-\infty}^t B(\tau) d\tau$  is the pulse area of the bichromatic field  $\mathbf{E}$ . The dark state and its population remain unchanged, i.e.,  $\alpha = \text{const}$  and  $x_d = B_{2m}/B_m$ . For our pulse train the dark state changes in time. If the atom follows adiabatically the time-dependent dark state, the excited state 3 remains empty throughout the excitation. Qualitatively, one can specify this condition as follows. The maximum value of the mixing parameter is  $\alpha_{\text{max}} = \tan^{-1}(B_{1m}/B_{2m})$ . The time interval during which  $\dot{\alpha} \neq 0$  determines the interaction time between the atom and the  $B_1$  pulse. During this time the atom has a small probability to leave the dark state with the rate  $\dot{\alpha}$  [see Eq. (10)]. The maximum value of  $|\dot{\alpha}|$  can be estimated roughly to be  $\alpha_{\text{max}}/t_p$ . Choosing  $\alpha_{\text{max}}$  small, the probability to find the atom in the bright state is small. According to [9], the probability to have the atom in the excited state  $c$  is small if  $B(t) \gg |\dot{\alpha}(t)|$ . Thus the second condition of the adiabatic following of the dark state (AFDS) is  $B(t)t_p \gg \alpha_{\text{max}}$ . This condition, together with the condition  $B_{2m} \gg B_{1m}$ , are the two requirements for AFDS. If both conditions hold, the atom comes back to the ground state after interacting with the pulse train.

This qualitative discussion can be made rigorous with the analytical solution of Eqs. (10). They remarkably coincide with the Bloch equations for the two-level system if one makes the substitution  $u = x_c$ ,  $v = x_b$ ,  $w = x_d$  for the variables and  $\Delta = B$ ,  $\chi = \dot{\alpha}$  for the parameters, where  $u, v, w$  are the Bloch-vector components and  $\Delta$  and  $\chi$  are the detuning and Rabi frequency of the field driving the two-level atom. The similarity between the equations stems from a similarity of the physical processes: large detuning of the driving field ( $\Delta \gg \chi$ ) prevents the excitation of the two-level atom. Similarly, the  $B_2(t)$ -driving field shifts the energy levels of the three-level atom out of the resonance with the probe field  $B_1(t)$  due to the ac Stark effect.

The adiabatic following approximation for the description of the nonlinear response of a two-level atom to a near resonant light pulse was developed by Crisp [10]. We use this approximation scheme to solve Eqs. (10), assuming that the change of the generalized Rabi frequency  $B(t)$  is negligibly small during the change of  $\dot{\alpha}$ . The solution is presented as a series expansion

$$x_d(t) = 1 - \frac{\dot{\alpha}^2}{2B^2} + \frac{3\dot{\alpha}^4 + 8\dot{\alpha}\ddot{\alpha} - 4\ddot{\alpha}^2}{8B^4} + O(B^{-6}), \quad (11)$$

$$x_b(t) = \frac{\ddot{\alpha}}{B^2} - \frac{2\dot{\alpha}\ddot{\alpha} + 3\dot{\alpha}^2\ddot{\alpha}}{2B^4} + O(B^{-6}), \quad (12)$$

$$x_c(t) = \frac{\dot{\alpha}}{B} - \frac{\dot{\alpha}^3 + 2\ddot{\alpha}}{2B^3} + O(B^{-5}). \quad (13)$$

For the  $B_2$  field to have a strong effect on the  $B_1$  excitation, it follows from (11)–(13) that one has to satisfy the condition  $B(t) \gg |\dot{\alpha}(t)|$  and then the adiabaticity parameter  $A \approx \dot{\alpha}/B \rightarrow 0$ . Figure 1(a) shows the result of a numerical simulation of Eqs. (10) for the AFDS of the atom in the point  $z_0$  obtained with the Gaussian-shaped pulses  $B_1(t) = B_{1m} \exp[-r_1^2(t - t_m)^2]$  and  $B_2(t) = B_{2m} \exp[-r_2^2(t - t_m)^2]$  where  $r_1 = 10r_2$ ,  $r_1 t_m = 10$ , and  $B_{1m} = 0.2B_{2m} = r_1 \sqrt{\pi}/2$ . The pulse area of the probe is  $\theta_1(+\infty) = 2 \int_{-\infty}^{+\infty} B_1(t) dt = \pi$ . Such a pulse, acting alone, would leave the atom in the excited state 3 at  $t = +\infty$ . The plot shows that the excitation of the common state is small and the atom is left in the dark state after the pulse train is gone. The difference between the numerical and analytical solutions is indistinguishable on this scale. Figure 1(b) shows the atomic evolution for the matched pulses with  $r_2 = r_1$  and the other parameters are the same as in the Fig. 1(a). The probability amplitude of the bright state  $x_b$  is completely transferred to the common, excited, state 3 [ $x_c(+\infty) = x_b(-\infty)$ ]. The atom is left excited. However, the excitation level is strongly reduced compared to the case without coupling field.

The wave equation for the probe field envelope [6] is

$$\left(\frac{\partial}{\partial z} + \frac{1}{c} \frac{\partial}{\partial t}\right) \vec{\mathcal{E}}_1(z, t) = -\frac{2\pi\Omega_1 N \mathbf{d}_{13}}{c} x_1 x_3, \quad (14)$$

where  $N$  is the concentration of the resonant atoms,  $x_1$  and  $x_3$  are the coefficients of the state vector  $|\Phi(t)\rangle = x_1 |1\rangle + x_2 e^{i(k_1 - k_2)z} |2\rangle + i x_3 e^{ik_1 z} |3\rangle$  expressed in the interaction representation. These coefficients are related to  $x_d, x_b$ , and  $x_c$  by the relations  $x_1 = x_d \cos\alpha + x_b \sin\alpha$ ,  $x_2 = x_b \cos\alpha - x_d \sin\alpha$ , and  $x_3 = x_c$ . For the adiabatic pulses one can obtain from Eqs. (11)–(13) the expression  $x_1 x_3 \approx \dot{B}_1/B_2^2$  which results from the approximation  $x_d \approx 1$ ,  $x_b \approx 0$ ,  $x_c \approx \dot{B}_1/B_2^2$ , and  $\cos\alpha \approx 1$ . Then Eq. (14) is reduced to

$$\left(\frac{\partial}{\partial z} + \frac{1}{c} \frac{\partial}{\partial t}\right) \vec{\mathcal{E}}_1(z, t) = -\frac{\eta}{c} \frac{\partial}{\partial t} \vec{\mathcal{E}}_1(z, t), \quad (15)$$

where  $\eta = Kc/B_2^2 T_2$ ,  $2K$  is the power attenuation coefficient of the probe if the atoms are not disturbed by the coupling field,  $T_2$  is the dephasing time of the optical coherence 1-3. The solution of Eq. (15) with constant coupling field amplitude shows that the probe is not absorbed. However, its group velocity is reduced as  $v_g = c/(1 + \eta)$ . The slowly varying amplitude of the polarization induced by the coupling pulse is  $\mathbf{P}_2(z, t) = iN \mathbf{d}_{23} x_2 x_3$ , where  $x_2 x_3 \approx -\dot{B}_1/B_1 B_2^3$ . This polarization is nonzero only during the probe pulse interaction with the atom and lasts  $t_p$ . One can find that the distortion

of the coupling pulse and the change of its group velocity are negligible compared to the probe. Therefore the probe pulse, having group velocity  $v_g$ , is delayed relative to the coupling pulse. When the probe pulse reaches the edge of the coupling pulse where  $B_1(t_r) = B_2(t_r)$ , strong reshaping of both pulses takes place. This happens at  $z_r = z_0 + ct_c v_g / (c - v_g)$  and  $t_r = t_m + ct_c / (c - v_g)$  if at  $z_0$  and  $t = t_m$  the pulse maxima were coincident, and  $t_c$  is defined as  $B_2(t_m + t_c, z_0) = B_{1m}$ . The distortion of both pulses if they have nearly equal Rabi frequencies was considered in [7]. These distortions are developed in adiabats.

The change of the group velocity is typical for EIT since in the EIT window anomalous dispersion takes place with the steep slope of the real part of the refractive index as a function of frequency (see, for example, Ref. [8]). The description in the spectral domain assumes that the probe pulse spectral width must be smaller than the transparency window given by the Rabi frequency of the coupling pulse, or  $1/t_p < B_{2m}$ . Our time domain consideration of AFDS gives another condition for EIT, i.e.,  $|B_1/B_2|_{\max} \approx R_1 R_2 < 1$ , where  $R_1 = B_{1m}/B_{2m}$  and  $R_2 = |\xi|_{\max}/B_{2m}$ . The parameter  $R_2$  is proportional to the inverse of  $B_{2m} t_p$ , if the probe pulse has a bell shape (for example, Gaussian). Then the condition that the probe pulse spectral width must be smaller than the EIT window is  $R_2 < 1$ . The product  $R_1 R_2$  is the maximum of  $|x_3|$  if the first term in the expansion Eq. (13) gives the main contribution. This occurs if the inequalities  $R_1 < 1$  and  $R_2 < 1$  are simultaneously verified. This can be shown by analyzing the second term in the expansion of  $x_c(t)$  [see Eq. (13)]. The term contains contributions proportional to  $(R_1 R_2)^3$  and  $R_1 R_2^3$  if we estimate roughly the third time derivative of the mixing parameter as  $\ddot{\alpha} \sim B_{1m}/B_{2m} t_p^3$ . Both components are small if  $R_1 < 1$  and  $R_2 < 1$  and, hence, the condition  $R_1 R_2 < 1$  alone is not sufficient for EIT.

It is worthwhile to mention that since matched pulses are nonadiabatic, the coherence between the states 1, 2, and 3 remains after the pulses. On the contrary, adiabatic pulse switching on and off leaves the atom unexcited. Therefore, only the simultaneous switching off of the coupling and probe pulses produces the remnant coherence—the imprinted information about the relative phases of the pulses. Figure 1(c) shows the evolution of the coherence 1-2 for the matched pulses  $(x_1 x_2)_m$ , and for the AFDS pulse sequence  $(x_1 x_2)_a$ . The matched pulses leave an appreciable 1-2 coherence which would reproduce the probe pulse if the coupling pulse alone would be applied again after some delay time.

In this Letter we have shown that, if the coupling pulse is much longer than the probe pulse, there will be a much stronger reduction of the absorption of the probe pulse compared to the case of matched pulses. The crucial point is the adiabatic switching on and off of the coupling pulse. Adiabatic following of the dark state does not leave any information in the atom about the pulse that interacted. Applying coupling pulses, one can induce a fast switching on the dark resonance (as fast as several  $t_p$ 's) for any particular pulse in a pulse train of  $B_1(t)$  pulses. Then all protected pulses of the probe will pass through the sample while the nonprotected pulses will be absorbed.

This work was supported by the Research Fund K. U. Leuven, INTAS, IAP Programme of the Belgian government, FWO-Vlaanderen, the Fonds National de la Recherche Scientifique, and the Belgian Federal Office for Scientific, Technical and Cultural Affairs.

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