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PHYSICS

Electromagnetically Induced Transparency and Controlling the Time Shape of Laser Pulses

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1. The development of new methods for controlling optical properties of a medium and the time shape of laser pulses is an important physical problem of fundamental and applied significance. The solution of this problem is urgently required in the areas of spectroscopy of fast processes, quantum control of atoms and molecules, optoelectronics and optical communication, etc. From this standpoint, the phenomenon of electromagnetically induced transparency (EIT) presents extremely interesting and rather unique opportunities [1–3]. Although the majority of studies on EIT have been carried out for atomic media (see reviews [3–6] and references therein), the relevant ideas have also undergone further development in the context of solid-state systems [7].

EIT is a quantum-interferential phenomenon arising as a result of the interaction between two laser fields and a three-level quantum system. The essence of the phenomenon is that one of the fields, called the controlling field, modifies the optical state of a medium at the frequency of the other (probe) field. In the general case, the intensity of the probe field may be comparable with that of the controlling field. Under conditions of EIT, the substance turns out to be in a coherent state with unusual properties [3–6]. For example, an optically dense medium becomes transparent for the probe field in the single-photon resonance region, whereas the dispersion of the refractive index strongly increases. Under these conditions, optical pulses can run to distances significantly exceeding the resonance length of the single-photon absorption (see, e.g., [8, 9]). The controlling field can control the group velocity of the probe-pulse propagation and can even reduce this velocity down to zero or to a negative value [5]. At group velocities on the order of 1–100 m s^{-1}, a spatial compression of the probe pulse occurs. As a result, it is completely localized in a medium. This phenomenon makes it possible to write down, store, and read out optical pulses [5, 6, 10] and opens the door to new approaches to the development of quantum memory [6].

EIT is used for controlling characteristics of optical radiation, for example, the generation of femtosecond and subfemtosecond pulses (see, e.g., [11]). At present, methods of controlling femtosecond pulses are developed sufficiently well and are applied in spectroscopy, microscopy, and optical monitoring [12]. However, the situation is not so good in the case of picosecond and nanosecond pulses.

In this paper, we discuss a new possibility for controlling the shape and duration of laser pulses on the basis of the EIT phenomenon. Using as an example the time compression of pulses, we now consider the principal concept of the control. Let the probe pulse propagate inside a three-level medium in the presence of a coupling pulse interacting with the adjacent transition (Fig. 1). The envelope of the latter pulse varies with time according to a certain law. Since, in the case of EIT, the propagation velocity of the probe pulse depends on the coupling-pulse intensity at a given instant of time, different parts of the probe pulse move at different velocities. It is important that one can control this velocity by variation of the coupling-pulse shape. For example, the envelope of the coupling pulse at the boundary of the medium can be chosen in a manner such that the propagation velocity of the trailing edge of the probe pulse will be higher than the velocity of its leading edge. As a result, the pulse is compressed with time. Varying the shape of the coupling pulse, it is possible to obtain various shapes for the probe-pulse envelope.

Thus, we here propose an efficient method for controlling the shape and duration of laser pulses, which is based on employing additional controlling radiation that interacts with the adjacent transition under conditions of EIT.

2. We now consider the interaction between a three-level medium and two optical pulses possessing envelopes $E_p(t)$ and $E_c(t)$ (Fig. 1). These pulses propagate in the same direction along the z axis. Probe pulse E_p resonantly interacts with the transition between the ground $|0\rangle$ and excited $|1\rangle$ states, whereas coupling pulse E_c interacts with levels $|2\rangle$ and $|1\rangle$. The dipole $|2\rangle$–$|0\rangle$

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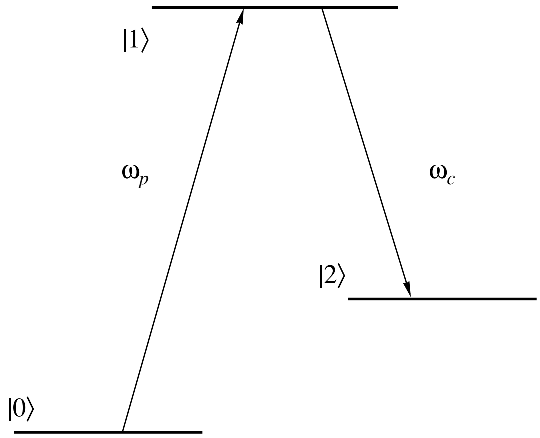

Fig. 1. Energy diagram for a three-level atom resonantly interacting with a probe (ω_p) and coupling (ω_c) pulse.

transition is forbidden. Further, we suggest that pulse durations T_p and T_c are much shorter than all relaxation times and that $T_p < T_c$.

The evolution of the probe and coupling pulses is described by the usual self-consistent set of Maxwell–Schrödinger equations. In the coordinate system with the local time $\tau = t - \frac{z}{c}$, this set is of the form

$$\frac{\partial a_0}{\partial \tau} = iG_p^* a_1, \quad \frac{\partial a_2}{\partial \tau} = iG_c^* a_1, \quad \frac{\partial a_1}{\partial \tau} = i(G_p a_0 + G_c a_2), \tag{1}$$

$$\frac{\partial G_p}{\partial z} = iK_p a_1 a_0^*, \quad \frac{\partial G_c}{\partial z} = iK_c a_1 a_2^*. \tag{2}$$

Here, $a_{0,1,2}$ are the amplitudes of atomic-state probabilities; $2G_p = \frac{d_{10}E_p}{\hbar}$, $2G_c = \frac{d_{12}E_c}{\hbar}$ are Rabi frequencies; $K_p = \pi\omega_p|d_{10}|^2\frac{N}{\hbar c}$, $K_c = \pi\omega_c|d_{12}|^2\frac{N}{\hbar c}$ are the propagation coefficients; d_{ij} are the matrix elements of the electric dipole moment for the $|i\rangle$–$|j\rangle$ transition ($i, j = 0, 1, 2$); $\omega_{p,c}$ and $k_{p,c}$ are the carrier frequencies and wave numbers (in vacuum), respectively; N is the atomic concentration; and c is the speed of light in vacuum.

Equations (1) for probability amplitudes are written for zero single-photon detunings $\omega_{10} - \omega_p = \omega_{12} - \omega_c = 0$. We consider all atoms to be in the ground state $|0\rangle$ at the initial instant of time; i.e., $a_0(-\infty, z) = 1$, $a_1(-\infty, z) = a_2(-\infty, z) = 0$, and the fields $E_{p,c}(t)$ being given at the medium boundary $z = 0$: $E_{p,c}(t, z = 0) = E_{0p,0c}(t)$.

Equations (1) and (2) must be solved by a self-consistent method. In the general case, this procedure can be realized numerically. An essential simplification is attained in the adiabatic approximation [13]. In this case, the solution to set (1) can be represented in the form (see, e.g., [8, 9])

$$a_0 = \cos[\theta(\tau)], \quad a_2 = -\sin[\theta(\tau)], \quad a_1 = -\frac{G_c\dot{G}_p - G_p\dot{G}_c}{G^3} = i\frac{\dot{\theta}}{G}. \tag{3}$$

The displacement angle θ is defined by the expression $\tan\theta = \frac{G_p}{G_c}$, where $G(\tau) = \sqrt{|G_p(\tau)|^2 + |G_c(\tau)|^2}$ is the generalized Rabi frequency. The dot from above denotes differentiation with respect to the local time τ: $\dot{\theta} = \frac{\partial\theta}{\partial\tau}$, etc. In the general case, envelopes G_p and G_c depend on the z coordinate.

The criterion of applicability for the adiabatic approximation can be written as

$$\left|\frac{\dot{G}_c G_p - \dot{G}_p G_c}{G^3}\right| \ll 1. \tag{4}$$

A detailed analysis of the adiabaticity condition with allowance for pulse propagation has been performed in [9].

It follows from formulas (3) and (4) that, in the adiabatic limit, the population of the intermediate state in the interaction process is close to zero, ($|a_1| \ll 1$). This implies that for transitions $|0\rangle$–$|1\rangle$ and $|2\rangle$–$|1\rangle$ the absorption is small. Therefore, the pulses run to a distance that considerably exceeds the length of the resonance linear absorption of probe radiation. This phenomenon is also interpreted in terms of coherent population trapping (CPT): atoms are excited into a coherent superposition of lower states $|0\rangle$ and $|2\rangle$, which is called the CPT state, or dark state [14]. In this state, atoms cease to interact with optical pulses. The EIT phenomenon arises as a result of this process.

Using expression (3), we can represent Eq. (2) in the form

$$\frac{\partial G_p}{\partial z} = -K_p\frac{\dot{\theta}}{G}\cos\theta, \quad \frac{\partial G_c}{\partial z} = K_c\frac{\dot{\theta}}{G}\sin\theta. \tag{5}$$

In the general case, the set of Eqs. (5) can be solved only numerically.

For $K_p = K_c$, it is easy to show from (5) that the generalized Rabi frequency G is independent of the z coordinate:

$$G(\tau, z) = G(\tau, z = 0) = G_0(\tau) = \sqrt{|G_{0p}(\tau)|^2 + |G_{0c}(\tau)|^2}. \tag{6}$$

From this, it follows that arbitrary variations occurring in the probe field are compensated by corresponding

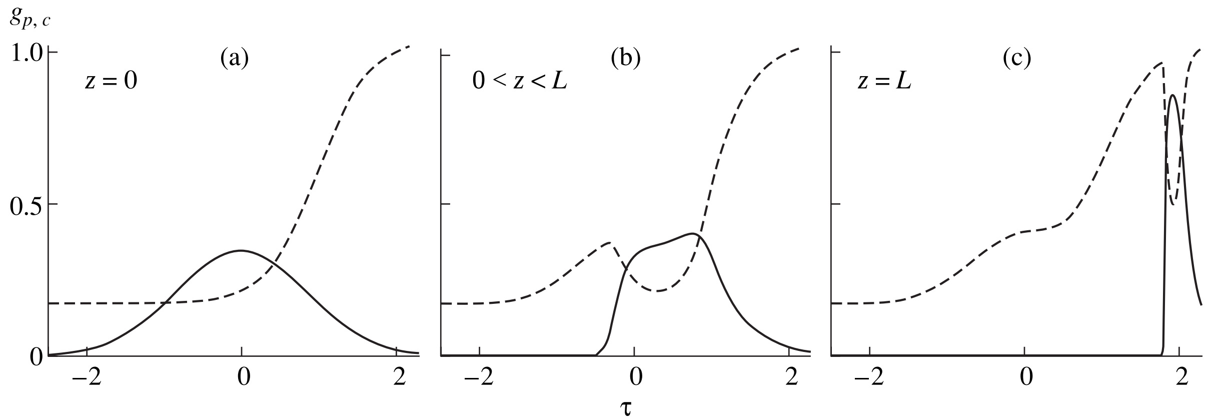

Fig. 2. Time profile of normalized Rabi frequencies for the probe pulse $g_p = \frac{G_p(\tau)}{G_{max}}$ (G_{max} is the maximum value of the effective Rabi frequency) and coupling pulse $g_c = \frac{G_c(\tau)}{G_{max}}$ (dashed curve) in the case of different values of z coordinates inside the medium: (a) at the input of the medium, $z = 0$, $G_p(\tau = 0, z = 0)\ T_p = 20$; (b) at a certain distance z; and (c) at the output of the medium, $z = L$.

changes in the controlling field. In this case, the set of Eqs. (5) is reduced to one equation for $\theta(\tau, z)$:

$$\frac{\partial\theta}{\partial\tau} + \frac{G_0^2(\tau)}{K}\frac{\partial\theta}{\partial z} = 0. \quad (7)$$

The solution to Eq. (7) can be written as

$$\theta(\tau, z) = \theta_0(Z^{-1}(Z(\tau) - z)) = 0, \quad (8)$$

where $Z(\tau) = K^{-1}\int_{\infty}^{\tau} G_0^2(\tau', 0)d\tau'$, Z^{-1} is the function inverse to Z, and $\theta_0 = \theta(\tau, z = 0)$.

The function $\theta(\tau, z)$ allows us to find

$$G_p = G_0(\tau)\sin[\theta(\tau, z)],$$
$$G_c = G_0(\tau)\cos[\theta(\tau, z)]. \quad (9)$$

Analysis of Eqs. (9) [with allowance for (8)] shows that the evolution of the probe pulse depends on the time shape of the coupling pulse at the boundary $z = 0$ of the medium. Figure 2 displays the evolution [described by solution (9)] of the Rabi frequency for the probe and coupling pulses as a function of the z coordinate. The shape of the pulses at the boundary $z = 0$ of the medium is shown in Fig. 2a. As is seen, a time compression of the probe pulse occurs under the conditions indicated. The probe-pulse duration significantly decreases at the output of the medium compared to the input duration. From a physical standpoint, this behavior is associated with the fact that under the indicated conditions in the medium, the propagation velocity of the probe-pulse trailing edge is higher than that of the leading edge. The constraints for the compression are stipulated by the finite width of the transparency window in which the probe pulse can propagate without absorption and also by the adiabaticity conditions.

The pattern of the space–time pulse evolution is similar to the propagation of adiabatons [8]. However, in our case, the pulse duration and the pulse shape vary with pulse propagation, the envelopes of both pulses changing consistently. Therefore, they may be called quasi-adiabatons. Thus, we can speak on the coherent control of the probe-pulse shape by the coupling pulse under the EIT conditions.

It is worth noting that the compression effect does not depend on details of the coupling-pulse time structure. It is sufficient to have a region in which the pulse amplitude (adiabatically) increases. The narrowing effect also arises in the case of a linear variation law for the coupling-pulse envelope. If the coupling-pulse amplitude becomes constant (with respect to time), then, as in [8], we have adiabatons at the output of the medium.

The results obtained are highly consistent with data obtained by the numerical solution of the self-consistent set of Maxwell–Schrödinger equations (1), (2) for the region of parameter in which the adiabaticity condition is fulfilled.

The compression of the probe pulse, which was demonstrated above, is a specific case in which coherent control of the probe pulse shape was achieved by means of EIT. Choosing the time shape of a coupling pulse, it is possible, e.g., to broaden the probe pulse and to form its flat-top or two-bump shape, etc.

3. Thus, it has been theoretically demonstrated that it is possible to control the envelope and duration of a probe pulse on the basis of the EIT phenomenon. This method is applicable to the control of pulses in a wide duration region and in a broad range of spectral wavelength.

The method proposed for the coherent control of laser-pulse time shape can be useful in optical-communication technologies, in processing optical signals, and in nonlinear optics.

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