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Adiabatic propagation of short pulses under conditions of electromagnetically induced transparency

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Abstract. The spatial and temporal dynamics of two short pulses propagating in an optically dense medium of resonant three-level Λ-atoms is investigated numerically and analytically. The maximum coherence for the Raman transition due to coherent population trapping. It is shown that, at the initial stage of propagation, the waveforms of such pulses only slightly change along the length of the medium, which may considerably exceed the length of linear absorption for a single weak pulse. As the length of the absorbing medium increases, the energy of the probe (first) pulse is completely transferred into the second (control) pulse.

1. Introduction

The interaction of two optical light beams with a three-level Λ-system when conditions of one- and two-photon resonances are simultaneously satisfied is one of the 'hot' topics of modern laser and optical physics. The interest of researchers in this problem is mainly due to the fact that the character of light – matter interaction under resonant conditions is determined to a considerable extent by atomic coherence and quantum interference, which may radically change the optical characteristics of a medium, allowing one to control these characteristics. Such phenomena include the nonlinear interference effect [1, 2], electromagnetically induced transparency (EIT) [3, 4], coherent population trapping (CPT) [5, 6], amplification without inversion and generation of coherent radiation [3, 7]. These effects are of fundamental importance for understanding the nature of resonant light – matter interactions and for various applications of quantum and nonlinear optics and laser physics in resonant laser photochemistry (for example, in isotope separation).

These effects have already been employed for controlling the absorption coefficient [1 – 3] and the refractive index [8], for ultrasensitive phase measurements and optical interferometry [9], in the measurement of weak magnetic fields [10], in laser frequency stabilisation [5], in isotope separation [11], and for the enhancement of the efficiency of resonant nonlinear frequency mixing [3, 12 – 14, 30, 31].

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Quantum interference effects also give rise to several unusual phenomena accompanying the propagation of light pulses in resonant three-level media. In particular, under definite conditions, EIT and CPT effects are observed in the temporal and spatial evolution of interacting light pulses. A pair of light pulses under these conditions may propagate in space over distances substantially exceeding the length of linear absorption for a single weak (probe) pulse without considerable changes in pulse waveforms. Such pulses are referred to as matched pulses [3], pulses dressed with a field [15], adiabatons [16], simultons, and solitons [17, 18]. When such pulses partially overlap in time and the 'Raman' pulse is switched on and off earlier than the probe pulse (a counterintuitive sequence [19]), the relevant population can be virtually completely transferred to the two-photon-excited state [19, 20]. Radiation may propagate over large distances in an absorbing medium under these conditions [21].

Coherent population trapping leads to efficient population of certain states of coherently phased atoms and induces a considerable atomic coherence (large off-diagonal elements of the density matrix) for a Raman transition. The modulus of the atomic coherence in such a situation may reach the maximum possible value of 1/2. Absorption of radiation involved in resonant interaction with allowed transitions decreases under these conditions. This effect can be observed for both cw radiation and light pulses, and may considerably enhance the efficiency of one- and two-photon-resonant laser frequency mixing and optical parametric oscillation [12, 13, 22, 30]. For pulsed lasers, this effect was experimentally observed by Jane et al. [13], who achieved an efficiency of four-wave mixing of about 40%. In this context it is of considerable importance and interest to investigate the coherent dynamics of an atomic system interacting with resonant pulsed laser fields and the specific features of the propagation of laser pulses under conditions when the medium has a noticeable influence on laser radiation.

Some aspects of the propagation of light pulses in resonant three-level media were investigated, for example, by the authors of Refs [15, 16, 23 – 27]. This problem is usually analysed for the cases when either both pulses have identical waveforms (identical pulses) and equal durations, exceeding the relaxation time of the intermediate resonant state, or the duration of one of the pulses (the control pulse) is much greater than the duration of the second (the probe pulse). In the majority of cases, the problem under study was simplified because of the adiabatic exclusion of the intermediate level.

In this paper we investigate the spatial and temporal evolution of two short laser pulses overlapping in time and resonantly interacting with an optically dense medium con-

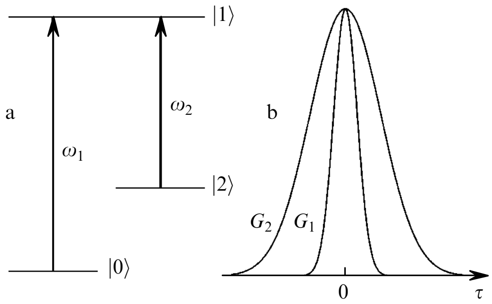

Figure 1. Diagram of atomic energy levels (a) and the envelopes of the Rabi frequencies of light pulses at the input of the medium (b). The arrows show dipole-allowed transitions.

sisting of three-level Λ-atoms. The pulses are assumed to have identical waveforms but different durations (see Fig. 1). We also assume that the pulse envelopes satisfy the adiabaticity criterion (e.g., see Ref. [19]):

$$|G_{1,2}|T_{1,2} \gg 1 \ , \tag{1}$$

where $G_{1,2}$ are the Rabi frequencies corresponding to the interacting pulses, and $T_{1,2}$ are the pulse durations.

As can be seen from expression (1), the adiabaticity condition can be satisfied for short (but sufficiently powerful) pulses that meet the requirement $\Gamma_{ij}T_1 \ll 1$ (Γ_{ij} are the relaxation rates of the atomic subsystem). For the second pulse, this condition is satisfied automatically, since the duration of this pulse satisfies the inequality $T_2 > T_1$. Physically, the adiabaticity criterion implies that the envelopes of interacting light pulses should vary slowly within a time interval on the order of the reciprocal of the Rabi frequency. When these conditions are satisfied, an effect similar to CPT occurs in the stationary regime. This effect considerably lowers the absorption of propagating resonant pulses. Under these conditions, a large coherence is induced for the Raman transition. As demonstrated below, the length inside the medium where this coherence remains large substantially exceeds the linear absorption length of a single probe pulse.

The theoretical model employed in our study involves a set of equations for the amplitudes of the probabilities to find the system under consideration in the 'relevant' levels and the wave equations for slowly varying pulse envelopes. These equations govern the self-consistent spatial and temporal dynamics of the atomic subsystem and radiation.

2. The basic equations and numerical results

Let us consider the propagation of two pulses in a medium consisting of three-level Λ-atoms (Fig. 1). We assume that the light pulses are characterised by the same linear polarisation and propagate along the z axis. Atomic energy levels 0 and 1, 1 and 2 have opposite parities, the 0 – 2 transition is dipole-forbidden, and level 0 is the ground state. The intermediate state $|1\rangle$ is coupled by one-photon resonances with each of the fields, which interact only with the relevant transitions. In what follows, the pulse with frequency ω_1 is considered as a probe pulse, and the pulse with frequency ω_2 is considered as the Raman, or control, pulse. The duration of the control pulse is assumed to be larger than the duration of the probe pulse. The durations of both pulses are assumed to be much less than all the relaxation times of the atomic subsystem.

The interaction of radiation with the atomic system will be described in terms of the Schrödinger equations for probability amplitudes in the resonant approximation. The spatial and temporal evolution of light pulses will be described with the use of a set of wave equations for slowly varying amplitudes [28]. The self-consistent set of equations for the probability amplitudes and the Rabi frequencies of interacting pulses in the moving frame of reference with a local time $\tau = t - z/c$ are written in a standard form:

$$\frac{\partial b_0}{\partial \tau} = \mathrm{i}G_1^* b_1 \exp(-\mathrm{i}k_1 z) \,, \qquad \frac{\partial b_2}{\partial \tau} = \mathrm{i}G_2^* b_1 \exp(-\mathrm{i}k_2 z) \ ,$$
$$\frac{\partial b_1}{\partial \tau} = \mathrm{i}G_1 b_0 \exp(\mathrm{i}k_1 z) + \mathrm{i}G_2 b_2 \exp(\mathrm{i}k_2 z) \ , \tag{2}$$

$$\frac{\partial G_1}{\partial z} = \mathrm{i}K_1 b_1 b_0^* \exp(\mathrm{i}k_1 z) \,, \qquad \frac{\partial G_2}{\partial z} = \mathrm{i}K_2 b_1 b_2^* \exp(\mathrm{i}k_2 z) \, . \tag{3}$$

In writing this set of equations, we have assumed that each carrier frequency of a pulse is resonant to the relevant transition, and the wave equations for slowly varying amplitudes are written for the frequencies $G_1 = d_{10}E_1(t)/2\hbar$ and $G_2 = d_{21}E_2(t)/2\hbar$. Here, $b_{0,1,2}$ are the probability amplitudes for the states 0, 1, and 2; $K_1 = \pi\omega_1|d_{10}|^2N/2\hbar = \alpha_1\Gamma_{10}/4$ and $K_2 = \pi\omega_2|d_{21}|^2N/2\hbar = \alpha_2\Gamma_{12}/4$ are the propagation coefficients; $\alpha_{1,2}$ are the linear absorption coefficients for the probe and control light beams when all the atoms reside in states 0 or 2, respectively; Γ_{ij} are the half-widths of transitions; N is the concentration of atoms; d_{ij} are the dipole matrix elements of transitions; and $k_{1,2}$ is the modulus of the wave vector of the interacting waves in a vacuum. We assume that, at the moment of time when the fields are switched on ($\tau = \infty$), all the atoms reside in the ground state 0, and both pulses have Gaussian envelopes at the input of the medium $z = 0$, $E_1(t) = E_1^0 \exp(-t^2/2T_1^2)$ and $E_2(t) = E_2^0 \exp[-t^2(2T_2^2)^{-1}]$, where the times $T_2 > T_1$ are much less than all the relaxation times of the atomic subsystem. The amplitudes of the pulses $E_{1,2}^0$ are assumed to be real. Parameters of the pulses are chosen in such a way as to satisfy the adiabaticity condition (1) at the input of the medium $z = 0$.

It is now convenient to introduce new variables: $a_0 = b_0 \exp(\mathrm{i}k_1 z)$, $a_1 = \mathrm{i}b_1$, and $a_2 = b_2 \exp(\mathrm{i}k_2 z)$. Using these variables, we can rewrite the set of equations (2) and (3) in the following form:

$$\frac{\partial a_0}{\partial \tau} = G_1^* a_1 \,, \quad \frac{\partial a_1}{\partial \tau} = -G_1 a_0 - G_2 a_2, \ \frac{\partial a_2}{\partial \tau} = G_2^* a_1 \ , \tag{4}$$

$$\frac{\partial G_1}{\partial z} = -K_1 a_1 a_0^* \,, \qquad \frac{\partial G_2}{\partial z} = -K_2 a_1 a_2^* \ . \tag{5}$$

The set of equations (4) and (5) was solved numerically. For time-domain equations, we employed the Adams method, whereas the spatial equations were solved by using the Euler method. Simulations were performed for the following parameters of the system: $T_2/T_1 = 3$, $G_1^0 T_1 = 10$, $G_2^0 T_1 = 10$ ($G_{1,2}^0$ are the maximum values of the Rabi frequencies $G_{1,2}$), $\Gamma_{10}T_1 = 0.1$, $\Gamma_{12}T_1 = 0.1$, and $\alpha_2/\alpha_1 = 1$ ($K_1 = K_2$).

We investigated the populations of atomic levels and the Rabi frequencies as functions of time at different points of the medium. Fig. 2 displays the time dependences of populations $\rho_{0,2} = |a_{0,2}|^2$ at different points of the medium, and Fig. 3 shows the modulus of the off-diagonal element of the density matrix $|\rho_{20}| = |a_2 a_0^*|$ as a function of time and the spatial

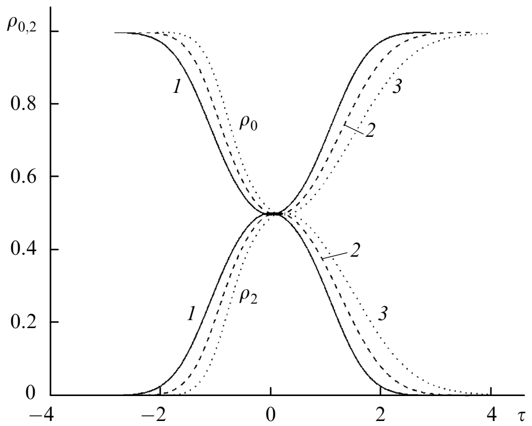

Figure 2. Temporal evolution of the populations $\rho_{0,2} = |a_{0,2}|^2$ at different points of the medium for $\xi = 0$ (1), 5×10^2 (2), and 10^3 (3). The time τ here and in all the other figures is measured in units of the pulse duration T_1, and the propagation length ξ of pulses in the medium (from the input plane $z = 0$) is measured in units of the length of linear absorption of radiation with the frequency ω_1 determined in accordance with the Bouguer law.

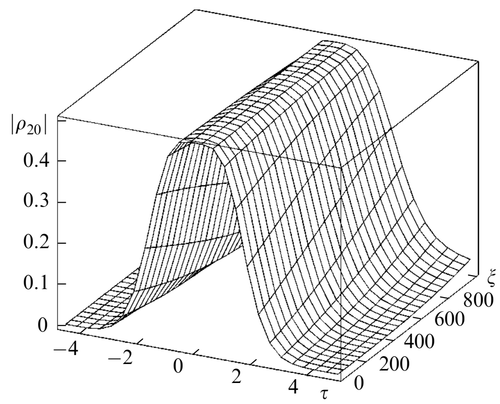

Figure 3. Atomic coherence $|\rho_{20}| = |a_2 a_0^*|$ as a function of time and the propagation length ξ of light pulses in the medium.

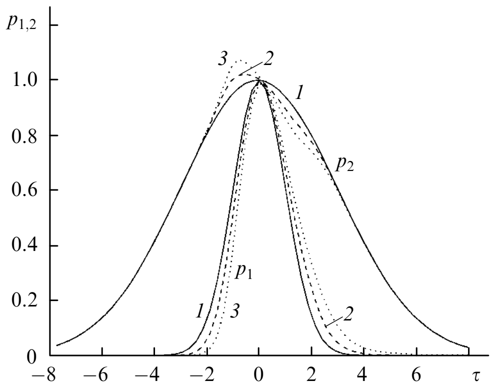

Figure 4. Time dependences of the normalised Rabi frequencies $p_{1,2} = G_{1,2}/G_{1,2}^0$ of the interacting light pulses at different points inside the medium along the z axis for $\xi = 0$ (1), 5×10^2 (2), and 10^3 (3).

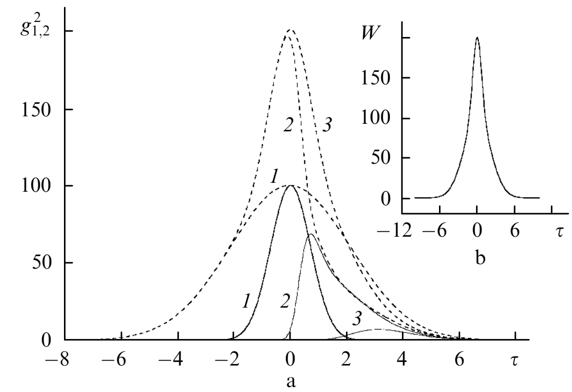

Figure 5. Time dependences of the envelopes of the Rabi frequencies squared, $g_{1,2}^2 = [G_{1,2}(\tau)T_1]^2$ (a), and the sum of these quantities, $W = g_1^2 + g_2^2$ (b), for the interacting pulses at different points inside the medium and $\xi = 0$ (1), 5×10^3 (2), and 10^4 (3).

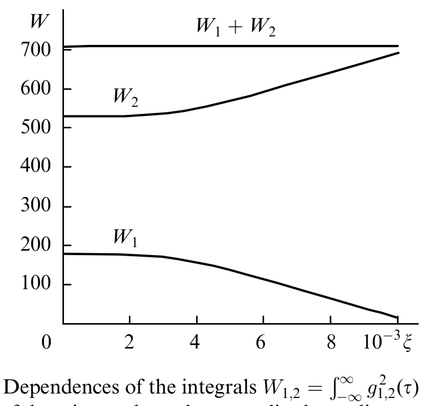

Figure 6. Dependences of the integrals $W_{1,2} = \int_{-\infty}^{\infty} g_{1,2}^2(\tau)\,\mathrm{d}\tau$ and the sum $W_1 + W_2$ of these integrals on the normalised coordinate.

coordinate for pulse parameters satisfying the adiabaticity criterion (1). As can be seen from these dependences, the atomic coherence $|\rho_{20}|$ may reach its maximum equal to 1/2. Under these conditions, the populations $\rho_{0,2}$ in atomic levels become virtually equal to each other when the field amplitudes reach their maxima, and the population ρ_1 of the intermediate state remains close to zero as the pulses propagate through the medium. The latter effect implies the absence of absorption for 0 – 1 and 1 – 2 transitions, which can be considered as an analogue of CPT in the case of cw radiation. As can be seen from Figs 2 and 3, this effect manifests itself within the length of a medium considerably exceeding the linear absorption length for a single probe pulse.

Fig. 4 shows the time dependences of the Rabi frequencies for propagating pulses at different points inside the medium. The results presented in these figures demonstrate that, under the CPT conditions, light pulses may propagate in an optically dense medium over distances substantially exceeding the linear absorption length for the probe pulse without noticeable changes in their waveforms.

Fig. 5 presents the time dependences of the envelopes of the Rabi frequencies squared and the sum of these quantities for the interacting pulses at different points inside the medium. As can be seen from these dependences, the sum of the envelopes of the Rabi frequencies squared is independent of the coordinate and is determined by the time dependence of its components at the input of the medium. Fig. 6 displays the time integrals of the quantities shown in Fig. 5 as functions of the coordinate. As can be seen from this figure, the integral of the sum of the Rabi frequencies squared remains constant, which is a manifestation of the self-consistent behaviour of the envelopes. Such a behaviour of the dependences studied is observed only when the propagation constants are equal to each other ($K_1 = K_2$).

3. Analysis of the results

The results presented above agree well with predictions of the analytical consideration in approximation (1). In the adiabatic approximation, derivatives in expressions (4) can be set equal to zero [29]. Hence we derive the following expressions for the probability amplitudes (see also [16]):

$$a_0 \approx \frac{G_2}{G}\,, \quad a_1 \approx \frac{1}{G_1}\frac{\partial(G_2/G)}{\partial\tau} \approx -\frac{1}{G_2}\frac{\partial(G_1/G)}{\partial\tau}\,, \quad a_2 \approx -\frac{G_1}{G}\,, \tag{6}$$

where $G = (G_1^2 + G_2^2)^{1/2}$.

The expression for a_1 can be reduced to the following form:

$$a_1 = \frac{G_2\dot{G}_1 - G_1\dot{G}_2}{G^3}. \tag{7}$$

In the adiabatic approximation, the inequality $|a_1| \ll 1$ is more precise adiabaticity condition than that given by expression (1) (e.g., see Ref. [19]). Under these conditions, the population of the intermediate state 1 is close to zero within the entire period of time corresponding to the interaction with laser pulses, as $|a_1(\tau)|^2 \ll 1$. Physically, this effect implies that the resonant absorption of light pulses is weak and that the population of levels in the process of interaction with the fields is distributed mainly between the initial state 0 and the final state 2. Thus we arrive at the approximate equality

$$|a_0|^2 + |a_2|^2 \approx 1\ . \tag{8}$$

The solutions for the probability amplitudes $a_{0,2}$ can be written as

$$a_0 = \cos\theta\,, \quad a_2 = -\sin\theta. \tag{9}$$

Here θ is understood as some angle whose meaning will become clear later.

The results obtained above can be interpreted in terms of the vector model where the vector variables $\boldsymbol{a} = \{a_0, a_1, a_2\}$ and $\boldsymbol{G} = \{G_2, 0, -G_1\}$ are introduced. Using these variables, we can rewrite the set of expressions (4) as:

$$\dot{\boldsymbol{a}} = \boldsymbol{G} \times \boldsymbol{a}. \tag{10}$$

The vector $\boldsymbol{a} = \{G_2/G, 0, -G_1/G\}$ defined as the solution to equation (10) lies in the plane $\boldsymbol{ik}$. Vector $\boldsymbol{G}$ lies in the same plane and makes an angle θ with the $\boldsymbol{i}$ axis, where $\cos\theta = G_2/G$ (Fig. 7). Components of vector $\boldsymbol{a}$ virtually coincide with the adiabatic solution (6), since $|a_1| \ll 1$ and we can neglect this quantity. Consequently, vector $\boldsymbol{a}$ corresponding to adiabatic solution (6) is virtually parallel to vector $\boldsymbol{G}$ and precesses around this vector with the frequency $G = (G_1^2 + G_2^2)^{1/2}$, adiabatically following the vector. Thus we deal with a complete analogy with adiabatic following observed in the case of interaction of a light pulse with a two-level atom. Since light pulses propagate in the medium, the angle θ is a function of the time and the coordinate. Note that the results obtained above are independent of

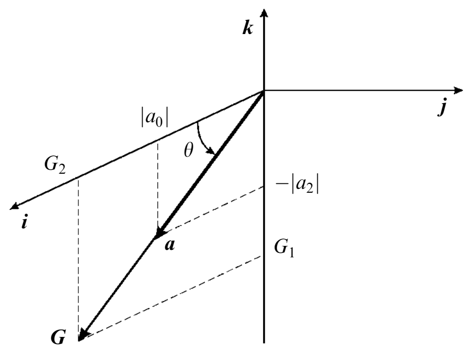

Figure 7. The vector model of the adiabatic interaction of two short pulses with a three-level Λ-system.

the waveform of light pulses if the adiabaticity condition $|a_1| \ll 1$ is satisfied.

Equality (8) also reflects the fact that atoms are trapped in the CPT state: $a_- = (G_2/G)a_0 - (G_1/G)a_2 = \cos(\theta)a_0 - \sin(\theta)a_2 = \text{const}$. This effect is responsible for the decrease in the resonant absorption of the propagating pulses.

Substituting solution (6) into field equations (5), we derive a set of coupled nonlinear equations:

$$\frac{\partial G_1}{\partial z} = -\frac{K_1}{G}\frac{\partial(G_1/G)}{\partial\tau}\,, \quad \frac{\partial G_2}{\partial z} = -\frac{K_2}{G}\frac{\partial(G_2/G)}{\partial\tau}\ . \tag{11}$$

Equations (11) can be conveniently reduced to

$$\frac{\partial G_1}{\partial z} = -\frac{K_1 G_1}{G}a_1\,, \quad \frac{\partial G_2}{\partial z} = -\frac{K_2 G_2}{G}a_1\ . \tag{12}$$

As can be easily shown with the use of expressions (12), the quantity $G_1^2(\tau, z) + G_2^2(\tau, z)$ with $K_1 = K_2$ is independent of coordinate z and is equal to $G^2(\tau, 0) = G_1^2(\tau, 0) + G_2^2(\tau, 0)$, and

$$\int_{-\infty}^{\infty} [G_1^2(\tau,\ 0) + G_2^2(\tau,\ 0)]\,\mathrm{d}\tau = \text{const}$$

(cf. the results of numerical simulations presented in Figs 5 and 6).

In general, the solution to the set of expressions (11) cannot be written in quadratures. However, with $K_1 = K_2 = K$, this solution can be found with the use of the method of characteristics (see also Ref. [16]), which yields

$$G_1 = G(0,\ \tau)\frac{G_1[0,\ Z^{-1}(x)]}{G[0,\ Z^{-1}(x)]}\,, \quad G_2 = G(0,\ \tau)\frac{G_2[0,\ Z^{-1}(x)]}{G[0,\ Z^{-1}(x)]}\,, \tag{13}$$

where $Z(\tau) = K^{-1}\int_{-\infty}^{\tau} G^2(0, \tau')\,\mathrm{d}\tau'$; $x = Z(\tau) - z$; and $Z^{-1}(x)$ is the inverse of $Z(x)$.

The results of numerical simulations agree very well with the results obtained with the use of formulas (6) and (13). Figs 8 and 9 present the populations $\rho_{0,2} = |a_{0,2}|^2$ and the envelopes of propagating pulses as functions of time and the penetration depth of radiation for the conditions of Figs 2–4. As can be seen from Figs. 8 and 9, light pulses can propagate in the medium over distances considerably exceeding the length of linear absorption for a single probe pulse. At the initial stage, the waveform of light pulses displays only small changes. However, at later stages, the

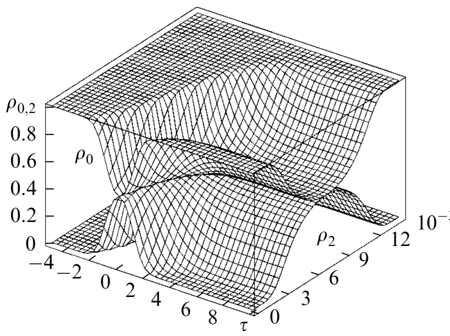

Figure 8. Populations of atomic levels $\rho_{0,2} = |a_{0,2}|^2$ as functions of time and the penetration depth of radiation in the medium.

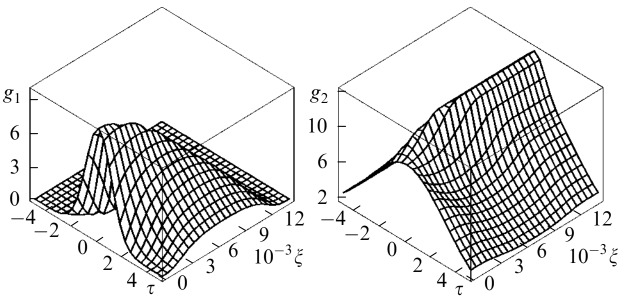

Figure 9. Envelopes of the Rabi frequencies $g_{1,2} = G_{1,2}T_1$ as functions of time and the penetration depth of radiation in the medium.

energy of the probe pulse is completely transferred into the second pulse. As the light pulses propagate through the medium, the populations of levels 0 and 2 become nonmonotonic functions of the coordinate (Fig. 8). The angle θ under these conditions reaches the value of $\pi/4$ and then decreases down to zero as the envelope of the probe field tends to zero owing to the transfer of the energy of this pulse into the second field.

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