Dressed-Field Pulses in an Absorbing Medium

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We have carried out a series of numerical experiments on the propagation of optical pulse pairs in absorbing media under one-photon and two-photon resonant lambda-system conditions. We report the first observation of the spatial evolution of "dressed-field" pulses, the exact analog of the temporal evolution of dressed-atom states.

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A variety of coherent nonlinear quantum optical phenomena are known in which a secondary optical pulse cooperates with or even controls a primary pulse. Recently this effect has been discussed theoretically in connection with electromagnetically induced transparency (EIT) [1,2], and other instances include simulton propagation [3], Raman solitons [4], laser-induced continuum structures [5], lasing without inversion [6], and efficient upper state excitation by counterintuitive pulse sequencing [7]. In most cases the pulses are to be injected into a medium of atoms that are well approximated as three-level lambda systems, such as sketched in Fig. 1. If the two cooperating fields are permitted to evolve dynamically, their interplay via the two-photon transition leads to significant mutual coherence [8].

The existence of a trapped state [9], or the presence of a "dark resonance," which is an exact (i.e., nonperturbative) single-atom two-photon effect, appears to be essential for EIT and several of the other phenomena mentioned. While the temporal evolution of trapped-state physics is well understood, no theoretical work has yet appeared to show how trapped states can evolve *spatially*, i.e., in the course of propagation through the necessary two-photon medium. Harris has noted [1] that even the details associated with the injection of the pulses into a medium raise interesting questions.

Our work was designed to engage these open questions and our results are obtained from "experiments" in which the coupled Maxwell-Schrödinger equations are solved exactly numerically. By means of these numerical solutions various input pulse pairs can be observed as they propagate in the medium over distances from a fraction of a Beers length to many Beers lengths. Our calculations are the first to show the following: (1) the mutual interplay and alteration of two cooperating injected pulses as they propagate, (2) changes in the secondary pulse as it "protects" the trailing portion of the primary pulse during propagation, and most interesting, (3) the rapid growth of a unique and extremely stable nondecaying composite pulse that is actually a "dressed-field" superposition of the pair of pulses. It is distinct from a two-photon pulse [10].

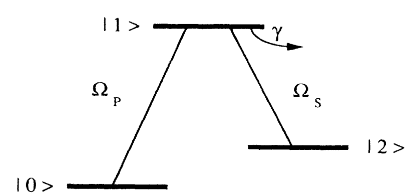

FIG. 1. A sketch of the atomic energy levels, connected by on-resonant pump and Stokes fields.

The slowly varying amplitude equations that describe the spatial and temporal evolution of two pulses in a three-level medium such as that shown in Fig. 1 are

$$\frac{\partial c_0}{\partial t} = \frac{1}{2}\Omega_P^*(ic_1)\,, \tag{1a}$$

$$\frac{\partial (ic_1)}{\partial t} = (-i\Delta-\gamma)(ic_1) - \frac{1}{2}\Omega_P c_0 - \frac{1}{2}\Omega_S c_2\,, \tag{1b}$$

$$\frac{\partial c_2}{\partial t} = \frac{1}{2}\Omega_S^*(ic_1)\,, \tag{1c}$$

$$\left(\frac{\partial}{\partial z}+\frac{\partial}{\partial ct}\right)\Omega_P = 2i\mu_P c_0^* c_1\,, \tag{2a}$$

$$\left(\frac{\partial}{\partial z}+\frac{\partial}{\partial ct}\right)\Omega_S = 2i\mu_S c_2^* c_1\,. \tag{2b}$$

We refer to the fields as pump (P) and Stokes (S) pulses and use the notation of Ackerhalt and Milonni [4], which has the advantage that the labels 0, 1, 2 indicate the number of photons required to reach a given level from the initial level, which is labeled 0. The pulses are represented by the corresponding Rabi frequencies in the usual way, $\Omega_P = 2d_P\mathcal{E}_P/\hbar$, etc., and the propagation coefficients are given by $\mu_P = 2\pi\mathcal{N}d_P^2\omega_P/\hbar c$, etc., where $\mathcal{N}$ is the density of three-level atoms in the medium [11]. All five of the dynamical variables are obviously functions of both z and t. We write the equations in terms of ic_1 instead of c_1 because in the results we will show c_1 is imaginary and all other variables are real if they are initially real, as we will assume. The existence of the usual "trapped" superposition state can be checked easily from Eqs. (1a) and (1c); in constant fields the combination of amplitudes

$$c_{\rm trap}(t) \equiv -c_2(t)\Omega_P^* + c_0(t)\Omega_S^* \tag{3}$$

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is independent of time.

Now let us consider the most interesting situation for spatial propagation, which occurs when the intermediate level $|1\rangle$ is one-photon resonant with both laser pulses. This is the case of most rapid decay of coherence, because the absorption coefficients are largest. For simplicity we will assume [as in Eqs. (1)] that the decay of level $|1\rangle$ occurs to "external" levels outside the lambda system, permitting the use of amplitude equations for the medium instead of Bloch variables, and that the two absorption coefficients are equal [12].

Our numerical experiments begin with the injection of pump and Stokes pulses into the medium, which is in its ground level $|0\rangle$ throughout. We then solve the coupled set of nonlinear equations (1) and (2) in the moving frame defined by the variables $\tau = t - z/c$ and $\zeta = z$. Many combinations of pump and Stokes pulses have been used in our experiments, but we will display here only relatively simple ones that turn on and off smoothly and have a constant amplitude in between, in order to concentrate on the most interesting results.

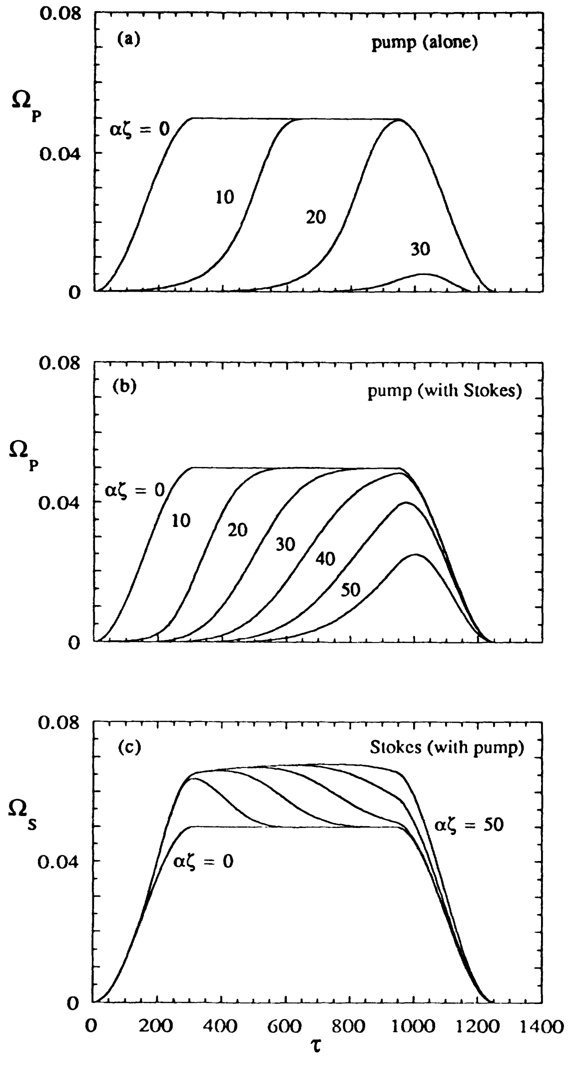

FIG. 2. Snapshots of long propagating pulses at different depths of penetration into the medium (measured in pump absorption lengths), given $\gamma = 0.040$ and the peak $\Omega = 0.050$. (a) Resonant pump pulse alone, showing rapid depletion; (b) resonant pump in the presence of resonant companion Stokes pulse, showing reduced depletion; and (c) companion resonant Stokes pulse, showing amplification.

Figure 2 shows snapshots of the pump pulse at different depths of penetration, measured in standard pump-field absorption lengths. Part (a) shows the behavior of the pump pulse alone, (b) shows the pump pulse if both pulses are injected, and (c) shows the accompanying Stokes pulse. This is done for a relatively lossy medium where the coefficient γ is approximately equal to the peak value of the Rabi frequencies: $\gamma = 0.040$ and $\Omega_P = \Omega_S = 0.050$ [13]. What happens during propagation is clear—the leading (early-time) edge of the pump pulse is continuously depleted as it penetrates the medium. If there is also a Stokes pulse present, it is correspondingly amplified, and *the presence of the Stokes pulse clearly retards the normal absorption of the pump pulse.* That is, note the contrast between the "protected" propagation of the pump pulse in (b) compared with the "unprotected" propagation shown in (a). Incidentally, this is the first observation of details of propagation implicit in EIT as proposed by Harris [2].

Our second example shows pulse behavior in a medium with much more coherence, i.e., 10 times smaller level decay, $\gamma = 0.004$, with the same input pulse parameters: $\Omega_P = \Omega_S = 0.050$. The results graphed in Fig. 3 show greater variation of both pump and Stokes pulse shapes during propagation. Despite this variation, one can still see protection of the pump pulse by the presence of the Stokes pulse.

What is much more interesting, however, is the discovery that certain superposition pulses, that we identify with dressed-field states in close analogy to the dressed-atom states of a two-level system [14], propagate essentially unchanged. We now introduce these dressed fields by a particular time-dependent combination in the two-dimensional space of field states:

$$\begin{pmatrix} \Omega_+ \\ \Omega_- \end{pmatrix} = \begin{pmatrix} c_0 & c_2 \\ -c_2 & c_0 \end{pmatrix} \begin{pmatrix} \Omega_P \\ \Omega_S \end{pmatrix} . \quad (4)$$

A similar combination, however, with time-independent coefficients, was identified by Harris [2]. In the limit of adiabatic excitation of the lambda system one has $c_1 \approx 0$. In this case c_0 and c_2 satisfy $|c_0|^2 + |c_2|^2 \approx 1$, and this makes the matrix in (4) a rotation matrix, and shows that Ω_+ is in this sense the unique partner of Ω_- [15].

Now note that the familiar trapped-state amplitude $c_{\rm trap}$ defined in (3), which is conventionally regarded as a *field-dressed combination of level amplitudes* c_0 and c_2, is in fact *exactly the same variable* as Ω_- in (4), which we see here in the context of field propagation should be

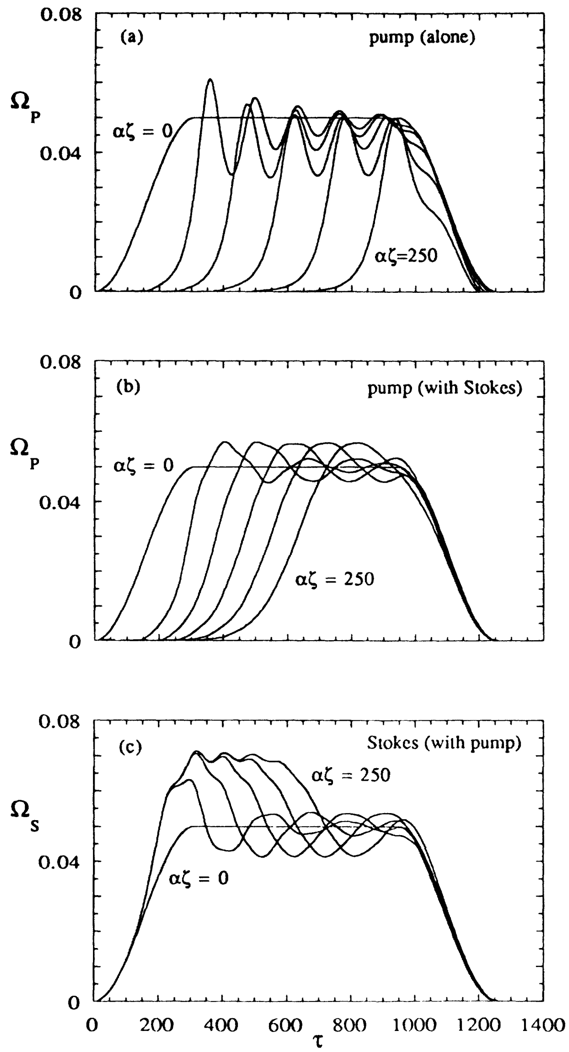

FIG. 3. Snapshots of propagating pulses identical to those of Fig. 2, in the case of greater atomic coherence. Here $\gamma=0.004$ (factor of 10 smaller than in Fig. 2). (a) Resonant pump pulse alone, showing strong rapid coherence oscillations and depletion; (b) pump in the presence of companion resonant Stokes pulse, showing reduced temporal variation and depletion; and (c) companion resonant Stokes pulse, showing coherence oscillations and amplification.

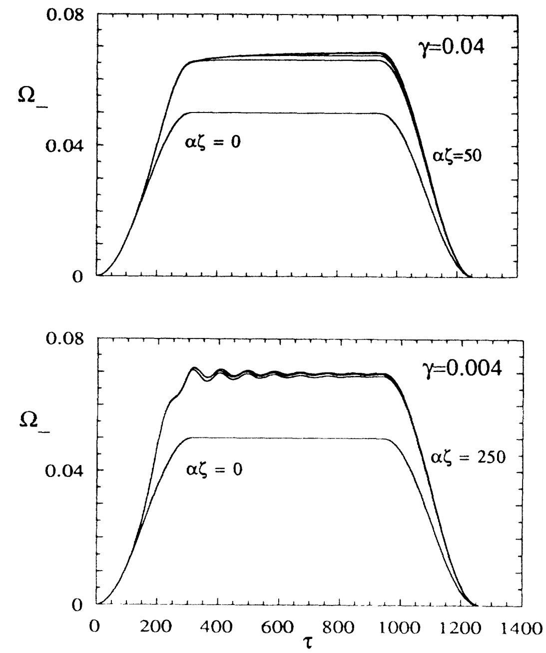

FIG. 4. Snapshots of the dressed-field pulse Ω_- defined in (4). Graphs for $\gamma=0.040$ and $\gamma=0.004$ were computed from the same data used for Figs. 2 and 3, respectively. Temporal variation is suppressed and spatial relaxation to steady state occurs rapidly and smoothly, before the second snapshot. In each case the stability of the dressed-field pulse is striking when compared to the variation and depletion shown for the individual pulses in Figs. 2 and 3.

regarded as an *atom-dressed combination of field amplitudes* Ω_P and Ω_S. It is well known in the atomic case that under appropriate circumstances the dressed superposition states have more convenient properties and are able to describe the physics more compactly than the bare states. The same is true for dressed fields, as we now demonstrate.

Using exactly the same computer data as in Figs. 2 and 3, we can replot the propagation physics in terms of Ω_- and Ω_+. Some of the results are shown in Fig. 4 for several depths of propagation, using the same graphical format as in Figs. 2 and 3. Comparisons clearly show the special character of the dressed state of the field denoted Ω_-. Recall that between Figs. 2 and 3 the atomic decay rate (and thus the absorption coefficient) differed by a factor of 10. This had a significant effect on the propagation of the bare fields (differences between Figs. 2 and 3), but Fig. 4 shows it has no significant effect on the dressed field. The dressed field quickly reaches a quasisteady state during propagation: For Ω_- the steady state has a practically constant amplitude that mimicks the input pulse shapes (as shown), and for Ω_+ the steady state is an almost zero value (not shown). This is a perfect parallel to the atomic case, where the trapped state quickly reaches a constant value while the orthogonal state, which is strongly coupled to the decay from the intermediate level, decays to zero.

We have carried our calculations over physically realistic distances, i.e., a few to many pump Beers lengths, and not to asymptotic distances. One can expect that eventually there will be modification of the dressed-field pulses, but it is important to note that over the realistic domain of our investigation we have found that our results apply to a wide variety of pulse pairs, even including temporally offset (but still overlapping) pulses and pulses with significantly different initial amplitudes. We can also say that the atomic level amplitudes, which we do not show here, follow the pulses in the sense that they also quickly reach quasisteady values that can be significantly different from their initial values. In particular, although c_2 is initially zero everywhere in the medium it adjusts quickly and can grow to more than 0.5 in the trailing edge of the excitation. Similarly, and importantly, we have found that c_1 can transiently exhibit values above 0.25 in magnitude.

In conclusion, we have presented the results of calculations of the spatial propagation of pairs of long optical pulses under simultaneous one-photon and two-photon resonance conditions. One consequence is a clear demonstration of the "protection" of a resonant pump pulse by a companion resonant Stokes pulse, as predicted by Harris [1,2] for EIT. In this sense our results provide additional information on EIT propagation. However, our main result is the discovery of dressed-field amplitudes. We have found that the *same dynamical variable* $c_{\rm trap}(t) \equiv -c_2(t)\Omega_P^* + c_0(t)\Omega_S^*$, which is traditionally interpreted as a field-dressed atomic amplitude with special properties in temporal evolution, can also be recognized as an *atom-dressed field amplitude* with special properties in spatial evolution. Our calculations show dressed fields that are generated during propagation without special attention to input pulse shapes and without the need for the preparation of inversion or coherence in the medium. They are relatively very stable during propagation, and stability is established rapidly and occurs even when propagation is accompanied by significant dynamical rearrangement of atomic level populations and significant degradation of "bare-field" pulse shapes. It is interesting to speculate on applications of dressed-field stability.

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[12] This is equivalent, given the main assumption about the principal role of the three-level system in the absorption process, to the simplification that the 0-1 and 1-2 oscillator strengths are approximately equal. For greater propagation distances than are contemplated in EIT, for example, more general assumptions about oscillator strengths, the nature of the detuning, the shape of the pulses and the type of relaxation can have interesting consequences. These will be described elsewhere in detail.
[13] All numerical quantities can be regarded as given in gigahertz for our lambda system. All frequencies scale with the size of the Rabi frequency, which sets the time-frequency scale for the problem.
[14] See B. W. Shore, *Theory of Coherent Atomic Excitation* (Wiley, New York, 1990), Vol. 1, Chap. 3.
[15] By their definition they are also distinct from the "normal modes" defined by Harris [2].