Knob for changing light propagation from subluminal to superluminal

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We show how the application of a coupling field connecting the two lower metastable states of a $\Lambda$ system can produce a variety of effects on the propagation of a weak electromagnetic pulse. In principle the light propagation can be changed from subluminal to superluminal. The negative group index results from regions of anomalous dispersion and gain in susceptibility.

A series of experiments have demonstrated both subluminal [1–5] and superluminal [6–8] propagation of light in a dispersive medium. The key to these successful demonstrations lies in one's ability to control optical properties of a medium with a laser field. Harris et al. [9] suggested how electromagnetically induced transparency (EIT) [10] can be used to obtain group velocities $v_g$ [11] much smaller than the velocity of light in vacuum. Early experiments [1,2] produced values of the group index $n_g = c/v_g$ in the range $10^2$–$10^3$. Hau et al. [3] could reduce the group velocity to 17 m/s in a Bose condensate. This was followed by an experiment in Rb vapor demonstrating reduction of the group velocity to 90 m/s [4] and to 8 m/s [5]. These experiments were based on the fact that EIT not only makes absorption zero at the line center but also leads to a dispersion profile [10,12] with a sharp derivative near the line center of the absorption line. In a different development Wang et al. [6] demonstrated superluminal propagation following the work of Chiao and co-workers [[13–16]; see also [17]]. Wang et al. used the stimulated Raman effect with the pumping beam replaced by a bichromatic field. This produces two regions of Raman gain with a region in between which has the right anomalous dispersion but with a negligible gain [18,19]. In this paper we propose a scheme where by changing a knob—an additional coupling field—one can switch the propagation of light from subluminal to superluminal. We use a $\Lambda$ system driven by a coherent control field, which has been extensively discussed in connection with subluminal propagation [1–5]. We apply, in addition, a field [referred to as the lower level (LL) coupling field] on the lower levels of the $\Lambda$ system and demonstrate how the application of the lower level coupling field can produce regions in the optical response with an appropriate dispersion profile. The dispersion can change from normal to anomalous depending on the intensity of the LL coupling field [20–22]. In addition, under suitable conditions the amplification of the light remains negligibly small.

We consider the scheme shown in Fig. 1(a). We consider propagation of a light pulse whose central frequency $\omega_1$ is close to the frequency of the atomic transition $|1\rangle \leftrightarrow |3\rangle$. We apply a control field on the optical transition $|1\rangle \leftrightarrow |2\rangle$. The transition $|2\rangle \leftrightarrow |3\rangle$ is generally an electric dipole forbidden transition. The states $|2\rangle$ and $|3\rangle$ are metastable states. We apply a field of frequency $\omega_3$ on the transition $|2\rangle \leftrightarrow |3\rangle$. The nature of this field will depend on the level structure. It could be a microwave field, say, in the case of Na, or an infrared field in the case of $^{208}$Pb. Moreover, it could be a dc field if one is considering transparency with Zeeman sublevels [5]. Let $2G = 2d_{12} \tilde{E}_c / \hbar$ and $2\Omega$ be the Rabi frequencies of the control field $\tilde{E}_c$ and the LL coupling field, respectively. The state $|1\rangle$ decays to the states $|3\rangle$ and $|2\rangle$ at the rates $2\gamma_1$ and $2\gamma_2$. For simplicity we ignore all collisional effects although these could easily be included. What is relevant for further consideration is the group velocity $v_g$ for the pulse applied to the transition $|1\rangle \leftrightarrow |3\rangle$. $v_g$ is related to the susceptibility $\chi'_{13}(\omega_1)$ for the transition $|1\rangle \leftrightarrow |3\rangle$:

$$v_g = \frac{c}{1 + 2\pi \chi'_{13}(\omega_1) + 2\pi \lambda \partial \chi'_{13}(\omega_1) / \partial \omega_1},$$

(1)

where $\chi'_{13}(\omega_1)$ is the real part of $\chi_{13}(\omega_1)$. We assume that we are working under conditions such that $\text{Im} \chi_{13}(\omega_1) = \chi_{13}''(\omega_1) = 0$. The susceptibility $\chi_{13}(\omega_1)$ will depend strongly on the intensities and the frequencies of the control laser and the LL coupling field. We concentrate on the group velocity although the actual pulse profiles could easily be simulated [23]. This susceptibility $\chi_{13}(\omega_1)$ is obtained by solving the density matrix equations for the $\Lambda$ system of Fig. 1(a), i.e., by calculating the density matrix element $\rho_{13}$ to first order in the applied optical field on the transition $|1\rangle \leftrightarrow |3\rangle$ but to all orders in the control field and the LL coupling field. By making a unitary transformation from the density matrix $\rho$ to $\sigma$ via

$$\rho_{12} = \sigma_{12} e^{-i\omega_2t}, \quad \rho_{13} = \sigma_{13} e^{-i(\omega_2 + \omega_3)t}, \quad \rho_{23} = \sigma_{23} e^{-i\omega_3t},$$

(2)

we have the relevant density matrix equations

$$\dot{\sigma}_{11} = iG\sigma_{21} + i\Omega \sigma_{32} - iG^*\sigma_{12} - i\Omega^*\sigma_{13} + 2\gamma_1 \sigma_{11},$$

$$\dot{\sigma}_{12} = iG^*\sigma_{12} + i\Omega \sigma_{23} - iG\sigma_{13} + 2\gamma_2 \sigma_{12},$$

$$\dot{\sigma}_{13} = -[\gamma_1 + \gamma_2 + \Gamma_{12} - i\Delta_2] \sigma_{12} + iG\sigma_{13} + i\gamma_2 e^{-i\Delta_3} \sigma_{13},$$

(3)

$$\dot{\sigma}_{22} = -[\gamma_2 + \gamma_3 + \Gamma_{23} + i\Delta_3] \sigma_{23} + i\gamma_3 e^{-i\Delta_3} \sigma_{23},$$

$$\dot{\sigma}_{23} = -[\Gamma_{23} - i\Delta_3] \sigma_{23} + iG^* \sigma_{13} + i\Omega \sigma_{33} - iG e^{-i\Delta_2} \sigma_{13} - i\Omega \sigma_{12},$$

We now consider the propagation of a pulse of light, which in the medium of the $\Lambda$ system, has frequency $\omega_1$, via the density matrix $\sigma_{11}$. The pulse is initially in a vacuum state $|0\rangle$ and at time $t = 0$ it is in a state $|\psi_0\rangle$ . The field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. Furthermore, the field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. The state of the field is initially in a vacuum state $|0\rangle$ and at time $t = 0$ it is in a state $|\psi_0\rangle$. The field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. Furthermore, the field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. The state of the field is initially in a vacuum state $|0\rangle$ and at time $t = 0$ it is in a state $|\psi_0\rangle$. The field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. Furthermore, the field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$. The state of the field is initially in a vacuum state $|0\rangle$ and at time $t = 0$ it is in a state $|\psi_0\rangle$. The field operator at time $t$ is $\hat{a}(t) = \int dt' \hat{a}(t) e^{i\omega_1(t-t')}$, where $\hat{a}(t)$ is an operator at time $t$ and $\hat{a}(t)'$ is a conjugate operator at time $t'$.