Electromagnetic-induced transparency and amplification of electromagnetic waves in photonic band-gap materials

Yuri V. Rostovtsev, Andrey B. Matsko, and Marlan O. Scully
Department of Physics, Texas A&M University, College Station, Texas 77843-4242

(Received 22 December 1997)

We investigate the propagation of an electromagnetic (EM) wave in a heterostructure formed by spatially modulated density of Λ three-level atoms. A new regime of propagation and amplification of EM waves with frequencies lying in the forbidden frequency range is found. An application of this phenomenon is discussed. [S1050-2947(98)08606-5]

PACS number(s): 42.50.Gy, 42.70.Qs, 42.25.Bs, 42.65.Tg

I. INTRODUCTION

There has been a growing interest in the study of the propagation of electromagnetic (EM) waves in periodic dielectric structures in the last decade [1]. This interest is strongly associated with an effect of photonic band-gap appearance in these structures. It is known that in photonic crystals a stop band with zero density of states for the propagation of EM modes in a given direction is formed. This phenomenon has been predicted theoretically and demonstrated experimentally [2–4].

The existence of a frequency gap where the propagation of EM waves is forbidden can lead to unusual quantum optical phenomena such as inhibition and enhancement of spontaneous emission [2,5,6], Anderson localization of light [3,7], photon-atom bound state [8–11], gap solitons [12,13], anomalous index of refraction [14], etc.

Many applications of photonic crystals utilizing these unique optical properties have been proposed [1]. For example, inhibition of spontaneous emission can be used to substantially enhance semiconductor laser operation.

Properties of photonic materials are usually investigated for weak EM fields, neglecting nonlinear properties of crystals. Meanwhile, photonic band-gap structures are often fabricated of dielectric materials with large nonlinearities, which can significantly change heterostructure properties in the presence of strong EM fields. It has been shown that the nonlinear periodic dielectrical structures admit solitary EM wave propagation in their band gaps [12,13] (self-induced transparency), shift the frequency of the band-gap position, and change the band-gap size [15]. Nonlinearity can lead to optical bistability [16] and chaotic behavior of the EM field [17] in the band-gap structures. A large amplitude EM wave propagating through different types of media is known to be subject to different nonlinear effects, such as Raman scattering, modulational instability, and self-focusing [13,18]. It is interesting to determine new features of these phenomena appearing due to the presence of a band gap.

The coherent effects in the atoms are known to give rise to new phenomena such as quenching of spontaneous emission, electromagnetic-induced transparency (EIT), lasing without population inversion, and high index of refraction without absorption [19]. Furthermore, recently it has been shown that in semiconductors [20] and in plasmas [21] there are similar coherent effects due to collective excitations of a medium. The EIT in cold overdense plasmas [21] and a new concept of the free-electron laser (FEL) [22] have been proposed.

In this paper we consider nonlinear EM wave propagation in the forbidden frequency region and possibility of control of such a gap via interaction with a strong EM field. For a model medium we use material doped by Λ three-level atoms (see Fig. 1). The gap arises as a result of spatial modulation of density of these dopants. For the sake of simplicity, we limit ourselves to the one-dimensional case.

We analyze the dispersion relations for the EM waves in the nonlinear photonic material in the presence of a strong driving field including both Stokes and anti-Stokes waves. We find EIT in such materials and the possibility of amplification for EM waves in band gap. These effects find possible application in Q switching, filtering, and others.

II. DISPERSION RELATIONS FOR BAND-GAP MATERIALS

Before we proceed to the dispersion relation for the nonlinear band-gap material, we consider a one-dimensional heterostructure consisting of a periodic array of two dielectric

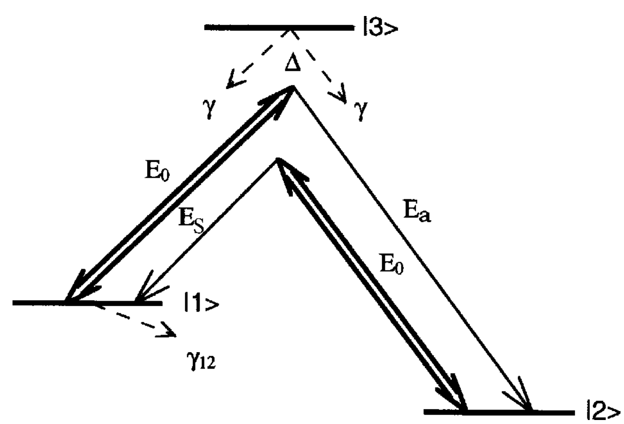

FIG. 1. A scheme of dopant levels. The strong coupling EM wave is detuned from the upper level $|3\rangle$ by Δ, so that a probability of one-photon processes is sufficiently small. It is possible to see the EIT phenomenon or amplification for Stokes sideband E_S in the photonic band-gap structure in the vicinity of two-photon resonance.

1050-2947/98/57(6)/4919(6)/$15.00

films with period $l=l_1+l_2$, where l_1 and l_2 are the thicknesses of the films with dielectric constants ϵ_1 and ϵ_2, respectively. For this structure, the simple dispersion equation can be written as

$$\frac{\partial^2 E(z,t)}{\partial z^2}-\frac{\epsilon(z)}{c^2}\frac{\partial^2 E(z,t)}{\partial t^2}=0. \tag{1}$$

Here $E(z,t)$ is an electric field of the EM wave inside a periodic structure, and $\epsilon(z)$ is a dielectric constant. A solution of the above equation is well known (see, for example, [13]). As a result, the dispersion relation for the EM wave with the frequency ω and the wave vector K has been obtained:

$$\cos(Kl)=\cos(K_1l_1+K_2l_2)-\frac{(\sqrt{\epsilon_1}-\sqrt{\epsilon_2})^2}{2\sqrt{\epsilon_1\epsilon_2}}\sin(K_1l_1)\sin(K_2l_2), \tag{2}$$

where $K_1=\omega\sqrt{\epsilon_1}/c$ and $K_2=\omega\sqrt{\epsilon_2}/c$.

To analyze the EM wave propagation through the nonlinear periodic structure, Eq. (1) should be solved by taking the dependence of $\epsilon(z)$ on the EM wave amplitude into account. To accomplish this we simplify the task by considering the dispersion relations (2) in the vicinity of the band gap. We assume that $Kl=\tilde{\kappa}l+\kappa l$, $K_1l_1=\tilde{k}_1l_1+k_1l_1$, $K_2l_2=\tilde{k}_2l_2+k_2l_2$. Here $\tilde{\kappa}l=\pi(1+2m)$, $\tilde{k}_1l_1=\pi(1+2m_1)$, and $\tilde{k}_2l_2=\pi(1+2m_2)$ (m, m_1, and m_2 are integer numbers), $\kappa l\ll 1$, $k_{1,2}l_{1,2}\ll 1$. Under these conditions the dispersion relation (2) can be approximated by

$$\kappa^2=\frac{\omega^2}{c^2}\left(\frac{\sqrt{\epsilon_1}l_1+\sqrt{\epsilon_2}l_2}{l}\right)^2-\frac{(\sqrt{\epsilon_1}-\sqrt{\epsilon_2})^2}{l^2\sqrt{\epsilon_1\epsilon_2}}. \tag{3}$$

Introducing

$$\sqrt{\epsilon_0}=\frac{\sqrt{\epsilon_1}l_1+\sqrt{\epsilon_2}l_2}{l},$$

and

$$\omega_g^2=\frac{c^2(\sqrt{\epsilon_1}-\sqrt{\epsilon_2})^2}{l^2\epsilon_0\sqrt{\epsilon_1\epsilon_2}},$$

we can rewrite Eq. (3) as

$$c^2\kappa^2=\epsilon_0(\omega^2-\omega_g^2). \tag{4}$$

This dispersion relation coincides with the dispersion relation for the EM waves in cold plasmas, where ω_g plays the role of the plasma frequency. On the other hand, the obtained relation is similar to that obtained in [9,10] if we consider small deviation of frequency ω from ω_g.

We are ready now to take into account the optical nonlinearities on the EM wave propagation in the periodic dielectrical structure. We consider the photonic band gap, which is produced by the spatial modulation of the density of optically active dopants incorporated into the host material with small linear dielectric susceptibility.

To analyze this problem we should solve the nonlinear wave equation

$$\hat{\mathcal{D}}_1(z,t)=\frac{\partial^2 E(z,t)}{\partial z^2}-\frac{\epsilon(z)}{c^2}\frac{\partial^2 E(z,t)}{\partial t^2}=\frac{4\pi}{c^2}\frac{\partial^2}{\partial t^2}P_{\rm nl}(z,t), \tag{5}$$

where $P_{\rm nl}(z,t)=(\chi_0^{(3)}+\Delta\chi^{(3)})E(z,t)^3$ in the one slide and $P_{\rm nl}(z,t)=\chi_0^{(3)}E(z,t)^3$ in the other. For linear parts of dielectric susceptibilities the expressions $\epsilon_1=\epsilon_0+\Delta\epsilon$ and $\epsilon_2=\epsilon_0$ are true. We also set $l_1=l_2$, and consider the small amplitude of the density modulation $\Delta\epsilon\ll\epsilon_0$ and $\Delta\chi^{(3)}\ll\chi_0^{(3)}$.

We use the amplitude of EM waves in the form

$$\begin{aligned}E(z,t)=&\{E_0+E_a\exp[i(kz-\omega t)]\\&+E_S\exp[-i(k^*z-\omega^*t)]\}\\&\times\exp[i(k_0z-\omega_0t)]\exp[i(\tilde{k}_0z-\tilde{\omega}_0t)]+\text{c.c.}\end{aligned} \tag{6}$$

Here a factor $\exp[i(\tilde{k}_0z-\tilde{\omega}_0t)]$ allows us to have the wave vector k in the first Brillouin zone, where we investigate the photonic band gap ($\tilde{k}_0$ consists of an integer number of reciprocal lattice vectors and $\tilde{\omega}_0$ is the corresponding frequency). We assume that the sideband amplitudes E_a (anti-Stokes) and E_S (Stokes) are much smaller in comparison with the carrier-wave amplitude E_0.

To analyze the possibility of a wave propagation we have to make the harmonic analysis of the strict equation (5), and obtain the dispersion relation for the heterostructure. Strict analysis of this equation is very complicated. To simplify it we have used the common method. For the linear part of the dispersion relation we can write a differential operator, which is an approximation of the strict operator $\hat{\mathcal{D}}_1(z,t)$ in the vicinity of the chosen band gap.

$$\hat{\mathcal{D}}_2(z,t)=\frac{\partial^2}{\partial z^2}-\frac{\epsilon_0}{c^2}\frac{\partial^2}{\partial t^2}-\frac{\epsilon_0}{c^2}\omega_g^2. \tag{7}$$

Then we replace the strict equation (5) by the equation

$$\begin{aligned}&\hat{\mathcal{D}}_2(z,t)\{E(z,t)\exp[-i(\tilde{k}_0z-\tilde{\omega}_0t)]\}\\&=\frac{4\pi}{c^2}\exp[-i(\tilde{k}_0z-\tilde{\omega}_0t)]\frac{\partial^2}{\partial t^2}P_{\rm nl}(z,t).\end{aligned} \tag{8}$$

In the approximation of the large coupling field E_0 for Fourier amplitudes of the Stokes and anti-Stokes fields the following set of equations can be written:

$$D_+E_a=\frac{16\pi^3}{\lambda_a^2}(\chi_{a,a}^{(3)}|E_0|^2E_a+\chi_{a,S}^{(3)}E_0^2E_S^*), \tag{9}$$

$$D_-E_S^*=\frac{16\pi^3}{\lambda_S^2}(\chi_{S,S}^{(3)}|E_0|^2E_S^*+\chi_{S,a}^{(3)}E_0^{*2}E_a). \tag{10}$$

Here λ_S and λ_a are the Stokes and anti-Stokes sideband wavelengths, respectively; the nonlinear susceptibilities $\chi^{(3)}$ depend on frequencies of fields;

$$D_{\pm}(\omega,k)=\frac{\epsilon_0}{c^2}\omega^2\pm 2\left(\frac{\epsilon_0}{c^2}\omega_0\omega-kk_0\right)-k^2$$

are the light-wave dispersion functions in the vicinity of the band gap;

$$\omega_0^2=\omega_g^2+\frac{c^2}{\epsilon_0}k_0^2 \tag{11}$$

is the dispersion relation for the sufficiently weak coupling field (i.e., $\epsilon_0 \gg 4\pi\chi_0^{(3)}|E_0|^2$).

In this case, the unstable solutions of Eqs. (9) and (10) automatically mean the instability of the sideband waves. From these equations we obtain the dispersion relation for EM waves in the band-gap heterostructure

$$\mathcal{D}(\omega,k)=\left(D_+-\frac{16\pi^3}{\lambda_a^2}\chi_{a,a}^{(3)}|E_0|^2\right)\left(D_--\frac{16\pi^3}{\lambda_S^2}\chi_{S,S}^{(3)}|E_0|^2\right)-\left(\frac{16\pi^3}{\lambda_a\lambda_S}\right)^2\chi_{a,S}^{(3)}\chi_{S,a}^{(3)}|E_0|^4=0. \tag{12}$$

This dispersion relation is similar to the dispersion relation of the EM waves in plasmas [23], but it has different dependence for nonlinear susceptibility of Stokes and anti-Stokes sidebands on the frequency. Analyzing the roots of Eq. (12) we investigate amplification and generation of Stokes and anti-Stokes waves in the medium. Also it is worth mentioning that many forbidden zones can exist in photonic band-gap materials, but only one zone can exist in plasmas.

III. ANALYSIS OF STABILITY FOR THE HETEROSTRUCTURE DOPED BY THREE-LEVEL ATOMS

To analyze the possible behavior of the dispersion relation (12) we consider dopants to be Λ three-level atoms (see Fig. 1). We assume that the field $E(z,t)$ couples to both $|1\rangle$-$|3\rangle$ and $|2\rangle$-$|3\rangle$ transitions and has a large detuning from one-photon resonance. Then the linear part of dielectric susceptibility for this field is a constant, but the nonlinear part has two-photon resonant behavior.

This scheme is thoroughly investigated in [24], and we use the results of those calculations for the nonlinear polarization. As a result Eqs. (9) and (10) take the form

$$D_+E_a=-\frac{16\pi^3}{\hbar^3\lambda_0^2}\frac{|\mathcal{P}|^4\mathcal{N}}{\Delta(\delta\omega+i\gamma)}\left(\frac{|E_0|^2E_a}{\Delta}+\frac{E_0^2E_S^*}{\Delta+\omega_{12}}\right)\times(\rho_{11}-\rho_{22}), \tag{13}$$

$$D_-E_S^*=-\frac{16\pi^3}{\hbar^3\lambda_0^2}\frac{|\mathcal{P}|^4\mathcal{N}}{(\Delta+\omega_{12})(\delta\omega+i\gamma)}\left(\frac{E_0^{*2}E_a}{\Delta}+\frac{|E_0|^2E_S^*}{\Delta+\omega_{12}}\right)\times(\rho_{11}-\rho_{22}). \tag{14}$$

Here Δ is the difference between the coupling field frequency and the frequency of the transition between level $|1\rangle$ and level $|3\rangle$, ρ_{11} and ρ_{22} are the populations of $|1\rangle$ and $|2\rangle$ levels, respectively, λ_0 is the wavelength of the coupling field, ω_{12} is the frequency of the transition $|1\rangle$-$|2\rangle$, $\delta\omega=\omega-\omega_{12}$ is the two-photon transition detuning, $\mathcal{P}$ is the dipole momentum of the transitions $|3\rangle$-$|1\rangle$ and $|3\rangle$-$|2\rangle$, and $\mathcal{N}$ is the dopant density. We assume that due to the large detuning Δ it is possible to omit the upper level decay rate γ in the expression $\Delta+i\gamma$ and set the population of the upper level $|3\rangle$ to zero. As a result, the sum of the lower level populations is equal to unity, $\rho_{11}+\rho_{22}=1$.

The dispersion relation in this case can be obtained from Eqs. (13) and (14),

$$D_+D_-+\frac{\Gamma_0}{\delta\omega+i\gamma}\frac{\omega_0^2}{c^2}\left[D_-+\left(\frac{\Delta}{\Delta+\omega_{12}}\right)^2D_+\right]=0, \tag{15}$$

where

$$\Gamma_0=\frac{4\pi|\mathcal{P}|^2|E_0|^2}{\hbar^2\Delta^2}\;\frac{|\mathcal{P}|^2\mathcal{N}}{\hbar}\;\;\frac{4\pi^2c^2}{\lambda_0^2\omega_0^2}(\rho_{11}-\rho_{22}).$$

Equation (15) is similar to the case of plasmas [23], if we set $\Delta \gg \omega_{12}$. It has been shown in [21] that EIT exists in cold homogeneous plasmas for frequencies below the plasma frequency. Therefore we can expect that the nonlinear photonic band-gap material has a similar property. Also from Eq. (15) we see that the amplification of the Stokes wave is possible under the usual condition of Raman inversion, i.e., $\rho_{22}>\rho_{11}$.

Equation (15), which is fourth order in the wave vector k and fifth order of frequency ω, has been solved numerically. There are two methods to investigate the dispersion relation. The first method is to investigate the behavior of the frequency ω as a function of the real wave vector. Such consideration is useful for an infinite uniform medium or a finite medium with periodic boundary conditions [18]. The second method is to consider the propagation of EM waves through nonlinear heterostructure. In this case it is necessary to find a spatial solution. To solve this problem we analyzed the behavior of the wave vector values obtained in Eq. (15) for the real frequency ω.

For parameters $\epsilon_0=1$, $\Delta=0.8\omega_0$, $\omega_{12}=0.3\omega_0$, and $\gamma_{12}=0$ we considered three cases. The first case was the linear heterostructure, where the coupling field was set equal to zero ($\omega_g=0.8\omega_0$, $\Gamma_0=0$). The solutions for this case are presented in Figs. 2 and 3 [for convenience we use the same line style for plotting the real and imaginary parts of every root of the solution to simplify tracing a particular solution of Eq. (15)]. The second case consisted of the nonlinear medium without photonic band gap ($\omega_g=0$, $\Gamma_0=0.1\omega_0$). The solutions for this case are presented in Figs. 4 and 5. Finally, the last case consisted of the nonlinear band-gap material ($\omega_g=0.8\omega_0$, $\Gamma_0=0.1\omega_0$). The solutions for this case are presented in Figs. 6 and 7. It should be mentioned here, that due to the shifting of our frame of reference, the frequencies in the dispersion relation are small ($\ll c/l\sqrt{\epsilon_0}$). To obtain the

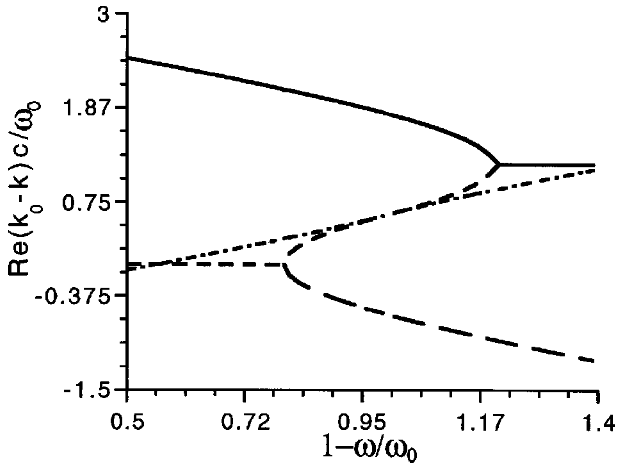

FIG. 2. The dispersion of Stokes sideband EM wave $\mathrm{Re}(k_0-k)c/\omega_0$ in the linear photonic heterostructure as a function of the real frequency $1-\omega/\omega_0$. It is easy to see the region in which the dispersion is constant and does not depend on the frequency. In this region the group and phase velocities of EM waves are indefinite. Physically it means that the EM waves with such frequencies are reflected from the heterostructure.

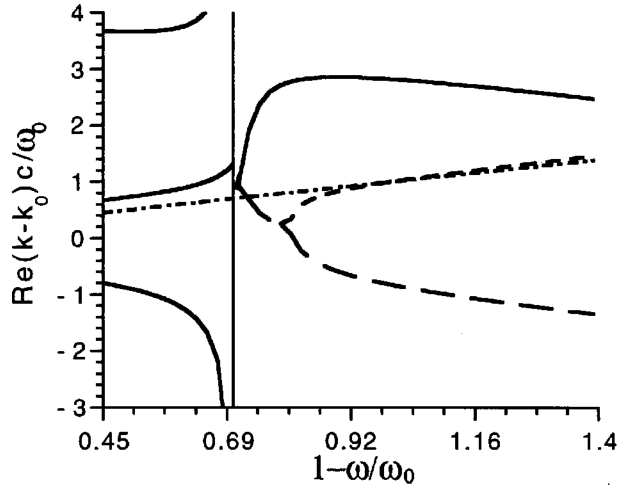

FIG. 4. The dispersion of EM waves in the nonlinear medium (without photonic band gap), $\mathrm{Re}(k_0-k)c/\omega_0$ as a function of the real frequency $1-\omega/\omega_0$.

actual values, one needs to add parts $\tilde{\omega}_0$ and $\tilde{k}_0$ to all the frequencies and wave vectors. Such transformation does not change the physical picture.

The results are analyzed next. There is an EIT gap close to the point of two-photon resonance (Fig. 7). The physics of EIT here is common. The nonlinear interaction of the sideband and coupling waves leads to a modulation of nonlinear refractive index. In the vicinity of the two-photon resonance the amplitude of the modulation increases and the obtained dynamical diffraction lattice compensates the phase shift due to the linear part of the refractive index of the heterostructure. This compensation prevents Bragg diffraction and allows the medium to appear transparent. This process is different from allowance of wave propagation which arises when dopants with one-photon transition are embedded into the photonic band gap [11,25], because in the one-photon transition case there are one-photon losses of radiation. In contrast, our system can be lossless. However, usually coherence decay rate between levels $|1\rangle$ and $|2\rangle$ is not equal to zero ($\gamma_{12}\neq 0$) and the EIT gap vanishes as this parameter increases.

In the linear photonic band-gap materials, spatial instabilities (nonzero proper imaginary part of the wave vector for the real frequency) are not useful for wave amplification because corresponding group velocity for unstable branches is undefined (see Fig. 2). The electromagnetic wave with the frequency lying in the gap reflects from such material due to Bragg diffraction. Nevertheless, in the nonlinear heterostructure this situation changes due to strong nonlinear wave interaction, which makes possible amplification (Fig. 6). The presence of the nonlinear dopants changes the properties of the media and makes possible the amplification even when the resonant frequencies for the anti-Stokes and Stokes waves $\tilde{\omega}_0+\omega_0\pm\omega_{12}$ lie outside of the band gap.

The behavior of the dispersion for EM waves (Figs. 6 and 7) in the nonlinear heterostructure has an essential singular-

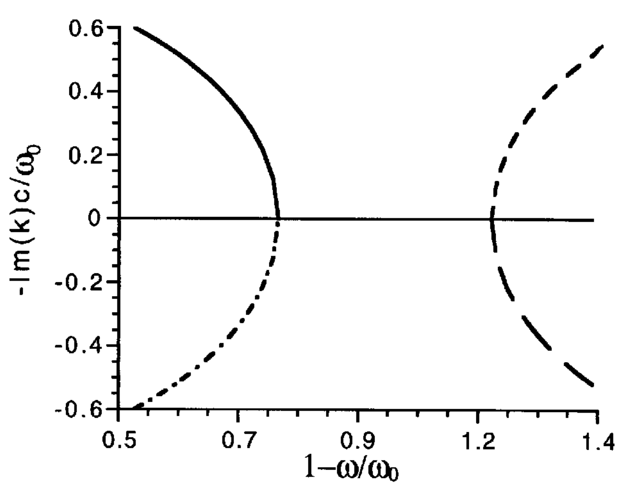

FIG. 3. The imaginary part of the wave vector for EM waves in the linear heterostructure, corresponding to Fig. 2, $\mathrm{Im}\,kc/\omega_0$ as a function of the real frequency $1-\omega/\omega_0$. The imaginary part is nonzero only in the region with indefinite phase and group velocities. As a result, it does not lead to amplification of the EM field.

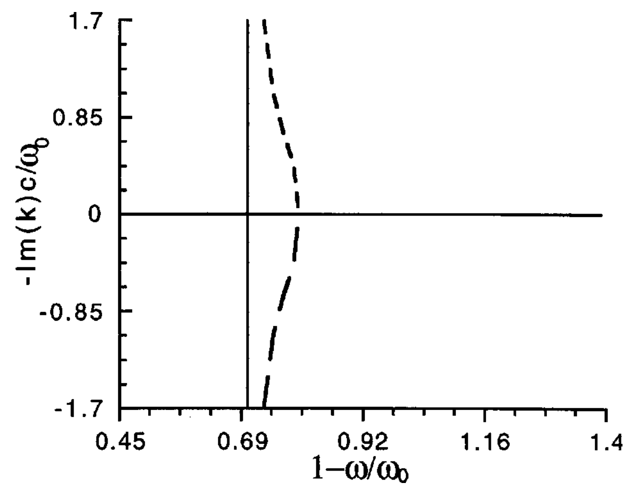

FIG. 5. The spatial growth rate of the EM wave instability in the nonlinear medium $\mathrm{Im}\,kc/\omega_0$, corresponding to Fig. 4, as a function of the real frequency. It is easy to see the two-photon resonance.

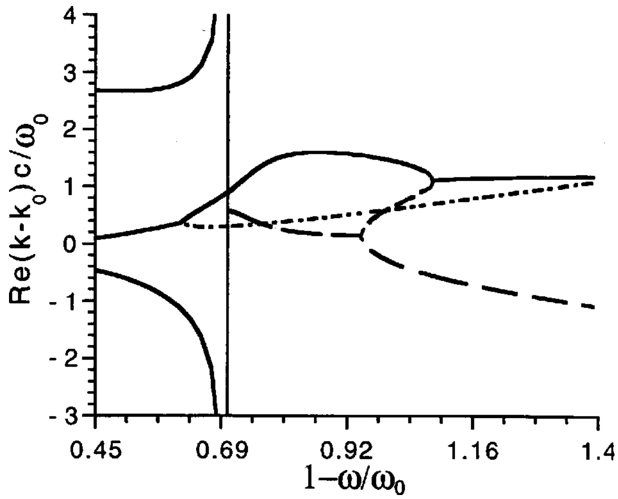

FIG. 6. The dispersion of EM waves in the nonlinear heterostructure, $\mathrm{Re}(k_0 - k)c/\omega_0$ as a function of the real frequency $1-\omega/\omega_0$. Due to nonlinear interaction of EM waves optical properties of the photonic band-gap material change and the propagation of EM waves in the vicinity of two-photon resonance becomes possible. This happens because the phase and group velocities are determined and nonzero.

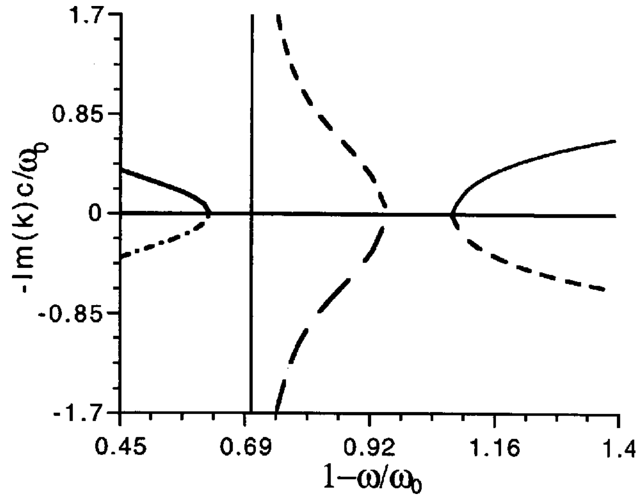

FIG. 7. The spatial growth rate of the EM wave instability in the nonlinear heterostructure $\mathrm{Im}\, kc/\omega_0$, corresponding to Fig. 4, as a function of the real frequency. The EIT gap is clearly seen in the vicinity of the two-photon resonance. The shape of the Stokes line is noticeably changed when compared with Fig. 5.

ity at the point $\omega = \omega_{12}$. This singularity corresponds to the two-photon resonance of the dopants (see Figs. 4 and 5). It is similar to the beam-plasma instability for one-dimensional electrostatic plasma oscillation at the high frequencies where the motion of ions is neglected [18] and an instability of multiwave propagation in cold overdense plasmas [26]. The essential singularity disappears for arbitrary nonzero γ_{12}.

To find out the amplification of the probe field E_S we used Figs. 6 and 7. Gain is equal to $\kappa = |\mathrm{Im}\, k|L$, where L is the length of the sample. Let us consider the situation when $L = 1$ cm, $\epsilon_0 = 9$, $\Delta = 8\times 10^9$ s^{-1}, $\omega_{12} = 3\times 10^9$ s^{-1}, $\gamma_{12} = 10^5$ s^{-1}, $\omega_g = 8\times 10^9$ s^{-1}, $\Gamma_0 = 10^9$ s^{-1}, $\omega_0 = 3\times 10^{10}$ s^{-1}, and $\omega = 2\times 10^9$ s^{-1}. We found the gain under these conditions to be approximately equal to $\kappa = 2.4$. However, for the simple nonlinear medium without photonic band gap (see Figs. 4 and 5) there is no gain under these conditions. This demonstrates gain line broadening in the nonlinear heterostructure.

In the present paper we have shown that propagation and amplification of EM waves inside the nonlinear photonic band-gap structure in the presence of the strong coupling field are possible. These processes are similar to propagation and amplification of EM waves in overdense plasmas. One can foresee that such results can be obtained for every nonlinear material with a forbidden gap, for example, for a nonlinear one-mode fiber.

Here we should stress that the photonic band gap does not restrict two-photon transitions of the dopants even when the frequency of the Stokes field lies in the forbidden gap. The presence of a sufficiently large number of dopants, in contrast to an atom embedded in the heterostructure, can significantly change the properties of the band-gap material. Therefore for an adequate description of the doped heterostructure a self-consistent problem taking into account changing of the index of refraction which appears due to the dopants should be undertaken.

Proposed EIT and amplification effects can be used for the construction of Q-switch devices. For example, the dielectric heterostructures are usually utilized for fabrication of multilayer dielectrical mirrors. Doping of selected layers of the mirror by multilevel atoms can allow us to change the mirror reflectance by changing the amplitude of the coupling field. Moreover, we can make this mirror absolutely transparent or amplified in the chosen bandwidth, according to the above results. This effect can be used in pulse generation and in the production of nonclassical states of light.

ACKNOWLEDGMENTS

The authors wish to thank D. Volchenkov for stimulating and helpful discussions, and they gratefully acknowledge support from the Office of Naval Research, the Welch Foundation, and the Texas Advanced Research and Technology Program. One of us (Yu.V.R.) wishes to thank the North Atlantic Treaty Organization for their financial support.

[1] For a review see articles in *Photonic Band Gaps and Localization*, Vol. 38 of *NATO Advanced Study Institute, Series B: Physics*, edited by C.M. Soukoulis (Plenum, New York, 1993); J. Mod. Opt. **41** (2) (1994), special issue on photonic band structures; *Photonic Band Gap Materials*, Vol. 315 of *NATO Advanced Study Institute, Series E: Applied Science*, edited by E.N. Ecomonou (Plenum, New York, 1996).

[2] E. Yablonovich, Phys. Rev. Lett. **58**, 2059 (1987); E. Yablonovich and T.J. Gmitter, *ibid.* **63**, 1950 (1989).

[3] S. John, Phys. Rev. Lett. **58**, 2486 (1987); S. John and R. Rangavajan, Phys. Rev. B **38**, 10 101 (1988).

[4] K.M. Ho, C.T. Chan, and C.M. Soukoulis, Phys. Rev. Lett. **65**, 3152 (1990); E. Ozbay *et al.*, Opt. Lett. **19**, 1155 (1994); E. Ozbay *et al.*, Phys. Rev. B **50**, 1945 (1994); E. Ozbay *et al.*,

Appl. Phys. Lett. **64**, 2059 (1994); M. Wada, K. Sakoda, and K. Inone, Phys. Rev. B **52**, 16 297 (1996); I.I. Tarhan and J.H. Watson, Phys. Rev. Lett. **76**, 315 (1996); Yu.A. Vlasov *et al.*, Phys. Rev. B **55**, R13 357 (1997); V.N. Bogomolov *et al.*, Phys. Rev. E **55**, 7619 (1997).

[5] J. Martorell and N.M. Lawandy, Phys. Rev. Lett. **65**, 1877 (1990); S. John and T. Quang, Phys. Rev. A **50**, 1764 (1994); G. Kweon and N.M. Lawandy, Opt. Commun. **118**, 388 (1995); Shi-Yao Zhu, Hong Chen, and Hu Huang, Phys. Rev. Lett. **79**, 205 (1997).

[6] A.G. Kofman, G. Kurizki, and B. Sherman, J. Mod. Opt. **41**, 345 (1994).

[7] M. Ohtaka *et al.*, Jpn. J. Appl. Phys., Part 1 **34**, 4502 (1995).

[8] S. John and T. Quang, Phys. Rev. Lett. **74**, 3419 (1995).

[9] S. John and J. Wang, Phys. Rev. Lett. **64**, 2418 (1990).

[10] S. John and J. Wang, Phys. Rev. B **43**, 12 772 (1991).

[11] R.D. Meade, K.D. Brommer, A.M. Rappe, and J.D. Joannopoulos, Phys. Rev. B **44**, 13 772 (1991).

[12] W. Chen and D.L. Mills, Phys. Rev. Lett. **58**, 160 (1987); D.L. Mills and S.E. Trullinger, Phys. Rev. B **36**, 917 (1987); S. John and N. Akozbek, Phys. Rev. Lett. **71**, 1168 (1993); M. Scalora, J.P. Dowling, C.M. Bowden, and M.J. Bloemer, *ibid.* **73**, 1368 (1994); A. Kozhekin and G. Kurizki, *ibid.* **74**, 5020 (1995); V.I. Rupasov and M. Singh, Phys. Rev. A **54**, 3614 (1996); Z. Cheng and G. Kurizki, *ibid.* **54**, 3576 (1996).

[13] D.L. Mills, *Nonlinear Optics: Basic Concepts* (Springer-Verlag, Berlin, 1991).

[14] J.P. Dowling and C.M. Bowden, J. Mod. Opt. **41**, 345 (1994).

[15] P. Tran, Phys. Rev. B **52**, 10 673 (1995); J. Opt. Soc. Am. B **52**, 2589 (1997).

[16] H.G. Winful, J.H. Marburger, and E. Garmire, Appl. Phys. Lett. **35**, 379 (1979).

[17] H.G. Winful and G.D. Cooperman, Appl. Phys. Lett. **40**, 298 (1982).

[18] *Handbook of Plasma Physics*, edited by M.N. Rozenbluth and R.Z. Sagdeev, Basic Plasma Physics Vol. I, edited by A.A. Galeev and R.N. Sudan (North-Holland, New York, 1983).

[19] O.A. Kocharovskaya and Ya.I. Khanin, Pis'ma Zh. Éksp. Teor. Fiz. **48**, 581 (1988) [JETP Lett. **48**, 630 (1988)]; S.E. Harris, Phys. Rev. Lett. **62**, 1033 (1989); M.O. Scully, S.-Y. Zhu, and A. Gavrielides, *ibid.* **62**, 2813 (1989); A.S. Zibrov *et al.*, *ibid.* **75**, 1499 (1995); G.G. Padmabandu *et al.*, *ibid.* **76**, 2053 (1996).

[20] A. Imamoglu and R.J. Ram, Opt. Lett. **19**, 1744 (1994).

[21] S.E. Harris, Phys. Rev. Lett. **77**, 5357 (1996).

[22] B. Sherman, G. Kurizki, D.E. Nikonov, and M.O. Scully, Phys. Rev. Lett. **75**, 4602 (1995); D. E. Nikonov, B. Sherman, G. Kurizki, and M.O. Scully, Opt. Commun. **123**, 363 (1996); D.E. Nikonov, M.O. Scully, and G. Kurizki, Phys. Rev. E **45**, 6780 (1996).

[23] C.J. McKinstre and R. Bingham, Phys. Fluids B **1**, 230 (1989); **4**, 2626 (1992); A.S. Sakharov and V.I. Kirsanov, Phys. Rev. E **49**, 3274 (1994).

[24] M. Fleischhauer *et al.*, Phys. Rev. A **46**, 1468 (1992); U. Rathe *et al.*, *ibid.* **47**, 4994 (1993).

[25] E. Yablonovitch *et al.*, Phys. Rev. Lett. **67**, 3380 (1991).

[26] A.B. Matsko, Yu.V. Rostovtsev, and M.O. Scully (unpublished).