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Probe spectrum of a four-level atom in a double-band photonic crystal*

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In this paper, the probe absorption spectrum of an atom in a double-band photonic crystal have been studied. In the modes, we assume that one of the two atomic transitions in a Λ-type atomic system is interacting with free vacuum modes, and another transition is interacting with free vacuum modes, isotropic photonic band gap (PBG) modes and anisotropic PBG modes, separately. The effects of the fine structure of the atomic lower levels on the probe absorption spectrum are investigated in detail in the three cases. The most interesting thing is that the two (four) transparencies at one (two) probe absorption peak(s), caused by the fine structure of the lower levels of an atom, are predicted in the case of isotropic PBG modes.

Keywords: photonic band gap, probe absorption spectrum
PACC: 3200, 3270, 4250

1. Introduction

The control of spontaneous emission and probe absorption of an atom will lead to a number of potential applications in high-precision measurements, quantum computation, teleportation, quantum information theory and lasing without population inversion. Therefore, the study of spontaneous emission and probe absorption has attracted much attention in recent years. It is well known that the spontaneous emission and probe absorption of an atom depend not only on the properties of the excited atomic system but also on the nature of the surrounding environment.[1] The quantum interference of spontaneous emission and probe absorption is widely studied. A classic example is the V-type three-level system consisting of two close excited levels and one ground state in the same environment reservoirs. Quantum interference phenomena will occur if the dipole moments of the two arms of transition are not orthogonal. The interference in such a system arises from indistinguishable spontaneous emission pathways. The interference, vacuum-induced or additionally driven by external coherent laser fields is widely studied, and leads to many remarkable phenomena, including very narrow absorption and fluorescence spectra,[2] population trapping in a degenerate system,[3] fluorescence quenching in the free space,[4] phase-dependent line shapes,[5,6] lasing without inversion,[7] etc.[8] From the viewpoint of the surrounding environment of atoms, photonic band gap (PBG) structures have different density of states (DOS) compared with a free-space vacuum field.[1,9,10] The study of quantum and nonlinear optical phenomena for types of atoms embedded in such PBG materials leads to the prediction of many interesting effects: for example, localization of light,[11,12] photon-atom bound states,[1,11] suppression and even complete cancellation of spontaneous emission,[13,14] enhancement of spontaneous emission,[15,16] electro-magnetically induced transparency and other phenomena.[17]

Many papers have reported on the population involution and distribution in upper levels [1,11] and the spectral components in the spontaneous emission fields [15,16] of an atom embedded in a photonic crystal. Of particular interest, Knight and his co-workers [17] investigated transparency phenomena of an atom in the Λ-type model where one transition interacts with free vacuum modes and another interacts with modes near the edges of photonic band gap. Focusing on the quantum interference of an atom without external driving coherent laser fields, Zhu [2] independently re-

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ported the quantum interference of spontaneous emission spectrum and probe absorption spectrum of an atom coupled to vacuum space reservoirs. In our previous papers,[18–20] the quantum interference in a spontaneous emission spectrum of an atom embedded in a double-band photonic crystal was reported. Only one lower level was considered in the above studies. From the point of view of practice the ground-state level of an atom with a superfine energy level structure is of versatility. In this paper, we adopt the model of a four-level atom with a superfine energy level structure in a double-band photonic crystal. We study the transition from an upper level to one lower level coupled to free vacuum modes, and the transitions from an upper level to two lower levels with fine structure coupled to the modes which are isotropic PBG, anisotropic PBG and free vacuum, respectively. The probe absorption spectrum with the transition coupled to free vacuum modes is investigated. In Section 2, the equations for describing the probe absorption spectrum of the four-level atom in a photonic crystal are deduced. The general calculated results and analysis are presented in Section 3. The major conclusions are summarized in Section 4.

2. Equations for the probe absorption spectrum

Consider a four-level atom with upper level $|3\rangle$ and lower levels $|2\rangle$, $|0\rangle$ and $|1\rangle$, as shown in Fig.1. The transitions from the upper level $|3\rangle$ to the lower level $|1\rangle$ are synchronously coupled by a free vacuum mode (ω_λ) and a weak probe field (ω_p), and to the lower levels $|2\rangle$ and $|0\rangle$ by the double-band isotropic PBG modes, anisotropic PBG modes and free vacuum modes, respectively. The dispersion relations near the photonic band edges are approximated by [18]

$$\begin{aligned} \omega_k &= \omega_{c1} - A_1(k-k_0)^2, \quad k < k_0 \\ \omega_k &= \omega_{c2} + A_2(k-k_0)^2, \quad k > k_0 \end{aligned} \tag{1}$$

$$\begin{aligned} \omega_k &= \omega_{c1} - A_1|(\boldsymbol{k}-\boldsymbol{k}_0)|^2, \quad |\boldsymbol{k}| < |\boldsymbol{k}_0| \\ \omega_k &= \omega_{c2} + A_2|(\boldsymbol{k}-\boldsymbol{k}_0)|^2, \quad |\boldsymbol{k}| > |\boldsymbol{k}_0| \end{aligned} \tag{2}$$

for isotropic and anisotropic dispersion relations, respectively. Here $A_n \approx \omega_{cn}/k_0^2$ $(n = 1, 2)$. ω_{c1} and ω_{c2} are the upper and lower frequencies at the edges of the forbidden gap, respectively. The Hamiltonian describing the dynamics of this system in the interaction picture and the rotating wave approximation can be written as

$$\begin{aligned} H_I =& i\hbar(\Omega_{31} b_1^+ b_3 e^{i\mu_p^{31} t} + C.C.) \\ &+ i\hbar \sum_{\lambda e} (g_{\lambda e}^{31} a_{\lambda e}^+ b_1^+ b_3 e^{i\mu_\lambda^{31} t} + C.C.) \\ &+ i\hbar \sum_{ke} (g_{ke}^{30} a_{ke}^+ b_0^+ b_3 e^{i\mu_k^{30} t} \\ &+ g_{ke}^{32} a_{ke}^+ b_2^+ b_3 e^{i\mu_k^{32} t} + C.C.), \end{aligned} \tag{3}$$

where b_i^+, b_i (i=0,1,2,3) are creation and annihilation operators for an atom in level i, respectively; $a_{\lambda e}^+$ ($a_{\lambda e}$) is the creation (annihilation) operator for the λth vacuum mode with frequency ω_λ; a_{ke}^+ (a_{ke}) is the creation (annihilation) operator for the kth PBG mode with frequency ω_k; Ω_{31} is the Rabi frequency of a probe field coupling with the transition from the upper level to the lower level $|1\rangle$; $g_{\lambda e}^{31}$ is the frequency-dependent coupling constant between the atomic transition $|3\rangle$-$|1\rangle$ and the free vacuum modes ω_λ, and g_{ke}^{3j} (j=0, 2) is another coupling constant between the atomic transition $|3\rangle$-$|j\rangle$ and PBG modes ω_k. More precisely,

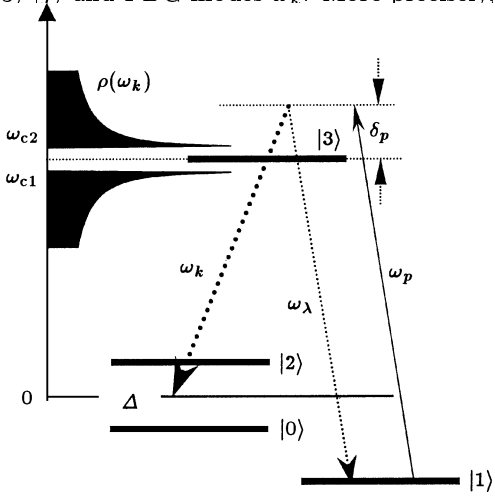


Fig.1. Schematic illustration of the four-level atom in an isotropic double-photonic-band structure. $\rho(\omega_k)$ denotes the density of state of the PBG modes. ω_{c2} and ω_{c1} are the upper and lower frequencies of the forbidden gap, respectively. It is assumed that the transitions from the upper level $|3\rangle$ to the lower levels $|0\rangle$ and $|2\rangle$ are coupled by PBG modes ω_k. The transition from the upper level $|3\rangle$ to the lower level $|1\rangle$ are coupled by both a weak probe laser field ω_p and the free vacuum modes ω_λ. δ_p is the detuning of frequency of probe field from the middle of the forbidden gap. Δ denotes the splitting width of the lower state levels.

$$
\begin{aligned}
\Omega_{31} =& d_{31}E/2\hbar = \sqrt{\gamma_{31}}E\Omega, \\
g_{\lambda e}^{31} =& (\omega_{31}d_{31}/\hbar)(\hbar/(2\varepsilon_0\omega_\lambda V))^{1/2}\boldsymbol{e}_{\lambda e}\cdot\boldsymbol{d}_{31}, \\
g_{ke}^{3j} =& (\omega_{3j}d_{3j}/\hbar)(\hbar/(2\varepsilon_0\omega_k V))^{1/2}\boldsymbol{e}_{ke}\cdot\boldsymbol{d}_{3j}, \\
& (j=0,2), \\
\mu_p^{31} =& \omega_p - \omega_{31}, \\
\mu_\lambda^{31} =& \omega_\lambda - \omega_{31}, \\
\mu_k^{3j} =& \omega_k - \omega_{3j}\ (j=0,2),
\end{aligned} \tag{4}
$$

where $\boldsymbol{d}_{32}$ and $\boldsymbol{d}_{30}$ are the atomic dipole moment unit vectors for the transitions $|3\rangle$-$|2\rangle$ and $|3\rangle$-$|0\rangle$, respectively; $\boldsymbol{e}_{\lambda e}$ and $\boldsymbol{e}_{ke}$ (e=1, 2) are the polarization unit vectors of the radiation fields from the double V-type transitions, E is the amplitude of the weak probe field, and V is the sample volume. The state vector of the system in the interaction picture can be written as

$$
\begin{aligned}
|\Psi(t)> =& A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}|3,\{0\}> + A_1(t)|1,\{0\}> \\
&+ \sum_{\lambda e} A_{1\lambda e}|1,\{\lambda e\}> \\
&+ \sum_{ke} A_{0ke}|0,\{ke\}> \\
&+ \sum_{ke} A_{2ke}|2,\{ke\}>,
\end{aligned} \tag{5}
$$

where the amplitude $A_iA_i^*$ (i=0,1,2,3) represent the probability of an atom occurring in level i. Substituting Eqs.(3) and (5) into the Schrödinger equation, we obtain

$$
\begin{aligned}
\frac{\mathrm{d}(A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t})}{\mathrm{d}t} =& -\mathrm{i}\Omega_{31}\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}A_1(t) \\
&- \mathrm{i}\sum_{\lambda e} g_{\lambda e}^{31}\mathrm{e}^{-\mathrm{i}\mu_\lambda^{31}t}A_{1\lambda e}(t) \\
&- \mathrm{i}\sum_{ke} g_{ke}^{30}\mathrm{e}^{-\mathrm{i}\mu_k^{30}t}A_{0ke}(t) \\
&- \mathrm{i}\sum_{ke} g_{ke}^{32}\mathrm{e}^{-\mathrm{i}\mu_k^{32}t}A_{2ke}(t), \\
\frac{\mathrm{d}A_1(t)}{\mathrm{d}t} =& -\mathrm{i}\Omega_{31}A_3(t), \\
\frac{\mathrm{d}A_{1\lambda e}(t)}{\mathrm{d}t} =& -\mathrm{i}g_{\lambda e}^{31}\mathrm{e}^{\mathrm{i}\mu_\lambda^{31}t}A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}, \\
\frac{\mathrm{d}A_{0ke}(t)}{\mathrm{d}t} =& -\mathrm{i}g_{ke}^{30}\mathrm{e}^{\mathrm{i}\mu_k^{30}t}A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}, \\
\frac{\mathrm{d}A_{2ke}(t)}{\mathrm{d}t} =& -\mathrm{i}g_{ke}^{32}\mathrm{e}^{\mathrm{i}\mu_k^{32}t}A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}.
\end{aligned} \tag{6}
$$

Formally integrating and eliminating $A_{1\lambda e}(t)$, $A_{0ke}(t)$ and $A_{2ke}(t)$ from Eq.(6), we obtain

$$
\begin{aligned}
\frac{\mathrm{d}(A_3(t)\mathrm{e}^{-\mathrm{i}\mu_p^{31}t})}{\mathrm{d}t} =& -\mathrm{i}\Omega_{31}\mathrm{e}^{-\mathrm{i}\mu_p^{31}t}A_1(t) \\
&- \int_0^t G_{33}^{\lambda 1}(t-t')\mathrm{e}^{-\mathrm{i}\mu_p^{31}t'}A_3(t')\mathrm{d}t' \\
&- \int_0^t G_{33}^{k0}(t-t')\mathrm{e}^{-\mathrm{i}\mu_p^{31}t'}A_3(t')\mathrm{d}t' \\
&- \int_0^t G_{33}^{k2}(t-t')\mathrm{e}^{-\mathrm{i}\mu_p^{31}t'}A_3(t')\mathrm{d}t', \\
\frac{\mathrm{d}A_1(t)}{\mathrm{d}t} =& -\mathrm{i}\Omega_{31}A_3(t),
\end{aligned} \tag{7}
$$

where

$$
\begin{aligned}
G_{33}^{\lambda 1}(t-t') =& \sum_{\lambda e} g_{\lambda e}^{31}g_{\lambda e}^{31}\mathrm{e}^{-\mathrm{i}\mu_\lambda^{31}(t-t')}, \\
G_{33}^{kj}(t-t') =& \sum_{ke} g_{ke}^{3j}g_{ke}^{3j}\mathrm{e}^{-\mathrm{i}\mu_k^{3j}(t-t')} \quad (j=0,2)
\end{aligned} \tag{8}
$$

are the delayed Green functions which can be turned into the following forms:[11]

$$
\begin{aligned}
G_{33}^{kj}(t-t') =& \frac{1}{2}\sqrt{\beta_{3j}^{3/2}\beta_{3j}^{3/2}}\left\{\frac{\mathrm{e}^{\mathrm{i}[\delta_{3jc2}(t-t')-\pi/4]}}{\sqrt{\pi(t-t')}}\right. \\
&\left. + \frac{\mathrm{e}^{\mathrm{i}[\delta_{3jc1}(t-t')+\pi/4]}}{\sqrt{\pi(t-t')}}\right\}, \\
G_{33}^{kj}(t-t') =& \frac{1}{2}\sqrt{\alpha_{3j}\alpha_{3j}}\left\{\frac{\mathrm{e}^{\mathrm{i}[\delta_{3jc2}(t-t')+\pi/4]}}{\sqrt{4\pi(t-t')^3}}\right. \\
&\left. + \frac{\mathrm{e}^{\mathrm{i}[\delta_{3jc1}(t-t')-\pi/4]}}{\sqrt{4\pi(t-t')^3}}\right\}, \\
G_{33}^{kj}(t-t') =& \frac{1}{2}\sqrt{\gamma_{3j}\gamma_{3j}}\delta(t-t') \quad (j=0,2), \\
G_{33}^{\lambda 1}(t-t') =& \frac{1}{2}\sqrt{\gamma_{31}\gamma_{31}}\delta(t-t'),
\end{aligned} \tag{9}
$$

for isotropic PBG, anisotropic PBG and free space reservoirs, respectively. δ_{3jc1} (δ_{3jc2}) (j=0,2) is the difference between the upper level and the lower (upper) edge of the forbidden gap. The definitions of the parameters are

$$
\begin{aligned}
\beta_{3j}^{3/2} =& \frac{1}{2\pi\varepsilon_0}\frac{\omega_{3j}^2 d_{3j}^2\omega_{c2}^{3/2}}{3\hbar c^3}, \\
\alpha_{3j} =& \frac{1}{2\pi\varepsilon_0}\frac{\omega_{3j}^2 d_{3j}^2\omega_{c2}^{1/2}}{3\hbar c^3}, \\
\gamma_{3j} =& \frac{2}{\pi\varepsilon_0}\frac{\omega_{3j}^3 d_{3j}^2}{3\hbar c^3} \quad (j=0,2), \\
\gamma_{31} =& \frac{2}{\pi\varepsilon_0}\frac{\omega_{31}^3 d_{31}^2}{3\hbar c^3}.
\end{aligned} \tag{10}
$$

Since the transition $|3> - |1>$ is coupled to free vacuum space modes, the steady-state linear susceptibility for the transitions can be given by [17]

$$\chi(\delta_p) \sim \frac{N}{E} d_{31} A_1(t \to \infty) A_3^*(t \to \infty), \tag{11}$$

where N is the atomic density. Therefore, the steady-state probe absorption spectrum is determined by the steady-state solution of Eqs.(7). We assume that the initial population is in the ground state $|1>$ and the interaction between the probe field and atom is very weak ($\Omega_{31},\ \Omega_{21} \ll \alpha,\ \beta^{3/2},\ \gamma$) so that $A_1(t) \approx 1$. If $\frac{\mathrm{d}A(t)}{\mathrm{d}t} = f(t)$, from the following relations:

$$A(\infty) = \int_0^\infty f(t')\mathrm{d}t' = \widetilde{f}(s) \mid_{s\to 0},$$

$$s\widetilde{A}(s) \mid_{s\to 0} - A(0) = \widetilde{f}(s) \mid_{s\to 0}, \tag{12}$$

we obtain

$$A(\infty) = s\widetilde{A}(s) \mid_{s\to 0} - A(0). \tag{13}$$

Therefore, by taking the Laplace transform of Eqs.(7), the steady state solution for $A_3(t)$ can be obtained:

$$\begin{aligned}(s - \mathrm{i}\mu_p^{31})\widetilde{A}_3(s) = &- \mathrm{i}\Omega_{31}\frac{A_1(0)}{s} - [\widetilde{G}_{33}^{\lambda 1}(s - \mathrm{i}\mu_p^{31}) \\ &+ \widetilde{G}_{33}^{k2}(s - \mathrm{i}\mu_p^{31}) \\ &+ \widetilde{G}_{33}^{k0}(s - \mathrm{i}\mu_p^{31})]\widetilde{A}_3(s)\end{aligned} \tag{14}$$

with the approximation of $A_1(t) \approx 1.0$. From Eqs.(13) and (14), we have

$$\widetilde{A}_3(\infty) = s\widetilde{A}_3(s) \mid_{s\to 0} = \frac{-\mathrm{i}\Omega_{31}A_1(0)}{-\mathrm{i}\mu_p^{31} + \widetilde{G}_{33}^{\lambda 1}(s \to -\mathrm{i}\mu_p^{31}) + \widetilde{G}_{33}^{k0}(s \to -\mathrm{i}\mu_p^{31}) + \widetilde{G}_{33}^{k2}(s \to -\mathrm{i}\mu_p^{31})}. \tag{15}$$

Here, $\widetilde{G}_{33}^{k(\lambda)j}(s)$ (j=0, 1, 2) is the Laplace transform of the delayed Green function in Eq.(9), which can be derived as follows:

$$\begin{aligned}\widetilde{G}_{33}^{kj}(s) =& \frac{1}{2}\sqrt{\beta_{3j}^{3/2}\beta_{3j}^{3/2}}\left[\frac{\mathrm{i}}{\sqrt{\mathrm{i}s + \delta_{3j\mathrm{c}1}}} + \frac{1.0}{\sqrt{\mathrm{i}s + \delta_{3j\mathrm{c}2}}}\right], \\ \widetilde{G}_{33}^{kj}(s) =& \frac{1}{2}\sqrt{\alpha_{3j}\alpha_{3j}}\left[-\mathrm{i}\sqrt{\mathrm{i}s + \delta_{3j\mathrm{c}1}} + \sqrt{\mathrm{i}s + \delta_{3j\mathrm{c}2}}\right], \\ \widetilde{G}_{33}^{kj}(s) =& \frac{1}{2}\sqrt{\gamma_{3j}\gamma_{3j}}, \quad (j = 0, 2) \\ \widetilde{G}_{3l}^{\lambda 1}(s) =& \frac{1}{2}\sqrt{\gamma_{31}\gamma_{31}},\end{aligned} \tag{16}$$

for isotropic PBG, anisotropic PBG and free space reservoirs, respectively. Here the definitions of the detuning parameters are

$$\begin{aligned}\delta_{30\mathrm{c}1} =& \delta_{3\mathrm{c}1} - \frac{1}{2}\Delta, \\ \delta_{32\mathrm{c}1} =& \delta_{3\mathrm{c}1} + \frac{1}{2}\Delta, \\ \delta_{30\mathrm{c}2} =& \delta_{30\mathrm{c}1} - \delta_{\mathrm{c}2\mathrm{c}1}, \\ \delta_{32\mathrm{c}2} =& \delta_{32\mathrm{c}1} - \delta_{\mathrm{c}2\mathrm{c}1},\end{aligned} \tag{17}$$

where δ_{3jcn} (j=0, 2, n=1, 2) is the difference between the upper level and the lower edge ($n = 1$) or the upper edge ($n = 2$) of the forbidden gap. The calculation of a probe absorption spectrum can be performed based on Eqs.(11), (15) and (16).

3. Results and discussion

For realistic simulation of the ground state of an atom, it has a superfine structure with splitting width Δ. In order to investigate the effects of the splitting on the probe absorption spectrum in the PBG and free vacuum reservoirs, the probe and dispersion spectra as a function of detuning δ_p are shown in Figs.2(a)–(c). From Figs.2(b) and 2(c), we see that the probe absorption peak is corresponding to the resonant transition from the upper level $|3\rangle$ to the lower level $|1\rangle$. This is expected and understandable as the splitting changes the other transitions coupled to the anisotropic PBG and free vacuum reservoirs. However, in the case of isotropic PBG reservoir (Fig.2(a)), the effect of the splitting width Δ on the probe absorption spectrum is quite different. Some additional and unexpected peaks appear besides the peaks corresponding to the resonant transitions from the upper level to the lower level. In order to investigate this new feature, we plot the probe absorption spectra for different splitting widths Δ in the case of isotropic PBG modes shown in Fig.3. From Fig.3, we see that the splitting of the lower levels has a complicated effect on probe spectrum in the case of isotropic PBG modes. The probe absorption spectra for different widths of forbid-

den gap are shown in Fig.4. From Fig.2(a), Figs.3 and 4. The detuning for the additional peaks in the case of isotropic PBG modes unconcerned with other parameters are covered in two ranges: from $\frac{1}{2}(\Delta - \delta_{c2c1})$ to $\frac{1}{2}(\Delta+\delta_{c2c1})$ and from $-\frac{1}{2}(\Delta-\delta_{c2c1})$ to $-\frac{1}{2}(\Delta+\delta_{c2c1})$. As a result, the four (two) transparencies with two (one) probe absorption peaks appear in the case of isotropic PBG modes. The origin of this feature can be traced back to Eqs.(15), and (16). The forms of Laplace transform of the delayed Green function involving in the probe absorption spectrum include the following factors:

$$\left(\frac{\mathrm{i}}{\sqrt{\delta_k + \frac{1}{2}(\delta_{c2c1} - \Delta)}} + \frac{1.0}{\sqrt{\delta_k - \frac{1}{2}(\delta_{c2c1} + \Delta)}}\right)$$

and

$$\left(\frac{\mathrm{i}}{\sqrt{\delta_k - \frac{1}{2}(\delta_{c2c1} - \Delta)}} + \frac{1.0}{\sqrt{\delta_k + \frac{1}{2}(\delta_{c2c1} + \Delta)}}\right)$$

They show clearly that the splitting of the lower levels will induce additional singularities of Laplace transform of the delayed Green function for isotropic PBG modes. These additional singularities occur at $\frac{1}{2}(\Delta - \delta_{c2c1})$, $\frac{1}{2}(\Delta + \delta_{c2c1})$, $-\frac{1}{2}(\Delta - \delta_{c2c1})$ and $-\frac{1}{2}(\Delta+\delta_{c2c1})$. This shows conclusively that the peaks resulting from the splitting of the atomic ground state stem from the contribution of the additional singularities of the Laplace transform of the delayed Green function of the isotropic PBG modes.

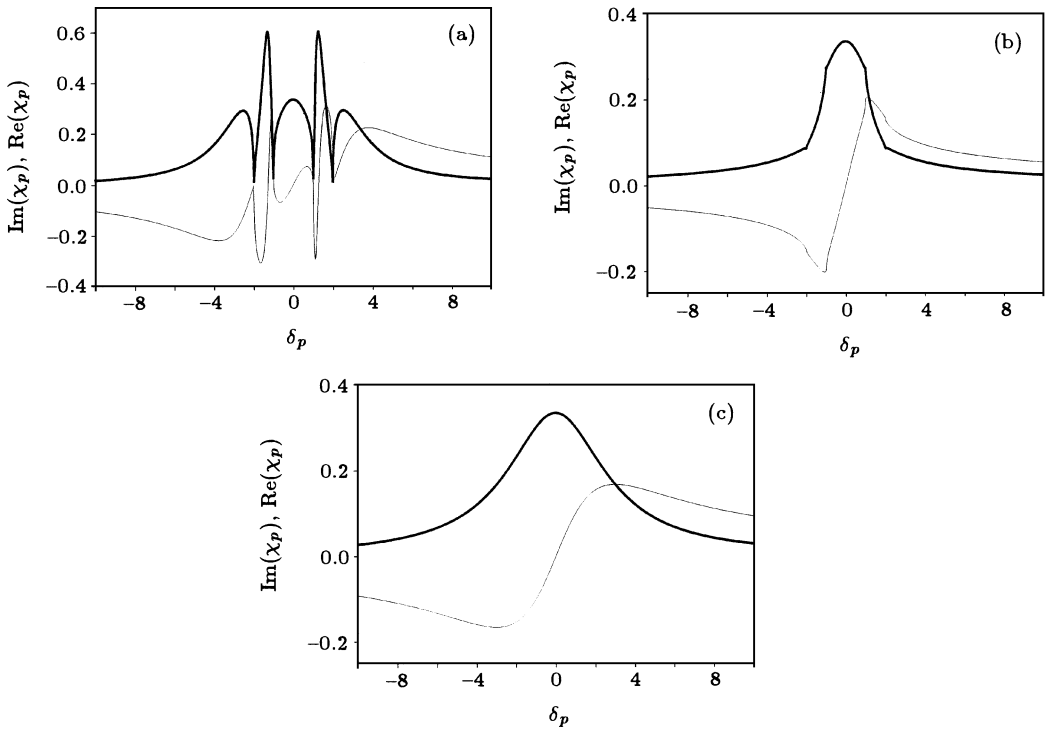


Fig.2. The absorption $\mathrm{Im}[\chi(\delta_p)]$ (thick solid lines) and dispersion $\mathrm{Re}[\chi(\delta_p)]$ (thin solid lines) spectra (in arbitrary units) as a function of detuning δ_p of probe field, (a) for double-band isotropic PBG reservoir and $\beta_{30} = \beta_{32} = 1.0$, (b) for double-band anisotropic PBG reservoir and $\alpha_{30}^2 = \alpha_{32}^2 = 1.0$, (c) for free vacuum reservoir and $\gamma_{30} = \gamma_{32} = 1.0$. The other parameters used are $\delta_{c2c1} = 1.0$, $\delta_{3c1} = 0.5$, $\Delta = 3.0$, $A_1(0) = 1.0$. All parameters in this paper are in units of γ_{31}.

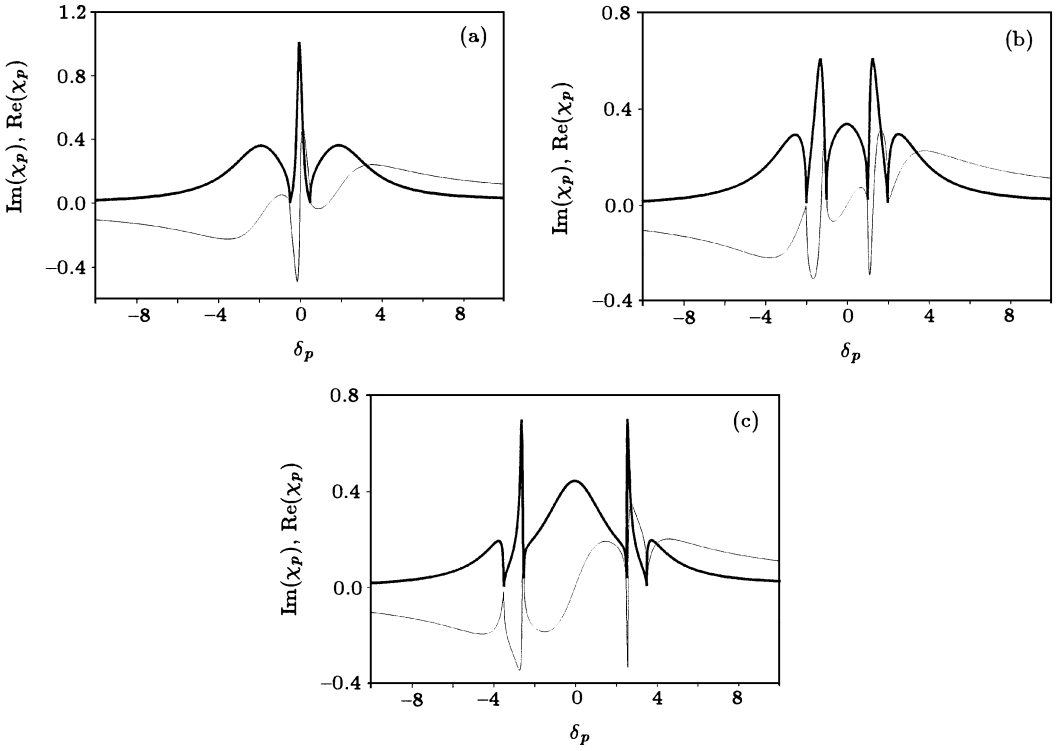


Fig.3. The absorption $\mathrm{Im}[\chi(\delta_p)]$ (thick solid lines) and dispersion $\mathrm{Re}[\chi(\delta_p)]$ (thin solid lines) spectra (in arbitrary units) as functions of detuning δ_p of probe field for different splitting widths of the lower state levels in the case of double-band isotropic PBG reservoir. (a) $\Delta = 0.0$, (b) $\Delta = 3.0$, (c) $\Delta = 6.0$. Other parameters used are the same as those in Fig.2.

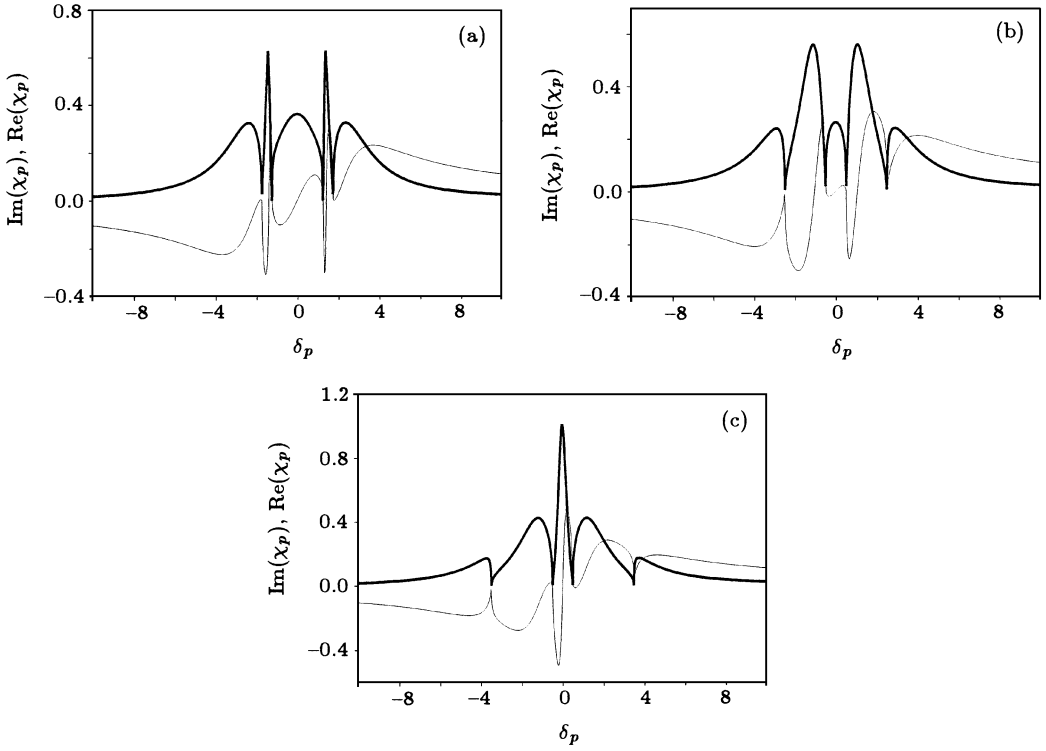


Fig.4. The absorption $\mathrm{Im}[\chi(\delta_p)]$ (thick solid lines) and dispersion $\mathrm{Re}[\chi(\delta_p)]$ (thin solid lines) spectra (in arbitrary units) as functions of detuning δ_p of probe field for different forbidden gap widths in the case of isotropic PBG reservoir. $\delta_{3c1} = \frac{1}{2}\delta_{c2c1}$, (a) $\delta_{c2c1} = 0.5$, (b) $\delta_{c2c1} = 3.0$, (c)$\delta_{c2c1} = 6.0$. Other parameters used are the those as that in Fig.2.

4. Conclusion

The probe absorption and dispersion spectra of a four-level atom embedded in a double-band photonic crystal have been discussed. The model used in this paper assumes atomic transitions from the one upper level to the three lower levels. One transition is assumed to interact with free vacuum modes, and the other two are assumed to interact respectively with isotropic PBG modes, anisotropic PBG modes and free vacuum modes. The effect of the splitting of the two lower levels on the probe absorption spectra of an atom are investigated in detail in the three cases. In the case of free vacuum modes, the splitting width has no effect on the probe and dispersion spectra. If the splitting width is larger in the case of anisotropic PBG modes, the width of the probe spectrum will be expanded. The two (four) transparencies with one (two) probe absorption peak(s) appear in the case of isotropic PBG modes. These new features result from the contribution of the additional singularities of Laplace transform of the delayed Green function of the isotropic PBG modes.

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