Subluminal to superluminal light propagation in a V-type three-level atomic system interacting with squeezed vacuum*

Zhang Qin(张 琴)[a)†], Qu Guang-Hui(屈光辉)[a)], Tang Yuan-He(唐远河)[a)], Jin Kang(金 康)[b)], and Ji Wei-Li(纪卫莉)[a)]

[a)] *Department of Physics , Xi'an University of Technology, Xi'an* 710054, *China*
[b)] *Department of Physics, Northwest University, Xi'an* 710069, *China*

(Received 14 July 2011; revised manuscript received 29 August 2011)

We study the light propagation properties in a V-type three-level atomic system interacting with two independent broadband squeezed baths. Using the graphic method, the dependences of the absorption and the dispersion of the system on the coupling field strength, the incoherent pumping field strength and the intensity of the squeezed vacuum are analysed. The result shows that the probe pulse propagation can be changed from subluminal to superluminal speed.

Keywords: squeezed vacuum, group velocity, superluminal propagation, subluminal propagation

PACS: 32.80.Qk, 37.10.Vz, 42.50.–p, 42.50.Nn **DOI:** 10.1088/1674-1056/21/2/023201

1. Introduction

In recent years, there has been great interest in the study of subluminal and superluminal light propagation in dispersive media.[1−6] This is because of the potential applications of subluminal and superluminal light propagation in nonlinear optics, high-speed optical switches and communication systems.[7] It is well known that the group velocity of a light pulse can be slowed down, or it can be greater than c (the speed of light in vacuum) or even become negative in a dispersive medium. Various schemes have been proposed to switch from subluminal to superluminal light propagation in an atomic medium.[8−14] Hau *et al.*[2] could reduce the group velocity to 17 m/s in a Bose–Einstein condensate. Later, the group velocity can be slowed down to 8 m/s,[15] and in a solid to 45 m/s.[16] In another experiment, Wang *et al.*[3] demonstrated superluminal light propagation in the region of lossless anomalous dispersion between two closely spaced gain lines. Akulshin *et al.*[17] and Kim *et al.*[7] reported on the experiments where negative group velocities of $v_{\rm g} = -c/14400$ in Cs vapour cells were obtained. Mikhailov *et al.*[18] have obtained the negative group velocities as low as −80 m/s in Rb vapour.

In view of the many potential applications of subluminal and superluminal light propagation, a natural question is whether one can have a controlling parameter for switching the group velocity from subluminal to superluminal. Generally, the control of the dispersion from normal to anomalous has been obtained by (i) applying a field connecting the two closed sublevels and (ii) considering spontaneously generated coherence (SGC). More recently, the effects of the incoherent pumping field and SGC on controlling the group velocity of the probe field were discussed.[19−21] The rate of an incoherent pumping field, as an additional parameter, dramatically changes the slope of the dispersion from normal to anomalous. The above proposals of applying an incoherent pumping field are associated only with the three-level Λ systems and the four-level systems in the normal heat baths. However, the squeezed vacuum plays an important role in the atomic dynamics since the seminal studies by Gardiner,[22] and Ficek *et al.*[23] have shown that a two-level atom damped by a broadband squeezed vacuum can exhibit a strong emission peak at the central frequency of the atomic levels. Akram *et al.*[24] have studied the index of refraction of a two-level atom by replacing the usually applied coherent field by a squeezed vacuum field.

*Project supported by the National Natural Science Foundation of China (Grant No. 10874138).
†Corresponding author. E-mail: zhangqin@xaut.edu.cn
© 2012 Chinese Physical Society and IOP Publishing Ltd

In this paper, considering the effect of the indirect incoherent pumping field of a V-type atomic system damped by two independent broadband squeezed baths, the absorption and the dispersion properties of the medium can be controlled with the rate of the incoherent pumping field and the intensity of the squeezed vacuum. Moreover, different from other proposals, the switching from normal to anomalous dispersion can be realized just by changing the intensity of the coupling field without using an extra lower-level coupling field.

In the following, we will discuss the relevant details of the atomic system. The rest of the paper is organized as follows. In Section 2, the model is described and the density-matrix equations of motion for the system are derived. In Section 3, we discuss the absorption and the dispersion properties of the atomic medium. A brief conclusion of this paper is presented in Section 4.

2. Theoretical mode and density-matrix equations

We consider a V-type three-level atom with levels $|1\rangle$, $|2\rangle$, and $|3\rangle$, having energies E_1, E_2, E_3. Transition $|3\rangle \rightarrow |2\rangle$ is driven by an external coupling field $E_{\rm c}$ and similarly, transition $|3\rangle \rightarrow |1\rangle$ is driven only by the probe field $E_{\rm p}$. Moreover, the upper levels $|1\rangle$ and $|2\rangle$ are coupled to lower level $|3\rangle$ by two independent broadband squeezed vacuum modes and there is no correlation between the two different modes. At the same time, the transition $|3\rangle \rightarrow |1\rangle$ is incoherently pumped with a rate Λ by an incoherent pumping field. The atomic transition frequencies between the ground level $|3\rangle$ and upper levels $|2\rangle$ and $|1\rangle$ are ω_{23} and ω_{13}, respectively.

The Hamiltonian of the system in the rotating wave approximation can be written as

$$\begin{aligned} H =\ & \hbar\sum_{m=1}^{3}\omega_m|m\rangle\langle m| + \hbar\sum_{l=1}^{2}\omega_l b_l^{\dagger}b_l \\ & - \hbar\alpha_{\rm p}\,{\rm e}^{-{\rm i}\,\omega_{\rm p}t-\phi_1}|1\rangle\langle 3| - \hbar\alpha_{\rm c}\,{\rm e}^{-{\rm i}\,\omega_{\rm c}t-\phi_2}|2\rangle\langle 3| \\ & - \hbar\sum_{m=1}^{2} g_m|m\rangle\langle 3|b_m - {\rm H.C.}, \end{aligned} \tag{1}$$

where $\hbar\omega_m$ is the energy of the atomic level; $b_l^{\dagger}$ and b_l are the creation operator and the annihilation operator of the squeezed vacuum field with angular frequency ω_l $(l = 1, 2)$; parameters $\alpha_{\rm p} = \boldsymbol{\mu}_{13}\cdot\boldsymbol{E}_{\rm p}/2\hbar$ and $\alpha_{\rm c} = \boldsymbol{\mu}_{23}\cdot\boldsymbol{E}_{\rm c}/2\hbar$ are the Rabi frequencies of the transitions $|3\rangle \rightarrow |1\rangle$ and $|3\rangle \rightarrow |2\rangle$, which are determined by the dipolar moment μ_{3m} $(m = 1, 2)$ and the electromagnetic field amplitudes E_i $(i = {\rm p}, {\rm c})$; φ_1 and φ_2 are the phases of the probe and coupling field; $\omega_{\rm p}$ and $\omega_{\rm c}$ are the angular frequencies of the two driving external field, respectively. Now we assume that the two squeezed vacuum modes are independent and they are δ correlated. The correlation function for the field operators b_l and $b_l^{\dagger}$ can be written as[22,25]

$$\begin{aligned} &\langle b_l^{\dagger}(t)b_m(t')\rangle = N(\omega_l)\delta(t-t')\delta_{lm}, \\ &\langle b_l(t)b_m^{\dagger}(t')\rangle = [N(\omega_l)+1]\delta(t-t')\delta_{lm}, \\ &\langle b_l(t)b_m(t')\rangle = M(\omega_l)\delta(t-t')\delta_{2\omega_{sl}-\omega_l-\omega_m}\delta_{lm}, \\ &\langle b_l^{\dagger}(t)b_m^{\dagger}(t')\rangle = M^*(\omega_l)\delta(t-t')\delta_{2\omega_{sl}-\omega_l-\omega_m}\delta_{lm}, \end{aligned} \tag{2}$$

where ω_{sl} $(l = 1, 2)$ are the carrier frequencies of the squeezed vacuum field; $N(\omega_l)$ and $M(\omega_l)$ are the frequency parameters that characterize the squeezing. In general, the squeezed field satisfies the inequality

$$|M(\omega_l)|^2 \leq N(\omega_l)(N(\omega_l)+1), \tag{3}$$

where M is the complex magnitude so that $M(\omega_l)$ can be written as $M(\omega_l) = |M(\omega_l)|\,{\rm e}^{{\rm i}\,\Phi_l}$ $(l = 1, 2)$, where $\Phi_l (l = 1, 2)$ are the phase of the l-th squeezed field. When $M(\omega_l) = 0$ the above equation describes the condition of the thermal reservoir, while $N(\omega_l)$ is the mean photon number of the squeezed vacuum field. In this work, we consider the bath is in the case of the maximum degree of squeezing. By using the density-matrix formalism, the master equation based on the Born–Markoff approximation is given by

$$\frac{\partial\rho}{\partial t} = -\,{\rm i}[H_0,\rho] + L_{\rm therm} + L_{\rm pump} - L_{\rm sqz}, \tag{4}$$

where

$$\begin{aligned} H_0 =\ & -\alpha_{\rm c}(\sigma_{23}+\sigma_{32}) - \alpha_{\rm p}(\sigma_{13}+\sigma_{31}) \\ & + \Delta_{\rm p}\sigma_{11} + \Delta_{\rm c}\sigma_{22}, \end{aligned} \tag{5}$$

$$\begin{aligned} L_{\rm therm} =\ & \gamma_1(N_1+1)(2\sigma_{31}\rho\sigma_{13} - \sigma_{11}\rho - \rho\sigma_{11}) \\ & + \gamma_1 N_1(2\sigma_{13}\rho\sigma_{31} - \sigma_{33}\rho - \rho\sigma_{33}) \\ & + \gamma_2(N_2+1)(2\sigma_{32}\rho\sigma_{23} - \sigma_{22}\rho - \rho\sigma_{22}) \\ & + \gamma_2 N_2(2\sigma_{23}\rho\sigma_{32} - \sigma_{33}\rho - \rho\sigma_{33}), \end{aligned} \tag{6}$$

$$\begin{aligned} L_{\rm pump} =\ & \Lambda(2\sigma_{31}\rho\sigma_{13} - \sigma_{11}\rho - \rho\sigma_{11}) \\ & + \Lambda(2\sigma_{13}\rho\sigma_{31} - \sigma_{33}\rho - \rho\sigma_{33}), \end{aligned}$$

$$\begin{aligned} L_{\rm sqz} =\ & \gamma_1|M_1|(2\sigma_{13}\rho\sigma_{13} - \sigma_{13}^2\rho - \rho\sigma_{13}^2)\,{\rm e}^{-{\rm i}\,\phi_1} \\ & + \gamma_1|M_1|(2\sigma_{31}\rho\sigma_{31} - \sigma_{33}^2\rho - \rho\sigma_{31}^2)\,{\rm e}^{{\rm i}\,\phi_1} \\ & + \gamma_2|M_2|(2\sigma_{23}\rho\sigma_{23} - \sigma_{22}^2\rho - \rho\sigma_{22}^2)\,{\rm e}^{-{\rm i}\,\phi_2} \\ & + \gamma_2|M_2|(2\sigma_{32}\rho\sigma_{32} - \sigma_{33}^2\rho - \rho\sigma_{32}^2)\,{\rm e}^{{\rm i}\,\phi_2}, \end{aligned} \tag{7}$$

with $\sigma_{ij} = |i\rangle\langle j|$ $(i, j = 1, 2, 3)$ being the atomic operators; γ_1 and γ_2 the decay rates for $|1\rangle \rightarrow |3\rangle$ and $|2\rangle \rightarrow |3\rangle$ transitions, respectively; $\Delta_{\rm p} = \omega_{13} - \omega_{\rm p}$ and

$\Delta_{\rm c} = \omega_{23} - \omega_{\rm c}$ the laser detunings ($\omega_{13} = (E_1 - E_3)/\hbar$ and $\omega_{23} = (E_2 - E_3)/\hbar$); phases $\phi_1 = 2\phi_{\rm p} - \Phi_1$ and $\phi_2 = 2\phi_{\rm c} - \Phi_2$ are the relative phases between the squeezed field and the probe field with phase $\phi_{\rm p}$, and between the squeezed field and the coupling field with phase $\phi_{\rm c}$ respectively. The evolution equation of the density-matrix elements for the atom in the rotating frame can be written as

$$\begin{aligned}
\frac{\partial \rho_{11}}{\partial t} &= {\rm i}\,\alpha_{\rm p}(\rho_{31} - \rho_{13}) - 2\gamma_1(N_1+1)\rho_{11} \\
&\quad + 2\gamma_1 N_1 \rho_{33} - 2\Lambda\rho_{11} + 2\Lambda\rho_{33}, \\
\frac{\partial \rho_{22}}{\partial t} &= {\rm i}\,\alpha_{\rm c}(\rho_{32} - \rho_{23}) - 2\gamma_2(N_2+1)\rho_{22} + 2\gamma_2 N_2\rho_{33}, \\
\frac{\partial \rho_{12}}{\partial t} &= {\rm i}\,\alpha_{\rm p}\rho_{32} - {\rm i}\,\alpha_{\rm c}\rho_{13} + {\rm i}(\Delta_{\rm c} - \Delta_{\rm p})\rho_{12} \\
&\quad - [\gamma_1(N_1+1) + \gamma_2(N_2+1)]\rho_{12} - \Lambda\rho_{12}, \\
\frac{\partial \rho_{13}}{\partial t} &= -{\rm i}\,\alpha_{\rm c}\rho_{12} + {\rm i}\,\alpha_{\rm p}(\rho_{33} - \rho_{11}) \\
&\quad - {\rm i}\,\Delta_{\rm p}\rho_{13} - 2\Lambda\rho_{13} - \gamma_1(2N_1+1)\rho_{13} \\
&\quad - \gamma_2 N_2\rho_{13} - 2\gamma_1|M_1|\rho_{31}\,{\rm e}^{-{\rm i}\,\phi_1}, \\
\frac{\partial \rho_{23}}{\partial t} &= -{\rm i}\,\alpha_{\rm c}(\rho_{33} - \rho_{22}) - {\rm i}\,\alpha_{\rm p}\rho_{21} \\
&\quad - {\rm i}\,\Delta_{\rm c}\rho_{23} - \Lambda\rho_{23} - \gamma_2(2N_2+1)\rho_{23} \\
&\quad - \gamma_1 N_1\rho_{23} - 2\gamma_2|M_2|\rho_{32}\,{\rm e}^{-{\rm i}\,\phi_2}, \\
\rho_{11} &+ \rho_{22} + \rho_{33} = 1. \qquad (8)
\end{aligned}$$

Using the fact that the probe field is weak compared with the coupling field, it is sufficient to treat the probe field as a perturbation and all of the atoms are initially in the ground state $|3\rangle$, i.e., $\rho_{11}^0 = 0$, $\rho_{22}^0 = 0$, and $\rho_{33}^0 = 1$. We are interested in the steady state behaviour of the probe transition $|1\rangle \to |3\rangle$, and the first order steady-state solution for ρ_{13} is given by

$$\rho_{13}^{(1)} = {\rm i}\,\frac{2\gamma_1 M_1^* + F^*}{|F|^2 - 4|M_1|^2\gamma_1^2}\alpha_{\rm p}, \qquad (9)$$

where

$$\begin{aligned}
F &= {\rm i}\,\Delta_{\rm p} + \gamma_1(2N_1+1) + \gamma_2 N_2 + 2\Lambda \\
&\quad + \frac{\alpha_{\rm c}^2}{\gamma_1(N_1+1) + \gamma_2(N_2+1) + {\rm i}(\Delta_{\rm p} - \Delta_{\rm c}) + \Lambda}, \qquad (10)
\end{aligned}$$

and the linear susceptibility $\chi^{(1)}$ is given by

$$\chi^{(1)} = {\rm i}\,\frac{2\gamma_1 M_1^* + F^*}{|F|^2 - 4|M_1|^2\gamma_1^2}. \qquad (11)$$

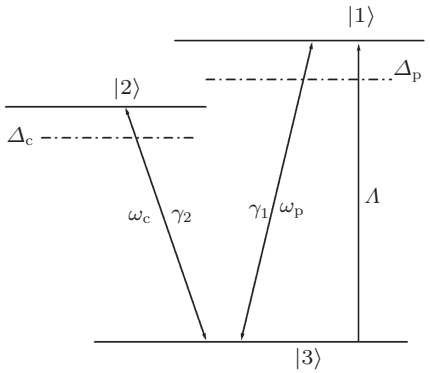


Fig. 1. Schematic diagram of a V-type atomic system driven by an incoherent pumping and a week probe field in a squeezed vacuum.

For simplicity, we assume that $\gamma_1 = \gamma_2 = \gamma$, $N_1 = N_2 = N$, and $M_1 = M_2 = M$. Separating the real part from the imaginary part of $\chi^{(1)} = \chi' + {\rm i}\,\chi''$, we obtain

$$\begin{aligned}
\chi' &= {\rm Re}\,\chi^{(1)} = \frac{\Delta_{\rm p} + 2|M|\sin\phi_1\gamma - \dfrac{(\Delta_{\rm p} - \Delta_{\rm c})\alpha_{\rm c}^2}{[2\gamma(N+1) + \Lambda]^2 + (\Delta_{\rm p} - \Delta_{\rm c})^2}}{|F|^2 - 4|M|^2\gamma^2}, \\
\chi'' &= {\rm Im}\,\chi^{(1)} = \frac{\gamma(3N+1) + 2|M|\cos\phi_1\gamma + \dfrac{[2\gamma(N+1) + \Lambda]\alpha_{\rm c}^2}{[2\gamma(N+1) + \Lambda]^2 + (\Delta_{\rm p} - \Delta_{\rm c})^2}}{|F|^2 - 4|M|^2\gamma^2}. \qquad (12)
\end{aligned}$$

The group index of the probe field is given by

$$\begin{aligned}
n_{\rm g} &= \frac{c}{v_{\rm g}} = 1 + 2\pi\chi'(\omega_{\rm p}) \\
&\quad + 2\pi\omega_{\rm p}\frac{\partial\chi'(\omega_{\rm p})}{\partial\omega_{\rm p}}, \qquad (13)
\end{aligned}$$

where $v_{\rm g}$ is the group velocity, and $\chi'(\omega_{\rm p})$ is the real part of the line susceptibility.

3. Results and discussion

For simplicity, we assume that all figures are plotted in the units of γ. We choose the modest values for frequency of the probe field $\omega_{\rm p} = 10^5\gamma$, where c is the speed of light in a vacuum and the refractive index is $n_{\rm r} = 1 + 2\pi\chi'(\omega_{\rm p})$. For normal dispersion $\partial\chi'(\omega_{\rm p})/\partial\omega_{\rm p} > 0$, the group velocity is smaller than

the phase velocity $v_g < c/n_r$, while for anomalous dispersion $\partial\chi'(\omega_p)/\partial\omega_p < 0$, the group velocity can be larger than the phase velocity $v_g > c/n_r$, or even becomes negative $v_g < 0$.[3,7,21,26] In the following, we present our numerical results.

In Fig. 2, we display the properties of absorption (a) and dispersion (b) of the probe field versus the normalized probe detuning Δ_p/γ for different values of the coupling field α_c. Figure 2(a) shows that for $\alpha_c = 0.5$, the probe absorption presents a sharp peak around $\Delta_p/\gamma = 0$, the phenomenon is called EIA.[22−24,27] As the value of the coupling field increases (i.e., $\alpha_c = 5.0$), the probe absorption exhibits two absorption peaks and a dip around $\Delta_p/\gamma = 0$, thia phenomenon is called EIT. The phenomenon can be considered as arising from the interference of the transitions between dressed states $|+\rangle = (|3\rangle + |2\rangle)/\sqrt{2}$ and $|-\rangle = (|3\rangle - |2\rangle)/\sqrt{2}$ to a third level $|1\rangle$. It is clear that the EIT dip becomes reduced by further increasing the strength of the coupling field α_c up to $\alpha_c = 8.0$. The effect of the coupling field on the dispersion is shown in Fig. 2(b). It is clearly seen that when $\alpha_c = 0.5$, the slope of dispersion is positive and superluminal light propagation occurs. With the strength of the coupling field increasing (i.e., $\alpha_c = 5.0$), the slope of dispersion becomes negative, corresponding to the subluminal light propagation. Note that the slope of the dispersion for $\alpha_c = 8.0$ becomes steep compared with the case $\alpha_c = 5.0$. Thus, the coupling field plays a major role in switching the light propagation from subluminal to superluminal speed.

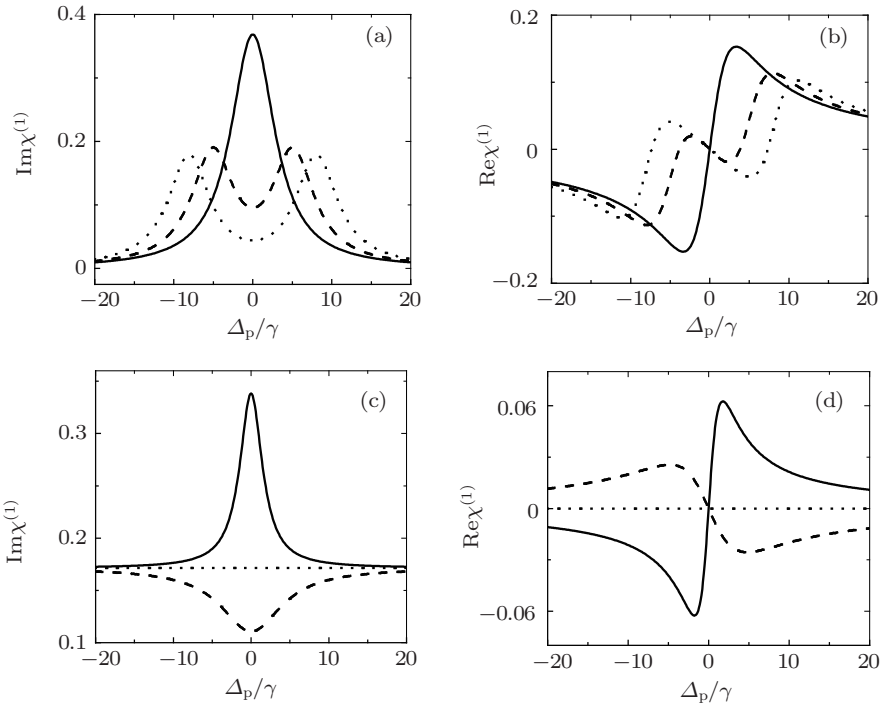


Fig. 2. Imaginary part (a) and real part (b) of $\chi^{(1)}$ each as a function of the probe field detuning for $\alpha_c = 0.5$ (solid line), 5.0 (dashed line), and 8.0 (dotted line); imaginary part (c) and real part (d) of $\chi^{(1)}$ of susceptibility versus probe detuning for $\alpha_c^2 = 0.1[\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ (solid line), $\alpha_c^2 = 2[\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ (dashed line), $\alpha_c^2 = [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ (dotted line). Other parameters are chosen as $\gamma = 1.0$, $N = 0.1$, $M = \sqrt{N(N+1)}$, $\Lambda = 1.0$, $\varphi = 0$, and $\Delta_c = 0$.

We come across an interesting situation by examining Eq. (11) for the $\mathrm{Re}\,\chi^{(1)}$. If we resort to the specific choice of the phase φ_1 (i.e., $\varphi_1 = 0, \pi$) with $\Delta_c = 0$, we have a simplified expression from Eq. (11)

$$\mathrm{Re}\,\chi^{(1)} = \frac{\Delta_p\left[1 - \dfrac{\alpha_c^2}{[2\gamma(N+1)+\Lambda]^2 + \Delta_p^2}\right]}{|F|^2 - 4|M|^2\gamma^2}. \quad (14)$$

It is evident that $\mathrm{Re}\,\chi^{(1)} = 0$, when we adjust the coupling field α_c according to the relation $\alpha_c^2 = [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$, where the probe detuning Δ_p is varied. This interesting result shows that we may obtain a flat refraction index profile for all the frequencies. We call this condition of $\alpha_c^2 = [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ a "critically condition"

for the coupling field strength. In Fig. 2, we once again display the values of absorption (Fig. 2(c)) and dispersion (Fig. 2(d)) versus probe field detuning for various intensities of the coupling field. It is seen that the normal absorption can be obtained under the condition of $\alpha_c^2 < [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ (i.e., $\alpha_c^2 = 0.1[\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$, denoted by the solid line in Fig. 2(c)), there is an absorption maximum at $\Delta_p/\gamma = 0$. Correspondingly, the dispersion is 0 at $\Delta_p/\gamma = 0$. On the contrary, there is an inverted and reduced absorption (dashed line in Fig. 2(c)) accompanied by opposite dispersion (dashed line in Fig. 2(d) under the condition of $\alpha_c^2 > [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$ (i.e., $\alpha_c^2 = 2[\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$)). Obviously, if we keep $\alpha_c^2 = [\Lambda + 2\gamma(N+1)]^2 + \Delta_p^2$, we will find that $\mathrm{Re}\,\chi^{(1)} = 0$ (this has been discussed above) always holds when the value of Δ_p is varied (dotted line of Fig. 2(d)). The absorption given by $\mathrm{Im}\,\chi^{(1)}$ is also flat at all the frequencies under the same conditions (dotted line in Fig. 2(c)). In addition, the slope of the index of refraction versus the probe detuning Δ_p/γ changes from positive to negative.

From Fig. 3, it is shown that the switching from subluminal to superluminal light propagation occurs for an appropriate value of the strength of the coupling field. It is obvious that the coupling field strength can be considered as a knob for changing the probe light propagation from subluminal to superluminal speed.

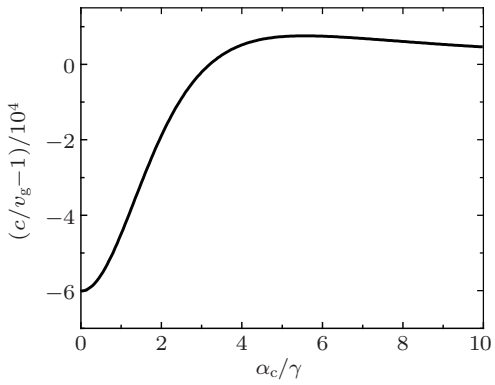


Fig. 3. Group index versus Rabi frequency of the coupling field for the parameters: $\gamma = 1.0$, $N = 0.1$, $M = \sqrt{N(N+1)}$, $\Lambda = 1.0$, $\varphi = 0$, $\Delta_c = 0$, and $\Delta_p = 0$.

In Fig. 4, we display the properties of absorption and dispersion of the probe field versus normalized probe detuning Δ_p/γ for different values of the incoherent pumping rate Λ. Figure 4(a) shows that for $\Lambda = 0.2$ there are two absorption peaks of the probe pulse. With the incoherent pumping rate Λ changing from $\Lambda = 0.2$ to $\Lambda = 1.0$, the probe absorption peaks and the dips all reduce. By further increasing the incoherent pumping rate Λ up to $\Lambda = 4.0$, the figures change from two absorption peaks near the Rabi frequencies and a dip near the centre Δ_p/γ to a single absorption peak near the centre. Thus, the EIA peak is established. Physically, as the incoherent pumping rate becomes larger, the incoherent pumping field brings more population from state $|3\rangle$ to state $|1\rangle$, which can result in an increase in the excited state population and an increase decay from the excited state to lower state. It should be noted that the appearance of EIA is due to the relatively large population in the excited states, which results in an increase of spontaneous transfer of the coherence from excited level to lower level. The effect of the incoherent pumping rate on the dispersion is shown in

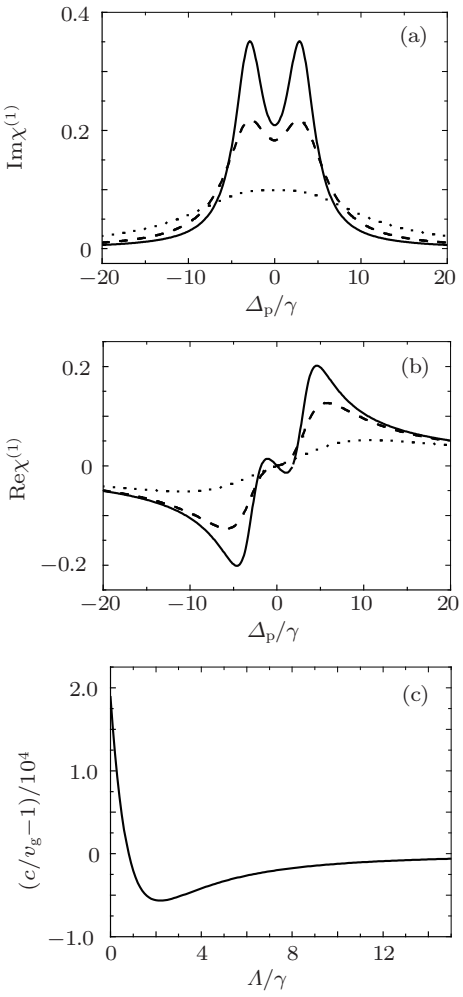


Fig. 4. Imaginary part (a) and real part (b) of $\chi^{(1)}$ each as a function of probe field detuning for $\gamma = 1.0$, $N = 0.1$, $M = \sqrt{N(N+1)}$, $\varphi = 0$, $\Delta_c = 0$, $\alpha_c = 3.0$, $\Lambda = 0.2$ (solid line), 1.0 (dashed line), and 4.0 (dotted line). Group index versus incoherent pump rate (c) for the parameters: $\gamma = 1.0$, $N = 0.1$, $M = \sqrt{N(N+1)}$, $\varphi = 0$, $\Delta_c = 0$, $\alpha_c = 3.0$, $\Delta_p = 0$.

Fig. 4(b). It is clearly seen that when $\Lambda = 0.2$, the slope of dispersion around zero probe field detuning is negative, corresponding to the subluminal light propagation. With the increase of the incoherent pumping rate (i.e., $\Lambda = 1.0$) , the slope of dispersion becomes positive. For such a case, superluminal light propagation occurs. Moreover, the slope of the dispersion with $\Lambda = 4.0$ has decreased compared with $\Lambda = 1.0$. It can be shown that the slope of the dispersion changes from negative to positive by increasing the intensity of the incoherent pump field. From Fig. 4(c) it follows that the switching from subluminal to superluminal light propagation occurs for an appropriate value of the incoherent pumping rate Λ. Thus, the incoherent pumping rate plays a major role in switching the light propagation from subluminal to superluminal speed.

In Fig. 5, we plot the values of group index $n_{\rm g}$ of the probe field versus the mean photon number of the squeezed vacuum N for different values of the Rabi frequency of the coupling field $\alpha_{\rm c}$. Here we assume that the incoherent pumping field is weak. We see that with the increase of N, the value of group index $n_{\rm g}$ changes from positive to negative. Thus the group velocity can be controlled by adjusting the photon number of the squeezed field and the propagation can change from subluminal to superluminal speed. Moreover, it is shown that there is an optimal intensity control for the largest negative value of $n_{\rm g}$. It is clear that as the Rabi frequency of the coupling field $\alpha_{\rm c}$ becomes larger, the largest negative value of $n_{\rm g}$ becomes smaller. There is an N value which makes the group index $n_{\rm g}$ reach the minimum [seen in Fig. 5 (dash–dot line)]. Correspondingly, the group velocity reaches the maximum value ($v_{\rm g} = -c/11004 = -2.73 \times 10^4$ m/s) under the appropriate parameters we choose. Meanwhile, in the positive region, the group velocity reaches the minimum value ($v_{\rm g} = c/27131 = 1.10 \times 10^4$ m/s) in the case of $N = 0$. Now, the system is damped by the standard vacuum (i.e., $N = M = 0$). Thus, the group velocity can be manipulated by the mean photon number of the squeezed vacuum. These results may be helpful for measuring the group velocity.

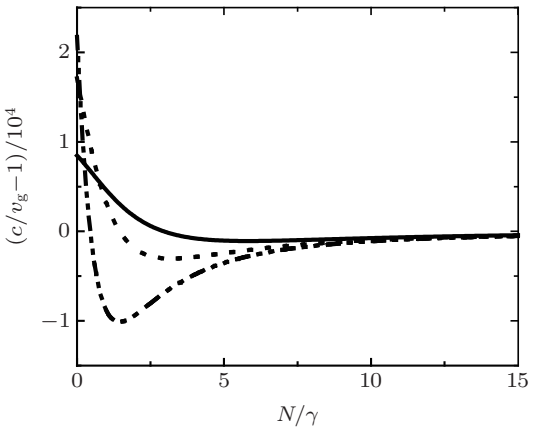


Fig. 5. Values of group index versus the average photon number N for $\gamma = 1.0$, $M = \sqrt{N(N+1)}$, $\varphi = 0$, $\Delta_{\rm c} = 0$, $\Delta_{\rm p} = 0$, $\Lambda = 0.1$, $\alpha_{\rm c} = 3.0$ (dash-dotted line), 5.0 (dotted line), and 8.0 (solid line).

4. Conclusion

We studied the properties of absorption and dispersion of a weak probe in a V-type three-level atomic configuration interacting with two independent broadband squeezed baths. It is shown that the slope of the dispersion changes from positive to negative by modulating the coupling field strength and the incoherent pumping field, which leads the probe pulse propagation to change from subluminal to superluminal speed in the medium. In addition, the effect of the intensity of the squeezed baths on the group index is discussed. By varying the squeezed bath parameters, one can control the light propagation from subluminal to superluminal speed. Therefore, one may use this system to realize the group velocity manipulation of a light pulse.

References

[1] Lukin M D and Imamoglu A 1990 *Nature* **413** 273
[2] Hau L V, Harris S E, Dutton Z and Behroozi C H 1999 *Nature* **397** 594
[3] Wang L J, Kuzmich A and Dogariu A 2000 *Nature* **406** 277
[4] Xu W H and Gao J Y 2005 *Chin. Phys.* **14** 2496
[5] Hu Z F, Deng J L and Wang Y Z 2008 *Chin. Phys.* B **17** 3358
[6] Agarwal G S and Dasgupta S 2004 *Phys. Rev.* A **70** 023802
[7] Kim K, Moom H S, Lee C, Kim S K and Kim J B 2003 *Phys. Rev.* A **68** 013810
[8] Kash M M, Sautenkov V A, Zirov A S, Hollberg L, Welch G R, Lukin M D, Rostovtsev Y, Fry E S and Scully M C 1999 *Phys. Rev. Lett.* **82** 229
[9] Agarwal G S, Dey T N and Menon S 2001 *Phys. Rev.* A **64** 0538809
[10] Hou B P, Wang S J, Yu W L and Sun W L 2006 *J. Phys.* B: *At. Mol. Opt. Phys.* **39** 2335
[11] Dutta S and Dasidar K R J 2007 *J. Phys.* B: *At. Mol. Opt. Phys.* **40** 4287
[12] Liu H Y, Meng Z M, Dai Q F, Wu L J, Lan S and Liu S H 2009 *Acta Phys. Sin.* **58** 4702 (in Chinese)
[13] Meng D D, Liu X D and Zhang S L 2011 *Acta Phys. Sin.* **60** 020305 (in Chinese)
[14] Lü C H, Tan L and Tan W T 2011 *Acta Phys. Sin.* **60** 024204 (in Chinese)

[15] Budker D, Kimball D F, Rochester S M and Yashchuk V V 1999 *Phys. Rev. Lett.* **83** 1767

[16] Turukhin A V, Sudarshanam V S, Shahriar M S, Musser J A, Ham B S and Hammer P R 2002 *Phys. Rev. Lett.* **88** 023602

[17] Akulshin A M, Barreiro S and Lezama A 1999 *Phys. Rev. Lett.* **83** 4277

[18] Mikhailov E E, Sautenkov V A, Novikova I and Welch G R 2004 *Phys. Rev.* A **69** 063808

[19] Mahmoudi M, Sahrai M and Tajalli H 2006 *J. Phys.* B: *At. Mol. Opt. Phys.* **39** 1825

[20] Wang L G, Qamar S, Zhu S Y and Zubairy M S 2008 *Phys. Rev.* A **77** 033833

[21] Mahmoudi M, Rabiei S W, Zohravi L E and Sahrai M 2008 *Opt. Commum.* **281** 4681

[22] Gardiner C W 1986 *Phys. Rev. Lett.* **56** 1917

[23] Ficek Z, Smyth W S and Swain S 1994 *Opt. Commun.* **110** 555

[24] Akram U, Wahiddin M R B and Ficek Z 1998 *Phys. Lett.* A **238** 117

[25] Dupertuis M A and Stenholm S J 1987 *J. Opt. Soc. Am.* B **4** 1094

[26] Bortman-Arbiv D, Wilson-Gordon A D and Friedmann H 2001 *Phys. Rev.* A **63** 043818

[27] Hu X M, Cheng G L, Zou J H, Li X and Du D 2005 *Phys. Rev.* A **72** 023803