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Raman scattering in a four-level atomic system with hyperfine structure*

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We propose and analyse an efficient Raman scheme for suppressing the absorption of a weak probe beam in a typical four-level atomic system with a nearly hyperfine doublet structure of two higher-lying excited levels for the two cases of transient regime and steady-state process. For the transient process, using the numerical calculations by a nice MATHEMATICA code, we find that the magnitude of the probe absorption at line centre of the probe transition is small compared to the standard three-level atomic system based on electromagnetically induced transparency (EIT). In particular, our results show that the probe absorption can be completely eliminated under the condition of Raman resonance, i.e. we only require that two-photon detuning is zero within the range of the hyperfine two-level frequency gap for the case of the steady state. In contrast to the standard three-level EIT scheme, one of the key advantages of our four-level Raman scheme is that under the Raman resonance condition we can observe one transparency window without the need of exact vanishing of one- and two-photon detuning. As a consequence, the atomic hyperfine structure cannot be a hindrance for obtaining EIT.

Keywords: transient absorption, steady state, dressed atomic state
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1. Introduction

Quantum coherence and interference have led to the appearance of many new effects and techniques in quantum optics and atomic physics, such as non-absorption resonances, electromagnetically induced transparency (EIT) and ultraslow optical solitons.[1−19] There have been many theoretical and experimental studies in the literature dealing with the absorption and amplification of light. Representative papers are cited in Refs.[5-16]. Agarwal and Harshawardhan[5] showed that in a four-level ladder-type system, two-photon absorption could be selectively suppressed or enhanced. In a recent experiment, Yan *et al* demonstrated that EIT in a standard Λ-type configuration could be used to suppress both single-photon and two-photon absorptions simultaneously.[6,7] Harris *et al* proposed the use of EIT to suppress absorption of the short-wavelength light generated in a four-wave mixing (FWM) scheme and showed that the FWM efficiency could be greatly enhanced.[8] Quite recently, Deng *et al* reported the first theoretical investigation of optical coherent four-wave mixing (OCFWM) with a weak coherent probe laser field based on EIT and showed that such a scheme could lead to many orders of magnitude enhancement in the amplitude of the generated wave in a typical four-level atomic system.[10] Later on, on the basis of the proposal of Deng *et al*, Wu *et al* analysed and discussed a FWM scheme in a five-level atomic system and hyper-Raman scattering (HRS) in resonant coherent media by the use of EIT, which led to suppressing both two-photon and three-photon absorptions in both FWM and HRS schemes and enabling the FWM to proceed through real, resonant intermediate states without absorption loss.[15,16] In particular, with the advent of Bose–Einstein condensate in an atomic gas there has been much interest in studying ultra-low light speed propagation with EIT,[2] coherent optical information storage based on EIT[3] and FWM with matter wave both experimentally and theoretically in the framework of nonlinear atomic optics.[20−26]

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In this paper, we analyse and discuss the probe-absorption properties in a typical four-level atomic system with a nearly hyperfine doublet structure of two higher-lying excited levels for the two cases of transient process and steady-state process. This paper is organized as follows. In Section 2, the model is described and the density-matrix equations of motion describing the atom–field interaction for the system under consideration are derived when a strong coherent coupling field and a weak probe field are applied. In Section 3, we investigate numerically the transient absorption properties via a simple MATHEMATICA code. It is shown that the magnitude of the probe absorption at the line centre is sufficiently decreased as compared with the usual three-level EIT scheme due to the fact that the additional hyperfine level is included. In Section 4, we examine the probe absorption under the steady-state condition. Our results show that the probe absorption can be completely eliminated under the condition of Raman resonance, i.e. we only require that the two-photon resonance be satisfied within the range of the hyperfine two-level frequency gap for the case of the steady state. In contrast to the standard three-level EIT scheme, our four-level Raman scheme can result in one EIT window without the need of exact vanishing of one- and two-photon detuning. To conclude it is worthwhile emphasizing that the atomic hyperfine structure cannot be a hindrance for obtaining EIT. Section 5 is a brief summary of our main results.

2. The atomic model and equations of motion

Let us consider a closed four-level atomic system coupled with two laser fields as depicted in Fig. 1, which includes a nearly hyperfine closely spaced doublet labelled $|3\rangle$ and $|4\rangle$ with a frequency difference of 2Δ. The transition $|2\rangle \leftrightarrow |3\rangle(|4\rangle)$ of frequency $\omega_{32}(\omega_{42})$ is driven by a strong-coupling laser of angular frequency $\omega_{\rm c}$ with Rabi frequency $\Omega_{32(42)} = \mu_{32(42)}E_{\rm c}/(2\hbar)$. A weak probe laser of angular frequency $\omega_{\rm p}$ with Rabi frequency $\Omega_{31(41)} = \mu_{31(41)}E_{\rm p}/(2\hbar)$ is applied to the transition $|1\rangle \leftrightarrow |3\rangle(|4\rangle)$ of frequency $\omega_{31(41)}$. Level $|4\rangle$ is the nearly hyperfine level with respect to the standard EIT-based three-level atomic configuration; μ_{ij} stands for the matrix element of the dipole moment between levels $|i\rangle$ and $|j\rangle$, and $2\gamma_{ij}$ is the spontaneous decay rate from levels $|i\rangle$ to $|j\rangle$.

In the present analysis we make the following assumptions. (1) Here we assume that the Rabi frequencies are real and that $\mu_{32} = \mu_{42} = \mu_c$ and $\mu_{31} = \mu_{41} = \mu_{\rm p}$. (2) As usual, we will ignore the dephasing rate γ_{21} between levels $|2\rangle$ and $|1\rangle$. (3) We assume that the excited levels $|3\rangle$ and $|4\rangle$ are closely spaced, but the corresponding frequency difference 2Δ is large as compared to the decay rates $2\gamma_{3j}$ and $2\gamma_{4j}(j = 1, 2)$. This approximation allows us to neglect the decay-induced coherence terms (proportional to $\sqrt{\gamma_{3j}\gamma_{4j}}$) in the following density matrix equations of motion.

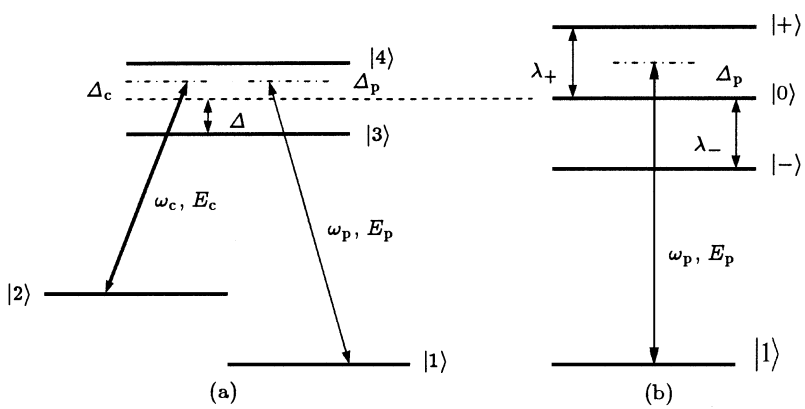


Fig.1. (a) Schematic diagram of the four-level atomic system under consideration. In the D_1 line of ^{87}Rb (nuclear spin $I = 3/2$) cold atoms, states $|1\rangle$ and $|2\rangle$ correspond to the ground-state Zeeman sublevels $|5\ {\rm S}_{1/2}$, F=1, $M_{\rm F} = \pm 1\rangle$, states $|3\rangle$ and $|4\rangle$ correspond to the excited-state Zeeman sublevels $|5\ {\rm P}_{1/2}, F = 1, M_F = 0\rangle$ and $|5\ {\rm P}_{1/2}$, F=2, $M_{\rm F} = 0\rangle$, respectively. In this case the *Rb* atom is driven by a right circularly polarized coupling light σ_+ with amplitude $E_{\rm c}$ and angular frequency $\omega_{\rm c}$ and a left circularly polarized probe light σ_- with amplitude $E_{\rm p}$ and angular frequency $\omega_{\rm p}$. (b) Corresponding dressed-state picture for the $\Delta_{\rm c} = 0$ case.

In Schrödinger's picture and in rotating-wave approximation, the semiclassical Hamiltonian describing the atom–field interaction for the system under study can be written as

$$\begin{aligned} H =& \sum_{j=1}^{4} E_j^{(a)} |j\rangle\langle j| - \hbar(\Omega_{\rm c} {\rm e}^{-{\rm i}\omega_{\rm c} t}|3\rangle\langle 2| \\ &+ \Omega_{\rm c} {\rm e}^{-{\rm i}\omega_{\rm c} t}|4\rangle\langle 2| + \Omega_{\rm p} {\rm e}^{-{\rm i}\omega_{\rm p} t}|3\rangle\langle 1| \\ &+ \Omega_{\rm p} {\rm e}^{-{\rm i}\omega_{\rm p} t}|4\rangle\langle 1| + {\rm h.c.}, \end{aligned} \tag{1}$$

where h.c. means Hermitian conjugation. The quantities $\Omega_n (n = c, p)$ denote one-half Rabi frequency of the coupling and probe lasers, i.e. $\Omega_{\rm p} = \mu_{\rm p} E_{\rm p}/(2\hbar)$ and $\Omega_{\rm c} = \mu_{\rm c} E_{\rm c}/(2\hbar)$, with $\mu_{\rm c,p} = \boldsymbol{\mu}_{\rm c,p} \cdot \boldsymbol{e}_{\rm L}$ ($\boldsymbol{e}_{\rm L}$ is the polarization unit vector of the corresponding laser field) denoting the dipole moment for the transition between levels $|i\rangle$ and $|j\rangle$, and $E_j^{(a)} = \hbar\omega_j$ is the energy of the atomic state $|j\rangle$. For simplicity of analysis, in what follows we will take $\omega_1 = 0$ (for ground state $|1\rangle$) as the energy origin. Turning to the interaction picture, the Hamiltonian can be rewritten as follows (taking $\hbar = 1$):

$$H_0 = (\omega_{\rm p} - \omega_{\rm c})|2\rangle\langle 2| + \omega_{\rm p}|3\rangle\langle 3| + \omega_{\rm p}|4\rangle\langle 4|, \tag{2}$$

$$\begin{aligned} H_{\rm int} =& (\Delta_{\rm c} - \Delta_{\rm p})|2\rangle\langle 2| - (\Delta_{\rm p} + \Delta)|3\rangle\langle 3| \\ &+ (\Delta - \Delta_{\rm p})|4\rangle\langle 4| \\ &- (\Omega_{\rm c}|3\rangle\langle 2| + \Omega_{\rm c}|4\rangle\langle 2| \\ &+ \Omega_{\rm p}|3\rangle\langle 1| + \Omega_{\rm p}|4\rangle\langle 1| + {\rm h.c.}), \end{aligned} \tag{3}$$

where

$$\Delta_{\rm c} = \omega_{\rm c} - \omega_0, \quad \Delta_{\rm p} = \omega_{\rm p} - \omega_0,$$

$$\Delta = \frac{1}{2}(\omega_4 - \omega_3) \quad \text{and} \quad \omega_0 = \frac{1}{2}(\omega_4 + \omega_3).$$

$\Delta_{\rm p}$ and $\Delta_{\rm c}$ are the related frequency differences between the driven-laser field and the two-level transition, as shown in Fig.1(a).

Under the electro-dipole and rotating-wave approximations, by the standard approach, we can easily obtain the following density matrix equations of motion:

$$\dot{\rho}_{44} = {\rm i}\Omega_{\rm c}(\rho_{24} - \rho_{42}) + {\rm i}\Omega_{\rm p}(\rho_{14} - \rho_{41}) - (2\gamma_{41} + 2\gamma_{42})\rho_{44}, \tag{4}$$

$$\dot{\rho}_{33} = {\rm i}\Omega_{\rm c}(\rho_{23} - \rho_{32}) + {\rm i}\Omega_{\rm p}(\rho_{13} - \rho_{31}) - (2\gamma_{31} + 2\gamma_{32})\rho_{33}, \tag{5}$$

$$\dot{\rho}_{22} = {\rm i}\Omega_{\rm c}(\rho_{32} - \rho_{23}) + {\rm i}\Omega_{\rm c}(\rho_{42} - \rho_{24}) + 2\gamma_{32}\rho_{33} + 2\gamma_{42}\rho_{44}, \tag{6}$$

$$\dot{\rho}_{43} = -(\gamma_{41} + \gamma_{42} + \gamma_{31} + \gamma_{32} + {\rm i}2\Delta)\rho_{43} + {\rm i}\Omega_{\rm c}\rho_{23} - {\rm i}\Omega_{\rm c}\rho_{42} + {\rm i}\Omega_{\rm p}\rho_{13} - {\rm i}\Omega_{\rm p}\rho_{41}, \tag{7}$$

$$\dot{\rho}_{42} = -(\gamma_{41} + \gamma_{42} + {\rm i}\Delta - {\rm i}\Delta_{\rm c})\rho_{42} + {\rm i}\Omega_{\rm p}\rho_{12} - {\rm i}\Omega_{\rm c}\rho_{43} + {\rm i}\Omega_{\rm c}(\rho_{22} - \rho_{44}), \tag{8}$$

$$\dot{\rho}_{41} = -(\gamma_{41} + \gamma_{42} + {\rm i}\Delta - {\rm i}\Delta_{\rm p})\rho_{41} + {\rm i}\Omega_{\rm c}\rho_{21} - {\rm i}\Omega_{\rm p}\rho_{43} + {\rm i}\Omega_{\rm p}(\rho_{11} - \rho_{44}), \tag{9}$$

$$\dot{\rho}_{32} = -(\gamma_{31} + \gamma_{32} - {\rm i}\Delta - {\rm i}\Delta_{\rm c})\rho_{32} + {\rm i}\Omega_{\rm p}\rho_{12} - {\rm i}\Omega_{\rm c}\rho_{34} + {\rm i}\Omega_{\rm c}(\rho_{22} - \rho_{33}), \tag{10}$$

$$\dot{\rho}_{31} = -(\gamma_{31} + \gamma_{32} - {\rm i}\Delta - {\rm i}\Delta_{\rm p})\rho_{31} + {\rm i}\Omega_{\rm c}\rho_{21} - {\rm i}\Omega_{\rm p}\rho_{34} + {\rm i}\Omega_{\rm p}(\rho_{11} - \rho_{33}), \tag{11}$$

$$\dot{\rho}_{21} = -{\rm i}(\Delta_{\rm c} - \Delta_{\rm p})\rho_{21} + {\rm i}\Omega_{\rm c}\rho_{31} + {\rm i}\Omega_{\rm c}\rho_{41} - {\rm i}\Omega_{\rm p}\rho_{23} - {\rm i}\Omega_{\rm p}\rho_{24}. \tag{12}$$

Closure of this atomic system requires that $\rho_{11} + \rho_{22} + \rho_{33} + \rho_{44} = 1$ and $\rho_{ij} = \rho_{ji}^*$. By a straightforward semiclassical analysis, the above matrix elements can be used to calculate the total linear complex susceptibility χ of the probe transition, i.e.

$$\chi = \frac{N_0 |\mu_{\rm p}|^2 (\rho_{31} + \rho_{41})}{2\hbar\varepsilon_0 \Omega_{\rm p}} \propto \rho_{31} + \rho_{41}, \tag{13}$$

where N_0 is the atomic number density in the medium and ε_0 is the permittivity of free space. Therefore, in the limit of a weak probe field, the gain-absorption coefficient for the probe laser coupled to the transition $|1\rangle \leftrightarrow |3\rangle(|4\rangle)$ is proportional to the term ${\rm Im}(\rho_{31} + \rho_{41})$. If ${\rm Im}(\rho_{31} + \rho_{41}) > 0$, the probe laser will be absorbed. In contrast, the probe laser will be amplified. In the following, in order to investigate the absorption properties of such a weak probe beam for the transient and steady-state processes, we begin by solving numerically the density matrix equations (4)–(12) using a nice MATHEMATICA code with the initial conditions $\rho_{11}(0) = 1$, $\rho_{22}(0) = \rho_{33}(0) = \rho_{44}(0) = 0$, and $\rho_{ij}(0) = 0$ for $i \neq j (i, j = 1 - 4)$. Note that, in this paper, parameters $\Omega_{\rm c,p}$, $\Delta_{\rm c,p}$, γ_{ij} and Δ are in units of γ.

3. The transient absorption of the probe field

In this section, we present several absorption spectra under the condition of using different sets of parameters, as shown in Figs.2–4. From Fig.2 one can conclude that at the line centre of the probe transition ($\Delta_p = 0$), the magnitude of the transient absorption of the probe field at point A corresponding to $\Delta = \gamma$ is much smaller than the magnitude of the transient absorption of the probe field at point B for the standard three-level EIT corresponding to $\Delta = 0$. From the above analysis, we can see that the probe absorption at the line centre ($\Delta_p = 0$) can be considerably suppressed due to the fact that the additional hyperfine level is introduced for the transient process. The calculated probe absorption coefficient [$\mathrm{Im}(\rho_{31} + \rho_{41})$] of such a weak probe laser versus the hyperfine two-level frequency gap Δ between levels $|4\rangle$ and $|3\rangle$ for the probe and coupling laser detunings $\Delta_p = \Delta_c = 0$ is shown in Fig.3. Note that, in Fig.3, the dot-dashed straight line stands for the magnitude of the probe absorption for the standard three-level EIT and the positions $A - F$ correspond to the minimum of the probe absorption. We observe that the magnitude of the transient absorption of the probe field with the inclusion of the additional hyperfine level is always lower than the magnitude of the probe absorption with the standard three-level EIT at $\Delta_p = \Delta_c = 0$. In Fig.4, we plot the probe absorption profile as a function of probe detuning Δ_p, under the condition of Raman resonance $\Delta_p = \Delta_c$. For this case, we find that the maximal absorption peak is always located below the probe absorption of the standard three-level EIT.

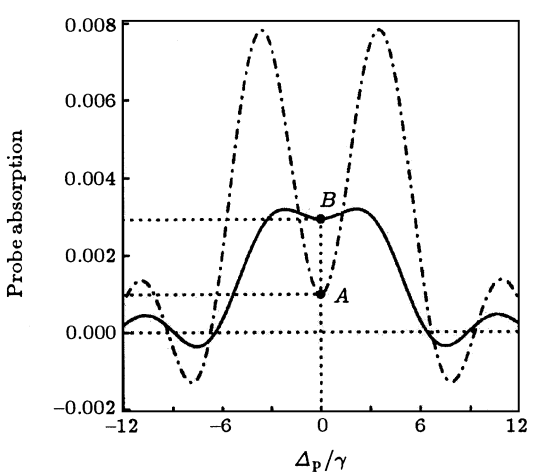


Fig.2. The probe absorption as a function of dimensionless probe detuning Δ_p/γ for the transition process. Solid curve: $\Delta = 0$; dot-dashed curve: $\Delta = \gamma$. Other fitting parameters used are $\Delta_c = 0$, $\Omega_p = 0.01\gamma$, $\Omega_c = 2\gamma$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma$.

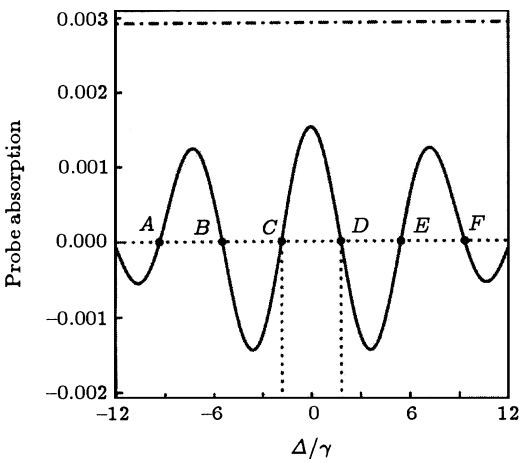


Fig.3. The probe absorption as a function of dimensionless hyperfine two-level frequency gap Δ/γ for the transient process. The dot-dashed straight line stands for the magnitude of the probe absorption for the standard three-level EIT at $\Delta_p = \Delta_c = 0$.The fitting parameters used are $\Delta_p = 0$, $\Delta_c = 0$, $\Omega_p = 0.01\gamma$, $\Omega_c = 2\gamma$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma$.

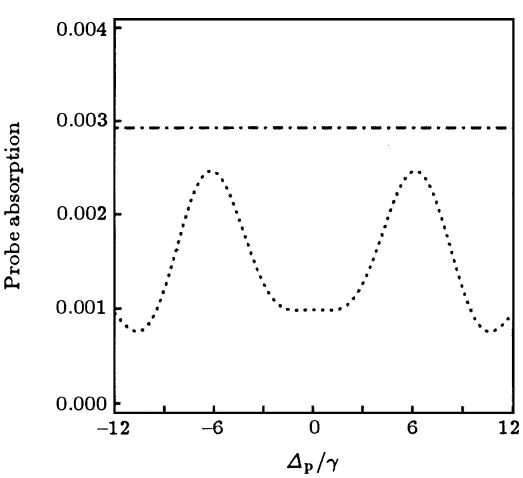


Fig.4. The probe absorption as a function of dimensionless probe detuning Δ_p/γ for the transient process under the Raman resonance condition $\Delta_p = \Delta_c$. The dot-dashed straight line stands for the magnitude of the probe absorption for the standard three-level EIT at $\Delta_p = \Delta_c = 0$. Other fitting parameters used are $\Omega_p = 0.01\gamma$, $\Omega_c = 2\gamma$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma$, $\Delta = \gamma$.

4. The probe absorption under the steady-state condition

The general steady-state solutions to Eqs.(4)–(12) can be derived analytically (for the analysis with two resonant coupling and probe lasers), but the tedious expressions for the solutions corresponding to the situation of an off-resonant coupling laser offers no clear physical insight. We therefore present a numerical calculations here. The present process is most clearly understood in terms of the so-called dressed states produced by the strong coupling field: that is, the $|2\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |4\rangle$ transitions together with the coupling field are treated as a global "atom+field" system forming the dressed states as shown in Fig.1(b). It is obvious that, under the action of the strong coherent coupling field, levels $|3\rangle$ and $|4\rangle$ should be split into three dressed-state sublevels $|\pm\rangle$ and $|0\rangle$. The energy eigenstates of the three semiclassical dressed states for the $\Delta_c = 0$ case are written as

$$\begin{aligned}|+\rangle =& -\frac{\Omega_c}{\sqrt{\Delta^2+2\Omega_c^2}}|2\rangle \\ & -\frac{1}{2}\left(\frac{\Delta}{\sqrt{\Delta^2+2\Omega_c^2}}-1\right)|3\rangle \\ & +\frac{1}{2}\left(\frac{\Delta}{\sqrt{\Delta^2+2\Omega_c^2}}+1\right)|4\rangle, \qquad (14)\\ |0\rangle =& \frac{\Delta}{\sqrt{\Delta^2+2\Omega_c^2}}|2\rangle - \frac{\Omega_c}{\sqrt{\Delta^2+2\Omega_c^2}}|3\rangle \\ & +\frac{\Omega_c}{\sqrt{\Delta^2+2\Omega_c^2}}|4\rangle, \qquad (15)\\ |-\rangle =& \frac{\Omega_c}{\sqrt{\Delta^2+2\Omega_c^2}}|2\rangle \\ & +\frac{1}{2}\left(\frac{\Delta}{\sqrt{\Delta^2+2\Omega_c^2}}+1\right)|3\rangle \\ & -\frac{1}{2}\left(\frac{\Delta}{\sqrt{\Delta^2+2\Omega_c^2}}-1\right)|4\rangle. \qquad (16)\end{aligned}$$

The corresponding energy eigenvalues are given by

$$\lambda_0 = 0, \quad \text{and} \quad \lambda_\pm = \pm\sqrt{\Delta^2+2\Omega_c^2}. \qquad (17)$$

With the above results, the dressed-state transition dipole moments are derived as follows:

$$\begin{aligned}\mu_{+1} &= \langle +|P|1\rangle = \mu_p,\\ \mu_{01} &= \langle 0|P|1\rangle = 0,\\ \mu_{-1} &= \langle -|P|1\rangle = \mu_p, \qquad (18)\end{aligned}$$

where $P = \mu_p(|3\rangle\langle 1| + |4\rangle\langle 1|)$ represents the atomic polarization operator of the probe transition.

On the basis of the above analysis, it is straightforwardly shown that the transition between dressed-levels $|0\rangle$ and $|1\rangle$ is not allowed because of the complete interference cancellation between the transitions of $|1\rangle \leftrightarrow |3\rangle$ and $|1\rangle \leftrightarrow |4\rangle$. From Fig.5, we can observe two symmetric absorption peaks with equal amplitude on the probe spectra in the case of $\Delta_c = 0$ because of the relation $\mu_{+1} = \mu_{-1}$. The two absorption peaks correspond respectively to the dressed-state transitions $|1\rangle \leftrightarrow |+\rangle$ and $|1\rangle \leftrightarrow |-\rangle$, so the peaks are always located at $\Delta_p = \pm\sqrt{\Delta^2+2\Omega_c^2}$ (see Fig. 5), and the difference between them becomes larger with increasing Ω_c. Due to the fact that the transition $|1\rangle \leftrightarrow |0\rangle$ in the dressed-state picture is interference-cancelled out and this interference always makes a negative contribution to the probe absorption, thus it is easy to understand that the probe absorption at the line centre can be approximately equal to zero, i.e. the probe absorption can be completely suppressed, such that the absorption peak at the line centre is too small to be observed. In Fig.6, under the condition of Raman resonance ($\Delta_p = \Delta_c$), we plot the probe absorption profile as a function of the probe detuning Δ_p. We find that the probe absorption is always approximately equal to zero within the range of the hyperfine two-level frequency gap 2Δ. This means that if the two-photon Raman resonance is satisfied, we can observe the transparency window without the need to make one- and two-photon detunings exactly vanish.

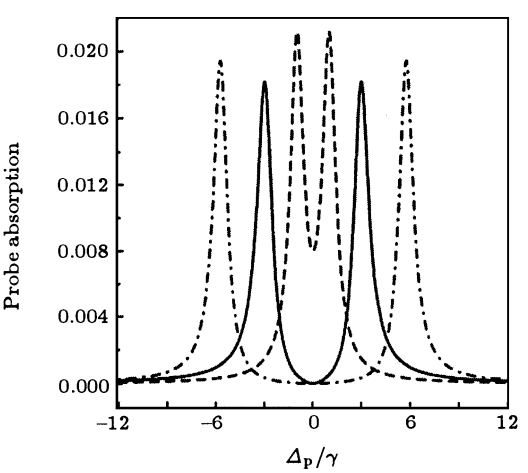


Fig.5. The probe absorption as a function of dimensionless probe detuning Δ_p/γ for the steady-state process. Dashed curve: $\Omega_c = 0$, solid curve: $\Omega_c = 2\gamma$. Dot-dashed curve: $\Omega_c = 4\gamma$. The corresponding peaks are located at $\Delta_p = \gamma$, $\Delta_p = 3\gamma$ and $\Delta_p \simeq 5.7\gamma$, respectively. Other fitting parameters used are $\Omega_p = 0.01\gamma$, $\Delta = \gamma$, $\Delta_c = 0$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma$.

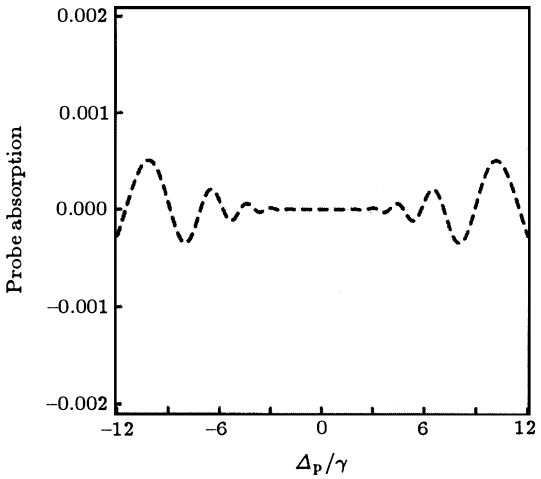


Fig.6. The probe absorption as a function of dimensionless probe detuning Δ_p/γ for the steady-state process under the Raman resonant condition $\Delta_p = \Delta_c$. Other fitting parameters used are $\Omega_p = 0.01\gamma$, $\Omega_c = 2\gamma$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma$, $\Delta = \gamma$.

5. Conclusion

In summary, in the present paper, we have analysed and discussed the probe-absorption properties in a typical four-level atomic system with a nearly hyperfine doublet structure of two higher-lying excited levels for the transient regime and the steady-state processes by using the detailed numerical simulations based on the density matrix equations (4)–(12) and analytical calculations in the dressed-state picture. For the former regime, it is shown that the magnitude of the probe absorption at the line centre of the probe transition $\Delta_p = 0$ is greatly decreased as compared with the usual three-level EIT scheme due to the fact that the additional hyperfine level is included. For the latter, our results show that the probe absorption can be completely eliminated under the condition of Raman resonance, i.e. we only require that two-photon resonance be satisfied within the range of the hyperfine two-level frequency gap 2Δ for the case of the steady state. In contrast to the standard three-level EIT scheme, one of the key advantages of our four-level Raman scheme is that under the Raman resonance condition we can observe one EIT window without the need to make one- and two-photon detunings exactly vanish. In conclusion it is worthwhile emphasizing that the atomic hyperfine structure cannot be a hindrance for obtaining EIT.

Acknowledgment

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