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Coupled atomic coherences induced by a standing wave

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ABSTRACT

Atomic coherence effects with a strong standing-wave coupling field on a four-level system of hot Rb atoms are investigated. Amplitude modulation and phase modulation of the strong standing-wave field are coherently coupled in linear and nonlinear atomic polarizations in various multi-wave mixing processes. Such coupled atomic coherence effect leads to one-photon triple-bandpass filters in an electromagnetically induced transparency structure and two-photon bandgaps at two different wavelengths under two-photon resonance in the hot four-level Rb atoms. Such system can be useful for dispersion compensation in multi-channel WDM systems, all-optical switching and optical filters at multiple wavelengths.

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1. Introduction

Recently, there are many works on coherent interactions between a strong standing-wave field and multi-level atoms for various applications, such as storage of light pulses [1], measurement of coherent properties of Bose–Einstein condensates [2], atomic interferometry [3,4], subwavelength localization of atoms [5], all-optical switching [6], diffraction-type electromagnetically induced grating (EIG) [7], dynamic controllable dispersion compensation [8], enhanced nonlinearities at low light intensities [9,10] and manipulation of light pulses by enhanced nonlinearities [11]. Here, an inverted Y-type four-level atomic system, as shown in Fig. 1, is employed to investigate fundamental physical mechanisms for forming photonic bandgaps using atomic coherence induced by a standing-wave driving field, which can help in designing appropriate coherent systems to achieve one-photon multi-channel filters or multi-photonic bandgaps.

It is well known that atomic coherence induced by a strong traveling-wave field can be used to manipulate the absorptive and dispersive properties of the probe field to generate the phenomenon of electromagnetically induced transparency (EIT) [12–15]. Taking a three-level ladder-type system as an example, by applying a strong electromagnetic field to dress the upper two bare states, an opaque transition can be made transparent to the weak probe field at its resonant frequency. The real and imaginary parts of the susceptibility for the atomic medium can be expressed as a function of $|\Omega_c|^2$ [14,15], where the Rabi frequency of the traveling-wave field $\Omega_c = g_{ij}E_c$, in which $2\hbar g_{ij}$ $(i,j = 1,2,3)$ are the dipole-matrix elements of atomic transitions and E_c is the slowly varying amplitude of the coupling field. When a standing-wave field is employed in place of the traveling-wave field, the expression for Rabi frequency is changed to be $\Omega_c = \Omega_+ e^{-ik_c \cdot z} + \Omega_- e^{ik_c \cdot z}$, where $\Omega_\pm$ are the forward and backward components which form the standing-wave field and k_c is the wave vector. With the standing-wave driving field, the refractive index of the medium is then periodically modulated along the propagating direction (z) for the probe field. Since $|\Omega_c|^2 = |\Omega_+|^2 + |\Omega_-|^2 + |\Omega_+||\Omega_-|e^{-2ik_c \cdot z} + |\Omega_+||\Omega_-|e^{2ik_c \cdot z}$ for the standing-wave driving field, atomic coherence induced by this standing-wave field has both amplitude (the former two terms) and phase (the last two terms) modulations. For the diffraction-type EIG [7], atomic coherence caused by the amplitude-modulation terms exerts over the probe field just as an amplitude grating does, and the effect of phase modulation leads to a dispersive relation for the probe field as modulated by a dispersive grating. For the transmission-type EIG to form photonic bandgaps [8,16], effects due to mixed amplitude and phase modulations on the probe field will be discussed.

Considering an inverted Y-type four-level system as shown in Fig. 1, the bare atomic energy levels $|1>$, $|2>$, $|3>$ and $|4>$ correspond to $|5S_{1/2}, F=1>$, $|5S_{1/2}, F=2>$, $5P_{3/2}$ and $5D_{3/2}$ states of the Rb atoms, respectively. The strong standing-wave field Ω_c is formed by two counter-propagating fields $\Omega_\pm$ in z direction, which keeps on resonance with the atomic transition $|3>-|2>$. Two probe fields Ω_{p1} and Ω_{p2} drive the atomic transitions $|3>-|1>$ and $|4>-|3>$, respectively. In order to utilize Doppler-free configuration for each component in the standing wave in hot atoms [15], two probe fields Ω_{p1} and Ω_{p2} are chosen to propagate along the z and $-z$ axis in the medium, respectively. Such four-level

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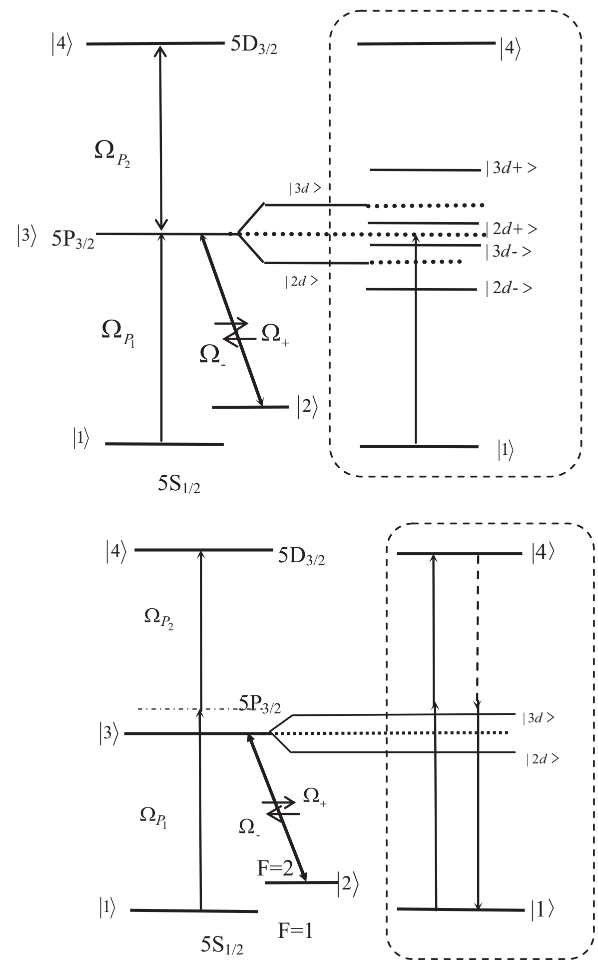

Fig. 1. An inverted Y-type four-level system in bare and dressed states (inside dotted boxes) for (a) EIT one-photon multi-channel structure and (b) two-photon resonance structure.

system in Fig. 1 has potential to generate one-photon triple-bandpass channels or two-photon bandgaps, and it can be made into two different configurations according to these specific situations. When Ω_{p2} is kept on resonance with the corresponding atomic transition and is as strong as the standing-wave field Ω_c, as shown in Fig. 1(a) (thereafter called the EIT structure), levels $|1>$, $|3>$ and $|4>$ form a traditional ladder-type EIT system [15]. However, when the probe beam Ω_{p2} scans into two-photon resonance with the weak probe field Ω_{p1} and keeping $\Delta_{p1}+\Delta_{p2}=0$, as shown in Fig. 1(b), atomic coherence induced by the strong standing-wave field Ω_c takes effect on the two resonant wavelengths, therefore we refer such system as a two-photon resonance structure. For both EIT and two-photon resonance structures, with the strong field Ω_c on, level $|3>$ is viewed as a coherent superposition of two split dressed states $|2\mathrm{d}>$ and $|3\mathrm{d}>$ and the frequency interval of the two dressed states is $2\Omega_c$. In the EIT structure, because a resonant strong Ω_{p2} is applied, $|2\mathrm{d}>$ and $|3\mathrm{d}>$ further split into four levels as shown in Fig. 1(a), where $|2d\pm\geq\cos{}^*\theta_-|2d>\pm\ \sin\theta_-|4>)$ and $|3d\pm\geq\cos{}^*\theta_+|3d>\pm\ \sin\theta_+|4>)$ with $tg(\theta_\pm)=\frac{\pm|\Omega_c|+\sqrt{|\Omega_{p2}|^2+|\Omega_c|^2}}{\Omega_{p2}}$ [17]. Therefore, in such inverted Y-type [18,19] or Y-type system [20,21], when Ω_c is used as a strong traveling-wave field, two-photon resonant absorption will be inhibited for the two probe beams Ω_{p1} and Ω_{p2} [18,20] (in two-photon resonance structure) or one-photon triple EIT pathways are produced [21] due to destructive interferences between three-photon and five-photon nonlinear processes (in the EIT structure). The scheme used in the current work is different from those in Refs. [18–21] since the traveling-wave coupling field is replaced by a standing-wave field Ω_c. If one neglects the effect of phase modulation terms in the standing-wave field, the amplitude modulation terms of the standing-wave field will have the same impact on the two probe fields as the traveling-wave field does since they have the same expressions (except double the strength). However, when both the amplitude and phase modulations are considered, a standing-wave grating is formed [8,11], and two new fields are generated by the phase modulation terms under phase-matching conditions. The frequencies of the generated fields are the same as those of the two probe fields Ω_{p1} and Ω_{p2}, but they propagate in the opposite directions relative to Ω_{p1} and Ω_{p2}, respectively. Meanwhile, with respect to the polarizations of Ω_{p1} and Ω_{p2} with the standing-wave drive, contributions from the generated fields are included in order to satisfy the momentum conservation or phase-matching conditions. This means that the coupled coherences between Ω_{p1}, Ω_{p2} and the generated fields are induced by the phase modulation of the standing-wave field. The coupled coherences result in destructive interferences between phase and amplitude modulations, and thus create narrow two-photon bandgaps at two resonant wavelengths in the two-photon resonance and one-photon triple-bandpass channels in the EIT structure, respectively.

In the interaction picture, the Hamiltonian of the system in bare states is given as follows:

$$H=-\hbar\Delta_{p1}|3><3|-\hbar\Delta_{p2}|4><4|-\hbar\Omega_c|3><2|\\ -\hbar\Omega_{p1}|3><1|-\hbar\Omega_{p2}|4><3|+c.c. \tag{1}$$

where $\Delta_{p1}=(\omega_{p1}-\omega_{31})$ and $\Delta_{p2}=(\omega_{p2}-\omega_{43})$ are frequency detunings of the probe fields Ω_{p1} and Ω_{p2} with corresponding atomic transitions and Rabi frequencies of the respective incident lasers as $\Omega_{p1}=\Omega_1 e^{-ik_{p1}\times r}$, $\Omega_{p2}=\Omega_2 e^{ik_{p2}\times r}$, $\Omega_1=g_{31}E_{p1}$, $\Omega_2=g_{43}E_{p2}$. The Rabi frequency of the strong standing-wave driving field is $\Omega_c=\Omega_+e^{-ik_c\times z}+\Omega_-e^{ik_c\times z}$; $\Omega_\pm=g_{32}E_\pm$. $2\hbar g_{ij}$ $(i,j=1,2,3,4)$ are the dipole-matrix elements for the atomic transitions; $E_{p1,p2,\pm}$ are slowly-varying field amplitudes, while $k_{p1,p2}$ and k_c are wave vectors of the two probe fields and standing-wave drive field, respectively.

The density–matrix equations of this system are governed by [22]:

$$\begin{aligned}
\dot\rho_{21}&=-(\gamma_{21}-i(\Delta_{p1}-\Delta_c))\rho_{21}-i\Omega_c^*\rho_{31}+i\Omega_{p1}\rho_{23}\\
\dot\rho_{31}&=-(\gamma_{31}-i\Delta_{p_1})\rho_{31}-i\Omega_c\rho_{21}-i\Omega_{p2}^*\rho_{41}+i\Omega_{p1}(\rho_{33}-\rho_{11})\\
\dot\rho_{41}&=-(\gamma_{41}-i(\Delta_{p_1}+\Delta_{p2}))\rho_{41}-i\Omega_{p2}\rho_{31}+i\Omega_{p1}\rho_{43}\\
\dot\rho_{32}&=-(\gamma_{32}-i\Delta_c)\rho_{32}-i\Omega_{p2}^*\rho_{42}-i\Omega_{p1}\rho_{12}+i\Omega_c(\rho_{33}-\rho_{22})\\
\dot\rho_{42}&=-(\gamma_{42}-i(\Delta p_2+\Delta_c))\rho_{42}+i\Omega_c\rho_{43}-i\Omega_{p2}\rho_{32}\\
\dot\rho_{43}&=-(\gamma_{43}-i\Delta_{p2})\rho_{43}+i\Omega_c^*\rho_{42}-i\Omega_{p1}^*\rho_{41}+i\Omega_{p2}(\rho_{44}-\rho_{33})\\
\dot\rho_{33}&=\Gamma_4\rho_{44}-\Gamma_3\rho_{33}-i\Omega_{p1}^*\rho_{31}-i\Omega_C^*\rho_{32}+cc.\\
\dot\rho_{44}&=-\Gamma_4\rho_{44}-i\Omega_{p2}^*\rho_{43}+cc.
\end{aligned} \tag{2}$$

Under the assumption of a weak probe field Ω_{p1} at all time, one can assume $\rho_{22}=0$, and therefore $\rho_{11}=1-\rho_{33}-\rho_{44}$. In the absence of collisions, $\gamma_{ij}=(\Gamma_i+\Gamma_j)/2$, where Γ_i is the natural decay rate of level $|i>$. Since Eq. (2) cannot be solved exactly, we apply the perturbation method to solve it since Ω_{p1} is assumed to be weak enough. We keep all terms of the strong fields Ω_c and Ω_{p2} in Eq. (2) and take only the lowest order of amplitude for the weak probe field Ω_{p1}. The steady-state solutions of Eq. (2) are then required to be in the form of [23]:

$$\rho_{ij}=\rho_{ij}^{(0)}+\Omega_{p1}\rho_{ij}^{(1)}+\Omega_{p1}^*\rho_{ij}^{(-1)}\ (i,j=1,2,3,4) \tag{3}$$

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