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Two-dimensional electromagnetically induced grating via nonlinear modulation in a five-level atomic system

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ABSTRACT

We theoretically propose a scheme for two-dimensional electromagnetically induced grating in a five-level atomic system interacting with two orthogonal standing-wave fields. In such atomic system, nonlinear absorption or refractivity can be significantly enhanced with nearly vanishing linear absorption under different resonant conditions. When applying two coupling fields with standing-wave pattern, absorption grating or phase grating, which efficiently diffracts a probe beam into high-order directions, can be formed in the media. The diffraction efficiencies of the gratings depend strongly on the interaction length, the intensities and detunings of the coupling fields. By investigating the third-order nonlinearity of the system, it is found that the spatial variations of amplitude and phase arise from the nonlinear modulation. The proposed gratings can be used as multi-channel beam splitter and have potential applications in optical information processing and optical networking.

1. Introduction

It is well known that the light–media interaction plays a significant role in modern quantum optics. Many fascinating phenomena based on atomic coherence and quantum interference have been actively investigated. Especially electromagnetically induced transparency (EIT) [1–3] is the most prominent one. In three-level Λ-type EIT system, the opaque medium is rendered transparent to a resonant probe field by means of a strong coupling field driving on the linked transition. Many applications based on EIT have been proposed, such as slow light [4,5], optical information storage [6,7], enhanced Kerr-nonlinear [8,9], optical switching [10,11], optical solitons [12,13], etc. In these schemes, a traveling-wave driving field is used to control the absorption and dispersion of light fields across a medium. When a standing-wave driving field substitutes for the traveling-wave driving field in EIT, the probe field, which propagates perpendicular or parallel to the standing-wave, experiences periodically modulated absorption and refraction. Consequently, a new type of all-optical device known as electromagnetically induced grating (EIG) can be formed in the media. The longitudinal EIG acts as a Bragg grating, which reflects the probe field within certain frequency range [14]. However, the transverse EIG is used as a diffraction grating, which distributes the probe energy into different directions [15]. Many potential applications has been proposed based on EIG, such as probing material optical properties [16], all optical switching and routing [17], beam splitting and fanning [18], optical bistability [19], electromagnetically induced Talbot effect [20,21], surface solitons of four-wave mixing [22,23], among others. A spatially dependent grating can also be used to enhance nonlinear multi-wave mixing processes [24]. In recent, the standing-wave driven EIG was applied to create parity-time (PT) symmetric optical potentials with periodical gain and loss profiles in a four-level N-type atomic system [25,26]. Additionally, the four-level N-type system can also create optical PT-antisymmetry in a standing-wave driven atomic lattice [27,28].

The most original transverse EIG is proposed in a three-level Λ system [15] and then experimentally demonstrated in cold and hot atomic samples [29,30]. However, the attainable diffraction efficiency is relatively small owing to the inevitable linear absorption. The absorption grating, which tends to gather light in the center maximal, has a limited diffraction efficiency due to the low average transmissivity. Phase gratings with nearly 100% transmissivity, however, resulted from the phase modulation are much more efficient. Thus, the high-order diffraction intensities of the probe beam can be extensively enhanced in the near-transparent medium [31–35]. In addition, nonlinear modulation have been also applied to realize EIG [36,37]. Generally, it can yield giant resonantly enhanced nonlinearities based on the EIT systems, while the linear susceptibilities of the media are almost zero for the probe field. Therefore, we can obtain nonlinear EIG under the

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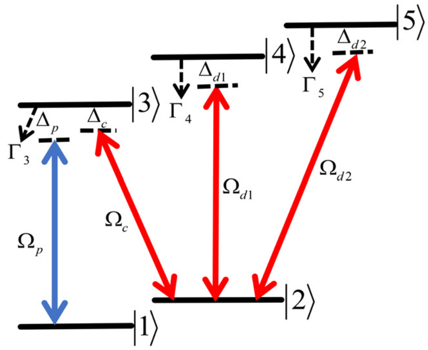

Fig. 1. Schematic diagram of a five-level atomic system.

appropriate parameters, which are distinct from other schemes founded on linear modulation.

Recently, the EIG schemes have been extended to two-dimensional (2D) for the best of our knowledge, which is demonstrated in different types of atomic systems [38–40]. In this paper, we propose a new scheme for two-dimensional electromagnetically induced grating in a five-level atomic system. By applying two standing-wave fields along the x and y axes to drive their corresponding atomic transitions, the probe susceptibility experiences a periodic variation, thereby leading to the formation of 2D EIG. It is shown that an amplitude grating can be converted into a phase grating by adjusting the detunings of two standing-wave fields, and thus improving the diffraction efficiency. Furthermore, we investigate the effect of third-order nonlinearity on the grating, the results show that the gratings result from the nonlinear absorption and phase modulations. Moreover, the controllability of the diffraction efficiency is also studied, and it is found that the diffracting power of the proposed gratings can be tuned by the intensities of the two standing-wave fields, Thus the gratings supply a further possibility of multi-channel beam splitter. We believe the proposed scheme may have potential applications in optical communications and optical information processing.

This paper is organized as follows: In Section 2, we present the atomic model and relevant equations. In Section 3, we respectively analysis and discuss the absorption grating and phasing grating. Finally, the conclusions are given in Section 4.

2. Atomic model and equations

We consider a five-level atomic system with two ground states $|1\rangle$, $|2\rangle$ and three excited states $|3\rangle$, $|4\rangle$, $|5\rangle$ as shown in Fig. 1. A weak probe field with Rabi frequency Ω_p interacts with the transition $|3\rangle-|1\rangle$. Three coupling fields with Rabi frequencies Ω_c, Ω_{d1} and Ω_{d2} drive the transitions $|3\rangle-|2\rangle$, $|4\rangle-|2\rangle$ and $|5\rangle-|2\rangle$, respectively. Here $\Delta_p = \omega_p - \omega_{31}$, $\Delta_c = \omega_c - \omega_{32}$, $\Delta_{d1} = \omega_{d1} - \omega_{42}$ and $\Delta_{d2} = \omega_{d2} - \omega_{52}$ denote the detunings of corresponding laser fields, where ω_{ij} $(i,j = 1-5)$ is the atomic resonant frequency of atomic transition $|i\rangle - |j\rangle$, and ω_p, ω_c, ω_{d1} and ω_{d2} denote the frequencies of laser fields. The Rabi frequencies of the optical fields are defined as $\Omega_p = \overrightarrow{E_p}\cdot\overrightarrow{\mu}_{13}/2\hbar$, $\Omega_c = \overrightarrow{E_c}\cdot\overrightarrow{\mu}_{23}/2\hbar$, $\Omega_{d1} = \overrightarrow{E}_{d1}\cdot\overrightarrow{\mu}_{24}/2\hbar$ and $\Omega_{d2} = \overrightarrow{E}_{d2}\cdot\overrightarrow{\mu}_{25}/2\hbar$, where $\overrightarrow{\mu}_{ij}$ represents the electric-dipole matrix element of transition $|i\rangle - |j\rangle$, E_p, E_c, E_{d1} and E_{d2} are the corresponding amplitudes of laser fields. The level structure can be realized in cold rubidium atoms for the reason that it can effectively reduce the impact of undesired effects such as Doppler effect and collisional broadening.

Under the electric-dipole approximation and rotating-wave approximation, the Hamiltonian of the system in the interaction picture is given by

$$\begin{aligned} H = & -\hbar\left(\Delta_p - \Delta_c\right)|2\rangle\langle 2| - \hbar\Delta_p|3\rangle\langle 3| - \hbar\left(\Delta_p - \Delta_c + \Delta_{d1}\right)|4\rangle\langle 4| \\ & -\hbar\left(\Delta_p - \Delta_c + \Delta_{d2}\right)|5\rangle\langle 5| \\ & -\hbar\left(\Omega_p|3\rangle\langle 1| + \Omega_c|3\rangle\langle 2| + \Omega_{d1}|4\rangle\langle 2| + \Omega_{d2}|5\rangle\langle 2| + H.C.\right). \end{aligned} \tag{1}$$

The system dynamics can be described by the motion equations for the probability amplitude of the states as follows

$$\dot{a}_1 = i\Omega_p a_3, \tag{2a}$$

$$\dot{a}_2 = i\left[\left(\Delta_p - \Delta_c\right)a_2 + \Omega_c a_3 + \Omega_{d1}a_4 + \Omega_{d2}a_5\right] - \gamma_2 a_2, \tag{2b}$$

$$\dot{a}_3 = i\left(\Omega_p a_1 + \Omega_c a_2 + \Delta_p a_3\right) - \frac{\Gamma_3}{2}a_3, \tag{2c}$$

$$\dot{a}_4 = i\left[\Omega_{d1}a_2 + \left(\Delta_p - \Delta_c + \Delta_{d1}\right)a_4\right] - \frac{\Gamma_4}{2}a_4, \tag{2d}$$

$$\dot{a}_5 = i\left[\Omega_{d2}a_2 + \left(\Delta_p - \Delta_c + \Delta_{d2}\right)a_5\right] - \frac{\Gamma_5}{2}a_5, \tag{2e}$$

where Γ_3, Γ_4, Γ_5 are the decay rates from the upper levels, and γ_2 represents the decoherence rate of the ground state. For simplicity, we assume $\Gamma_3 = \Gamma_4 = \Gamma_5 = \gamma$ and the other parameters take γ as unit.

The atomic-optical response of the probe field is described by the susceptibility χ_p. By solving the above equations under the steady-state condition ($\dot{a}_i = 0$) and weak probe field approximation ($|a_1|^2 = 1$), and according to the polarization of the medium $\overrightarrow{P}_p = \varepsilon_0\chi_p\overrightarrow{E}_p = \mathcal{N}\overrightarrow{\mu}_{13}\rho_{31} = \mathcal{N}\overrightarrow{\mu}_{13}a_3a_1^*$, the susceptibility of probe field can be written as

$$\chi_p = \frac{\mathcal{N}\mu_{13}^2}{2\varepsilon_0\hbar}\chi, \tag{3}$$

where $\mathcal{N}$ denotes the atomic density, ε_0 is dielectric constant in vacuum, and χ is given by

$$\chi = -\frac{d_2d_4d_5 - d_5\Omega_{d1}^2 - d_4\Omega_{d2}^2}{d_4d_5\left(d_2d_3 - \Omega_c^2\right) - d_3\left(d_5\Omega_{d1}^2 + d_4\Omega_{d2}^2\right)}, \tag{4}$$

where $d_2 = \left(\Delta_p - \Delta_c\right) + i\gamma_2$, $d_3 = \Delta_p + i\frac{\Gamma_3}{2}$, $d_4 = \left(\Delta_p - \Delta_c + \Delta_{d1}\right) + i\frac{\Gamma_4}{2}$, $d_5 = \left(\Delta_p - \Delta_c + \Delta_{d2}\right) + i\frac{\Gamma_5}{2}$. In order to analyze the nonlinear modulation induced by the controlling fields with standing-wave pattern, we expand χ into its third-order parts which is expressed as

$$\chi = \chi^{(1)} + \chi_{d1}^{(3)}\Omega_{d1}^2 + \chi_{d2}^{(3)}\Omega_{d2}^2, \tag{5}$$

where $\chi^{(1)}$ ($\chi_{d1}^{(3)}$ and $\chi_{d2}^{(3)}$) denotes the first-order linear (third-order Kerr-nonlinear) part of the probe susceptibility given by

$$\chi^{(1)} = -\frac{d_2}{d_2d_3 - \Omega_c^2}, \tag{6a}$$

$$\chi_{d1}^{(3)} = -\frac{\Omega_c^2}{d_4\left(d_3d_2 - \Omega_c^2\right)^2}, \tag{6b}$$

$$\chi_{d2}^{(3)} = -\frac{\Omega_c^2}{d_5\left(d_3d_2 - \Omega_c^2\right)^2}. \tag{6c}$$

As we know, owing to the intensity-dependent susceptibility, the controlling fields with standing-wave pattern can lead to the spatially modulated absorption and refractivity for the probe field, thus the atoms act as a grating which can diffract the probe into different directions. In the following, we will analyze the physical mechanism of the grating. Two orthogonal standing-wave fields $\Omega_{d1}(x) = \Omega_1\sin\left(\pi x/\Lambda_x\right)$ and $\Omega_{d2}(y) = \Omega_2\sin\left(\pi y/\Lambda_y\right)$ are applied to replace the original two controlling fields Ω_{d1} and Ω_{d2}, which can make the formation of 2D EIG. Here, Λ_x and Λ_y are the spatial periods of corresponding standing-wave fields. Ω_1 and Ω_2 are the Rabi frequencies and we take $\Omega_1 = \Omega_2 = \Omega$ in the numerical calculations.

In order to obtain the diffraction pattern of the probe field, we begin with the Maxwell's equations. Here, we follow the method in Ref. [15]. In the slowly varying envelope approximation and steady state regime, the propagation equation of the probe reduces to

$$\frac{\partial E_p}{\partial z} = i\frac{\pi}{\lambda_p\varepsilon_0}P_p, \tag{7}$$

where λ_p is the wavelength of the probe field. By using $\overrightarrow{P}_p = \varepsilon_0\chi_p\overrightarrow{E}_p$, Eq. (7) can be written as

$$\frac{\partial E_p}{\partial z'} = i\chi E_p, \tag{8}$$