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Enhanced Kerr nonlinearity via quantum interference from spontaneous emission

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ABSTRACT

A novel atom configuration is proposed for a giant Kerr nonlinearity in zero linear and nonlinear probe absorption. It is shown that without coherent control field and just by quantum interference of spontaneous emission, a giant Kerr nonlinearity can be obtained.

1. Introduction

The interference of spontaneous emission channels refers to spontaneously generated coherence (SGC). Spontaneous emissions from a single excited state to two lower closely spaced levels V-type [1] or from two closely spaced upper levels to a common atomic ground state V-type [2] can interfere. In a cascade three-level system, SGC can also be created in a nearly spaced atomic levels case [3,4]. The existence of such coherence relies on the two atomic dipole matrix elements be non-orthogonal when the atom is placed in a free space. Some efforts have been made experimentally in the past two decades to generate SGC [5]. The first experimental investigation of spontaneous emission interference by using the sodium molecule carried out by Xia et al. [5]. It is believed that this type of coherence can alter some optical properties of atomic media. Although it is difficult to realize SGC experimentally in atomic systems, there are several proposals to investigate this effect. The effect of SGC on lasing without population inversion (LWI) [6], electromagnetically induced transparency (EIT) [7], optical bistability (OB) and optical multistability (OM) [8], and population inversion [9] have been extensively studied in a three-level or multilevel systems. Recently, it is shown that EIT has opened up a completely new route to achieving large optical nonlinearity [10,11]. EIT is also the mechanism underlying the recent experiments in ultra-slow group velocity of a probe pulse and therefore greatly increases the effective interaction time of the pulse with the medium [12]. These features enable one to use an EIT medium to achieve nonlinear optical processes at very low light intensities [13]. Moreover, EIT can significantly enhance the nonlinear interaction strength in multilevel atomic systems.

Recent studies have shown that Kerr nonlinearity can be used for quantum nondemolition measurements [14], quantum bit regeneration [15], quantum state teleportation [16], and the generation of the optical solitons [17]. Thus, it is suitable to have large third-order nonlinear susceptibilities under condition of slow light level and high sensitivities. This requires that the linear susceptibilities should be as small as possible for all pump and signal fields in order to minimize absorption loss. Note that Kerr nonlinearity corresponds to the refraction part of the third-order susceptibility in an optical medium. The large nonlinear susceptibility with small absorption causes the nonlinear optics to be studied at low light levels [13,18,19]. Since the ideal EIT regime dose not interact with the probe light, it cannot produce any nonlinear effects. However, it is shown that a large Kerr nonlinearity with vanishing absorption can be occurred in general three-level systems with spontaneously generated coherence (SGC) [20]. The effect of SGC on Kerr nonlinearity in a four-level atomic system has also been proposed [21]. The four-level EIT media are the convenience optical media to enhance the Kerr nonlinearity. Schmidt et al. [22] proposed a four-level N-type system to enhance the third order susceptibilities and, at the same time, completely suppressing the linear susceptibilities. Nakajima [23] found that the autoionizing off-resonance could lead to enhanced third-order susceptibilities. In another proposal, Niu et al. [24] predicted that in a four-level double-dark resonance system, the interacting double-dark resonance can effectively be enhanced the Kerr nonlinearity. The experimental observation of a large Kerr nonlinearity at low light intensities in four-level rubidium atoms has also been reported [25]. To the best of our knowledge, the giant Kerr nonlinearity in the absence of coherent control field is not reported for any atomic media.

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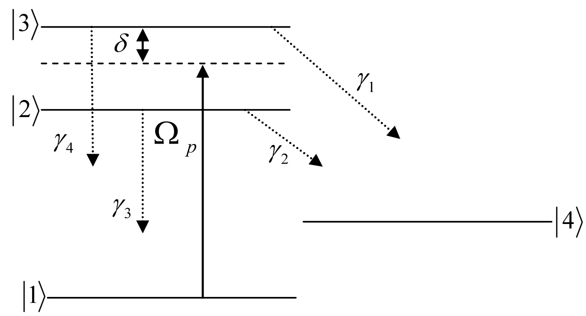

Fig. 1. Schematic diagram of proposed four-level atomic system is shown. In ^{87}Rb-D_1 line (nuclear spin $I = 3/2$) cold atoms, states $|1\rangle$ and $|4\rangle$ correspond to $|5S_{1/2}, F = 1, M_F = \pm 1\rangle$, states $|2\rangle$ and $|3\rangle$ correspond to $|5P_{1/2}, F = 1, M_F = 0\rangle$ and $|5P_{1/2}, F = 2, M_F = 0\rangle$, respectively.

In this Letter, we consider a double V-type of a four-level atom interacting by a weak probe field in the absence of coherent control field. The effects of SGC on linear and nonlinear susceptibility are investigated. It is shown that the SGC can dramatically reduce linear absorption and at the same time enhance the Kerr nonlinearity.

2. Model and equations

Consider a four-level atomic medium interacting by a probe laser field as depicted in Fig. 1. A probe laser field couple ground level $|1\rangle$ to upper levels $|2\rangle$ and $|3\rangle$. The spontaneous decay rates from levels $|3\rangle$ and $|2\rangle$ to level $|1\rangle$ denotes by γ_4 and γ_3, respectively. Also the spontaneous emission from levels $|3\rangle$ and $|2\rangle$ to level $|4\rangle$ are denotes by γ_1 and γ_2. The other decay rates are ignored. There are two major dynamical processes occurring in the system: (i) interaction with the reservoir governing the decay processes from levels $|3\rangle$ and $|2\rangle$ to the levels $|1\rangle$ and $|4\rangle$, (ii) interaction with the weak probe field. The processes are described by two interaction Hamiltonian terms H_1 and H_2 respectively. Thus including to the free energy term, the total Hamiltonian can be written as [26]

$$H = H_0 + H_1 + H_2, \tag{1}$$

with

$$H_0 = \hbar\omega_1|1\rangle\langle 1| + \hbar\omega_2|2\rangle\langle 2| + \hbar\omega_3|3\rangle\langle 3| + \hbar\omega_4|4\rangle\langle 4|, \tag{2a}$$

$$H_1 = -\hbar\big(\Omega_{p1}e^{-i\nu_p t}|2\rangle\langle 1| + \Omega_{p2}e^{-i\nu_p t}|3\rangle\langle 1|\big) + \text{h.c.}, \tag{2b}$$

$$\begin{aligned} H_2 = -\sum g_k^{(1)} e^{-i(\nu_k-\omega_{43})t}|4\rangle\langle 3|b_k + g_k^{(2)} e^{-i(\nu_k-\omega_{42})t}|4\rangle\langle 2|b_k \\ + g_k^{(3)} e^{-i(\nu_k-\omega_{21})t}|2\rangle\langle 1|b_k \\ + g_k^{(4)} e^{-i(\nu_k-\omega_{31})t}|3\rangle\langle 1|b_k + \text{h.c.} \end{aligned} \tag{2c}$$

where $\hbar\omega_i$ gives the energy of state $|i\rangle$ $(i = 1, 2, 3, 4)$. Here, Ω_{p1} and Ω_{p2} are the Rabi-frequency of the weak probe field, corresponding to the transitions $|3\rangle \to |1\rangle$ and $|2\rangle \to |1\rangle$. $b_k(b_k')$ is the annihilation (creation) operator for the kth vacuum mode with frequency ν_k; k here represents both the momentum and polarization of the vacuum mode. The density matrix equations of motion in the rotating wave approximation and in rotating frame can be written as

$$\begin{aligned} \frac{\partial\rho_{31}}{\partial t} &= -\left(\frac{1}{2}(\gamma_4+\gamma_1) + i\delta\right)\rho_{31} - i\Omega_{p1}(\rho_{33}-\rho_{11}) \\ &\quad - i\Omega_{p2}\rho_{32} - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})\rho_{21}, \end{aligned} \tag{3a}$$

$$\begin{aligned} \frac{\partial\rho_{21}}{\partial t} &= -\left(\frac{1}{2}(\gamma_3+\gamma_2) + i(\delta-\omega_{23})\right)\rho_{21} - i\Omega_{p1}\rho_{23} \\ &\quad + i\Omega_{p2}(\rho_{11}-\rho_{22}) - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})\rho_{31}, \end{aligned} \tag{3b}$$

$$\frac{\partial\rho_{41}}{\partial t} = -\left(\frac{1}{2}(\gamma_1+\gamma_2) + i\delta\right)\rho_{41} - i\Omega_{p1}\rho_{43} + i\Omega_{p2}\rho_{42}, \tag{3c}$$

$$\begin{aligned} \frac{\partial\rho_{42}}{\partial t} &= -\left[\frac{1}{2}(\gamma_3+\gamma_2) - i\Delta_2\right]\rho_{42} + i\Omega_2(\rho_{22}-\rho_{44}) \\ &\quad - i\Omega_{p2}\rho_{41} - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})\rho_{43}, \end{aligned} \tag{3d}$$

$$\begin{aligned} \frac{\partial\rho_{43}}{\partial t} &= -\left[\frac{1}{2}(\gamma_4+\gamma_1)\right]\rho_{43} - i\Omega_{p1}\rho_{41} \\ &\quad - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})\rho_{42}, \end{aligned} \tag{3e}$$

$$\frac{\partial\rho_{44}}{\partial t} = +\gamma_1\rho_{33} + \gamma_2\rho_{22} + \frac{1}{2}\beta\sqrt{\gamma_1\gamma_2}(\rho_{23}+\rho_{32}), \tag{3f}$$

$$\begin{aligned} \frac{\partial\rho_{32}}{\partial t} &= -\left[\frac{1}{2}(\gamma_3+\gamma_4+\gamma_1+\gamma_2) + i\omega_{23}\right]\rho_{32} + i\Omega_{p1}\rho_{12} \\ &\quad - i\Omega_{p2}\rho_{31} - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})(\rho_{33}+\rho_{22}), \end{aligned} \tag{3g}$$

$$\begin{aligned} \frac{\partial\rho_{33}}{\partial t} &= i\Omega_{p1}(\rho_{13}-\rho_{31}) - (\gamma_4+\gamma_1)\rho_{33} \\ &\quad - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})(\rho_{23}+\rho_{32}), \end{aligned} \tag{3h}$$

$$\begin{aligned} \frac{\partial\rho_{22}}{\partial t} &= i\Omega_{p2}(\rho_{12}-\rho_{21}) - (\gamma_3+\gamma_2)\rho_{22} \\ &\quad - \frac{1}{2}(P\sqrt{\gamma_3\gamma_4} + \beta\sqrt{\gamma_1\gamma_2})(\rho_{23}+\rho_{32}), \end{aligned} \tag{3i}$$

$$\begin{aligned} \frac{\partial\rho_{11}}{\partial t} &= i\Omega_{p2}(\rho_{12}-\rho_{21}) + i\Omega_{p1}(\rho_{13}-\rho_{31}) \\ &\quad + (\gamma_3+\gamma_4)\rho_{11} + \frac{1}{2}P\sqrt{\gamma_3\gamma_4}(\rho_{23}+\rho_{32}). \end{aligned} \tag{3j}$$

Here, detuning parameters are defined as $\delta = \Delta + \frac{\omega_{32}}{2}$, where parameter $\Delta = \nu_p - \frac{\omega_{21}+\omega_{31}}{2}$ measures the common detuning of probe field from the middle of doublet states. In addition, we assume that the dipole matrix elements for both transitions $|1\rangle \to |2\rangle$ and $|1\rangle \to |3\rangle$ are both equal, i.e. $|\vec{\wp}_{12}| = |\vec{\wp}_{13}| = \wp$. Thus, the Rabi-frequency of probe field is defined as $\Omega_{p1} = \Omega_{p2} = \Omega_p = \frac{\varepsilon_p\wp}{2\hbar}$. The parameter P $(= \frac{\vec{\wp}_{13}\cdot\vec{\wp}_{12}}{|\wp_{13}||\wp_{12}|})$ and β $(= \frac{\vec{\wp}_{34}\cdot\vec{\wp}_{24}}{|\wp_{34}||\wp_{24}|})$ denote the alignment of the matrix elements of four dipole moments. If the dipole moments of the two transitions are parallel or antiparallel, we have $P(\beta) = \pm 1$, while for the orthogonal case we have $P(\beta) = 0$. As pointed out above, the term $P\sqrt{\gamma_3\gamma_4}$ and $\beta\sqrt{\gamma_1\gamma_2}$ represent the effect of quantum interference between spontaneous emission pathways from $|3\rangle$ to $|1\rangle$ $(|4\rangle)$ and from $|2\rangle$ to $|1\rangle$ $(|4\rangle)$.

As is known, the response of the atomic medium to the weak probe field is governed by its polarization $P = \varepsilon_0(E_p\chi + E_p^*\chi^*)/2$ with χ being the susceptibility of the atomic medium. By performing a quantum average of the dipole moments over an ensemble of N atoms, it is found that $P = N(\wp_{13}\rho_{31} + \wp_{31}\rho_{13} + \wp_{12}\rho_{21} + \wp_{21}\rho_{12})$. In order to derive the linear and nonlinear susceptibility it is needed to obtain the steady state solution of density matrix equations. The density matrix elements can be expanded as $\rho_{ij} = \rho_{ij}^{(0)} + \rho_{ij}^{(1)} + \rho_{ij}^{(2)} + \cdots$. The zeroth order solution of ρ_{11} will be identical, i.e., $\rho_{11}^{(0)} = 1$, and other elements are set to be zero. The first and third order susceptibilities $\chi^{(1)}$ and $\chi^{(3)}$ of the medium can be determined by coherence terms $\rho_{31}^{(1)}$ and $\rho_{21}^{(1)}$ in Eq. (5a) and $\rho_{31}^{(3)}$, $\rho_{21}^{(3)}$ in Eq. (5b), respectively. For the analytical solution we set $\gamma_1 = \gamma_2 = \gamma_3 = \gamma_4 = \gamma$. The matrix elements ρ_{21} and ρ_{31} up to third order is obtained: