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Phase control of Kerr nonlinearity due to quantum interference in a four-level N-type atomic system

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ABSTRACT

Linear and nonlinear response of a four-level N-type atomic system for a weak probe field is investigated. It is demonstrated that the giant Kerr nonlinearity with reduced absorption can be achieved by the spontaneously generated coherence. In addition, the effect of a relative phase between coupling fields on linear and nonlinear absorption as well as Kerr nonlinearity is then discussed.

1. Introduction

Quantum coherence and interference are the basic mechanisms for controlling the optical properties of the medium. It is well known that the Kerr nonlinearity is determined by the real part of the third order susceptibility. All materials exhibit this effect, but the main difference is their extend of the nonlinear response to the applied fields. For many years, experimental research on quantum nonlinear optics has been limited because of the weak nonlinear response of even the best materials. Fortunately, the field of electromagnetically induced transparency (EIT) [1] leads to achieving large optical nonlinearity [2–4]. Schmidt and Imamoglu [5] proposed a scheme to achieve a large Kerr nonlinearity in an EIT medium. Several methods were described to enhancement the Kerr nonlinearity with reduced linear absorption in an EIT medium [6–12]. It has been observed that the optical properties of an atomic system can be controlled by coherent driving fields. Also, it is usually believed that the spontaneous emission destroy the atomic coherence. However, spontaneous emission can be used to produce atomic coherence as long as there exist two closely-lying levels with non-orthogonal dipoles in an atomic system. Atomic coherence based on spontaneous emission is usually referred to as vacuum-induced coherence or spontaneously generated coherence (SGC) [13]. One important application of quantum coherence and interference created by SGC is modification of linear and nonlinear absorption and dispersion [14]. The enhanced Kerr nonlinearity with zero absorption via SGC has also been investigated [15,16]. SGC along with EIT generally can increase the effective interaction time of the pulse with the medium leading to achieving of nonlinear optical processes at low light intensities. Thus, it is desirable to have large third order nonlinear susceptibilities under the condition of low light power and high sensitivities [17]. Giant Kerr nonlinearity may lead to many important applications in nonlinear phenomena, for example; multiple usages of giant Kerr nonlinearity in quantum information process (QIP) enable us to detect and resolve individual optical number states, such as quantum bit regeneration [18], long-distance quantum teleportation [19], Bell-state measurements [20], optical Fock state synthesis [21] and so on. In the present study, theoretical investigation is carried out on the third order nonlinear susceptibility of a four level atomic system. Recently, we (with collaborates) investigated the optical bistability and multi-stability of N-type atomic system in the presence of SGC [22]. Now, we intend to study the Kerr nonlinearity of the same atomic system via quantum interference arising from SGC. In fact, the effect of SGC on the enhancement of Kerr nonlinearity of the four-level N-type system is proposed. By proper choice of atomic parameters, the giant Kerr nonlinearity with reduced absorption can be achieved. It is found that in the presence of SGC the system completely becomes phase dependent, thus phase control of Kerr nonlinearity is discussed.

The paper is organized as follows: we present the model and equation of motions in Section 2. Analytical solution of density matrix equation of motion for the first and third order susceptibility is given in Section 3. Result and discussion are presented in Section 4, and the conclusion can be found in Section 5.

2. Model and equations of motion

Consider a four-level N-type atomic system as shown in Fig. 1. The intermediate level $|2\rangle$ is coupled to upper levels $|3\rangle$ and $|4\rangle$ by

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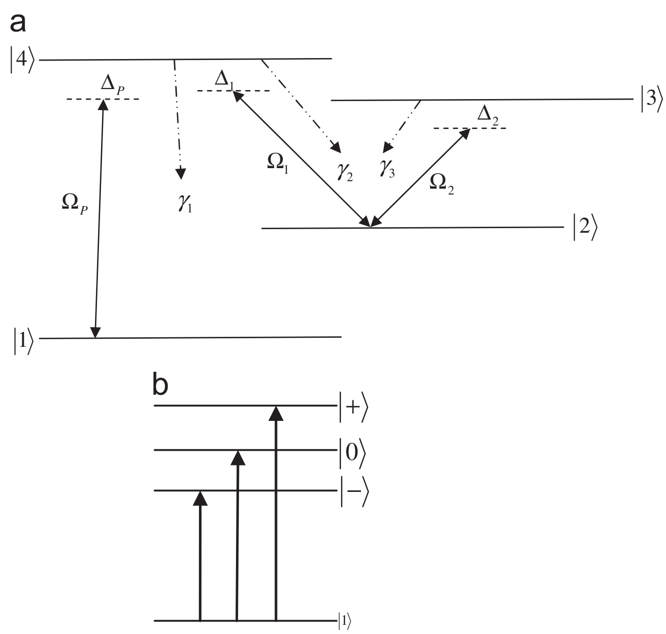

Fig. 1. (a) Configuration of a four level N-type atomic system with two coherent coupling fields, and a weak tunable probe field. (b) Dressed state configuration of four level N-type atomic system.

two coherent laser fields. Levels $|2\rangle$ and $|4\rangle$ are coupled by a control field with frequency υ_1 and Rabi-frequency $\Omega_1 = \frac{E_1 \cdot \vec{\wp}_{42}}{\hbar}$, while another control field with frequency υ_2 and Rabi-frequency $\Omega_2 = \frac{E_2 \cdot \vec{\wp}_{32}}{\hbar}$ couples levels $|2\rangle$ and $|3\rangle$. The ground level $|1\rangle$ is coupled to the exited level $|4\rangle$ by a weak tunable probe field with frequency υ_p and Rabi-frequency $\Omega_p = \frac{E_p \cdot \vec{\wp}_{41}}{\hbar}$, where $E_{1,2}(E_p)$ are the amplitudes of the coupling (probe) fields and $\vec{\wp}_{ij}$ denotes the atom dipole moments associated with the transition$|i\rangle \to |j\rangle$. The Rabi-frequencies are considered to be real, so we can define them as $\Omega_i \equiv \Omega_i e^{i\varphi_i}(i=1,2,p)$. The spontaneous decay rates from upper level $|4\rangle$ to lower levels $|1\rangle$ and $|2\rangle$ are denoted by γ_1 and γ_2, respectively. The corresponding decay rate from level $|3\rangle$ to level $|2\rangle$ is also denoted by γ_3. The transition between levels$|4\rangle$ and $|3\rangle$, levels $|2\rangle$ and $|1\rangle$ are electrically forbidden.

The dynamics of the system is described by the density matrix equation of motion

$$\dot{\rho} = -\frac{i}{\hbar}[H,\rho] + L_\rho \tag{1}$$

Here, H represents the interaction Hamiltonian of the system. Under the rotating wave approximation and in the rotating frame and L_ρ represents the atomic damping to background modes. The density matrix equations of motion can be written as [22]

$$\dot{\tilde{\rho}}_{21} = -i(\Delta_p - \Delta_1)\tilde{\rho}_{21} + \frac{i}{2}\Omega_2\tilde{\rho}_{31} - \frac{i}{2}\Omega_p\tilde{\rho}_{24} + \frac{i}{2}\Omega_1\tilde{\rho}_{41}.$$

$$\dot{\tilde{\rho}}_{22} = \gamma_2\tilde{\rho}_{44} + \gamma_3\tilde{\rho}_{33} + \frac{i}{2}\Omega_2\tilde{\rho}_{32} + \frac{i}{2}\Omega_1\tilde{\rho}_{42} - \frac{i}{2}\Omega_1\tilde{\rho}_{24} - \frac{i}{2}\Omega_2\tilde{\rho}_{23} + P\sqrt{\gamma_2\gamma_3}(\tilde{\rho}_{34}e^{i\phi}e^{i\omega_{43}t} + \tilde{\rho}_{43}e^{-i\phi}e^{-i\omega_{43}t}).$$

$$\dot{\tilde{\rho}}_{11} = \gamma_1\tilde{\rho}_{44} + \frac{i}{2}\Omega_p\tilde{\rho}_{41} - \frac{i}{2}\Omega_p\tilde{\rho}_{14}.$$

$$\dot{\tilde{\rho}}_{41} = -\left(i\Delta_p + \frac{1}{2}(\gamma_1+\gamma_2)\right)\tilde{\rho}_{41} + \frac{i}{2}\Omega_1\tilde{\rho}_{21} + \frac{i}{2}\Omega_p(\tilde{\rho}_{11} - \tilde{\rho}_{44}) - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}\tilde{\rho}_{31}e^{i\phi}e^{i\omega_{43}t}$$

$$\dot{\tilde{\rho}}_{42} = -\left(i\Delta_1 + \frac{1}{2}(\gamma_1+\gamma_2)\right)\tilde{\rho}_{42} + \frac{i}{2}\Omega_p\tilde{\rho}_{12} + \frac{i}{2}\Omega_1(\tilde{\rho}_{22} - \tilde{\rho}_{44}) - \frac{i}{2}\Omega_2\tilde{\rho}_{43} - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}\tilde{\rho}_{32}e^{i\phi}e^{i\omega_{43}t}.$$

$$\dot{\tilde{\rho}}_{43} = \left(i(\Delta_2 - \Delta_1) - \frac{1}{2}(\gamma_1+\gamma_2+\gamma_3)\right)\tilde{\rho}_{43} + \frac{i}{2}\Omega_p\tilde{\rho}_{13} + \frac{i}{2}\Omega_1\tilde{\rho}_{23} - \frac{i}{2}\Omega_2\tilde{\rho}_{42} - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}(\tilde{\rho}_{33} + \tilde{\rho}_{44})e^{i\phi}e^{i\omega_{43}t}.$$

$$\dot{\tilde{\rho}}_{31} = -\left(i(\Delta_p + \Delta_2 - \Delta_1) + \frac{1}{2}\gamma_3\right)\tilde{\rho}_{31} + \frac{i}{2}\Omega_2\tilde{\rho}_{21} - \frac{i}{2}\Omega_p\tilde{\rho}_{34} - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}\tilde{\rho}_{41}e^{-i\phi}e^{-i\omega_{43}t}.$$

$$\dot{\tilde{\rho}}_{32} = -\left(i\Delta_2 + \frac{1}{2}\gamma_3\right)\tilde{\rho}_{32} - \frac{i}{2}\Omega_1\tilde{\rho}_{34} + \frac{i}{2}\Omega_2(\tilde{\rho}_{22} - \tilde{\rho}_{33}) - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}\tilde{\rho}_{42}e^{-i\phi}e^{-i\omega_{43}t}.$$

$$\dot{\tilde{\rho}}_{33} = -\gamma_3\tilde{\rho}_{33} + \frac{i}{2}\Omega_2\tilde{\rho}_{23} - \frac{i}{2}\Omega_2\tilde{\rho}_{32} - \frac{1}{2}P\sqrt{\gamma_2\gamma_3}(\tilde{\rho}_{43}e^{-i\phi}e^{-i\omega_{43}t} + \tilde{\rho}_{34}e^{i\phi}e^{i\omega_{43}t}) \tag{2}$$

The above equations are constrained by $\tilde{\rho}_{11} + \tilde{\rho}_{22} + \tilde{\rho}_{33} + \tilde{\rho}_{44} = 1$ and $\tilde{\rho}^*_{ji} = \tilde{\rho}_{ij}$.

The detuning parameters are defined as $\Delta_1 = \omega_{42} - \upsilon_1$, $\Delta_2 = \omega_{32} - \upsilon_2$, and $\Delta_p = \omega_{41} - \upsilon_p$, where ω_{ij} is the frequency deference between level $|i\rangle$ and level $|j\rangle$. The parameter $P(=((\vec{\wp}_{42} \cdot \vec{\wp}_{32})/(|\vec{\wp}_{42}||\vec{\wp}_{32}|)) = \cos\theta)$ arises due to the quantum interference between two decay channels $|4\rangle \to |2\rangle$ and$|3\rangle \to |2\rangle$. In fact, the parameter P denotes the alignment of the two dipole moments $\vec{\wp}_{42}$ and $\vec{\wp}_{32}$, which represent the strength of the interference in spontaneous emission. Here θ is the angle between two induced dipole moments $\vec{\wp}_{42}$ and $\vec{\wp}_{32}$. If two dipole moments are orthogonal, there is no interference due to spontaneous emission and $P=0$, whereas for the parallel dipole moments, the effect of quantum interference is maximal and $P=1$. By considering $\Omega_p \equiv \Omega_p e^{i\varphi_p}$, $\Omega_1 \equiv \Omega_1 e^{i\varphi_1}$, $\Omega_2 \equiv \Omega_2 e^{i\varphi_2}$,we can obtain new equations for the redefined density matrix elements which are found to be identical to Eq. (2), with the interference parameter P replaced by $P = Pe^{i\phi}$. It means the change of phase difference between couplings fields may change the direction of the dipole moments; thus it changes the parameter P. Therefore, the linear and nonlinear behavior of medium can be changed via relative phase ϕ.

Note that in the case of closely two upper levels $|3\rangle$ and$|4\rangle$, the frequency difference should be zero, i.e. $\omega_{43} \simeq 0$. So, the time dependent exponential terms should be identical. Eq. (2) is obtained by using the transmissions$\rho_{43} = \tilde{\rho}_{43}e^{i(\Delta_2 - \Delta_1)t}e^{-i\phi}$, $\rho_{32} = \tilde{\rho}_{32}e^{-i\Delta_2 t}e^{-i\varphi_2}$, $\rho_{42} = \tilde{\rho}_{42}e^{-i\Delta_1 t}e^{-i\varphi_1}$, $\rho_{31} = \tilde{\rho}_{31}e^{i(\Delta_1 - \Delta_2 - \Delta_p)t}e^{i(\phi - \varphi_p)}$ and $\rho_{41} = \tilde{\rho}_{41}e^{-i\Delta_p t}e^{-i\varphi_p}$ and assuming the frequency difference $(v_3 - v_2)$ is small, so that $e^{\pm i(v_3 - v_2)} \simeq 1$.

3. Analytical solution

Now, we derive analytical expressions for the first and third order susceptibilities. In order to obtain the linear and nonlinear susceptibilities, we need to solve the steady state solution of the density matrix equations. Assuming that the probe field is much weaker than the other fields, density matrix elements are expanded as $\rho_{ij} = \tilde{\rho}^{(0)}_{ij} + \tilde{\rho}^{(1)}_{ij} + \tilde{\rho}^{(2)}_{ij} + \tilde{\rho}^{(3)}_{ij} + \ldots$. The 0th order solution of $\tilde{\rho}^{(0)}_{11}$ will be identical, i.e. $\tilde{\rho}^{(0)}_{11} = 1$, and other elements are set to be zero. We consider the resonance condition, i.e. $\Delta_1 = \Delta_2 = 0$. Therefore, the first and third-order susceptibilities $\chi^{(1)}$ and $\chi^{(3)}$