Light controlling light in a coupled atom-cavity system

Bichen Zou and Yifu Zhu

Department of Physics, Florida International University, Miami, Florida 33199, USA

(Received 7 November 2012; revised manuscript received 17 February 2013; published 2 May 2013)

We present a theoretical study of a multiatom cavity quantum electrodynamics system consisting of three-level atoms confined in a cavity and interacting with a control laser from free space. A probe laser is coupled into the cavity, and the transmitted cavity field and reflected probe field from the cavity are analyzed. We show that the free-space control laser induces large nonlinearities for the intracavity probe field, which can be used to control the amplitude and phase of the reflected probe light and the transmitted probe light. Under appropriate conditions, large cross-phase modulation for the reflected and transmitted light fields can be obtained with a weak control field, which renders the system useful for studies of all-optical switching and quantum phase gates.

DOI: 10.1103/PhysRevA.87.053802 PACS number(s): 42.50.Pq, 42.50.Gy, 32.80.Qk, 42.65.−k

I. INTRODUCTION

Optical cavity can be used to boost the interaction strength of the matter and the light field and has been used in diverse areas of research and applications such as ultrasensitive spectroscopic measurements, optical frequency metrology, and nonlinear optical devices [1]. Fundamental studies of the two-level atom and photon interactions in a cavity has developed into a research field of cavity quantum electrodynamics (cavity QED), which is useful in a variety of applications in quantum physics and quantum electronics [2–4].

It is interesting to extend studies of cavity QED to the interactions of the cavity mode and multilevel atoms. There were earlier studies of optical bistability in three-level atoms confined in an optical cavity [5,6]. Recent studies of atom-cavity interactions have been directed to a composite system of an optical cavity and coherently prepared multilevel atoms, in which the atomic coherence and interference such as electromagnetically induced transparency (EIT) [7] can be used to manipulate quantum states of the coupled cavity and atom system [8–12]. It has been shown that in a coherently coupled cavity and multiatom system, the interplay of the collective coupling of the atoms and the cavity mode, and the atomic coherence and interference manifested by EIT may lead to interesting linear and nonlinear optical phenomena [13–21].

Here we report a theoretical study of an atom-cavity system consisting of N three-level atoms confined in an optical cavity and coherently coupled from free space by a control laser. The system forms a Λ-type standard EIT configuration with the cavity mode. We show that the free-space control laser induces large nonlinearities for the intracavity probe field, which can be used to control both the amplitude and phase of the cavity-reflected probe light and the cavity-transmitted probe light. Under appropriate conditions, large cross-phase modulation (XPM) for the reflected and transmitted light fields can be obtained with a weak control field. The cavity-atom system provides an example of resonant nonlinear optics at low light intensities and can be used for fundamental studies of light control light in such applications as all-optical switching and quantum phase gates.

II. THEORETICAL ANALYSIS

A composite atom-cavity system that consists of a single mode cavity containing N Λ-type three-level atoms interacting with a control laser from free space is shown in Fig. 1. The cavity mode couples the atomic transition $|1\rangle - |3\rangle$ and the classical control laser drives the atomic transition $|2\rangle - |3\rangle$ with Rabi frequency 2Ω.$\Delta = \nu - \nu_{23}$ is the control frequency detuning and $\Delta_c = \nu_c - \nu_{13}$ is the cavity-atom detuning. A probe laser is coupled into the cavity mode and its frequency is detuned from the atomic transition $|1\rangle - |3\rangle$ by $\Delta_p = \nu_p - \nu_{13}$. The interaction Hamiltonian for the coupled cavity-atom system is

$$H = -\hbar\left(\sum_{i=1}^{N}\Omega\hat{\sigma}_{32}^{(i)} + \sum_{i=1}^{N} g\hat{a}\hat{\sigma}_{31}^{(i)}\right) + \text{H.c.}, \tag{1}$$

where $\hat{\sigma}_{lm}^{(i)}$ $(l,m = 1\text{–}3)$ is the atomic operator for the ith atom, $g = \mu_{13}\sqrt{\omega_c/2\hbar\varepsilon_0 V}$ is the cavity-atom coupling coefficient, and $\hat{a}$ is the annihilation operator of the cavity photons. The resulting operator equations of motion for the intracavity light field (two-sided cavity, one input) is given by [22]

$$\dot{\hat{a}} = -\frac{i}{\hbar}[\hat{a},H] - \frac{\kappa_1+\kappa_2}{2}\hat{a} + \sqrt{2\kappa_1/\tau}\,\hat{a}_p^{\text{in}} + \tilde{F}(t), \tag{2}$$

where $\hat{a}_p^{\text{in}}$ is the input probe field, $\kappa_i = T_i/2\tau (i = 1\text{–}2)$ is the loss rate of the cavity field of the mirror i (T_i is the mirror transmission and τ is the photon round-trip time inside the cavity). $\tilde{F}(t)$ is the noise term. The equation of motion for the expectation value of the intracavity probe field then is

$$\dot{a} = -[(\kappa_1+\kappa_2)/2 - i\Delta_c]a + \sum_{i=1}^{N} ig\sigma_{31}^{(i)} + \sqrt{2\kappa_1/\tau}\,a_p^{\text{in}}. \tag{3}$$

For simplicity, we consider a symmetric cavity such that $\kappa_1 = \kappa_2 = \kappa/2$. With $g \ll \Omega$, the atomic population is concentrated in $|1\rangle$ and the steady-state solution of the intracavity probe field is given by

$$a = \frac{\sqrt{\kappa/\tau}\,a_p^{\text{in}}}{\kappa - i\Delta_c - i\chi}, \tag{4}$$

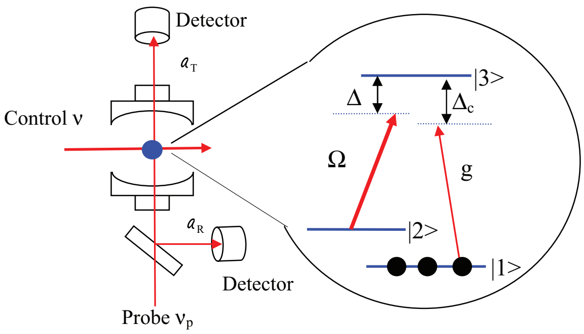

FIG. 1. (Color online) Schematic coupling scheme of coherently coupled four-level atoms in a cavity. A control laser drives the $|2\rangle-|3\rangle$ transition with Rabi frequency 2Ω. Δ is the control detuning. A cavity is coupled to the atomic transition $|1\rangle-|3\rangle$ with the collective coupling coefficient $\sqrt{N}g$ (Δ_c is the cavity-atom detuning). A probe laser is coupled into the cavity and Δ_p is its frequency detuning from the atomic transition. The cavity-transmitted probe light and the cavity-reflected probe light are collected by two detectors.

where χ is the atomic susceptibility given by $\chi = ig^2N/\{\Gamma_3 - i\Delta_p + \Omega^2/[\gamma_{12} - i(\Delta_p - \Delta)]\}$. Here Γ_3 is the spontaneous decay rate of the excited state $|3\rangle$ and γ_{12} is the decoherence rate between the ground states $|1\rangle$ and $|2\rangle$. The transmitted probe field is then given by $a_T = \sqrt{\kappa\tau}a$ and the reflected probe field from the cavity is $a_R = \sqrt{\kappa\tau}a - a_p^{\text{in}}$. We are interested in the regime of parameters near cavity EIT in which the laser fields are near or at resonance with the atomic transitions, and under the conditions of low light intensities in which the intracavity field is very weak and the control field is well below the saturation level. We show that a weak control light can induce large nonlinearities in the cavity-atom system, which then may be used to control the amplitude and phase of the cavity-transmitted and the cavity-reflected weak probe field. The primary control parameters of the light-control-light system are the frequency and intensity of the control laser that are characterized by the control detuning Δ and control Rabi frequency Ω, respectively. Therefore we apply the resonance condition for the probe laser and the cavity ($\Delta_c = \Delta_p = 0$), and derive the analytical results with $\gamma_{12} = 0$, neglecting the decoherence. (The effect of decoherence with $\gamma_{12} \neq 0$ will be considered in the numerical calculations.) Specifically, the light field transmitted through the cavity is

$$a_T = \frac{\kappa(\Omega^2 + i\Delta\Gamma_3)a_p^{\text{in}}}{\kappa\Omega^2 + i(g^2N + \kappa\Gamma_3)\Delta} = \alpha_T e^{i\Phi_T} a_p^{\text{in}}. \tag{5}$$

The amplitude ratio of the transmitted probe field to the input probe field is given by

$$\alpha_T = \frac{\kappa[\Omega^4 + (\Delta\Gamma_3)^2]^{1/2}}{\{(\kappa\Omega^2)^2 + [(g^2N + \kappa\Gamma_3)\Delta]^2\}^{1/2}}, \tag{6}$$

and the phase shift Φ_T is given by

$$\Phi_T = \tan^{-1}\frac{(\Gamma_3\kappa - g^2N)\Omega^2\Delta}{g^2N\Delta^2\Gamma_3 + \kappa\Omega^4}. \tag{7}$$

Under the condition of the strong collective coupling ($g^2N \gg \kappa\Gamma_3$) and for a weak control field ($\Omega^2 \ll \Delta\Gamma_3$ with $\Delta < \Gamma_3$), the phase shift of the probe field is $\Phi_T = -\Omega^2/\Delta\Gamma_3$, and the amplitude ratio of the transmitted field to the input probe field is $\alpha_T = \kappa\Gamma_3/g^2N$. The phase shift is proportional to the control laser intensity and corresponds to the standard Kerr nonlinearity. Under the strong collective coupling condition, a large phase shift of the transmitted field can be obtained but the transmitted field amplitude is very small.

The reflected probe field from the cavity is

$$a_R = \frac{g^2N\Delta a_p^{\text{in}}}{i\kappa\Omega^2 - (g^2N + \kappa\Gamma_3)\Delta} = \alpha_R e^{i\Phi_R} a_p^{\text{in}}. \tag{8}$$

The amplitude ratio of the reflected probe field to the input probe field is given by

$$\alpha_R = \frac{g^2N\Delta}{[(\kappa\Omega^2)^2 + (g^2N + \kappa\Gamma_3)^2\Delta^2]^{1/2}}, \tag{9}$$

and the phase shift Φ_R of the reflected probe field is

$$\Phi_R = \tan^{-1}\frac{\kappa\Omega^2}{(g^2N + \kappa\Gamma_3)\Delta}. \tag{10}$$

As an example, with a weak, near resonant control ($\Omega = 0.5\Gamma_3$ and $\Delta = 0.01\Gamma_3$), $g\sqrt{N} = 10\Gamma_3$, and $\kappa = 2\Gamma_3$, the amplitude ratio of the reflected probe field to the input probe field is ~79% and the XPM phase shift of the reflected field from the control laser is ~0.5 rad. In the limit of $\kappa\Omega^2 \ll |(g^2N + \kappa\Gamma_3)\Delta|$, one derives $\alpha_R = g^2N/(g^2N + \kappa\Gamma_3)$ and $\Phi_R \approx \kappa\Omega^2/(g^2N + \kappa\Gamma_3)\Delta$, which is the phase shift from the Kerr nonlinearity induced by the weak control field on the reflected cavity field. The reflected field amplitude can be nearly equal to the input probe field when the condition of strong collective coupling is satisfied ($g^2N \gg \kappa\Gamma_3$), but the phase shift Φ_R approaches zero. Therefore the optimal performance of the cavity system for the cross-phase modulation is achieved for moderate g^2N values. As will be further clarified in the numerical calculations that by inducing cavity EIT in the coherently coupled cavity-atom system, the XPM can be obtained near the atomic resonance with the control detuning $\Delta \ll \Gamma$, which leads to a larger Φ_R value, but at the same time still maintains a sufficiently large amplitude of the transmitted and reflected fields. This is in sharp contrast with the conventional XPM systems in which the laser fields have to be tuned far away from the atomic resonance to minimize the absorption loss.

III. NUMERICAL CALCULATIONS

In order to provide a detailed picture for the study of light control light in the coherently coupled cavity-atom system, we present numerical calculations in Figs. 2–5 and show that practical parameters can be identified, in which the transmitted or reflected fields with large phase shifts and sufficiently large amplitudes can be obtained.

Figures 2(a) and 2(b) plot the intensity ratio of the cavity-transmitted field $I_{\text{trans}}/I_{\text{in}} = \alpha_T^2$ and the phase shift Φ_R versus the probe frequency detuning $\Delta p/\Gamma_3$, respectively. Figures 2(c) and 2(d) plot the intensity ratio of the cavity-reflected field $I_{\text{Ref}}/I_{\text{in}} = \alpha_R^2$ and the phase shift Φ_T versus $\Delta p/\Gamma$, respectively. The relevant parameters are $g\sqrt{N} = 10\Gamma_3$,$\kappa = 2\Gamma_3$, $\gamma_{12} = 0.001\Gamma_3$, and $\Delta_c = \Delta = 0$. The transmission spectrum in Figs. 2(a) and 2(b) exhibit the standard cavity EIT spectral peak at $\Delta p = 0$ with a peak linewidth that

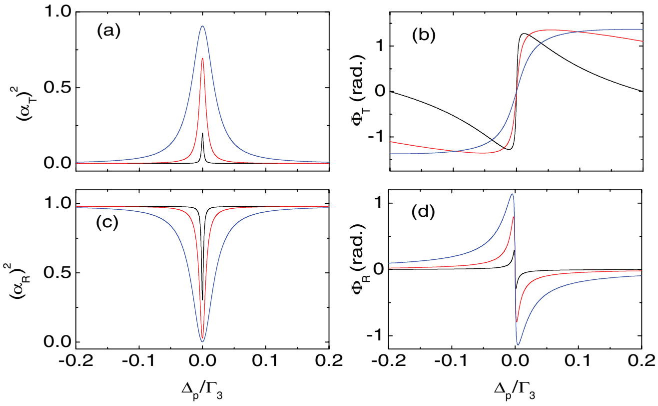

FIG. 2. (Color online) (a) The intensity ratio α_T^2 and (b) phase shift Φ_T of the cavity-transmitted probe field versus the probe frequency detuning $\Delta p/\Gamma_3$. (c) The intensity ratio α_R^2 and (d) phase shift Φ_R of the cavity-reflected probe field versus the probe frequency detuning $\Delta p/\Gamma_3$. The spectral features at $\Delta p = 0$ represent the cavity EIT. The control Rabi frequency $\Omega = 0.2\Gamma_3$, $0.5\Gamma_3$, and Γ_3 for the black lines, red lines, and blues lines, respectively.

is substantially smaller than the decay width Γ_3 and increases with the increasing Ω of the control laser. The phase shift of the transmitted field varies rapidly across the resonance at $\Delta p = 0$ and has a line shape of the normal dispersion. Concomitantly, the reflected field from the cavity exhibits a dip with the linewidth matching the peak linewidth of the transmitted field and its phase shift has a line shape of the anomalous dispersion with a steep slope across $\Delta p = 0$.

We next show that the coupled cavity-atom system can be used for studies of light control light. Specifically, a weak control light may induce large optical nonlinearities for the intracavity probe field near the frequency of cavity EIT, which can be used to control the amplitude and phase of both the transmitted and reflected light fields from the cavity. Figures 3(a) and 3(b) plot the intensity ratio α_T^2 and the phase shift Φ_T of the cavity-transmitted probe field versus the control frequency detuning Δ/Γ_3, respectively. Figures 3(c) and 3(d) plot the intensity ratio α_R^2 and the phase shift Φ_R of the cavity-reflected probe field versus Δ/Γ_3, respectively. The calculations of Fig. 3 are done with the probe laser tuned to the resonance $\Delta p = 0$ and show that the transmitted field and the reflected field can be controlled by varying the frequency of the control laser. For example, Figs. 3(a) and 3(c) show that with a resonant control laser ($\Delta = 0$), the transmitted light field can be turned on or off by the turning on or off of the control light, and at the same time and in contrast, the reflected light field is turned off or on. Thus the coupled cavity-atom system effectively performs all-optical switching with reciprocal functions for the transmission and reflection with a weak control laser. When the control laser is tuned slightly away from the resonance ($\Delta \neq 0$ but $\Delta \ll \Gamma_3$), large phase shifts are induced on the transmitted field and the reflected field [Figs. 3(b) and 3(d)], which effectively performs the XPM on the transmitted and reflected fields with a weak control light near the atomic resonance.

It is necessary to quantify the XPM performance of the cavity-atom system by comparing the phase shift and the field amplitude. A useful system has to be able to produce large

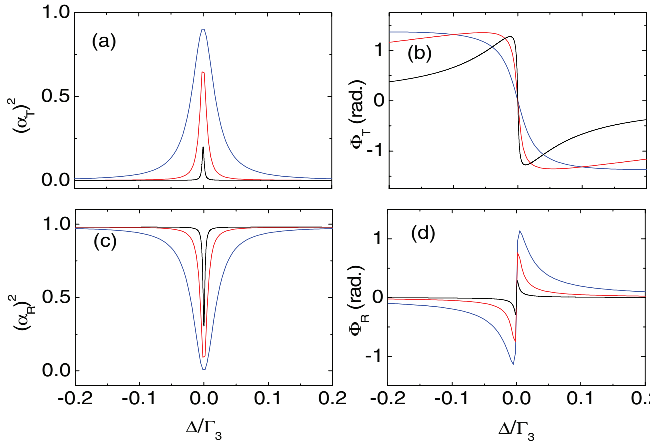

FIG. 3. (Color online) (a) The intensity ratio α_T^2 and the phase shift Φ_T of the cavity-transmitted field versus the control frequency detuning Δ/Γ_3. (c) The intensity ratio α_R^2 and (d) the phase shift Φ_R of the cavity-reflected field versus the control frequency detuning Δ/Γ_3. The parameters used in the calculations are $g\sqrt{N} = 10\Gamma$, $\kappa = 2\Gamma_3$, $\gamma_{12} = 0.001\Gamma$, and $\Delta_c = \Delta_p = 0$. The control Rabi frequency $\Omega = 0.2\Gamma$, 0.5Γ, and Γ for the black lines, red lines, and blues lines, respectively.

XPM phase shifts and at the same time, provides sufficiently large field amplitudes. For studies of nonlinear optics at low light intensities, such requirements have to be obtained at the condition of low control field intensities. Figure 4 plots α_T^2,α_R^2, Φ_T, and Φ_R versus the square of the control Rabi frequency Ω^2, which is proportional to the control light intensity [the control light intensity is given by $I = c\varepsilon_0(h\Omega/\mu_{23})^2$]. It can be seen that for the reflected field, the XPM phase shift increases with the control light intensity and exhibits saturation at high control intensities, while the reflected light intensity decreases with the control intensity and approaches zero at high control intensities. For comparison, the green curve in Fig. 4(c)

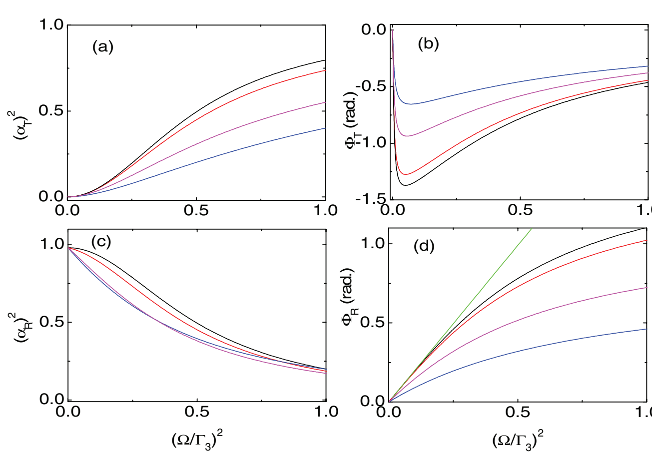

FIG. 4. (Color online) (a) The intensity ratio α_T^2 and (b) phase shift Φ_T of the cavity-transmitted probe field versus $(\Omega/\Gamma_3)^2$. (c) The intensity ratio α_R^2 and (d) phase shift Φ_R of the cavity-reflected field versus $(\Omega/\Gamma_3)^2$. The parameters used in the calculations are $g\sqrt{N} = 10\Gamma_3$, $\kappa = 2\Gamma_3$, $\Delta = 0.01\Gamma_3$, and $\Delta_c = \Delta_p = 0$. The decoherence rate $\gamma_{12} = 0$, $0.001\Gamma_3$, $0.005\Gamma_3$, and $0.01\Gamma_3$ for the black lines, red lines, purple lines, and blues lines, respectively.

plots the phase shift given by $\Phi_R = \kappa\Omega^2/(g^2N + \kappa\Gamma_3)\Delta$ [Eq. (10) with $\kappa\Omega^2 \ll |(g^2N + \kappa\Gamma_3)\Delta|$]. For the transmitted field, the XPM phase shift exhibits rapid change at low control intensities near zero and also saturates at high control intensities, while the transmitted intensity increases with the control intensity.

As shown in Fig. 4 the ground state decoherence degrades the system performance of the light control light. γ_{12} ultimately determines the minimum linewidth of the cavity EIT and also the lower limit of the control Rabi frequency Ω. As the decoherence rate γ_{12} increases, the amplitudes of the transmitted field and the reflected field decrease; the phase shifts of the transmitted field and the reflected field also decrease. It is preferable to select an atomic system with a small γ_{12}. We note that in cold Rb atoms, the decoherence rate $\gamma_{12} = 10^{-4}\Gamma_3$ has been observed [23]. Therefore, a possible experimental system may consist of cold Rb atoms ($\Gamma_3 \approx 6$ MHz) and a cavity with a moderate finesse (with $\kappa = 2\Gamma_3$ and a cavity length of 5 cm, the required finesse is $F = 124$). With a weak control laser, the linewidth of the cavity EIT is much smaller than Γ_3 (with $\Omega = 0.5\Gamma_3$ and $\gamma_{12} = 0.001\Gamma_3$, the linewidth of the cavity EIT is $\sim 0.01\Gamma_3$), it is necessary to have the control laser and the probe laser with a linewidth much smaller than $0.01\Gamma_3$. For the Rb atoms, the required laser linewdth is in the KHz range, which can be obtained with the frequency stabilized diode lasers.

Figures 4(a) and 4(c) show that at low control intensities, the reflected field amplitude is greater than the transmitted field amplitude. To achieve a large field amplitude, it is desirable to operate the cavity-atom system with the reflected light field. At low control intensities, the XPM phase shift is proportional to the control intensity, which is in the Kerr nonlinearity regime. A large Φ_R value can be obtained at a sufficiently large α_R^2 value with a weak control light. As a numerical example, in a cavity-atom system with $\gamma_{12} = 0.001\Gamma_3$, the XPM phase shift $\Phi_R \sim 0.5$ rad, and the reflected intensity ratio $\alpha_R^2 \sim 70\%$ can be obtained with practical parameters $\Omega = 0.5\Gamma_3$, $g\sqrt{N} = 10\Gamma_3$, $\kappa = 2\Gamma_3$, $\Delta = 0.01\Gamma_3$, and $\Delta_c = \Delta_p = 0$.

Figure 5 presents the parametric plotting of the phase shift versus the intensity ratio for both the transmitted probe field

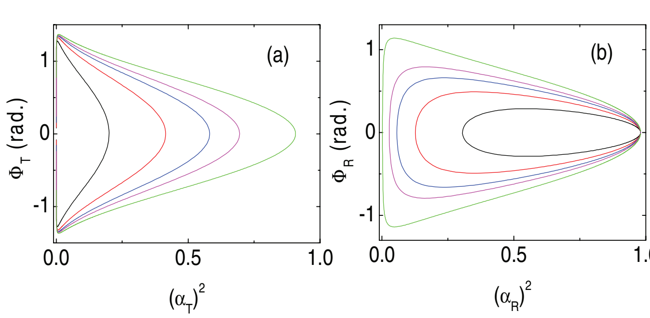

FIG. 5. (Color online) (a) Phase shift Φ_T of the cavity-transmitted field versus the intensity ratio α_T^2 of the transmitted field and (b) phase shift Φ_R of the cavity-reflected field versus the intensity ratio α_R^2 of the cavity-reflected field (the results are obtained versus the control laser detuning Δ). The parameters used in the calculations are the same as in Fig. 3. The control Rabi frequencies are $\Omega = 0.2\Gamma$, 0.3Γ, 0.4Γ, 0.5Γ, and Γ for the black lines, red lines, blue lines, purple lines, and green lines, respectively.

and the reflected probe field. The calculations are done by varying the control detuning Δ with the same set of parameters as in Fig. 3. It shows that with a weak control laser, the reflected probe field amplitude typically has larger amplitude than the transmitted probe field.

The nonlinear optics at low light intensities requires a weak control field below the saturation intensity ($\Omega < \Gamma$). The presented calculations are valid under the condition $\rho_{11} \sim 1$, i.e., the atomic population is concentrated in $|1\rangle$, which requires $g \ll \Omega$. Therefore, the cavity-atom coupling coefficient $g \ll \Gamma$; consequently, this is the weak-coupling regime of the cavity QED with single atoms (the bad cavity regime).

It is interesting to compare the light-control-light scheme based on the cavity EIT here with the earlier published scheme based on the cavity polariton interference [19,20]. The cavity polariton scheme also works with three-level atoms, but the free-space control laser and the probe laser are tuned to the polariton resonances (the normal mode) that are separated from the respective atomic resonances by the vacuum Rabi splitting. The control laser induces the destructive interference for the polariton excitation and renders the cavity-atom system opaque to the probe light [20]. For the cavity EIT scheme presented here, the control laser and the probe laser are tuned to the atomic resonant transitions of the three-level system and create the EIT condition. EIT renders the medium transparent to the probe laser. Therefore, the scheme in Refs. [19] and [20] is based on electromagnetically induced opaque in the normal mode excitation, while the scheme here is based on electromagnetically induced transparency. The two schemes result in comparable phase shifts and amplitudes for the transmitted field under the conditions of a weak control laser and moderate optical density.

Our earlier study of the cavity-atom polariton system only analyzed the cavity-transmitted field. Here the analysis is carried out for both the cavity-transmitted field and the cavity-reflected field. The phase shift and amplitude of the cavity-reflected field as well as the cavity-transmitted field are analyzed and compared. The results show that the cavity-reflected field and cavity-transmitted field are complementary in their performance merits. The cavity system provides the versatility of choosing either the reflected field or transmitted field for the studies of all-optical switching and the cross-phase modulation according to their respective performance merits in a given parameter regime.

There have been several studies of the EIT enhanced Kerr nonlinearities and XPM at low light intensities in multilevel atomic systems in recent years [24–31]. Most of these studies were done in four-level N-type systems coupled by three laser fields in free space. They can be divided into four types: a four-level N type coupled by three laser fields (Refs. [23–25]), a five-level M type coupled by four laser fields (Ref. [26]), a four-level tripod type coupled by three laser fields (Refs. [27,28]), and a modified five-level tripod type coupled by four laser fields (Ref. [29]). Reference [25] proposed to inject a signal light (that induces the Kerr nonlinearity on a free-space probe light) into a Michelson-type double cavity containing the four-level atoms and therefore increase its coupling time with the free-space probe light. The double cavity is a passive device used solely for the signal light. Reference [31] explores basically the N type in the time domain through the EIT light storage process. Reference [32] deals with the self-Kerr nonlinearity (self-phase modulation) in the three-level Λ-type system confined in a traveling-wave cavity. In contrast with the EIT schemes in free space with multilevel atoms coupled by at least three laser fields, our scheme requires only two laser fields coupled with three-level atoms, but needs to operate in a cavity (it differs from Ref. [26] in that the probe light is coupled into the cavity and the cavity QED effect measured by the vacuum Rabi splitting plays an essential role). Due to the cavity feedback enhancement, our scheme produces a large XPM phase shift ($\sim$0.5 rad and with a transmitted light intensity about 50% of the input light intensity under practical conditions). In the fee-space EIT schemes, in order to produce the comparable XPM phase shift, a much greater optical depth of the atomic medium is required (a few orders of magnitude greater than our scheme; see Refs. [24,26,30]).

The XPM phase shift in our scheme is very small when the medium optical depth becomes very large ($\sqrt{N}g \gg \kappa\Gamma_3$). Therefore, our scheme performs better with an atomic medium of low to moderate optical depths while the free-space EIT schemes perform better with an atomic medium of high optical depths. But at very high optical depth, the nonvanishing decoherence γ_{12} reduces the transmitted light intensity and places the upper limits on the practical phase shift obtainable with the free-space EIT schemes. Recently, large phase shifts have been observed in a gain-assisted XPM scheme that overcomes the limitation of the light attenuation in the absorption based XPM schemes [33].

IV. CONCLUSION

In conclusion, we have shown that a coherently coupled cavity-atom system can be used to study resonant nonlinear optics at low light intensities. The system is well suited for studies of all-optical switching and XPM at low control light intensities. In the weak-coupling regime of the cavity QED, the interplay of the EIT and the collective coupling of atoms with the cavity mode enable the nonlinear control of the intracavity light field through a free-space control laser. The analytical results and numerical calculations show that large optical nonlinearities can be induced by the control laser near the narrow resonance of cavity EIT, which can be used to control the amplitude and phase of the cavity-transmitted and -reflected fields.

ACKNOWLEDGMENT

This work is supported by the National Science Foundation under Grant No. 1205565.

[1] J. Ye and T. W. Lynn, *Advances in Atomic, Molecular, and Optical Physics*, edited by B. Bederson and H. Walther (Academic Press, San Diego, CA, 2003), Vol. 49, pp. 1–19.

[2] *Cavity Quantum Electrodynamics*, edited by P. R. Berman (Academic, San Diego, 1994).

[3] J. M. Raimond, M. Brune, and S. Haroche, Rev. Mod. Phys. **73**, 565 (2001).

[4] H. Tanji-Suzuki, W. Chen, R. Landig, J. Simon, and V. Vuletić, Science **333**, 1266 (2011).

[5] G. Rempe, R. J. Thompson, R. J. Brecha, W. D. Lee, and H. J. Kimble, Phys. Rev. Lett. **67**, 1727 (1991).

[6] P. Grangier, J. F. Roch, J. Roger, L. A. Lugiato, E. M. Pessina, G. Scandroglio, and P. Galatola, Phys. Rev. A **46**, 2735 (1992).

[7] S. E. Harris, Phys. Today **50**(7), 36 (1997).

[8] M. D. Lukin, M. Fleischhauer, M. O. Scully, and V. L. Velichansky, Opt. Lett. **23**, 295 (1998).

[9] H. Wang, D. J. Goorskey, W. H. Burkett, and M. Xiao, Opt. Lett. **25**, 1732 (2000).

[10] G. Hernandez, J. Zhang, and Y. Zhu, Phys. Rev. A **76**, 053814 (2007).

[11] H. Wu, J. Gea-Banacloche, and M. Xiao, Phys. Rev. Lett. **100**, 173602 (2008).

[12] M. Mücke, E. Figueroa, J. Bochmann, C. Hahn, K. Murr, S. Ritter, C. J. Villas-Boas, and G. Rempe, Nature (London) **465**, 755 (2010).

[13] A. Joshi and M. Xiao, Phys. Rev. Lett. **91**, 143904 (2003).

[14] T. Kampschulte, W. Alt, S. Brakhane, M. Eckstein, R. Reimann, A. Widera, and D. Meschede, Phys. Rev. Lett. **105**, 153603 (2010).

[15] J. K. Thompson, J. Simon, H. Loh, and V. Vuletic, Science **313**, 74 (2006).

[16] J. Sheng, H. Wu, M. Mumba, J. Gea-Banacloche, and M. Xiao, Phys. Rev. A **83**, 023829 (2011).

[17] M. Albert, A. Dantan, and M. Drewsen, Nat. Photonics **5**, 633 (2011).

[18] A. E. B. Nielsen and J. Kerckhoff, Phys. Rev. A **84**, 043821 (2011).

[19] Y. Zhu, Opt. Lett. **35**, 303 (2010).

[20] X. Wei, J. Zhang, and Y. Zhu, Phys. Rev. A **82**, 033808 (2010).

[21] Y. Wang, J. Zhang, and Y. Zhu, Phys. Rev. A **85**, 013814 (2012).

[22] D. F. Walls and G. J. Milburn, *Quantum Optics* (Springer-Verlag, Berlin, Heidelberg, 1994).

[23] R. Zhao, Y. O. Oudin, S. D. Jenkins, C. J. Campbell, D. N. Matsukevich, T. A. B. Kenndedy, and A. Kuzmich, Nat. Phys. **5**, 100 (2009).

[24] H. Schmidt and A. Imamoglu, Opt. Lett. **21**, 1936 (1996).

[25] H. Kang and Y. Zhu, Phys. Rev. Lett. **91**, 093601 (2003).

[26] T. Opatrny and D. G. Welsch, Phys. Rev. A **64**, 023805 (2001).

[27] A. B. Matsko, I. Novikova, G. R. Welch, and M. S. Zubairy, Opt. Lett. **28**, 96 (2003).

[28] D. Petrosyan and G. Kurizki, Phys. Rev. A **65**, 033833 (2002).

[29] D. Petrosyan and Y. P. Malakyan, Phys. Rev. A **70**, 023822 (2004).

[30] Z. B. Wang, K. P. Marzlin, and B. C. Sanders, Phys. Rev. Lett. **97**, 063901 (2006).

[31] Y. F. Chen, C. Y. Wang, S. H. Wang, and I. A. Yu, Phys. Rev. Lett. **96**, 043603 (2006).

[32] H. Wang, D. Goorskey, and M. Xiao, Phys. Rev. Lett. **87**, 073601 (2001).

[33] R. B. Li, L. Deng, and E. W. Hagley, Phys. Rev. Lett. **110**, 113902 (2013).