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Phase control of squeezed state in double electromagnetically induced transparency system with a loop-transition structure

Yuan Li [a], Yusheng Zhou [a], Yong Wang [a], Qiang Ling [a], Bing Chen [a], Yan Dou [a], Wei Zhang [a], Weiqing Gao [a], Zhiqiang Guo [a,*], Junxiang Zhang [b,**]

[a] *School of Electronic Science and Applied Physics, School of Resources and Environmental Engineering or School of Electrical Engineering and Automation, Hefei University of Technology, Hefei 230009, China*
[b] *Department of Physics, Zhejiang University, Hangzhou 310027, China*

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A B S T R A C T

We theoretically study the squeezed probe light passing through a double electromagnetically induced transparency (DEIT) system, in which a microwave field and two coupling lights drive a loop transition. It is shown that the output squeezing can be maintained in both two transparency windows of DEIT, and it can also be manipulated by the relative phase of the three driving fields. The influence of the intensity of applied fields and the optical depth of atoms on the squeezing is also investigated. This study offers possibilities to manipulate the squeezing propagation in atomic media by the phase of electromagnetic fields.

1. Introduction

As a quantum interference effect in light–atom interaction, electromagnetically induced transparency (EIT) has been introduced and discussed for three decades [1,2]. Under the action of a strong resonant coupling light, an opaque medium gets transparent for a weak probe light. Meanwhile, there is a steep normal dispersion for the probe light. Namely, a resonant probe light can propagate subluminally through the medium without absorption [3]. Based on this, EIT can be used to the optical and quantum memory [4,5], which is an indispensable part of the quantum information processing [6].

In order to extend quantum communication protocols from finite to infinite dimensions, it is necessary to deal with continuous quantum variables in quantum information processing [7–10]. Squeezed state light, which exhibits less noise in one quadrature component than coherent states, is one of the most commonly used nonclassical lights in the continuous variables (CV) regime [11]. Consequently, the ultraslow propagation, storage and retrieval of a squeezed vacuum under the condition of EIT have been realized experimentally and theoretically [12–17]. Akamatsu et al. firstly demonstrated the preservation of the quadrature squeezing for the probe light passing through the EIT medium [12], and they also observed the ultraslow propagation of squeezed vacuum pulses [13]. Lam et al. and Dantan et al. theoretically examined these phenomena even in the presence of ground state dephasing and atomic noise [14,15]. Appel et al. and Honda et al. observed the storage and retrieval of a squeezed vacuum in atomic vapor and trapped atoms [16,17]. Arikawa et al. developed a quantum memory of a squeezed vacuum for arbitrary frequency sidebands using bichromatic EIT [18].

The majority of researches on the above-list phenomena are based on three-level EIT media, and the physical contents have been comprehensively studied. Although multilevel systems do not provide new physical principles, they can contribute to extensive applications [19]. In a four-level atomic system such as an inverted-Y, N or tripod-type configuration, we can obtain a double-EIT (DEIT) [20]. It is found that DEIT can be applied to laser cooling of atoms [21], microwave electrometry [22,23], storage of multiple images [24], controllable splitting and modulation of single-photon-level pulses [25], matched group velocity and enhanced nonlinear effects [26,27]. DEIT has also been proposed to realize the two-channel quantum memory for a squeezed probe light [28,29].

As we known, the relative phase of electromagnetic fields plays a very important role in the interaction between fields and a multilevel system with a loop-transition structure [30]. The DEIT can be manipulated by the relative phase of fields in an inverted-Y [31]

* Corresponding author at: Hefei University of Technology, Tunxi Road 193, Hefei 230009, China.
** Corresponding author.
E-mail addresses: guozq@hfut.edu.cn (Z. Guo), junxiang_zhang@zju.edu.cn (J. Zhang).

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and a tripod-type system [32,33] with a microwave field. These phenomena have been used to obtain a giant Kerr nonlinearity [34], optical bistability and multistability [35], three-dimensional atom localization [36] and the dynamic control of coherent pulse propagation and switching [37]. In this paper, we study the phase control of a squeezed-state probe light passing through a tripod-type DEIT system with a closed interaction contour for a more realistic model, which is the basis for the phase-dependent quantum memory.

2. Theoretical model and basic equations

Consider a four-level tripod-type atomic system [32] which can be experimentally obtained in Rb atoms [38], as depicted in Fig. 1. Two classical strong coupling lights Ω_1 and Ω_2 and a microwave field Ω_m with angular frequencies ω_1, ω_2 and ω_m drive the transitions $|1\rangle \leftrightarrow |3\rangle$, $|2\rangle \leftrightarrow |3\rangle$ and $|1\rangle \leftrightarrow |2\rangle$, respectively; a weak quantum probe light $\hat{a}(z,t)$ of angular frequencies ω_p couples with the transition $|4\rangle \leftrightarrow |3\rangle$. The decay rate of the excited state $|3\rangle$ to the ground states $|1\rangle$, $|2\rangle$ and $|4\rangle$ is expressed as $\Gamma = \Gamma_{31} + \Gamma_{32} + \Gamma_{34}$. $\delta_p = \omega_p - \omega_{34}$, $\delta_1 = \omega_1 - \omega_{31}$, $\delta_2 = \omega_2 - \omega_{32}$ and $\delta_m = \omega_m - \omega_{21}$ are the detunings of the four fields, where $\omega_{\mu\nu}$ is the atomic transition frequency of $|\mu\rangle \leftrightarrow |\nu\rangle$. Under the condition of the three-photon resonance $\omega_1 = \omega_2 + \omega_m$, the Hamiltonian of the system under the rotating-wave and dipole approximations can be represented as:

$$\begin{aligned}\hat{H} &= \hbar(\delta_1 - \delta_2)|1\rangle\langle 1| - \hbar\delta_2|3\rangle\langle 3| + \hbar(\delta_p - \delta_2)|4\rangle\langle 4| \\ &\quad - \hbar\Omega_m\big(|2\rangle\langle 1| + |1\rangle\langle 2|\big) - \hbar\Omega_1\big(e^{i\phi}|3\rangle\langle 1| + e^{-i\phi}|1\rangle\langle 3|\big) \\ &\quad - \hbar\Omega_2\big(|3\rangle\langle 2| + |2\rangle\langle 3|\big) - \hbar\big(g\hat{a}|3\rangle\langle 4| + g^*\hat{a}^+|4\rangle\langle 3|\big),\end{aligned} \tag{1}$$

where g is the coupling constant between the atomic transition $|4\rangle \leftrightarrow |3\rangle$ and the quantized probe light $\hat{a}(z,t)$, $\phi = \phi_2 + \phi_m - \phi_1$ is the relative phase of the three driving fields.

Under the condition that the intensity of the quantized probe light is much weaker than intensities of the three driving fields, all the atoms are initially in the lower level $|4\rangle$ with $\langle\sigma_{44}\rangle \approx 1$. The evolution equations for both the atomic operators and the annihilation operator of the probe light in the slowly varying envelope approximation are given by

$$\dot{\hat{\sigma}}_{41} = \big[i(\delta_p - \delta_1) - \gamma_{14}\big]\hat{\sigma}_{41} + i\Omega_1 e^{-i\phi}\hat{\sigma}_{43} + i\Omega_m\hat{\sigma}_{42} + \hat{F}_{41}, \tag{2a}$$

$$\dot{\hat{\sigma}}_{42} = \big[i(\delta_p - \delta_2) - \gamma_{24}\big]\hat{\sigma}_{42} + i\Omega_2\hat{\sigma}_{43} + i\Omega_m\hat{\sigma}_{41} + \hat{F}_{42}, \tag{2b}$$

$$\dot{\hat{\sigma}}_{43} = (i\delta_p - \gamma_{34})\hat{\sigma}_{43} + i\Omega_1 e^{i\phi}\hat{\sigma}_{41} + i\Omega_2\hat{\sigma}_{42} + ig\hat{a} + \hat{F}_{43}, \tag{2c}$$

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{a}(z,t) = ig^* N\hat{\sigma}_{43}(z,t), \tag{2d}$$

where $\hat{F}_{\mu\nu}$ is the associated Langevin noise operator resulting from the coupling of atoms to all the vacuum field modes, N is the number of atoms and γ_{ij} are the decay rates of the atomic dipole operators with $\gamma_{34} = \Gamma/2 \gg \gamma_{14} \approx \gamma_{24}$. γ_{14} and γ_{24} denote the decay rates of the ground-state coherence due to atomic collisions and atomic transit outside the interaction region. The decay rate of the ground-state coherence is much smaller than other decay rates [14].

In order to solve above equations, we convert them into ones in the frequency domain by the Fourier transform. Then the annihilation operator of the probe light at the exit of the medium can be obtained:

$$\hat{a}(L,\omega) = \hat{a}(0,\omega)e^{-\Lambda(\omega)L} + \frac{g^* N}{c}\int_0^L \hat{F}(s,\omega)e^{-\Lambda(\omega)(L-s)}ds \tag{3}$$

with

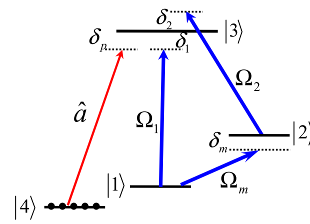

Fig. 1. A four-level tripod-type atomic system.

$$\Lambda(\omega) = \frac{|g|^2 N}{c} \times \frac{r_b r_c + \Omega_m^2}{Z(\omega)} - \frac{i\omega}{c} \tag{4}$$

and

$$\begin{aligned}\hat{F}(s,\omega) &= \big(-(r_b\Omega_1 e^{i\phi} + i\Omega_2\Omega_m)\hat{F}_{41}(s,\omega) \\ &\quad + i(r_b r_c + \Omega_m^2)\hat{F}_{43}(s,\omega) \\ &\quad - (r_c\Omega_2 + i\Omega_1\Omega_m e^{i\phi})\hat{F}_{42}(s,\omega)\big)/Z(\omega),\end{aligned} \tag{5}$$

where $Z(\omega) = r_a r_b r_c + r_b\Omega_1^2 + r_c\Omega_2^2 + r_a\Omega_m^2 + 2i\Omega_1\Omega_2\Omega_m\cos\phi$, $r_a = \gamma_{34} + i(\delta_p - \omega)$, $r_b = \gamma_{24} + i(\delta_p - \delta_2 - \omega)$, $r_c = \gamma_{14} + i(\delta_p - \delta_1 - \omega)$ and ω is the detection frequency in experiments.

According to the definition of the field quadrature operator, the amplitude and phase quadratures of the output probe light are as follows:

$$\hat{X}(L,\omega) = \hat{a}(L,\omega) + \hat{a}^+(L,-\omega), \tag{6a}$$

$$\hat{Y}(L,\omega) = -i\big[\hat{a}(L,\omega) - \hat{a}^+(L,-\omega)\big]. \tag{6b}$$

To study the preservation of squeezing after passing through the EIT medium, the normalized quadrature amplitude spectrum of the output probe light can be calculated by the following definition [39]:

$$S_X(L,\omega) = \frac{c}{L}\big\langle \hat{X}(L,\omega)\hat{X}(L,-\omega)\big\rangle. \tag{7}$$

To compute the equation (7), the correlation functions of Langevin noise operators must be calculated. By using the quantum regression theorem [14], the nonzero terms are as follows:

$$\big\langle \hat{F}_{41}(s,\omega)\hat{F}_{41}^+(s',-\omega')\big\rangle = 2\gamma_{14}L\delta(s-s')\delta(\omega-\omega')/N, \tag{8a}$$

$$\big\langle \hat{F}_{42}(s,\omega)\hat{F}_{42}^+(s',-\omega')\big\rangle = 2\gamma_{24}L\delta(s-s')\delta(\omega-\omega')/N, \tag{8b}$$

$$\big\langle \hat{F}_{43}(s,\omega)\hat{F}_{43}^+(s',-\omega')\big\rangle = 2\gamma_{34}L\delta(s-s')\delta(\omega-\omega')/N. \tag{8c}$$

Substitute Eqs. (3), (6) and (8) into Eq. (7), and the amplitude noise spectrum of the output probe can be expressed as:

$$S_X(L,\omega) = S_1 + S_2 + S_3 \tag{9a}$$

with

$$\begin{aligned}S_1 &= \frac{S_X(0,\omega)}{4}\big[\exp\big(-\big[\Lambda(\omega) + \Lambda(-\omega)\big]L\big) \\ &\quad + \exp\big(-\big[\Lambda(\omega) + \Lambda^*(\omega)\big]L\big) + \exp\big(-\big[\Lambda^*(-\omega) + \Lambda(-\omega)\big]L\big) \\ &\quad + \exp\big(-\big[\Lambda^*(-\omega) + \Lambda^*(\omega)\big]L\big)\big],\end{aligned} \tag{9b}$$

$$\begin{aligned}S_2 &= -\frac{S_Y(0,\omega)}{4}\big[\exp\big(-\big[\Lambda(\omega) + \Lambda(-\omega)\big]L\big) \\ &\quad - \exp\big(-\big[\Lambda(\omega) + \Lambda^*(\omega)\big]L\big) - \exp\big(-\big[\Lambda^*(-\omega) + \Lambda(-\omega)\big]L\big) \\ &\quad + \exp\big(-\big[\Lambda^*(-\omega) + \Lambda^*(\omega)\big]L\big)\big],\end{aligned} \tag{9c}$$

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