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RESEARCH ARTICLE

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Quantum noise property in coherent atomic system

Abstract The coherent superposition of atomic states leads to the characteristic change of interacting lights because of the coupling between the lights and atoms. In this paper, the noise spectrum of the quantified light interacting with the atoms is studied under the condition of electromagnetically induced transparency (EIT). It is shown that the noise spectrum displays a double M-shape noise profile resulted from the conversion of phase noise of probe beam. A squeezing of 0.3 dB can be observed at the detuning of probe light at the proper parameters of atoms and coupling beam.

Keywords quantum coherence, quantum fluctuation, noise spectrum

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1 Introduction

Quantum coherence in atomic system is responsible for a large number of important effects such as electromagnetically induced transparency (EIT) [1,2], coherent population trapping (CPT) [3], lasing without inversion (LWI) [4－5], and refractive index enhancement [6]. The motivation of such studies lies in the potential wide range of applications in quantum information system, for examples, storage of quantum information in a coherent media [7], quantum computation [6], quantum logic gate [8], quantum switches [9] and quantum interferometric optical lithography [10].

Recently many interesting investigations about application of atomic coherence in quantum information science have received a lot of public attention [11－12]. In quantum information, a "qubit" is the quantum equivalent of a classical "bit" of information. The qubit with photon states is usually called flying qubit since the states of photons are easily controlled, their coherence can be preserved over long distance with optical fibers, thus it is the useful resources for quantum communication application. But photons can not be stored for long, and manipulation collective photon state present considerable difficulties even when photons are confined in the same Microcavity. And the creation of long lived information processing with particles such as atoms, on the other hand, relatively attracts pursuit in recent years. Because the motional state of material particles is stable, the qubit of it is called static qubit. So with the quantum coherence between atoms or between atoms and light fields, the atoms may enable the conversion of "flying qubits" into "storage qubits". The further advance is made on issues such as register and exchange of quantum information from atom to photon or from photon to atom et al. And the most far reaching is the neutral atom quantum computer, for which quantum theory predicts a new type of computing logic, able in some cases, to outrun the present classical computers by many orders of magnitude in processing time.

Using atoms for quantum information processing requires efficient means for the preparation, manipulation, and storage of qubits inscribed into atoms as well as for the readout of the information of quantum state from atomic memory. One of the ways to verify the faithful atomic memory preparation and readout is the measurement of quantum noise or correlations of retrieved states [13]. The experiment of EIT with squeezed vacuum showed that the squeezing of output probe light was maintained under the condition of two-photon resonance [14]. And the quantum mechanical treatment for both atomic and field fluctuation in the situation of EIT or Raman process were thus presented to clarify the result that the quantum state can be well preserved under the condition of EIT [15, 16].

Regard the atoms as a storage medium with pleasure, we

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may meet rebarbative noises which affected the readout of the information of the field state from the storage atomic medium. For a well-stabilized laser beam, the noise behavior is determined by the presence of quantum noise of both the intensity and phase of the lasing field, especially the phase noise, which changes with time in a random fashion, will converts into amplitude noise when the light pass through a dispersion medium.

In the present paper, we concentrate on the interaction of two single-mode fields and atomic medium under the conditions of electromagnetically induced transparency (EIT). The essence of EIT is that an atomic coherence is induced in a multilevel system by a strong control laser field, which alters the response of system to a probe laser field. After the two light beams passed through the atomic medium, both the quantum and absorption properties of probe beam are affected by the atomic coherence, we calculate the atom-light interaction and show the spectrum of the quantum probe laser field noise in the present of detuning of probe beam in lambda atomic system, and we show that the dispersion characteristic of Raman scheme converts the laser phase noise into amplitude noise in the transmitted probe beam when the phase noise is larger than the for amplitude noise. Thus on the other hand, it provides a possible method to measure the phase noise of laser beam. The amplitude noise spectrum of output probe beam resulting from the phase noise conversion is found to display double M-shape structure, which shows that besides in the case of two-photon resonance EIT, the quantum noise of probe beam can also be maintained even in the case that the two-photon detuning exists. This phenomenon may be referenced to characterize the quantum state propagating extremely slowly or superluminally in the transfer process with atoms.

2 Theoretical model

Consider a closed three-level Λ-type system, as shown in Fig. 1. A strong coupling laser of frequency ω_c is coupled to the $|c\rangle \leftrightarrow |a\rangle$ transition, while a weak quantum probe beam labeled as $\hat{\varepsilon}$ with frequency ω_p interacts with transition $|a\rangle \leftrightarrow |b\rangle$, the positive frequency part of the electric field is defined by

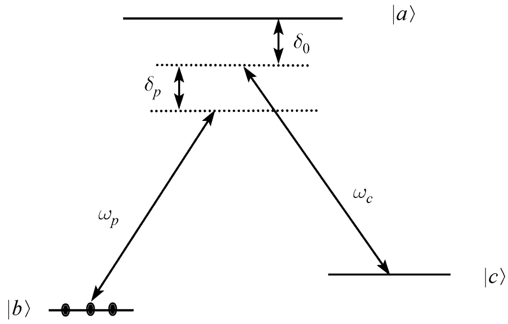

Fig. 1 Schematic of Λ -type system.

$$\hat{E}^{(+)}=\sqrt{\frac{\hbar\omega_{ab}}{2\varepsilon_0 V}}\hat{\varepsilon}(z,t)\exp\left[\mathrm{i}\omega_{ab}\left(\frac{z}{c}-t\right)\right] \qquad (1)$$

where V is the quantization volume of the electromagnetic field, which for simplicity is taken to be the interaction volume of the electromagnetic field and atomic medium. $\delta_0 = \omega_{ac}-\omega_c$ is one one-photon detuning, $\delta_p = \omega_{ab}-\omega_p-\delta_0$ is two-photon detuning. For this system, we make the usual assumption that all atoms are initially in the ground state $|b\rangle$ and the quantum probe light is much weaker than the classical coupling field. The evolution of the quantum probe field operator can be depicted in the slowly varying amplitude approximation by the propagation equation:

$$\left(\frac{\partial}{\partial t}+c\frac{\partial}{\partial t}\right)\hat{\varepsilon}=\mathrm{i}gN\hat{\sigma}_{ba} \qquad (2)$$

where N is the total numble of atoms interacting with the field. $g=\wp\sqrt{\omega_{ab}/2\varepsilon_0 V\hbar}$ is the corresponding coupling strength, with $\wp$ represents the atomic dipole moment .

As shown in the propagation Eq. (2) of the quantum probe field, an important operator $\hat{\sigma}_{ba}$ is a space- and time-dependant collective, slowly-varying operator denoting the atomic dipole which describes the quantum properties of the atomic medium .The collective, slowly-varying atomic operators are related to the single atom operators by the following definition. The atomic medium of total length L is divided into $2M$+1 slices, each of thickness $L/(2M+1)$ with the center at $Z_l = lL/(2M+1)$, ($l=-M, \cdots, M$). Then the space- and time- dependent collective, slowly-varying atomic operators are described

$$\hat{\sigma}_{\mu\nu}(z,t)=\lim_{M\to\infty}\frac{2M+1}{N}\sum_j \hat{\sigma}^j_{\mu\nu}(t)\Big|_{Z_j\to z}\exp\left[\mathrm{i}\omega_{\mu\nu}\left(\frac{z}{c}-t\right)\right] \qquad (3)$$

where $\hat{\sigma}^j_{\mu\nu}(t)=\left|\mu^j(t)\right\rangle\left\langle\nu^j(t)\right|$ is the j th single atom operator, j is to be taken over all atoms in the slice l around z . $\omega_{\mu\nu}=(E_\mu-E_\mu)/\hbar$ is the frequency of the $|\mu\rangle \leftrightarrow |\nu\rangle$ transition. The collective, slowly-varying atomic operators should comply with a set of Heisenberg-Langevin equations [17] and be governed by the interaction Hamiltonian between light and atoms and δ- correlated Langevin noise operators $\hat{F}_{\mu\nu}$:

$$\frac{\partial\hat{\sigma}_{\mu\nu}(z,t)}{\partial t}=-\gamma_{\mu\nu}\hat{\sigma}_{\mu\nu}(z,t)+\frac{\mathrm{i}}{\hbar}[\hat{H},\hat{\sigma}_{\mu\nu}(z,t)]+\hat{F}_{\mu\nu}$$

$$\hat{H}=-\int\frac{N\hbar}{L}[g\hat{\sigma}_{ab}(z,t)\hat{\varepsilon}(z,t)+\Omega\ (z,t)\hat{\sigma}_{ac}(z,t)+Hc]\mathrm{d}z$$

$$\hat{F}_{\mu\nu}(z,t)=\lim_{M\to\infty}\frac{2M+1}{N}\sum_j \hat{F}^j_{\mu\nu}(t)\Big|_{Z_j\to z}\exp\left[\mathrm{i}\omega_{\mu\nu}\left(\frac{z}{c}-t\right)\right] \qquad (4)$$

where $\gamma_{\mu\nu}$ is the decays of the atomic dipole operators and $\hat{F}_{\mu\nu}$ is the Langevin noise operator which describe the effect of spontaneous decay caused by the coupling of atoms and all the vacuum field modes. The random decay process adds noise to the single atom Langevin noise operators $\hat{F}_{\mu\nu}^{j}$. The collective, slowly-varying Langevin noise operators $\hat{F}_{\mu\nu}$ are related to the single-atom Langevin noises operators by the same relation like Eq. (3). For the Λ-type 3-level atoms shown in Fig. 1, the evolution equations of the atomic operators are expanded by:

$$\begin{aligned}
\dot{\hat{\sigma}}_{aa} &= -2\gamma\hat{\sigma}_{aa} - \mathrm{i}\left(g^{*}\hat{\varepsilon}^{+}\hat{\sigma}_{ba} - g\hat{\varepsilon}\hat{\sigma}_{ab}\right) - \mathrm{i}\left(\Omega^{*}\hat{\sigma}_{ca} - \Omega\hat{\sigma}_{ac}\right) + \hat{F}_{aa} \\
\dot{\hat{\sigma}}_{bb} &= \gamma\hat{\sigma}_{aa} + \gamma_0\left(\hat{\sigma}_{cc} - \hat{\sigma}_{bb}\right) + \mathrm{i}\left(g^{*}\hat{\varepsilon}^{+}\hat{\sigma}_{ba} - g\hat{\varepsilon}\hat{\sigma}_{ab}\right) + \hat{F}_{bb} \\
\dot{\hat{\sigma}}_{cc} &= \gamma\hat{\sigma}_{aa} + \gamma_0\left(\hat{\sigma}_{bb} - \hat{\sigma}_{cc}\right) + \mathrm{i}\left(\Omega^{*}\hat{\sigma}_{ca} - \Omega\hat{\sigma}_{ac}\right) + \hat{F}_{cc} \\
\dot{\hat{\sigma}}_{ba} &= -[\gamma - \mathrm{i}(\delta_0 + \delta_p)]\hat{\sigma}_{ba} + \mathrm{i}g\hat{\varepsilon}\left(\hat{\sigma}_{bb} - \hat{\sigma}_{aa}\right) + \mathrm{i}\Omega\hat{\sigma}_{bc} + \hat{F}_{ba} \\
\dot{\hat{\sigma}}_{ac} &= -(\gamma - \mathrm{i}\delta_0)\hat{\sigma}_{ac} + \mathrm{i}\Omega^{*}\left(\hat{\sigma}_{aa} - \hat{\sigma}_{cc}\right) - \mathrm{i}g^{*}\hat{\varepsilon}^{+}\hat{\sigma}_{bc} + \hat{F}_{ac} \\
\dot{\hat{\sigma}}_{bc} &= -\left(\gamma_0 - \mathrm{i}\delta_p\right)\hat{\sigma}_{bc} - \mathrm{i}g\hat{\varepsilon}\hat{\sigma}_{ac} + \mathrm{i}\Omega^{*}\hat{\sigma}_{ba} + \hat{F}_{bc}
\end{aligned} \tag{5}$$

where γ is the atomic dipole decay rate, and γ_0 is the dephasing rate for the two ground states of the Λ-type system, g is the atom-field coupling constants, Ω is the Rabi frequency of coupling field, N refers to the number of atoms per unit volume. The correlation functions of the Langevin noises operators $\hat{F}_{\mu\nu}$ can be calculated via the quantum regression theorem [18, 19]:

$$\begin{aligned}
\left\langle \hat{F}_{\mu\nu}(z_1,t_1)\hat{F}_{\alpha\beta}(z_2,t_2)\right\rangle = \frac{L}{N}\Big\langle D(\hat{\sigma}_{\mu\nu}\hat{\sigma}_{\alpha\beta}) - D(\hat{\sigma}_{\mu\nu})\hat{\sigma}_{\alpha\beta} \\
- \hat{\sigma}_{\mu\nu}D(\hat{\sigma}_{\alpha\beta})\Big\rangle \delta(z_2 - z_1)\delta(t_2 - t_1)
\end{aligned} \tag{6}$$

Here $D(\hat{\sigma}_{\mu\nu})$ is the expression for $\dot{\hat{\sigma}}_{\mu\nu}$ obtained from the Heisenberg-Langevin equation for $\hat{\sigma}_{\mu\nu}$ but without the Langevin force term. The Dirac delta function in Eq. (6) represents the short memory of the reservoir of vacuum modes responsible for the Langevin forces. L is the length of medium interacting with lights. Using the definition of the collective, slowly-varying Langevin noise operators in terms of their single atom counterpart, we derive the following for the nonzero correlations which suffice Eq. (6), these are

$$\begin{aligned}
\left\langle \hat{F}_{ba}(z_1,\ \omega_1)\hat{F}_{ba}^{+}(z_2,\omega_2)\right\rangle &= \frac{L\delta(z_1 - z_2)\delta(\omega_1 - \omega_2)}{N}(2\gamma - \gamma_0) \\
\left\langle \hat{F}_{bc}(z_1,\ \omega_1)\hat{F}_{bc}^{+}(z_2,\omega_2)\right\rangle &= \frac{L\delta(z_1 - z_2)\delta(\omega_1 - \omega_2)}{N}\gamma_0 \\
\left\langle \hat{F}_{bc}^{+}(z_1,\omega_1)\hat{F}_{bc}(z_2,\ \omega_2)\right\rangle &= \frac{L\delta(z_1 - z_2)\delta(\omega_1 - \omega_2)}{N}\gamma_0
\end{aligned} \tag{7}$$

We now assume that the quantum probe field intensity is much less than that of the classical control field and all the atoms are initially in the ground state $|b\rangle$, we have

$$\hat{\sigma}_{bb}^{(0)} = 1,\ \hat{\sigma}_{aa}^{(0)} = \hat{\sigma}_{cc}^{(0)} = \hat{\sigma}_{ac}^{(0)} = 0 \tag{8}$$

and then we can solve Eqs. (5), (7) and (8) perturbatively to first order in $\frac{g\hat{\varepsilon}}{\Omega}$ to obtain a set of three closed equations:

$$\begin{aligned}
\dot{\hat{\sigma}}_{ba} &= -[\gamma - \mathrm{i}(\delta_p + \delta_0)]\hat{\sigma}_{ba} + \mathrm{i}g\hat{\varepsilon} + \mathrm{i}\Omega\hat{\sigma}_{bc} + \hat{F}_{ba} \\
\dot{\hat{\sigma}}_{bc} &= -\left(\gamma_0 - \mathrm{i}\delta_p\right)\hat{\sigma}_{bc} + \mathrm{i}\Omega^{*}\hat{\sigma}_{ba} + \hat{F}_{bc} \\
\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{\varepsilon} &= \mathrm{i}g^{*}N\hat{\sigma}_{ba}
\end{aligned} \tag{9}$$

In order to solve those equations, we convert them into Fourier form in the frequency domain via

$$\hat{F}(z,\omega) = \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{+\infty}\hat{F}(z,t)\exp(\mathrm{i}\omega t)\,\mathrm{d}t \tag{10}$$

then we solve Eq. (8) and perform formal integration over z to obtain the output field of the EIT medium. We obtain

$$\begin{aligned}
\hat{\varepsilon}(L,\omega) = \mathrm{e}^{-\Lambda(\omega)L}\hat{\varepsilon}(0,\omega) + \frac{g^{*}N}{c}\int_0^L \mathrm{e}^{-\Lambda(\omega)(L-s)} \\
\left\{\frac{-\Omega\hat{F}_{bc}(s,\omega) + \hat{F}_{ba}(s,\omega)\left(\omega + \delta_p + \mathrm{i}\gamma_0\right)}{\left[\gamma - \mathrm{i}\left(\omega + \delta_0 + \delta_p\right)\right]\left[\gamma_0 - \mathrm{i}\left(\omega + \delta_p\right)\right] + \left|\Omega\right|^2}\right\}\mathrm{d}s
\end{aligned} \tag{11}$$

where

$$\begin{aligned}
\Lambda(\omega) = \frac{|g|^2 N}{c} \\
\frac{\gamma_0 - \mathrm{i}(\omega + \delta)}{[\gamma - \mathrm{i}(\omega + \delta_p + \delta_0)]\left[\gamma_0 - \mathrm{i}(\omega + \delta_p)\right] + |\Omega|^2} - \frac{\mathrm{i}\omega}{c}
\end{aligned} \tag{12}$$

Equation (11) can be interpreted as follows: The first term on the right-hand side describes the amplitude of the probe field operator is attenuated and phase shifted according to the function (12), and the second term represents vacuum noise added to the probe field as it interacts with the atoms.

To obtain the amplitude noise spectrum of the probe field, we introduce the correlation of quadrature amplitude and phase components:

$$\begin{aligned}
\left\langle X^{\mathrm{out}}(\omega)X^{\mathrm{out}}(\omega')\right\rangle &= \frac{2\pi L}{c}\delta(\omega + \omega')S_X^{\mathrm{out}}(\omega) \\
\left\langle Y^{\mathrm{out}}(\omega)Y^{\mathrm{out}}(\omega')\right\rangle &= \frac{2\pi L}{c}\delta(\omega + \omega')S_Y^{\mathrm{out}}(\omega)
\end{aligned} \tag{13}$$

where $X^{\mathrm{out}}(\omega)$ and $X^{\mathrm{out}}(\omega)$ are the Fouirer transforms of the output quadrature amplitude and phase operators X^{out} and Y^{out}, which are defined as $\hat{X}^{\mathrm{out}}(\omega) = \hat{\varepsilon}^{\mathrm{out}}(L,\omega) + \hat{\varepsilon}^{\mathrm{out}+}(L,-\omega)$, $\hat{Y}^{\mathrm{out}}(\omega) = \mathrm{i}\left[\hat{\varepsilon}^{\mathrm{out}}(L,\omega) - \hat{\varepsilon}^{\mathrm{out}+}(L,-\omega)\right]$. And $S_X^{\mathrm{out}}(\omega), S_Y^{\mathrm{out}}(\omega)$, are the output noise spectra of quadratures X^{out}, Y^{out}. ω is the analysis frequency.

Then one gets the normalized output amplitude spectrum:

$$S_X^{\mathrm{out}}(\omega) = S_1(\omega) + S_2(\omega) + S_3(\omega) \tag{14}$$

with

$$\begin{aligned}
S_1(\omega) = (S_X^{\mathrm{in}}(\omega)/4) \times \{\exp\{-[\Lambda(\omega) + \Lambda(-\omega)]\} \\
+ \exp\{-[\Lambda^{*}(\omega) + \Lambda^{*}(-\omega)]\} + \exp\{-[\Lambda(\omega) \\
+ \Lambda^{*}(\omega)]\} + \exp\{-[\Lambda(-\omega) + \Lambda^{*}(-\omega)]\}\}
\end{aligned}$$

$$S_2(\omega) = (S_Y^{\text{in}}(\omega)/4) \times \{-\exp\{-[\Lambda(\omega)+\Lambda(-\omega)]\} - \exp\{-[\Lambda^*(\omega)+\Lambda^*(-\omega)]\} + \exp\{-[\Lambda(\omega) + \Lambda^*(\omega)]\} + \exp\{-[\Lambda(-\omega)+\Lambda^*(-\omega)]\}\}$$

$$S_3(\omega) = \frac{|g|^2 N}{c}\left\{\left[\frac{1-\exp[-2\operatorname{Re}\Lambda(\omega)]}{2\operatorname{Re}\Lambda(\omega)}\right] \times B1 + \left[\frac{1-\exp[-2\operatorname{Re}\Lambda(-\omega)]}{2\operatorname{Re}\Lambda(-\omega)}\right]\times B2\right\}$$

$$B1 = (|\Omega|^2 \gamma_0 + [(\omega+\delta_p)^2+\gamma_0^2])/ (\gamma_0[\gamma\gamma_0 - (\delta_p+\delta_0+\omega)\times(\delta_p+\omega)+|\Omega|^2] + (\delta_p+\omega)\times[\gamma_0(\delta_p+\delta_0+\omega)+\gamma(\delta_p+\omega)])$$

$$B2 = (|\Omega|^2\gamma_0)/(\gamma_0[\gamma\gamma_0 - (\delta_p+\delta_0-\omega)\times(\delta_p-\omega)+|\Omega|^2] + (\delta_p-\omega)\times[\gamma_0(\delta_p+\delta_0-\omega)+\gamma(\delta_p-\omega)])$$

Note that the output amplitude noise spectrum consists of three components as seen for the result in Eq. (13), the noise contributed by the first part $S_1(\omega)$ is related to the amplitude noise of input probe light $S_X^{\text{in}}(\omega)$, the second part $S_2(\omega)$ is related to the phase noise of input probe light $S_Y^{\text{in}}(\omega)$, and the third part $S_3(\omega)$ arises from the continuous Langevin atomic noise resulted from the random decay process of atoms.

Normally the laser interacting with atoms is diode laser, for which the phase noise of input light is larger than the amplitude noise, thus we assume $S_Y^{\text{in}}=1$, $S_Y^{\text{in}}=30$ in Fig. 2. As shown in Fig. 2, the phase noise of the input light and the Langevin noise of the atomic medium would be primarily responsible for the output noise spectrum of the probe field. Thus it may provide a method for phase noise measurement. The M-type noise shape like the comb in the wings of the one-photon resonance of the probe field is caused by the phase noise of the input light, and the noise intensity is found to be suppressed below the shot noise level at the probe field detuning. For the sake of seeing the noise squeezing clearly, we plot the noise behaviors in Figs. 3, 4. We can see that the noise squeezing range becomes deep and wide with increasing of the dephasing between two ground states and keeps away from the one-photon resonance

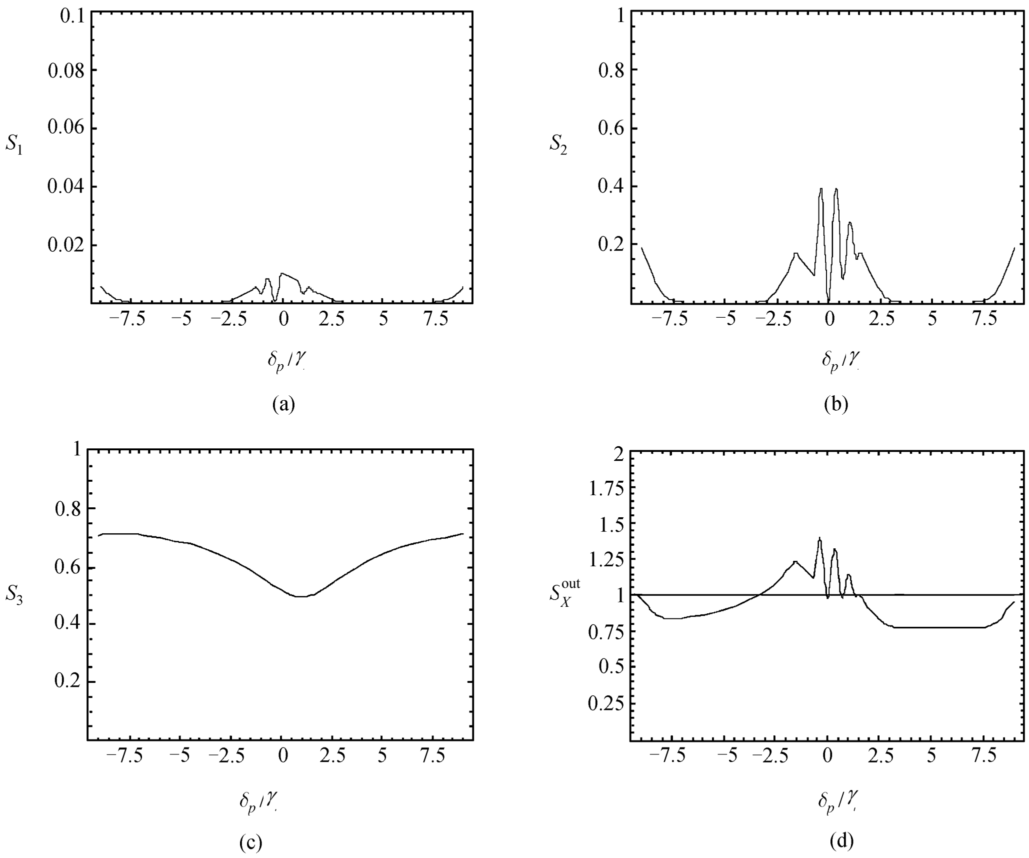

Fig. 2 The output amplitude noise spectrum: (**a**) the noise spectrum resulted from the amplitude fluctuation of the input field; (**b**) the noise spectrum resulted from the phase fluctuation of the input field; (**c**) the noise spectrum resulted from Langevin atomic noise; (**d**) the total amplitude noise spectrum and $S_X^{\text{out}}(\omega)=1$ represents shot noise level (SNL) of the output field. $|g|^2 N/c=100$, $\gamma_0/\gamma=0.5$, $\omega/\gamma=1$, $\Omega/\gamma=5$.

of the probe field with the increasing of the control field Rabi frequency Ω, If the deeper and wider noise squeezing is acquired as one wishes, we must increase the atomic collisions in principle. However, the atomic collisions will cause the atoms drifting out of the interaction region. The result of it is that the noise squeezing range does not largen and widen illimitably.

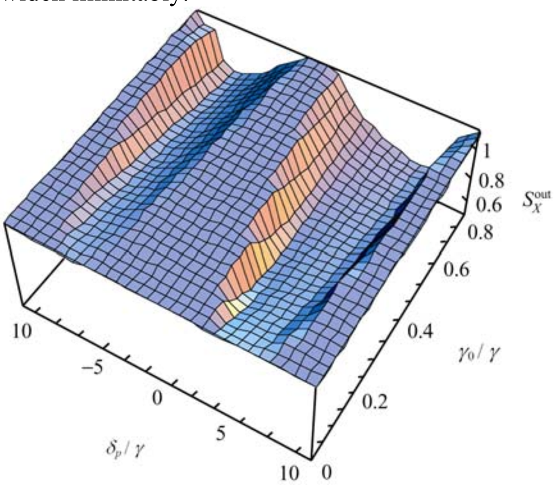

Fig. 3 The output noise spectrum of the probe field as the functions of normalized probe field detuning and the normalized dephasing of the ground-states. $|g|^2 N/c = 100,\ \omega/\gamma = 1,\ \Omega/\gamma = 5$.

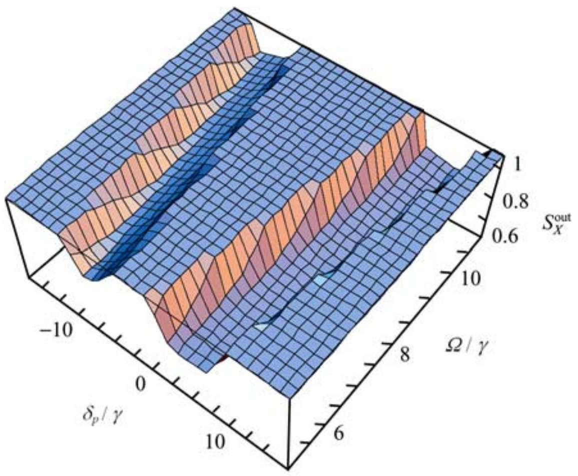

Fig. 4 The output noise spectrum of the probe field as the functions of normalized probed field detuning and normalized coupling field Rabi frequency. $|g|^2 N/c = 100,\ \gamma_0/\gamma = 0.5,\ \omega/\gamma = 1$.

3 Experiment

The atomic level used in this experiment is the Λ-type three-level system within the D_2 line of Cs atoms. A strong laser light called coupling beam couples the transition $6S_{1/2}\ (F_g = 4) \leftrightarrow 6P_{3/2}\ (F_e = 4)$. While a very weak laser called quantum probe beam scanned across the transition $6S_{1/2}\ (F_g = 3) \leftrightarrow 6P_{3/2}\ (F_e = 4)$. The frequency of coupling and probe laser lights are ω_c and ω_p, and they propagate in the same direction (i.e., Doppler-free configuration) through the Cs atoms cell. In this case the system is considered as a closed system since there is no more possibility for atoms in excited state decay to the other ground levels because of selection rules for electric dipole transition.

The experimental set-up is shown in Fig. 5. Two external cavity diode lasers (TOPTICA; DL100) with less than 1MHz line width are used as coupling and probe laser lights respectively. Both of the two beams have been spatially filtered to produce a quasi-Gaussian beam profile with a beam diameter of 2 mm. The coupling laser is locked to the resonant frequency of $6S_{1/2}\ (F_g = 4) \leftrightarrow 6P_{3/2}\ (F_e = 4)$ transition; it propagates through the optical isolator and half-wave plate, and then passes through a 5-cm-long Cs cell at room temperature. The probe light is scanned by a piezoelectric ceramic in the vicinity of the transition $6S_{1/2}\ (F_g = 3) \leftrightarrow 6P_{3/2}\ (F_e = 4)$, the beam passes through the other isolator and half-wave plate, and overlaps with the coupling laser by a polarizing beam splitter before the Cs cell. The extinction ratios of the polarizing beam splitter is greater than 25 dB, thus the polarization of two beams are well linearly polarized and their polarizations are orthogonal to each other. The power of two beams are 3.2 mW and 260 μW. The frequencies of two lasers are monitored using Doppler free saturated absorption spectroscopy. After passing through the Cs cell, a polarizing beam splitter separates the probe and coupling beams, only the probe beam is detected by photodiode for EIT effect, and the noise spectrum is measured by low noise detector and amplifier.

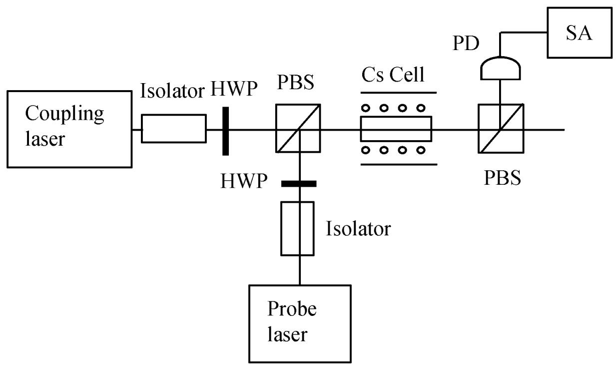

Fig. 5 Experimental setup. PBS: polarized beam splitter; PD: low-noise detector; HWP: half-wave plate; SA: the spectrum analyzer; Cs Cell: Cs sample cell.

Figure 6 shows the noise spectra of probe beam as a function of the probe detuning. The total Rabi frequencies of the coupling and probe lights are fixed to be $\Omega_c = 2\pi \times 9.8$ MHz and $\Omega_p = 2\pi \times 2.8$ MHz respectively, which is deduced from the definition about the Rabi frequency of light $\Omega = \Gamma \times \sqrt{\frac{I}{2I_0}}$, here $I = \frac{P}{\pi r^2}$ $\mathrm{mW/cm^2}$ is power density of the field, and the saturation intensity is $I_0 = 1.65$ mW/cm^2. The black line in Fig. 6 is the measured M-type noise spectrum, which prove that the phase noise of probe beam is converted into amplitude noise as analyzed for theoretical results. The red line is shot-noise limit, it is seen that at a bit large detuning, the noise of output probe beam is below shot-noise level 0.3 dB.

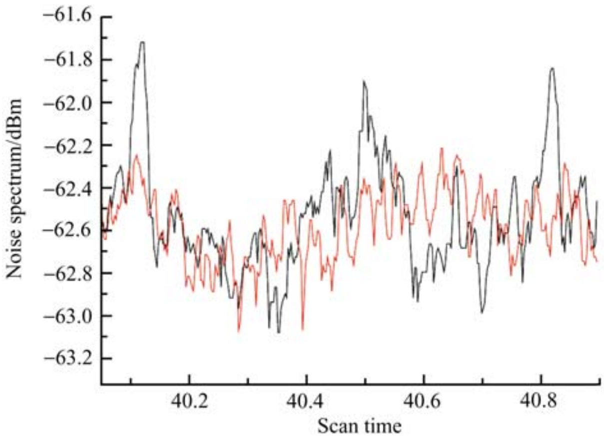

Fig. 6 The noise spectrum of the probe field vs probe detuning at the analyzing frequency of 5 MHz.

In conclusion, we have investigated the noise spectrum of a weak quantum probe field interacting with the atomic medium under the condition of electromagnetically induced transparency (EIT). The resulting noise root in the amplitude and phase fluctuations of the input field and the atomic noise which adds to the output field. In addition, the vital point is that if the proper parameters of the atom and light are selected, the noise of the output field is suppressed below the shot noise level at the detuning away from the resonance . What is more, the remotion of the noise squeezing range is controlled by the control field Rabi frequency Ω.

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