The difference in noise property between the Autler–Townes splitting medium and the electromagnetically induced transparent medium*

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The quantum noise of squeezed probe light passing through an atomic system with different electromagnetically induced transparency and Autler–Townes splitting effects is investigated theoretically. It is found that the optimal squeezing preservation of the outgoing probe beam occurs in the strong-coupling-field regime rather than in the weak-coupling-field regime. In the weak-coupling-field regime, which was recently recognized as the electromagnetically induced transparency regime (Abi-Salloum T Y 2010 *Phys. Rev.* A **81** 053836), the output amplitude noise is affected mainly by the atomic noise originating from the random decay process of atoms. While in the strong-coupling-field regime, defined as the Autler–Townes splitting regime, the output amplitude noise is affected mainly by the phase-to-amplitude conversion noise. This is useful in improving the quality of the experiment for efficient quantum memory, and hence has an application in quantum information processing.

Keywords: Autler–Townes splitting, electromagnetically induced transparency, quantum noise

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1. Introduction

Electromagnetically induced transparency (EIT) is a laser induced quantum interference effect originally proposed by Harris *et al.* in 1990,[1] and has received great attention in connection to the phenomena of slow light propagation and light storage.[2–4] The storage and retrieval of a weak light pulse in an EIT medium[3,5,6] and quantum memory in mapping a quantum state of light onto a long-lived atomic state[7–14] have been theoretically and experimentally demonstrated. Though the theoretical calculations showed that the quantum noise of the quantum state could be well preserved throughout an EIT medium or double EIT medium,[15–17] the undesirable excess noise was still introduced into the delayed output state after the interaction between the light and the atoms.[18] As a result, the quantum noise of the delayed output state in the EIT medium may not be kept the same as that of the input state, owing to its phase noise induced fluctuation in electric susceptibility, which in turn causes the fluctuation in transmitted intensity.[19] It has recently been revealed that the two similar but distinct phenomena of EIT and Autler–Townes (AT) splitting are always involved in the interaction process between light and atoms,[20] and the discussion over these two effects in a three-level system shows that EIT and AT splitting are both characterized by a reduction in the absorption of a weak field in the presence of a stronger field. EIT is a result of destructive interference between two competing excitation pathways, while AT splitting is just a result of splitting in the absorption line induced by a relatively strong coupling field. As described in Ref. [20], EIT and AT splitting can be distinguished by the threshold factor, which is defined as the Rabi frequency of the coupling field divided by the decay rate of the probe transition. If the Rabi frequency of the coupling field is much greater than the polarization decay rate, which is called the strong-coupling-field regime, the absorption spectrum of the probe light exhibits AT splitting.[21] While in the weak-coupling-regime, where the Rabi frequency of the coupling field is less than the polarization decay rate, the absorption spec-

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trum of the probe light is shown by EIT. Recently, AT-based slow light has been investigated in inhomogeneously broadened quantum dot media, where the EIT effect cannot be generated.[22] Sheremet *et al.*[23] considered the quantum memory in an AT resonance structure created by a strong control field, and found that it was more efficient to control the AT splitting to obtain more a effective delay of the signal pulse and high memory efficiency. The distinctions between EIT and AT splitting inspire us to explore the role of quantum interference in the noise property for the two effects. In this paper, we study the quantum noise spectrum of the outgoing probe light through a medium showing the EIT effect in the weak-coupling-field regime and the AT splitting effect in the strong-coupling-field regime, with an initial 3 dB amplitude squeezed vacuum state. It is shown that the optimum condition for quantum state preservation occurs in the AT splitting regime.

2. Theoretical model

Consider a closed three-level Λ-type system as shown in Fig. 1. The atoms have two metastable lower states $|b\rangle$ and $|c\rangle$, and one excited state $|a\rangle$. $\hat{a}$ is the annihilation operator for the weak probe light of the quantum field that couples the transition $|b\rangle \leftrightarrow |a\rangle$, while Ω is the Rabi frequency of the strong coupling field that couples the transition $|c\rangle \leftrightarrow |a\rangle$. ν is the frequency of the probe light and ν_c is the frequency of the coupling light. The one-photon detuning is given by $\Delta_p = \omega_{ab} - \nu$. Here, we have set the coupling field resonant with its corresponding transition. We give the evolution equations for both the slowly varying atomic operators and the slowly varying annihilation operator of the quantum probe field $\hat{a}$ as

$$\frac{\partial}{\partial t}\hat{\sigma}_{ba} = -(i\Delta_p + \gamma_{ba})\hat{\sigma}_{ba} + ig\hat{a}(\hat{\sigma}_{bb} - \hat{\sigma}_{aa}) + i\Omega\hat{\sigma}_{bc} + \hat{F}_{ba}, \tag{1a}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{bc} = -(i\Delta_p + \gamma_{bc})\hat{\sigma}_{bc} - ig\hat{a}\hat{\sigma}_{ac} + i\Omega^*\hat{\sigma}_{ba} + \hat{F}_{bc}, \tag{1b}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{ca} = -\gamma_{ca}\hat{\sigma}_{ca} + ig\hat{a}\hat{\sigma}_{cb} - i\Omega(\hat{\sigma}_{aa} - \hat{\sigma}_{cc}) + \hat{F}_{ca}, \tag{1c}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{bb} = \gamma_b\hat{\sigma}_{aa} + i(g^*\hat{a}^+\hat{\sigma}_{ba} - g\hat{a}\hat{\sigma}_{ab}) + \hat{F}_{bb}, \tag{1d}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{cc} = \gamma_c\hat{\sigma}_{aa} + i(\Omega^*\hat{\sigma}_{ca} - \Omega\hat{\sigma}_{ac}) + \hat{F}_{cc}, \tag{1e}$$

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{a}(z,t) = ig^*N\hat{\sigma}_{ba}(z,t), \tag{1f}$$

where we have included the decays γ_{ba} and γ_{ca} of the atomic dipole operators, the spontaneous decays γ_b and γ_c, and the associated Langevin noise operators describing the effect of spontaneous decay caused by the coupling of atoms to all the vacuum field modes. γ_{bc} is the dephasing rate for the two ground states of this system. $g = \wp_{ab}\sqrt{\nu/2\varepsilon_0 V\hbar}$ is the atom-field coupling constant with $\wp_{ab}$ being the atomic dipole moment for the $|b\rangle \leftrightarrow |a\rangle$ transition and V the interaction volume. N is the number of atoms. In Eqs. (1d) and (1e), we did not consider the population transfer between the two metastable lower states.

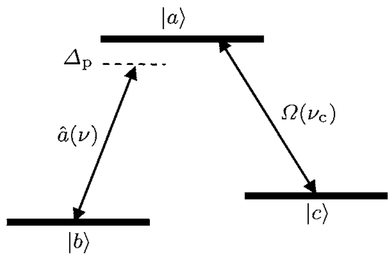

Fig. 1. Schematic of the Λ-type system.

On the assumption that the intensity of the quantum probe field is much less than that of the classical coupling field, and all the atoms are initially in the ground state $|b\rangle$, which satisfies $\langle\hat{\sigma}_{bb}\rangle = 1$, we can solve Eqs. (1a) and (1b) perturbatively in $\hat{a}$, and obtain the first order equations

$$\frac{\partial}{\partial t}\hat{\sigma}_{ba} = -(i\Delta_p + \gamma_{ba})\hat{\sigma}_{ba} + ig\hat{a} + i\Omega\hat{\sigma}_{bc} + \hat{F}_{ba}, \tag{2a}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{bc} = -(i\Delta_p + \gamma_{bc})\hat{\sigma}_{bc} + i\Omega^*\hat{\sigma}_{ba} + \hat{F}_{bc}. \tag{2b}$$

Make the Fourier transform of Eqs. (2a) and (2b), and solve them for $\hat{\sigma}_{ba}(z,\omega)$, substitute it into Eq. (1f) and perform formal integration over z,, then we will obtain the output probe field at the exit of the vapour cell with the length L as

$$\hat{a}(L,\omega) = e^{-\Lambda(\omega)L}\hat{a}(0,\omega) + \frac{g^*N}{c}\int_0^L e^{-\Lambda(\omega)(L-s)} \times \frac{i\gamma_2\hat{F}_{ba}(s,\omega) - \Omega\hat{F}_{bc}(s,\omega)}{\gamma_1\gamma_2 + |\Omega|^2}ds, \tag{3}$$

where ω is the detection frequency,

$$\Lambda(\omega) = \frac{|g|^2 N}{c} \times \frac{\gamma_2}{\gamma_1\gamma_2 + |\Omega|^2} - \frac{i\omega}{c}, \tag{4}$$

$\gamma_1 = \gamma_{ba} + i(\Delta_p - \omega)$, and $\gamma_2 = \gamma_{bc} + i(\Delta_p - \omega)$.

Define the amplitude and phase quadratures $\hat{X}(z,t)$ and $\hat{Y}(z,t)$ of the probe light as

$$\hat{X}(z,t) = \hat{a}(z,t) + \hat{a}^{+}(z,t), \quad (5a)$$

$$\hat{Y}(z,t) = -\mathrm{i}[\hat{a}(z,t) - \hat{a}^{+}(z,t)]. \quad (5b)$$

In order to calculate the noise spectrum of the output probe, we use the quadrature flux spectrum

$$\left\langle \hat{X}(L,\omega)\hat{X}(L,\omega')\right\rangle = \frac{2\pi L}{c}\delta(\omega+\omega')S_X(L,\omega), \quad (6a)$$

$$\left\langle \hat{Y}(L,\omega)\hat{Y}(L,\omega')\right\rangle = \frac{2\pi L}{c}\delta(\omega+\omega')S_Y(L,\omega), \quad (6b)$$

and the correlation functions of the Langevin noise operators, which can be calculated according to the quantum regression theorem[15,19]

$$\begin{aligned}&\left\langle \hat{F}_{\mu\nu}(z_1,t_1)\hat{F}_{\alpha\beta}(z_2,t_2)\right\rangle \\ &= \frac{L}{N}\left\langle D(\hat{\sigma}_{\mu\nu}\hat{\sigma}_{\alpha\beta}) - D(\hat{\sigma}_{\mu\nu})\hat{\sigma}_{\alpha\beta} - \hat{\sigma}_{\mu\nu}D(\hat{\sigma}_{\alpha\beta})\right\rangle \\ &\quad\times \delta(z_2-z_1)\delta(t_2-t_1), \end{aligned} \quad (7)$$

where $D(\hat{\sigma}_{\mu\nu})$ denotes the deterministic part of the Heisenberg–Langevin equation of motion for $\hat{\sigma}_{\mu\nu}$ without the Langevin noise term. The Dirac delta function in Eq. (7) represents the short memory of the reservoir of vacuum modes. Using the quantum regression theorem to calculate the correlation function of $\hat{F}(s,\omega)$ according to Eq. (7), then we will obtain

$$\begin{aligned}&\left\langle \hat{F}_{\mathrm{ba}}(z_1,\omega)\hat{F}^{+}_{\mathrm{ba}}(z_2,\omega')\right\rangle \\ &= \frac{L\delta(z_1-z_2)\delta(\omega+\omega')}{N}2\gamma_{\mathrm{ba}},\end{aligned} \quad (8a)$$

$$\begin{aligned}&\left\langle \hat{F}_{\mathrm{bc}}(z_1,\omega)\hat{F}^{+}_{\mathrm{bc}}(z_2,\omega')\right\rangle \\ &= \frac{L\delta(z_1-z_2)\delta(\omega+\omega')}{N}2\gamma_{\mathrm{bc}}.\end{aligned} \quad (8b)$$

Using Eqs. (6)–(8), one obtains the normalized quadrature amplitude spectrum of the output quantum probe related to the input via the relation

$$S_X(L,\omega) = S_1(\omega) + S_2(\omega) + S_3(\omega), \quad (9)$$

with

$$\begin{aligned} S_1(\omega) = {} & \frac{S_X^{\mathrm{in}}(\omega)}{4}(\exp\{-[\Lambda(\omega)+\Lambda(-\omega)]L\} \\ & + \exp\{-[\Lambda(\omega)+\Lambda^*(\omega)]L\} \\ & + \exp\{-[\Lambda^*(-\omega)+\Lambda(-\omega)]L\} \\ & + \exp\{-[\Lambda^*(-\omega)+\Lambda^*(\omega)]L\}), \\ S_2(\omega) = {} & -\frac{S_Y^{\mathrm{in}}(\omega)}{4}(\exp\{-[\Lambda(\omega)+\Lambda(-\omega)]L\} \\ & - \exp\{-[\Lambda(\omega)+\Lambda^*(\omega)]L\} \\ & - \exp\{-[\Lambda^*(-\omega)+\Lambda(-\omega)]L\} \\ & + \exp\{-[\Lambda^*(-\omega)+\Lambda^*(\omega)]L\}), \\ S_3(\omega) = {} & 1 - \exp[-2\mathrm{Re}(\Lambda(\omega)L)]. \end{aligned}$$

It can be seen that the output amplitude noise (9) is composed of three parts. The first part, $S_1(\omega)$, represents the contribution of the amplitude noise spectrum of the input probe $S_X^{\mathrm{in}}(\omega)$; the second part, $S_2(\omega)$, represents the contribution of the phase noise spectrum of the input probe $S_Y^{\mathrm{in}}(\omega)$ via the phase-to-amplitude noise conversion;[24–26] and the last part, i.e. $S_3(\omega)$, arises from the Langevin atomic noise due to the random decay process of atoms. The three parts have different influences on the output amplitude noise of the probe.

3. Results and discussion

In our calculations and discussion, the parameters, Ω, Δ_{p}, γ_{bc}, $|g|^2 NL/c$ and ω have been normalized to the decay γ_{ba} of the atomic transition $|b\rangle \leftrightarrow |a\rangle$. The spectral component of the probe beam at $\omega = 0$ for a 3 dB squeezed input beam passing through the system is plotted in Fig. 2 as a function of probe detuning Δ_{p} with $\gamma_{\mathrm{bc}} = 0.01$, $|g|^2 NL/c = 25$, $S_X^{\mathrm{in}}(\omega) = 0.5$, and $S_Y^{\mathrm{in}}(\omega) = 2$. Figure 2(a) is for the strong-coupling-field regime $\Omega = 3$ for AT splitting and Fig. 2(b) is for the weak-coupling-field regime $\Omega = 0.6$ for EIT. The green thick solid lines represent the corresponding transparency windows, the red thin solid lines denote the total output amplitude noise $S_X(L,0)$ values, the black dashed lines, blue dotted lines and magenta dash-dotted lines refer to the noise spectral components of $S_1(0)$, $S_2(0)$ and $S_3(0)$, respectively. Note that the vertical coordinates for $S_1(0)$, $S_2(0)$, $S_3(0)$ and $S_X(L,0)$ are different in the case of the weak-coupling-field with $\Omega = 0.6$. $S_X(L,0) = 1$ represents the shot noise level (SNL) of the output field.

It can be clearly seen that the total amplitude noises of the output probe exhibit resonant line shapes for the two cases of AT splitting and EIT regimes, but with different amplitudes. Comparing the noise spectra for AT splitting and EIT regimes, we find that the bandwidths of the transparency window and the squeezing below SNL in the AT splitting regime (Fig. 2(a)) are much wider than those in the EIT regime (Fig. 2(b)), and meanwhile the squeezing at $\Delta_{\mathrm{p}} = 0$ in the AT splitting regime is much larger than that in the EIT regime.

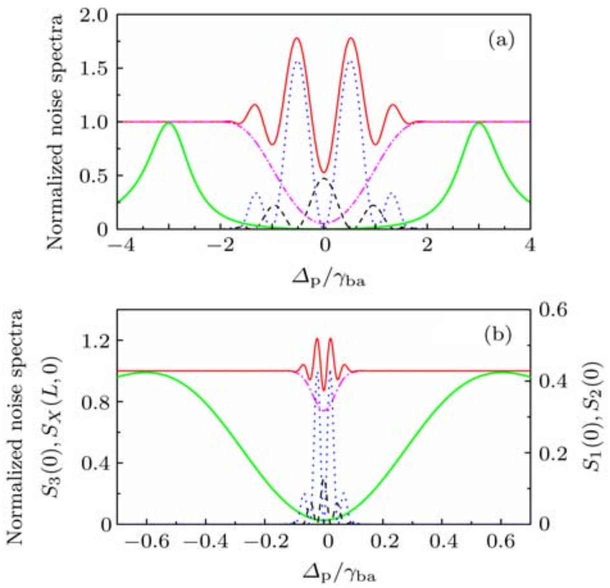

Fig. 2. (colour online) Curves for output amplitude noise and probe absorption versus probe detuning. The green thick solid lines refer to the probe absorption. The red thin solid lines denote the total output amplitude noise. $S_X(L, 0)$, the black dashed, blue dotted and magenta dash-dotted lines represent the noises of $S_1(0)$, $S_2(0)$ and $S_3(0)$,, respectively. (a) $\Omega = 3$,, (b) $\Omega = 0.6$.

For the AT splitting in Fig. 2(a), the output noise at $\Delta_p = 0$ is mainly from the contribution of the input amplitude noise $S_1(0)$, while the contributions of phase-to-amplitude conversion noise $S_2(0)$ and atomic noise $S_3(0)$ are almost zero. Therefore, the noise of the output probe beam nearly equals that of the input one, that is, the squeezing of the probe beam passing through the atomic medium can be well preserved. In the case of off resonance ($\Delta_p \neq 0$), the contribution of the input amplitude noise $S_1(0)$ decreases, but the phase-to-amplitude noise $S_2(0)$ and the atomic noise $S_3(0)$ increase monotonically with probe detuning, leading to the increase in output noise.

For the case of EIT in Fig. 2(b), although the phase-to-amplitude conversion noise is zero and the contribution of the input amplitude noise $S_1(0)$ is much smaller than 0.5 at $\Delta_p = 0$,, the total amplitude noise becomes larger because of the high atomic noise $S_3(0)$, which is related to the random decay process of atoms, and adds noise to the probe field as it interacts with the atoms. When the probe detuning becomes larger but still remains in the transparency window, the probe absorption and atomic noise increase gradually, leading to the increase in output noise of the probe light.

Note that at and near resonance $\Delta_p = 0$, the probe absorption becomes close to zero in the AT splitting regime, which results in the small Langevin noise from the atoms. However, in the EIT regime, due to the imperfect transparency accompanied with some absorption, input squeezing is contaminated by the atomic noise. Thus, we can conclude that it is better to manipulate squeezing preservation in the AT splitting regime rather than in the EIT regime. Accordingly, we deduce that the well squeezing preservation in the atomic medium will happen in the AT splitting regime.

In the physically realistic scheme, though the best squeezing preservation occurs at zero detection frequency ($\omega = 0$), the relaxation oscillation of the laser at low frequency[27] prevents us from detecting the squeezing at zero frequency in the process of squeezing measurement. In what follows, we will focus on the effect of detection frequency on the probe beam, as illustrated in Fig. 3. According to the detection frequency usually used in the squeezing measurement experiment, we take $\omega = 0.3$, which corresponds to about 4.5 MHz. In the AT splitting regime, though the output amplitude noise line (Fig. 3(a)) shows a similar oscillation line shape to that in the case of zero detection frequency, we see that two squeezing points exist, with the squeezing values being close to each other. It is also clearly seen that the two squeezing points correspond to the positions of the minimum phase-to-amplitude conversion noises. In terms of expression (9), the noise $S_1(\omega)$ which is related to the input amplitude noise and the phase-to-amplitude conversion noise $S_2(\omega)$ are always symmetrical about the probe detuning Δ_p for different nonzero detection frequencies, while the atomic noise $S_3(\omega)$ changes with detection frequency. In the EIT regime (Fig. 3(f)), however, the output amplitude noise $S_X(L, \omega)$ is determined mainly by atomic noise $S_3(\omega)$, so the minimum squeezing position is also the position of the dip of $S_3(\omega)$. It is also important to note that for the case of the intermediate Rabi frequency of coupling light, which is between the strong and weak coupling regimes (see Figs. 3(b)–3(e)), a hybrid regime is comprised of EIT and AT splitting simultaneously as described in Refs. [28] and [29]. If we reduce the Rabi frequency of the coupling light, then the atomic noise $S_3(0)$ will increase and the phase-to-amplitude conversion noise $S_2(0)$ will decrease gradually at off resonance. We can also see the changes from one to two squeezing windows if the Rabi frequency of the coupling light increases. From the above analyses, we can conclude that the result with two squeezing points will not occur in the EIT regime.

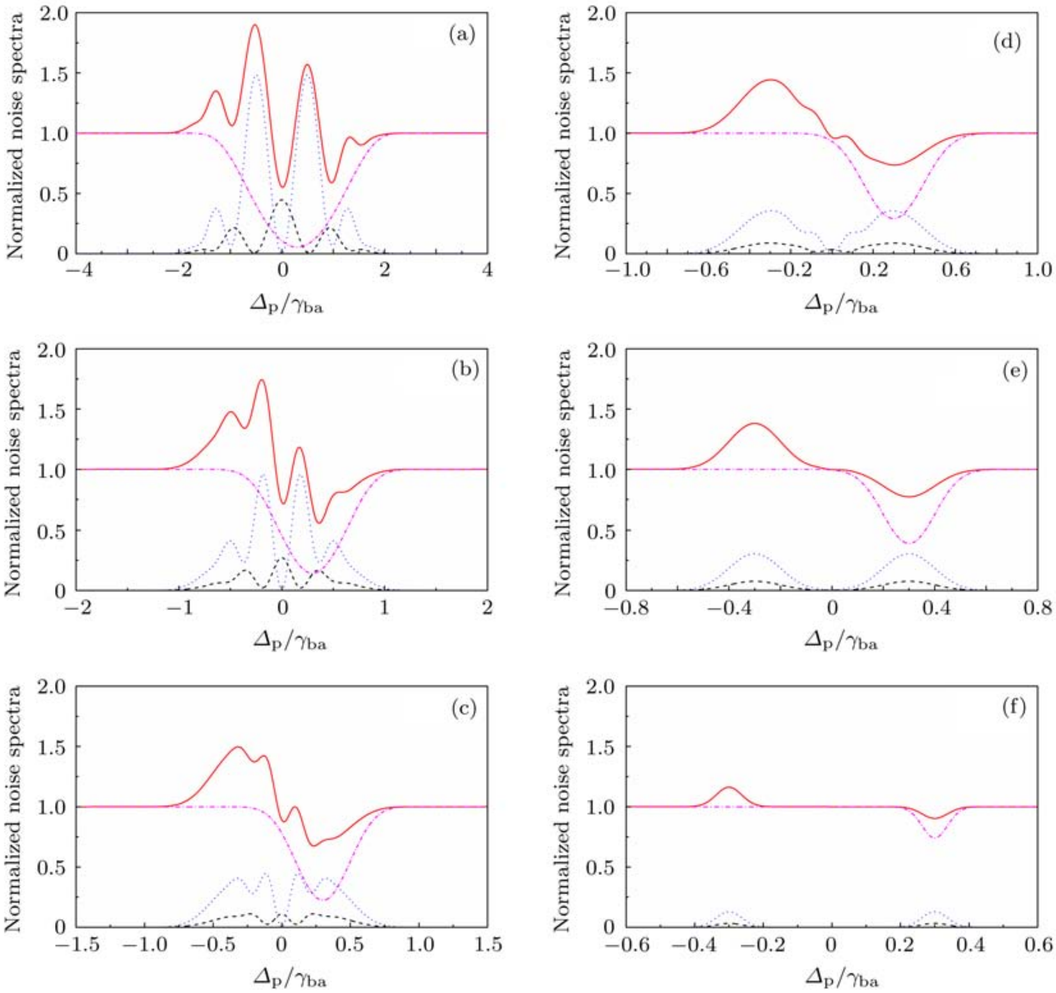

Fig. 3. (colour online) Curves for output amplitude noise versus probe detuning with $\omega = 0.3$. The red solid lines are the total output amplitude noise $S_X(L, 0.3)$, the black dashed, blue dotted, and magenta dash-dotted lines represent the noises of $S_1(0.3)$, $S_2(0.3)$ and $S_3(0.3)$,, respectively. Panels (a)–(f) show the amplitude noise changing from the case of AT splitting to EIT at (a) $\Omega = 3$, (b) $\Omega = 1.8$, (c) $\Omega = 1.4$, (d) $\Omega = 1.2$, (e) $\Omega = 1$ and (f) $\Omega = 0.6$. The other parameters are the same as those in Fig. 2.

The dependences of the output noise on detection frequency for different probe detunings are plotted in Figs. 4(a) and 4(b). As discussed above, there exist two squeezing points corresponding to the minimum phase-to-amplitude conversion noises in the AT splitting regime. One exists at $\Delta_{\rm p} = 0$ and the other at $\Delta_{\rm p} \approx 1$. The two points give us two channels for quantum state preservation with low noise. When the probe beam is resonant with its corresponding optical transition, the output amplitude noise exhibits the same dip as an EIT window, as described in Refs. [16], [17] and [19]. With the increase in detection frequency, the maximum squeezing decreases. Beyond the squeezing window the probe field displays shot noise. When the probe detuning is set as $\Delta_{\rm p} = 1$, which indicates that the phase-to-amplitude conversion noise becomes close to its minimum, the output amplitude noise exhibits a wider squeezing window than the previous one. Thus, we can use the two squeezing windows occurring at different probe detunings to keep the initial input squeezing, though the maximum squeezings are below the input one. While for the case of EIT, the output amplitude noise exhibits a similar dip to that in the case of AT splitting, but the dip is less squeezing and has a narrower squeezing bandwidth in the former than the latter case. If we set probe detuning to be 0.3, then we will see that the squeezing window varies with probe detuning, and that minimum squeezing occurs at $\omega = 0.3$. The reason for this is that we can view $\Delta_{\rm p} - \omega$ as the detuning of the probe field in the frequency domain,[17] and the minimum squeezing determined by the atomic noise occurs at the transparent point satisfying $\Delta_{\rm p} - \omega = 0$ for the case of the weak-coupling-field.[19] It can also be seen that the minimum squeezing at nonzero probe detuning is less than that at resonance, which corresponds to large absorption.

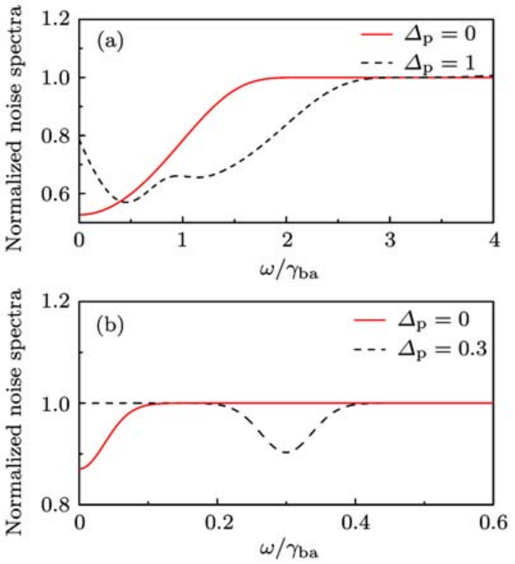

Fig. 4. (colour online) Curves for output amplitude noise versus detection frequency for different probe detunings (a) at $\Omega = 3$ with the red solid line for $\Delta_p = 0$ and the black dashed line for $\Delta_p = 1$, and (b) at $\Omega = 0.6$ with the red solid line for $\Delta_p = 0$ and the black dashed line for $\Delta_p = 0.3$. The other parameters are the same as those in Fig. 2.

4. Conclusion

We theoretically compared the output amplitude noise passing through an atomic medium in the EIT regime and that in the AT splitting regime. It is found that the Λ-type atomic system with a strong coupling for only the AT splitting effect is more suitable for the experiments of quantum state preservation and retrieval than the system with weak coupling for the EIT effect.

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