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Quantum interference effects on the probe amplification without and with inversion in a four-level system*

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A four-level system driven by two coherent fields is considered. It is shown that in the presence of an incoherent pump, the probe gain at a short wavelength can be achieved due to the quantum interference. Our density matrix calculation provides the conditions for probe amplification from different origins, including gain without population inversion on any state basis, gain with population inversion on the dressed-state basis, and gain with population inversion on the bare-state basis. Also, by controlling the Rabi frequency of the coupling field a total change from non-inversion to inversion can be achieved which does not depend on the intensity of the incoherent pump.

Keywords: quantum interference, probe gain, inversion, non-inversion
PACC: 4250, 3150

1. Introduction

Quantum interference resulting from indistinguishable quantum transition pathways renders many atomic coherence effects, such as coherent trapping of atomic population in a ground state or in an excited state,[1−5] electromagnetically induced transparency,[6−9] amplification (laser) without inversion (AWI)[10−21] and so on. During the last two decades, substantial efforts have been made to study atomic coherence effects in laser physics and quantum optics.[10] Lasing without inversion (LWI) is an interesting and potentially useful manifestation of these effects for it could provide a new tool in pursuit of the short-wavelength laser. Several experiments in the visible spectral region[21−24] have confirmed the initial ideas of LWI. The origin of the inversionless gain can be attributed to either inversion between dressed states or coherence among these states resulting from a strong coherent field.[25,26] Recently, Dong[27] and Wu *et al*[28] have shown that the probe gain without inversion can be achieved in a four-level model. In this model, population inversion or non-inversion depends on both coherent field and incoherent pumping.

In this paper, we propose a new four-level atomic model for achieving the probe gain at a short wavelength, where a metastable state is included; atoms in a ground state are pumped into a middle level by a coherent field and then into the top state by an incoherent pumping.[27] In such a case, population distribution in the upper level of the probe transition, i.e. the top level, becomes much easier to achieve in practical cases. From analytical as well as numerical calculations, we find that the probe amplification displays a transition from inversion to non-inversion which depends on Rabi frequency of the coherent field and the spontaneous decay rates, but not on intensity of the incoherent pumping. This characteristic is different from that in the models presented by Dong and Wu. Here, by showing the population distribution, the explicit gain process, which has been divided into three catalogues—gain without population inversion on any state basis, gain with population inversion

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on the dressed-state basis, and gain with population inversion on the bare-state basis—gives us a fundamental understanding of the amplification origin in the system. In particular, we distinguish the different gain mechanisms in the model and derive conditions in the steady state under which the different gain mechanisms may take place.

2. The atomic model and the density-matrix equation

We consider a closed four-level atomic system, shown in Fig.1. In this model, a coherent coupling field E_c with Rabi frequency Ω_c drives transition $|3\rangle \leftrightarrow |1\rangle$, and populates state $|2\rangle$ via a decay from $|3\rangle$to $|2\rangle$. An incoherent pump at a rate Λ is applied to the transition from metastable state $|2\rangle$ to excited state $|4\rangle$ and prepares the atoms in order to generate a gain on the transition labelled with $|4\rangle - |1\rangle$. The transition $|2\rangle \leftrightarrow |1\rangle$ is forbidden. A weak probe field E_p with Rabi frequency Ω_p is used to probe the possible gain on the transition $|1\rangle \leftrightarrow |4\rangle$.

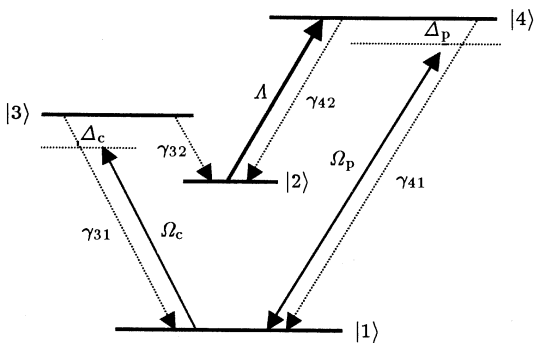


Fig.1. Schematic diagram of a four-level atomic system driven by a coherent field and an incoherent pumping.

In the framework of the semiclassical theory, by the standard density-matrix formalism in the dipole approximation and the rotating wave approximation, the interaction Hamiltonian H_I can be represented in an interaction picture

$$\begin{aligned} H_I =& \Delta_c|3\rangle\langle 3| + \Delta_p|4\rangle\langle 4| \\ &+ (\Omega_c|3\rangle\langle 1| + \Omega_p|4\rangle\langle 1| + \text{c.c.}), \end{aligned} \tag{1}$$

where the detunings of the probe and the coherent field are defined as $\Delta_p = \omega_{41} - \omega_p$, and $\Delta_c = \omega_{31} - \omega_c$, respectively; the Rabi frequencies Ω_p and Ω_c are defined as $\Omega_p = E_p\mu_{14}/2\hbar$ and $\Omega_c = E_c\mu_{13}/2\hbar$, respectively.

Assuming $\hbar = 1$ for simplicity, the master equation of motion for the density operator in an arbitrary multilevel atomic system in an interaction picture can be written as

$$\frac{\partial \rho}{\partial t} = -\mathrm{i}[H_I, \rho] + \Lambda\rho. \tag{2}$$

Expanding Eq.(2), we can easily arrive at their equation of motion

$$\begin{aligned} \dot\rho_{11} =& \Gamma_{21}\rho_{22} + \Gamma_{31}\rho_{33} + \Gamma_{41}\rho_{44} + \mathrm{i}\Omega_p\rho_{14} \\ &+ \mathrm{i}\Omega_c\rho_{13} - \mathrm{i}\Omega_p^*\rho_{41} - \mathrm{i}\Omega_c^*\rho_{31}, \\ \dot\rho_{22} =& -\Gamma_{21}\rho_{22} + \Gamma_{32}\rho_{33} + \Gamma_{42}\rho_{44} + \Lambda(\rho_{44} - \rho_{22}), \\ \dot\rho_{33} =& -\Gamma_{31}\rho_{33} - \Gamma_{32}\rho_{33} + \mathrm{i}\Omega_c^*\rho_{31} - \mathrm{i}\Omega_c\rho_{13}, \\ \dot\rho_{41} =& -\gamma_{41}\rho_{41} - \mathrm{i}\Delta_p\rho_{41} + \mathrm{i}\Omega_p(\rho_{44} - \rho_{11}) + \mathrm{i}\Omega_c\rho_{43}, \\ \dot\rho_{43} =& -\gamma_{43}\rho_{43} + \mathrm{i}\Delta_c\rho_{43} - \mathrm{i}\Delta_p\rho_{43} + \mathrm{i}\Omega_c^*\rho_{41} - \mathrm{i}\Omega_p\rho_{13}, \\ \dot\rho_{31} =& -\gamma_{31}\rho_{31} + \mathrm{i}\Omega_c(\rho_{33} - \rho_{11}) - \mathrm{i}\Delta_c\rho_{31} + \mathrm{i}\Omega_p\rho_{34}, \\ \dot\rho_{21} =& -\gamma_{21}\rho_{21} + \mathrm{i}\Omega_p\rho_{24} + \mathrm{i}\Omega_c\rho_{23}, \\ \rho_{11} &+ \rho_{22} + \rho_{33} + \rho_{44} = 1, \\ \rho_{ij} =& \rho_{ji}^*, \end{aligned} \tag{3}$$

where Γ_{ij} designates the spontaneous damping of population from i to j; γ_{ij} is referred to as coherence decay rate between i and j. To simplify the calculation, we have set the driving laser detuning as zero. In the radiative limit, the decay rates of the coherence are given by

$$\begin{aligned} \gamma_{41} =& \frac{1}{2}(\Gamma_{41} + \Gamma_{42} + \Gamma_{43} + \Lambda), \\ \gamma_{42} =& \frac{1}{2}(\Gamma_{41} + \Gamma_{42} + \Gamma_{43} + 2\Lambda), \\ \gamma_{43} =& \frac{1}{2}(\Gamma_{41} + \Gamma_{42} + \Gamma_{43} + \Gamma_{31} + \Gamma_{32} + \Lambda), \\ \gamma_{32} =& \frac{1}{2}(\Gamma_{31} + \Gamma_{32} + \Lambda), \\ \gamma_{31} =& \frac{1}{2}(\Gamma_{31} + \Gamma_{32}), \\ \gamma_{21} =& \frac{1}{2}\Lambda. \end{aligned} \tag{4}$$

In the limit of a weak probe and under the steady-state condition, the solution of Eq.(1) for ρ_{41} to the first order and for the other ρ_{ij} to the order in the probe field are

$$\begin{aligned} \rho_{11} =& \frac{\gamma_{31}}{2|\Omega_c|^2}\left(\Gamma_{31} + \Gamma_{32} + \frac{2|\Omega_c|^2}{\gamma_{31}}\right) \\ &\times \frac{\Gamma_{41}\Lambda}{\Gamma_{32}(\Gamma_{41} + \Gamma_{42} + \Lambda)}\rho_{22}, \\ \rho_{33} =& \frac{\Gamma_{41}\Lambda}{\Gamma_{32}(\Gamma_{41} + \Gamma_{42} + \Lambda)}\rho_{22}, \\ \rho_{44} =& \frac{\Lambda}{\Gamma_{41} + \Gamma_{42} + \Lambda}\rho_{22}, \\ \rho_{13} =& \frac{\mathrm{i}\Omega_c}{\gamma_{31} + \mathrm{i}\Delta_c}(\rho_{33} - \rho_{11}), \end{aligned} \tag{5}$$

and

$$\rho_{41} = \frac{\mathrm{i}\Omega_{\mathrm{p}}(\rho_{44}-\rho_{11}) + \Omega_{\mathrm{c}}\Omega_{\mathrm{p}}\rho_{13}/(\gamma_{43}+\mathrm{i}\Delta_p)}{\gamma_{41}+\mathrm{i}\Delta_{\mathrm{p}}+\dfrac{|\Omega_{\mathrm{c}}|^2}{\gamma_{43}+\mathrm{i}\Delta_{\mathrm{p}}}}$$
$$= \frac{\mathrm{i}\Omega_{\mathrm{p}}(\rho_{44}-\rho_{11}) - \mathrm{i}\Omega_{\mathrm{p}}|\Omega_{\mathrm{c}}|(\rho_{33}-\rho_{11})/\gamma_{31}(\gamma_{43}+\mathrm{i}\Delta_p)}{\gamma_{41}+\mathrm{i}\Delta_{\mathrm{p}}+\dfrac{|\Omega_{\mathrm{c}}|^2}{\gamma_{43}+\mathrm{i}\Delta_{\mathrm{p}}}}. \quad (6)$$

3. Solutions and analysis

The indices of refraction and absorption for the probe field are governed by the real and imaginary parts of the complex polarization ρ_{41}. The gain coefficient for the probe coupled to the transition $|4\rangle \leftrightarrow |1\rangle$ is proportional to $\mathrm{Im}(\rho_{41})$. If $\mathrm{Im}(\rho_{41}) > 0$, the system exhibits gain for the probe field, while for $\mathrm{Im}(\rho_{41}) < 0$ the probe field is attenuated. For this model, the bare-state basis consists of the states $|1\rangle, |2\rangle, |3\rangle$, and $|4\rangle$, while the dressed-state basis consists of the states $|+\rangle, |-\rangle, |2\rangle$ and $|4\rangle$. $|+\rangle(|-\rangle)$ is the superposition of states $|1\rangle$ and $|3\rangle$: $|+\rangle = \frac{1}{\sqrt{2}}(|3\rangle + |1\rangle)$ and $|-\rangle = \frac{1}{\sqrt{2}}(|3\rangle - |1\rangle)$. It is seen from Eq.(6) that the probe gain ($\propto \mathrm{Im}(\rho_{41}/\Omega_{\mathrm{p}})$) is contributed by two terms: the first term is proportional to $\rho_{44}-\rho_{11}$, and the second term is due to the induced coherence ρ_{13}. In the following, we derive the analytical results under the resonant excitation with a strong driving field, and then present the numerical calculations for more general situation. For the convenience of the calculation, Ω_{p} and Ω_{c} are chosen to be real. When $\Delta_{\mathrm{p}} = \Delta_{\mathrm{c}} = 0$, the solution of $\mathrm{Im}(\rho_{41}/\Omega_{\mathrm{p}})$ is

$$\mathrm{Im}\left(\frac{\rho_{41}}{\Omega_{\mathrm{p}}}\right) = \frac{(\rho_{44}-\rho_{11}) - \dfrac{|\Omega_{\mathrm{c}}|^2}{\gamma_{43}\gamma_{31}}(\rho_{33}-\rho_{11})}{\gamma_{41}+\dfrac{|\Omega_{\mathrm{c}}|^2}{\gamma_{43}}}. \quad (7)$$

The steady-state population difference between state $|3\rangle$ and $|1\rangle$ is

$$\rho_{33}-\rho_{11} = -\frac{\gamma_{31}(\Gamma_{31}+\Gamma_{32})\Gamma_{41}\Lambda}{2|\Omega_{\mathrm{c}}|^2\Gamma_{32}(\Gamma_{41}+\Gamma_{42}+\Lambda)}\rho_{22} < 0. \quad (8)$$

The steady-state population difference between state $|4\rangle$ and $|1\rangle$ is

$$\rho_{44}-\rho_{11} = \frac{(\Gamma_{32}-\Gamma_{41})\Lambda - \gamma_{31}(\Gamma_{31}+\Gamma_{32})\Gamma_{41}\Lambda/2|\Omega_{\mathrm{c}}|^2}{\Gamma_{32}(\Gamma_{41}+\Gamma_{42}+\Lambda)}. \quad (9)$$

It is easy to see from Eqs.(7) and (8) that the second term in $\mathrm{Im}(\rho_{41}/\Omega_{\mathrm{p}})$ is always positive. When the negative contribution from the first term is less than the second positive contribution, or the first term is positive, the probe could be amplified under the steady-state condition. Without the incoherent pump Λ, there will be no gain $\mathrm{Im}(\rho_{41}/\Omega_{\mathrm{p}}) = 0$, which is commonly known as the electromagnetically induced transparency (EIT) phenomenon.

With the incoherent pump Λ, if $\rho_{44} > \rho_{11}$, and $\rho_{44} > \rho_{++} = \rho_{--}$, where $\rho_{++}(\rho_{--})$ is the population probability in the dressed-state $|+\rangle(|-\rangle)$, $\rho_{++} = \rho_{--} = \frac{1}{2}(\rho_{11}+\rho_{33})$, the gain with population inversion in both the bare states and dressed states can be established when the following condition is satisfied:

$$\Gamma_{32} > \Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{4|\Omega_{\mathrm{c}}|^2}. \quad (10)$$

If $\rho_{++} < \rho_{44} < \rho_{11}$, the gain with population inversion in the dressed state and without population inversion in the bare state can be established. The condition is

$$\Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{8|\Omega_{\mathrm{c}}|^2} < \Gamma_{32} < \Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{4|\Omega_{\mathrm{c}}|^2}. \quad (11)$$

If $\rho_{44} < \rho_{++}$, and $\rho_{41} > 0$, the system exhibits gain without population inversion on any state basis (the dressed states or the bare states) if the following condition is satisfied:

$$\Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{4|\Omega_{\mathrm{c}}|^2} - \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})}{2\gamma_{43}} < \Gamma_{32} < \Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{8|\Omega_{\mathrm{c}}|^2}. \quad (12)$$

When the atomic parameters are chosen to satisfy the condition $\Gamma_{32} < \Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{4|\Omega_{\mathrm{c}}|^2} - \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})}{2\gamma_{43}}$ or $\Lambda = 0$, the probe amplification behaviour will never occur.

In Fig.2 are shown the population distribution and the calculated $\mathrm{Im}(\rho_{41}/\Omega_{\mathrm{p}})$ versus Λ for various

system parameters. Different parameters are chosen in Figs.2(a)–(c) in order that different amplification conditions expressed by Eqs.(9)–(11) could be found. From Eqs.(9)–(11) and Fig.2, it is easy to find that the necessary condition for probe amplification in this model is $\Gamma_{32} > \Gamma_{41} + \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})^2}{4|\Omega_c|^2} - \frac{\Gamma_{41}(\Gamma_{31}+\Gamma_{32})}{2\gamma_{43}}$. Λ and Γ_{42} do not essentially affect the onset of non-inversion or inversion of the populations in this system. (Also note that zero threshold value for probe amplification in this model under a certain parameter condition is very convenient for experimental realization.)

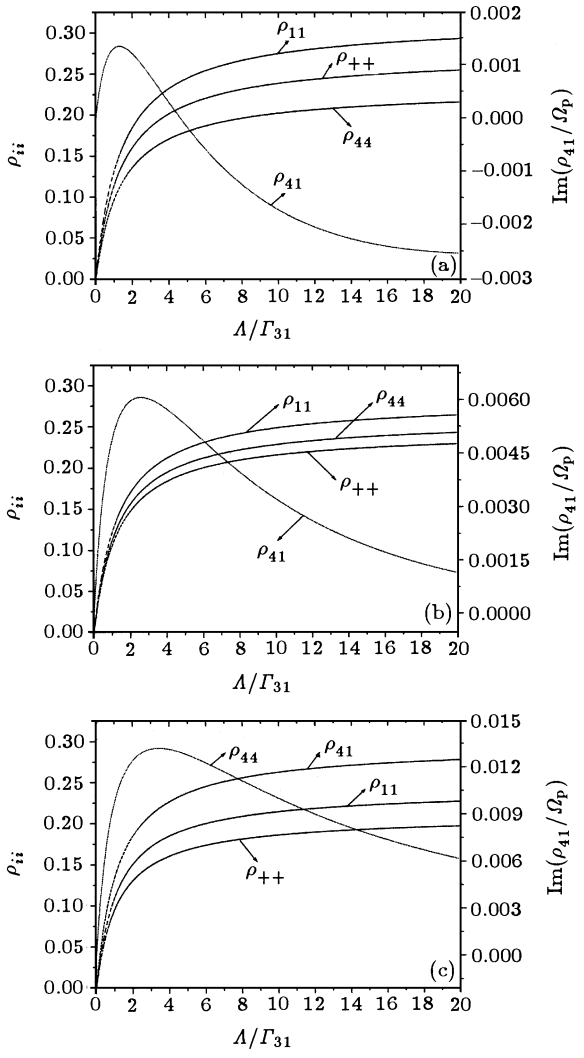


Fig.2. The dimensionless gain coefficient at $\Delta_p = 0$ and population distribution as a function of the incoherent pump rate Λ for (a) $\Gamma_{41} = 5\Gamma_{31}$, (b) $\Gamma_{41} = 4\Gamma_{31}$, (c) $\Gamma_{41} = 3\Gamma_{31}$. Other parameters are set at $\Gamma_{42} = \Gamma_{31}$, $\Gamma_{32} = 5\Gamma_{31}$, $\Omega_c = 5\Gamma_{31}$.

Since the population distribution and the probe gain depend on the coupling Rabi frequency, it is necessary to show graphically the light amplification mechanisms changing with the coupling Rabi frequency. We have calculated numerically the induced polarization ρ_{41} versus Ω_c and plotted the results in Fig.3. The population distribution is also drawn in Fig.3, so that we could discern which amplification mechanism dominates the probe amplification behaviour. The normalized parameters used are $\Gamma_{41} = 4\Gamma_{31}, \Gamma_{42} = \Gamma_{31}$, $\Gamma_{32} = 5\Gamma_{31}$, $\Gamma_{21} = 0$ and $\Lambda = \Gamma_{31}$. It is easy to find that the threshold value of Ω_c for probe amplification is Ω_{c1}. For $\mathrm{Im}(\rho_{41}) > 0$, there are three regions of population distribution: first, for $\Omega_{c1} < \Omega_c < \Omega_{c2}$, $\rho_{44} < \rho_{++} < \rho_{11}$, the gain exists without population inversion on any state basis (bare state and dressed state); second, for $\Omega_{c2} < \Omega_c < \Omega_{c3}$, $\rho_{++} < \rho_{44} < \rho_{11}$, the gain exists without population inversion in the bare state, but with population inversion in the dressed state; third, for $\Omega_c > \Omega_{c3}$, $\rho_{44} > \rho_{11} \geq \rho_{++}$, the gain exists with population inversion in the bare state as well as in the dressed state. It is clear that in this system the dressed-state inversion gain could never be operating alone, which is a conclusion in Ref.[24]. In Fig.3 the inversion gain region emerges along with the larger coupling Rabi frequency. As the coupling Rabi frequency becomes larger, the driving field brings more population from $|1\rangle$ to $|3\rangle$, the metastable state $|2\rangle$ obtains more population by spontaneous decay from $|3\rangle$ to $|2\rangle$. On the other hand, the incoherent pump makes up a large portion of population distribution of state $|4\rangle$. With increasing population in state $|4\rangle$, the system changes from state without a population inversion to that with a population inversion. The three gain regions become wider and wider with increasing coupling field frequency as we have checked. It is seen in Fig.3 that the gain is peaked at a moderate Ω_c in the dressed-state inversion gain region.

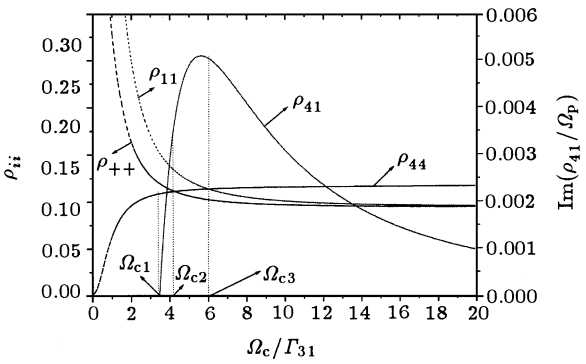


Fig.3. The dimensionless gain coefficient at $\Delta_p = 0$ and population distribution as a function of the coupling Rabi frequency Ω_c. The atomic parameters are given in the text.

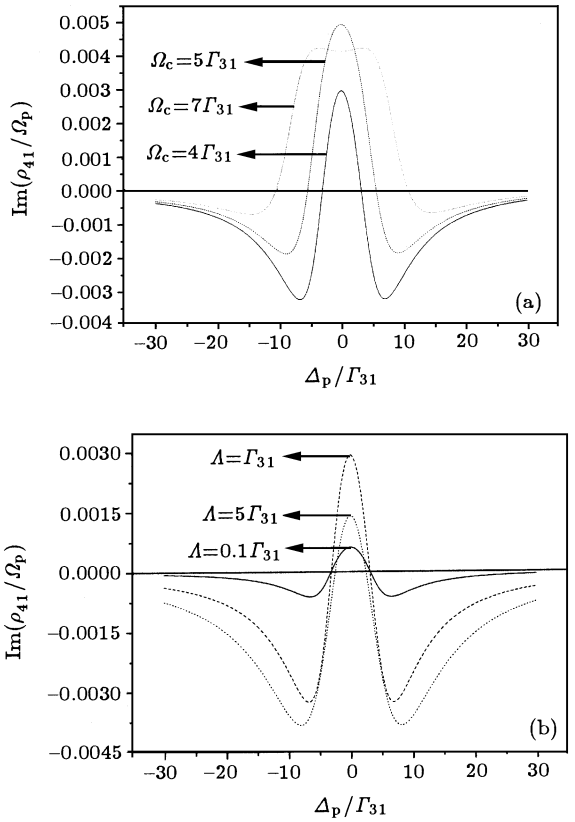


Fig.4. The probe gain as a function of probe detuning from resonance under different conditions. The atomic parameters are the same as those in Fig.3. (a) $\Lambda = \Gamma_{31}$, (b) $\Omega_c = 4\Gamma_{31}$.

In Fig.4(a), three different lines represent the probe spectra in three different amplification regions discussed above. We can see that the maximal probe gain with inversion in any state ($\Omega_c > \Omega_{c3}$) is located at $\Delta_p = \pm\Omega_c$ as shown by the curve labelled with $\Omega_c = 7\Gamma_{31}$, the maximal probe gain without inversion in the bare state ($\Omega_{c1} < \Omega_c < \Omega_{c2}$) and in the dressed state ($\Omega_{c2} < \Omega_c < \Omega_{c3}$) is located at $\Delta_p = 0$ as shown by the curves labelled with $\Omega_c = 4\Gamma_{31}$ and $\Omega_c = 5\Gamma_{31}$ respectively. The amplification range for the detuning of probe beam is mainly dominated by the coupling Rabi frequency, and becomes larger and larger with the coherent field increasing. But the amplification range does not depend on the intensity of the incoherent pumping. This is be demonstrated by Fig.4 (b), which shows the probe curve at different incoherent pump rates but the same coupling Rabi frequency in the region of gain without inversion.

The fundamental mechanism responsible for the AWI in this system can be attributed to quantum interference caused by the quantum indistinguishable transition pathways. In the dressed-state picture, the strong coupling laser produces a pair of dressed states $|+\rangle$ and $|-\rangle$. The atoms will not absorb the probe laser at $\Delta_p = 0$ due to the fact that probability amplitudes interfere destructively for transition of $|4\rangle - |+\rangle$ and $|4\rangle - |-\rangle$. This interference clearly results from the fact that we cannot tell where ($|+\rangle$ or $|-\rangle$) the electron has gone. However, for the stimulated emission, only the transition pathway $|4\rangle - \frac{1}{\sqrt{2}}(|+\rangle + |-\rangle)$ i.e. $|4\rangle - |1\rangle$ can occur. Thus the lack of destructive interference, in comparison with the absorption, gives rise to probe amplification.

4. Conclusions

In this paper, we have presented a four-level model that may exhibit either LWI or lasing with inversion. The amplification origin is analysed and three amplification mechanisms are distinguished. We have emphasized that quantum interference plays an important role in amplification without inversion. The advantage of this model is to display the transition from non-inversion to inversion depending on intensity of the coherent field but not on the incoherent pumping. We can achieve the amplification with or without inversion even by a very weak incoherent pumping. This model thus provides a new experimental scheme to realize LWI and a short-wavelength laser.

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