Influences of ac Stark shifts on coherent population trapping in the atom-field coupling system via Raman two-photon processes

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By means of adiabatically eliminating the upper level or levels of a Λ-configuration atom coupling to a two-mode quantized field, we have studied the influences of ac Stark shifts on the atomic coherent population trapping in two important atom-field coupling systems via nondegenerate Raman two-photon processes. The states of the field that trap the atom in its two nondegenerate lower levels are obtained, and the important roles of ac Stark shifts and properties of the atomic upper level or levels in the states of the trapping field in two Raman-coupling systems are also analyzed.

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I. INTRODUCTION

The coherent population trapping phenomena in the system of a Λ-configuration atom interacting with two coherent fields have been extensively studied because of applications in a number of different contexts such as laser cooling [1], lasing without inversion [2], and electromagnetically induced transparency [3]. Recently, Agarwal [4] investigated coherent population-trapping states in the system of a Λ-configuration three-level atom interacting with a two-mode quantized field, and showed that in the resonant case the quantized field, which keeps the atom initially in its semiclassical coherent population-trapping state from evolving in time, must be a two-mode photon statistically matched field. Eliminating adiabatically the upper level in a Λ-type three-level atom-field coupling system and neglecting the effects of ac Stark shifts, Deb, Gangopadhyay, and Ray [5] pointed out that, for an atom prepared in a coherent superposition of two nondegenerate lower states, population trapping occurs when the two-mode field is either anticorrelated or correlated, depending on the initial atomic state. In the limit of large total photon number, this field can be identified as an eigenstate of a two-mode phase difference operator.

On the other hand, under certain conditions, a multilevel atom interacting with a quantized field can be identified as a two-level system with ac Stark shifts by means of an adiabatically eliminating method [6,7]. Zhu and Scully [8] and Boone and Swain [9] have studied the properties of nondegenerate and degenerate two-photon lasers. They found that the photon distribution, the linewidth, and the frequency shift depend strongly on the detailed atomic structure because of the effects of ac Stark shifts. Brune *et al.* [10] showed that the ac Stark shifts may be proposed to realize the quantum-nondemolition scheme to measure the number of photons stored in a high-Q cavity. Also, the quantum properties of an atom-field coupling system, such as the collapses and revivals of atomic inversion [11–15], the atomic transition line [14], atomic emission spectra and atomic dipole squeezing [16,17], the squeezing of field [15], and the phase properties of the field [18], can be changed drastically due to the influences of ac Stark shifts. In particular, Cirac and Sanchez-Soto [19] pointed out that, in the degenerate two-photon Jaynes-Cummings model in the presence of ac Stark shifts, the field which traps the atom in the linear superposition of its two levels must be a single-mode squeezed vacuum field.

In order to reveal the trapping-field dependence on ac Stark shifts and detailed atomic structure, here we examine atomic coherent population trapping in two popular but important atom-field coupling systems via nondegenerate Raman two-photon processes. In Sec. II, our attention is focused on the influence of ac Stark shifts on atomic coherent population trapping in a system of a Λ-type three-level atom interacting with a two-mode quantized field. We find that the ac Stark shifts play an important role in the trapping field. For the case including the ac Stark shifts, the trapping field is a two-mode SU(2) coherent field [20], whose partition parameter is dependent on the initial atomic state and the atom-coupling constants, and the maximum possible photon number N is arbitrary. When the ac Stark shifts are neglected, the probability amplitude of the trapping field is not related to the atom-field coupling constants. Section III is devoted to investigating population trapping in a system of a two-mode quantized field coupled to a Λ-type atom with two nondegenerate bound levels and a flat continuum of levels. If the intensities of ac Stark shifts are chosen appropriately, then under the interaction of an appropriate two-mode SU(2) coherent field with a similar formula as obtained in Sec. II, except that N is fixed by the atom-field coupling parameters, the atom initially in a linear superposition of its two bound levels can be trapped completely and not ionized. But if the effects of ac Stark shifts are ignored, the trapping field cannot be found. The above results indicate that the properties of the de-

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tailed atomic structure have strong effects on the trapping field. Finally, our conclusions are summarized in Sec. IV.

II. STATE OF TRAPPING FIELD IN A TWO-LEVEL SYSTEM VIA RAMAN TWO-PHOTON PROCESSES IN THE PRESENCE OF STARK SHIFTS

First we consider a three-level atom of Λ type (Fig. 1). The three atomic levels are denoted as $|1\rangle$, $|2\rangle$, and $|3\rangle$ with corresponding energies ω_1, ω_2, and ω_3, respectively. The dipole-allowed transitions $|1\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |3\rangle$ are mediated by photons of the two different modes of the two-mode quantized field, which are characterized by the photon operators a and b with corresponding frequencies ω_a and ω_b. Following Refs. [6,7,11,12,15], one can adiabatically remove the upper level $|3\rangle$ when the one-photon detuning $|\Delta|$ is much larger than the Rabi frequency of the oscillation between $|3\rangle$ and $|1\rangle$ ($|2\rangle$), where $\Delta = \omega_3 - \omega_1 - \omega_a = \omega_3 - \omega_2 - \omega_b$, so that the effective Hamiltonian of the system including the ac Stark shifts is written in the rotating-wave approximation as [12–18]

$$H = \omega_a a^\dagger a + \omega_b b^\dagger b + (\omega_2 - \omega_1) S_z + \frac{g_1^2}{\Delta} a^\dagger a |1\rangle\langle 1| + \frac{g_2^2}{\Delta} b^\dagger b |2\rangle\langle 2| + \frac{g_1 g_2}{\Delta}(a b^\dagger S_+ + a^\dagger b S_-) \quad (\hbar = 1) \ . \tag{1}$$

Here g_1 (g_2) represents the coupling strength for the dipole-allowed transition $|3\rangle \leftrightarrow |1\rangle$ ($|3\rangle \leftrightarrow |2\rangle$), and S_+ and S_- are the atomic raising and lowering operators between levels $|2\rangle$ and $|1\rangle$.

If the atom is initially in the coherent superposition of its two nondegenerate lower levels $|1\rangle$ and $|2\rangle$ as

$$|\Psi_A(0)\rangle = \cos(\theta/2)|2\rangle + \sin(\theta/2) e^{i\varphi}|1\rangle \qquad (0 \le \theta < \pi, \ 0 < \varphi \le 2\pi) \tag{2}$$

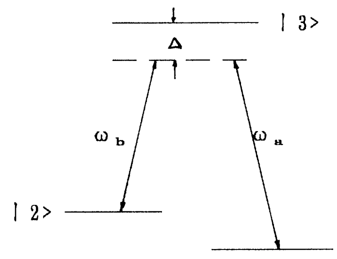

FIG. 1. Diagram of a Λ-configuration three-level atom coupling to a two-mode quantized field.

and the field is

$$|\Psi_F(0)\rangle = \sum_{n_1, n_2} F_{n_1 n_2} |n_1, n_2\rangle \ , \tag{3}$$

then the initial state of the system can be written as

$$\begin{aligned}|\Psi(0)\rangle &= |\Psi_A(0)\rangle \otimes |\Psi_F(0)\rangle \\ &= \sum_{n_1, n_2} [\cos(\theta/2) F_{n_1 - 1, n_2 + 1} |2, n_1 - 1, n_2 + 1\rangle \\ &\qquad + \sin(\theta/2) e^{i\varphi} F_{n_1, n_2} |1, n_1, n_2\rangle] \ .\end{aligned} \tag{4}$$

In the interaction picture, the state vector of the system at time t evolves into

$$\begin{aligned}|\Psi_I(t)\rangle = &\sum_{n_2=0}^{\infty} \sin(\theta/2) e^{i\varphi} F_{0, n_2} |1, 0, n_2\rangle \\ &+ \sum_{n_1=0}^{\infty} \cos(\theta/2) F_{n_1, 0} |2, n_1, 0\rangle \\ &+ \sum_{\substack{n_1=1 \\ n_2=0}} [A_{n_1, n_2} |1, n_1, n_2\rangle \\ &\qquad + B_{n_1 - 1, n_2 + 1} |2, n_1 - 1, n_2 + 1\rangle]\end{aligned} \tag{5}$$

with

$$A_{n_1, n_2}(t) = \frac{2 g_1 g_2 \sqrt{n_1(n_2+1)}}{\Delta} \left\{ C_1 \frac{\exp[i(\delta_{n_1, n_2} - \beta_{n_1, n_2}) t/2]}{\beta_{n_1, n_2} - \delta_{n_1, n_2}} - C_2 \frac{\exp[i(\delta_{n_1, n_2} + \beta_{n_1, n_2}) t/2]}{\beta_{n_1 n_2} + \delta_{n_1, n_2}} \right\} \exp\left[-i \frac{g_1^2}{\Delta} n_1 t \right] , \tag{6}$$

$$B_{n_1 - 1, n_2 + 1} = \{ C_1 \exp[-i(\delta_{n_1, n_2} + \beta_{n_1, n_2}) t/2] + C_2 \exp[i(\beta_{n_1, n_2} - \delta_{n_1, n_2}) t/2] \} \exp\left[-i \frac{g_2^2}{\Delta}(n_2 + 1) t \right] , \tag{7}$$

$$C_1 = \frac{1}{2\beta_{n_1, n_2}} \left[\cos(\theta/2) F_{n_1 - 1, n_2 + 1} (\beta_{n_1, n_2} - \delta_{n_1, n_2}) + \frac{2 g_1 g_2}{\Delta} \sqrt{n_1(n_2+1)} \sin(\theta/2) e^{i\varphi} F_{n_1, n_2} \right] , \tag{8}$$

$$C_2 = \frac{1}{2\beta_{n_1, n_2}} \left[\cos(\theta/2) F_{n_1 - 1, n_2 + 1} (\beta_{n_1, n_2} + \delta_{n_1, n_2}) - \frac{2 g_1 g_2}{\Delta} \sqrt{n_1(n_2+1)} \sin(\theta/2) e^{i\varphi} F_{n_1, n_2} \right] , \tag{9}$$

$$\delta_{n_1, n_2} = \frac{g_1^2}{\Delta} n_1 - \frac{g_2^2}{\Delta}(n_2 + 1), \quad \beta_{n_1, n_2} = \sqrt{\delta_{n_1, n_2}^2 + 4 g_1^2 g_2^2 n_1 (n_2 + 1)/\Delta^2} \ . \tag{10}$$

Explicitly, β_{n_1,n_2} is the generalized Rabi oscillating frequency of atom and δ_{n_1,n_2} represents the detuning due to the ac Stark shifts.

Let us now study population trapping in the two-level system described by Eq. (1), i.e., a persistent probability of finding the atom in its level $|1\rangle$ or $|2\rangle$ in spite of the existence of the two-mode quantized field. Evidently, if the state vectors $\{|1,n_1,n_2\rangle\}$ and $\{|2,n_1-1,n_2+1\rangle\}$ are stationary vectors, atomic coherent population trapping occurs. This demands that C_1 or C_2 in the probability amplitudes $A_{n_1,n_2}(t)$ and $B_{n_1-1,n_2+1}(t)$ must be equal to zero. Hence, according to Eq. (8) or (9) we can obtain the trapping condition as, for $C_1=0$,

$$F_{n_1-1,n_2+1}\cos(\theta/2)(\beta_{n_1,n_2}-\delta_{n_1,n_2}) = -\frac{2g_1g_2}{\Delta}\sqrt{n_1(n_2+1)}\sin(\theta/2)e^{i\varphi}F_{n_1,n_2}\ , \tag{11}$$

or for $C_2=0$,

$$F_{n_1-1,n_2+1}\cos(\theta/2)(\beta_{n_1,n_2}+\delta_{n_1,n_2}) = \frac{2g_1g_2}{\Delta}\sqrt{n_1(n_2+1)}\sin(\theta/2)e^{i\varphi}F_{n_1,n_2}\ . \tag{12}$$

From the above Eqs. (11) and (12), we can find that the recurrence formula of the probability amplitude F_{n_1,n_2} of the initial two-mode field is only related to F_{n_1-1,n_2+1}, i.e., the total photon number $N=n_1+n_2$ of each two-mode number state $|n_1,n_2\rangle$ is a constant. This means that the initial state of the driven field [described by Eq. (3)] must be expressed as

$$|\Psi_F(0)\rangle=\sum_{n=0}^{N}F_{n,N-n}|n,N-n\rangle\ . \tag{13}$$

Here, we assume $n_1=n$ and $n_2=N-n$. The reason is that the two-mode photon sum operator $N=a^\dagger a+b^\dagger b$ is a motion constant of the system, i.e., $[N,H]=0$. In order to exhibit the effects of the ac Stark shifts on the atomic coherent population trapping, we first consider the case of neglecting the ac Stark shifts.

A. Neglecting the influence of ac Stark shifts

When the influence of ac Stark shifts is neglected, i.e., $\delta_{n_1,n_2}=0$, Eq. (11) is simplified as

$$F_{n,N-n}=-\cot(\theta/2)e^{-i\varphi}F_{n-1,N-n+1}\ . \tag{14}$$

Solving the above equation, we obtain the normalized state vector of the two-mode field as

$$|\Psi_{F1}(0)\rangle=\left[\frac{1-\cot^2(\theta/2)}{1-[\cot^2(\theta/2)]^{N+1}}\right]^{1/2}\times\sum_{n=0}^{N}[-\cot(\theta/2)e^{-i\varphi}]^n|n,N-n\rangle\ . \tag{15}$$

The properties of the initial field described by Eq. (15) are given in Ref. [5]. From Eq. (15), we known that the total photon number N is an arbitrary constant which is not related to the atom-coupling system, and the probability amplitude $F_{n,N-n}$ is only decided by the parameters θ and φ describing the initial atomic state and is not connected with the parameters describing the atom-field coupling constants, such as g_1 and g_2. That is to say, when the ac Stark shifts are neglected in Raman-coupling two-level systems with different coupling strengths, if the atoms are prepared in the same state, then the fields trapping the atoms are in the same state.

B. Considering the effects of ac Stark shifts

If we consider the effects of ac Stark shifts in Eq. (11), the recurrence formula is written as

$$F_{n,N-n}=\cot(\theta/2)e^{-i\varphi}\frac{g_2}{g_1}\left[\frac{N-n+1}{n}\right]^{1/2}F_{n-1,N-n+1}\ . \tag{16}$$

Using the normalization condition, we obtain the state function of the driven field as

$$|\Psi_{F2}(0)\rangle=F_{0,N}\sum_{n=0}^{N}\left[-\frac{g_2}{g_1}\cot(\theta/2)e^{-i\varphi}\right]^n\times\left[\frac{N!}{n!(N-n)!}\right]^{1/2}|n,N-n\rangle\ . \tag{17}$$

Here

$$F_{0,N}=\left[1+\frac{g_2^2}{g_1^2}\cot^2(\theta/2)\right]^{-N/2}\ .$$

It is clear that $|\Psi_{F2}(0)\rangle$ is the two-mode coherent SU(2) state [20] whose maximum possible number of photons is N and the parameter describing the partition of photons is $-(g_2/g_1)\cot(\theta/2)e^{-i\varphi}$. It is easy to verify that the two modes are anticorrelated. From Eq. (17), we find that the probability amplitude $F_{n,N-n}$ is dependent not only on the initial atomic state but also on the atom-field coupling constants g_1 and g_2. This means that for Raman-coupling two-level systems with different coupling strengths, the fields which trap the atoms are in different two-mode SU(2) coherent states even if the atoms are prepared in the same state.

In this section, by means of adiabatically eliminating the upper level $|3\rangle$ of a Λ-configuration three-level atom coupling to a two-mode quantized field, we obtain the states of the trapping field in a Raman-coupling two-level system in the absence and presence of the ac Stark shifts. As we see, the influences of the ac Stark shifts play an important role in the states of the trapping field. However, if the upper levels of the Λ-type atom are continuum states, for example, in the Raman-type photoionization system of a Λ-configuration atom with two lower bound levels and a flat continuum of levels driven by a two-mode quantized field, then what is the trapping field? We answer this question in the following section.

III. EFFECTS OF THE AC STARK SHIFTS ON THE STATE OF TRAPPING FIELD IN A RAMAN-TYPE PHOTOIONIZATION SYSTEM DRIVEN BY A TWO-MODE QUANTIZED FIELD

As we know, a stronger laser can embed a low-lying atomic state into a flat atomic continuum to produce a tunable resonance of adjustable width. This process is termed the laser-induced continuum structure (LICS), [21–24] and has drawn much attention both theoretically [21,23] and experimentally [24], due to its possible applications in laser isotope separation [21] and new photodetecting devices for nonclassical fields [22]. Describing the two driven lasers classically in a LICS system involving Raman-transition process and neglecting the effects of ac Stark shifts, Knight and his co-workers [21] found that, if the intensities of two driven laser fields and the nondegenerate two-photon detuning satisfy an appropriate condition, the atom initially in the semiclassical population-trapping state cannot be ionized, and in this case the atomic coherent population trapping of two bound states occurs. Because this result is obtained based on semiclassical theory, it cannot completely reveal the coherent population-trapping properties in a Raman-type photoionization system. Since the atomic coherent population phenomenon in this atom-field coupling system plays an important role in the electromagnetically induced transparency proposed by Harris and others [3], it is necessary to investigate this phenomenon in detail. Thus we study this by means of the full quantum theory in the following.

The model to be considered here, shown in Fig. 2, consists of a two-mode quantized field coupling to a Λ-configuration atom with two bound states $|1\rangle$ and $|2\rangle$ and a flat atomic continuum of states $\{|\varepsilon\rangle\}$. The one-photon transition processes $|1\rangle \leftrightarrow |\varepsilon\rangle$ and $|2\rangle \leftrightarrow |\varepsilon\rangle$ are driven by the mode a and the mode b, respectively. In the rotating-wave approximation, the total Hamiltonian of this system is given by

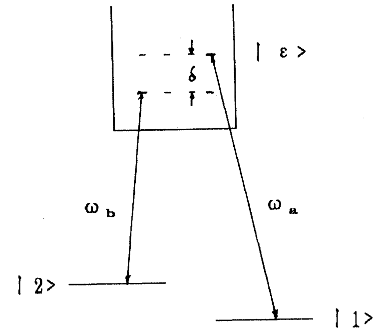

FIG. 2. Diagram of the Λ-type photoionization system driven by a two-mode quantized field.

$$H = H_0 + V \quad (\hbar = 1) , \tag{18}$$

$$H_0 = \omega_1 |1\rangle\langle 1| + \omega_2 |2\rangle\langle 2| + \int d\varepsilon\, \varepsilon |\varepsilon\rangle\langle\varepsilon| + \omega_a a^\dagger a + \omega_b b^\dagger b , \tag{19}$$

$$V = \left[\int d\varepsilon\, g_{1\varepsilon} a^\dagger |1\rangle\langle\varepsilon| + \int d\varepsilon\, g_{2\varepsilon} b^\dagger b |2\rangle\langle\varepsilon| + \text{H.c.} \right] , \tag{20}$$

where $g_{1\varepsilon}$ ($g_{2\varepsilon}$) is the atom-field coupling constant which is related to the one-photon processes $|\varepsilon\rangle \leftrightarrow |1\rangle$ ($|\varepsilon\rangle \leftrightarrow |2\rangle$).

If the system is initially in the state described by Eq. (4), then in the interaction picture the state vector of the system at time t develops into

$$|\Psi_I(t)\rangle = \sum_{\substack{n_1=1 \\ n_2=0}} \left[A_1(t)|1,n_1,n_2\rangle + A_2(t)|2,n_1-1,n_2+1\rangle + \int d\varepsilon\, A_\varepsilon(t)|\varepsilon,n_1-1,n_2\rangle \right] + \sin(\theta/2) e^{i\varphi} \sum_{n_2=0} F_{0,n_2}|1,0,n_2\rangle + \cos(\theta/2) \sum_{n_1=0}^{\infty} F_{n_1,0}|2,n_1,0\rangle . \tag{21}$$

Bringing the Schrödinger equation into the interaction picture, we obtain

$$i\frac{d}{dt} A_1(t) = -\sqrt{n_1} \int d\varepsilon\, g_{1\varepsilon} A_\varepsilon(t) e^{-i\Delta_{1\varepsilon} t} , \tag{22}$$

$$i\frac{d}{dt} A_2(t) = -\sqrt{n_2+1} \int d\varepsilon\, g_{2\varepsilon} A_\varepsilon(t) e^{-i\Delta_{2\varepsilon} t} , \tag{23}$$

$$i\frac{d}{dt} A_\varepsilon(t) = -g_{\varepsilon 1} \sqrt{n_1} A_\varepsilon(t) e^{i\Delta_{1\varepsilon} t} - g_{\varepsilon 2} \sqrt{n_2+1} A_\varepsilon(t) e^{i\Delta_{2\varepsilon} t} . \tag{24}$$

Here $\Delta_{1\varepsilon} = \varepsilon - \omega_a - \omega_1$ and $\Delta_{2\varepsilon} = \varepsilon - \omega_b - \omega_2$ are the one-photon detunings. Following a method similar to that given in Ref. [21], Eqs. (22) and (23) are modified as

$$\frac{d}{dt} A_1 = -\frac{\gamma_1 + i\delta_1}{2} n_1 A_1 + i\frac{\gamma}{2} \sqrt{n_1(n_2+1)} A_2 (q+i) e^{-i\delta t} , \tag{25}$$

$$\frac{d}{dt} A_2 = -\frac{\gamma_1 + i\delta_2}{2} (n_2+1) A_2 + i\frac{\gamma}{2} \sqrt{n_1(n_2+1)} A_1 (q+i) e^{i\delta t} , \tag{26}$$

where $\delta = \Delta_{1\varepsilon} - \Delta_{2\varepsilon}$ is the nondegenerate two-photon detuning, and γ_i and δ_i ($i=1,2$) are the parameters related to the decay rates and the ac Stark shifts, due to the transitions shown in Fig. 2, for levels $|1\rangle$ and $|2\rangle$, i.e.,

$$\gamma_1=2\pi\int d\varepsilon|g_{1\varepsilon}|^2\delta(\varepsilon-\omega_a-\omega_1)=2\pi|g_{1a}|^2,\quad \gamma_2=2\pi|g_{2b}|^2 , \tag{27}$$

$$\delta_i=-P\int d\varepsilon\frac{2|g_{i\varepsilon}|^2}{\Delta_{i\varepsilon}} . \tag{28}$$

$\gamma=\sqrt{\gamma_1\gamma_2}$ represents the nondegenerate two-photon coupling constant and q is the constant related to the Fano parameter [21–24],

$$q=\left[\frac{\gamma}{2}\right]^{-1}\mathrm{P}\int_0^\infty d\varepsilon\frac{g_{1\varepsilon}g_{2\varepsilon}}{\varepsilon-(\omega_a+\omega_b+\omega_1+\omega_2)/2} . \tag{29}$$

It is clear that the Raman-coupling system described by Eqs. (25) and (26) can be described by the effective interaction Hamiltonian in the interaction picture,

$$\begin{aligned}V^I_{\text{eff}}=&-\frac{i}{2}\gamma_1a^\dagger a|1\rangle\langle 1|-\frac{i}{2}\gamma_2b^\dagger b|2\rangle\langle 2|\\&+\frac{\delta_1}{2}a^\dagger a|1\rangle\langle 1|+\frac{\delta_2}{2}b^\dagger b|2\rangle\langle 2|\\&+i\frac{\gamma}{2}(1-iq)(a^\dagger b|1\rangle\langle 2|e^{-i\delta t}+ab^\dagger|2\rangle\langle 1|e^{i\delta t}).\end{aligned} \tag{30}$$

Here we see, if the upper levels of a Λ-type atom which are a flat continuum are adiabatically eliminated, the effective Hamiltonian for the Raman-coupling two-level system is non-Hermitian and different from the case in which the upper level is a bound state.

Considering the initial condition Eq. (4), Eqs. (25) and (26) are solved analytically to yield

$$A_1=e^{\alpha t}(B_1e^{\beta t}+B_2e^{-\beta t}) , \tag{31}$$

$$A_2=\frac{1}{F}e^{\alpha t}e^{i\delta t}[B_1(\beta-\alpha')e^{\beta t}-B_2(\beta+\alpha')e^{-\beta t}] , \tag{32}$$

where

$$\alpha'=\frac{i}{2}\left[\delta+\frac{\delta_2(n_2+1)-\delta_1n_1}{2}\right]+\frac{\gamma_2(n_2+1)-\gamma_1n_1}{4} , \tag{33a}$$

$$\alpha=-\frac{i}{2}\left[\delta+\frac{\delta_2(n_2+1)+\delta_1n_1}{2}\right]-\frac{\gamma_2(n_2+1)+\gamma_1n_1}{4} , \tag{33b}$$

$$F=-\gamma\frac{1-iq}{2}\sqrt{n_1(n_2+1)},\quad \beta^2=\alpha'^2+F^2 , \tag{33c}$$

$$\begin{aligned}B_1=\frac{1}{2\beta}[&F\cos(\theta/2)F_{n_1,-1,n_2+1}\\&+(\beta+\alpha')\sin(\theta/2)e^{i\varphi}F_{n_1,n_2}] ,\end{aligned} \tag{34a}$$

$$\begin{aligned}B_2=\frac{-1}{2\beta}[&F\cos(\theta/2)F_{n_1-1,n_2+1}\\&-(\beta-\alpha')\sin(\theta/2)e^{i\varphi}F_{n_1,n_2}] .\end{aligned} \tag{34b}$$

In the general case, from Eqs. (31)–(33) we see that with the time development the atom which is initially in the state $|\Psi_A(0)\rangle$ can be ionized because the exponential functions $e^{(\alpha\pm\beta)t}$ in $A_1(t)$ and $A_2(t)$ contain decay factors. In order to trap the atom in its two bound levels,the decay factors in $e^{(\alpha\pm\beta)t}$ must disappear. This means that the complex functions $\alpha\pm\beta$ must be reduced to pure imaginary functions. Fortunately, when the detuning δ and the photon number n_1,n_2 satisfy

$$\frac{2\delta}{q\gamma_1+\delta_1}=n_1-\frac{q\gamma_2+\delta_2}{q\gamma_1+\delta_1}(n_2+1) , \tag{35}$$

$\alpha\pm\beta$ are simplified as

$$\alpha+\beta=-\frac{i}{2}(q\gamma_1+\delta_1)n_1 , \tag{36a}$$

$$\begin{aligned}\alpha-\beta=&-\tfrac{1}{2}[\gamma_1n_1+\gamma_2(n_2+1)]\\&-\frac{i}{2}\left[(q\gamma_1+\delta_1)(n_1+n_2+1)\right.\\&\left.-\frac{q\gamma_1-\delta_1}{2}n_1-\frac{q\gamma_2-\delta_2}{2}(n_2+1)\right] .\end{aligned} \tag{36b}$$

It is evident that the function $e^{(\alpha+\beta)t}$ in $A_1(t)$ and $A_2(t)$ does not decay with time development. So in this case the atom initially in $|\Psi_A(0)\rangle$ cannot be completely ionized and can be partially trapped in its two bound levels. This means that when the photons between the mode a and the mode b of the driven field have the correlation shown in Eq. (35) the atomic coherent population trapping may occur in the Raman-type photoionization system. Furthermore, we discuss the effects of ac Stark shifts on the states of the trapping field in a Raman-type photoionization system when $q\gamma_1+\delta_1=\pm(q\gamma_2+\delta_2)$.

(1) When $q\gamma_1+\delta_1=-(q\gamma_2+\delta_2)$, Eq. (35) is reduced to

$$n_1+n_2+1=\frac{2\delta}{q\gamma_1+\delta_1}=N+1 . \tag{37}$$

Because the photon numbers n_1,n_2 must be $0,1,2,\ldots$, N must be a positive integer. So the total number of photons for both modes is a constant N, which is fixed by δ, γ_1, δ_1, and q. In this case, the initial state of the driven field can be written as Eq. (13).

However, in order that the atom in the initial state $|\Psi_A(0)\rangle$ can be completely trapped in its two bound states, $A_1(t)$ and $A_2(t)$ should not contain the function $e^{(\alpha-\beta)t}$. This requires that $B_2=0$, i.e.,

$$F\cos(\theta/2)F_{n_1-1,n_2+1}-(\beta-\alpha')\sin(\theta/2)e^{i\varphi}F_{n_1,n_2}=0 . \tag{38}$$

Substituting Eqs. (13) and (37) into the above equation, the probability amplitude $F_{n,N-n}$ must satisfy

$$F_{n,N-n}=-\cot(\theta/2)e^{-i\varphi}\times\frac{g_{2b}}{g_{1a}}\left[\frac{N-n+1}{n}\right]^{1/2}F_{n-1,N-n+1} . \tag{39}$$

Here we have assumed $(\gamma_2/\gamma_1)^{1/2}=g_{2b}/g_{1a}$. Comparing Eq. (16) with Eq. (39), the only difference is that in Eq. (39) the total photon number N is fixed by Eq. (37), but in Eq. (16) N is arbitrary. The cause leading to this difference is that the Λ-type atom in the Raman-type photoionization system has a flat continuum of levels, so that the atom driven by the field may be ionized. Although under the interaction with a two-mode radiation field, the flat continuum of atomic levels can be adiabatically eliminated and the atom-field coupling system can be identified as the two-level system described by Eq. (30), the atom in its two bound levels can be damped (ionized). In order to avoid this damping,the total photon number N of the driven field must be fixed by Eq. (37). So, if the parameter related to the atom-field coupling character satisfy $q\gamma_1+\delta_1=-(q\gamma_2+\delta_2)$, the field which traps the atom in its two bound levels is the two-mode SU(2) coherent field, in which the maximum possible number of photons is N, fixed by δ, γ_1, δ_1, and q, and the parameter describing the partition of photons is $-(g_{2b}/g_{1a})\cot(\theta/2)e^{-i\varphi}$. That is to say, the phenomenon of the atom being trapped completely in its two bound levels occurs only in some special Raman-type photoionization systems, and only in these systems may the electromagnetically induced transparency phenomenon [3] happen.

(2) If the amplitudes of ac Stark shifts in the system satisfy $q\gamma_1+\delta_1=q\gamma_2+\delta_2$, Eq. (35) reduces to

$$n_1=n_2+N+2 \ . \tag{40}$$

Inserting Eq. (40) into Eq. (38), the probability amplitude of the driven field obeys

$$F_{n+N+2,n}=-\cot(\theta/2)e^{-i\varphi}\frac{g_{2b}}{g_{1a}} \times\left[\frac{n+1}{n+N+2}\right]^{1/2} F_{n+N+1,n+1} \ . \tag{41}$$

Therefore the initial state [Eq. (4)] of the atom-field coupling system is modified as

$$|\Psi(0)\rangle=\sum_{n=0}^{\infty}\cos(\theta/2)F_{n+N+1,n+1}\left[|2,n+N+1,n\rangle-\frac{g_{2b}}{g_{1a}}\left[\frac{n+1}{n+N+2}\right]^{1/2}|1,n+N+2,n\rangle\right] \ . \tag{42}$$

It is evident that Eq. (42) cannot be expressed as $|\Psi_A(0)\rangle\otimes|\Psi_F(0)\rangle$. In this case, the trapping field which traps the atom completely in its two bound levels cannot be found. This means that when the intensities of ac Stark shifts satisfy $q\gamma_1+\delta_1=q\gamma_2+\delta_2$, the atom driven by an arbitrary field can certainly be ionized and be partially trapped. Meanwhile, there appear a series of peaks that are very sharp in the steady photoelectron spectrum. We refer to this phenomenon as "comb coherent confluence" [25] as Leonski and Buzek [26] show in Fano autoionization system driven by a single-mode quantized field.

Employing the semiclassical theory and neglecting the ac Stark shifts, Knight and co-workers [21] pointed out that in a Raman-type photoionization system, if the atom is initially in its semiclassical trapping state and the intensities of two classical fields are chosen appropriately, then the atom cannot be ionized and can be completely trapped. But, if the ac Stark shifts are ignored, i.e., $\delta_2=\delta_1=0$, and the two-photon detuning $\delta=0$, then according to Eq. (42) the initial state of the atom-field coupling system is reduced to

$$|\Psi(0)\rangle=\frac{1}{\sqrt{2}}e^{i\varphi}\sum_{n=0}^{\infty}F_{n+1,n}[|1\rangle|n+1,n\rangle -|2\rangle|n,n+1\rangle] \ . \tag{43}$$

Here we have assumed $\gamma_1=\gamma_2$ and have chosen $\theta=\pi/4$ according to the normalization condition. Explicitly, although the atomic initial state

$$|\Psi_A(0)\rangle=\frac{1}{\sqrt{2}}e^{i\varphi}(|2\rangle-|1\rangle)$$

is the semiclassical trapping state [21], the trapping field which can completely trap the atom in its two bound levels cannot be found. This result is clearly different from Ref. [21]. The reason inducing this difference is that in Ref. [21] the condition leading to atomic coherent population trapping only restricts the relation of the intensities of the classical driven fields, but here the photons between two modes in the quantized driven field have a strong correlation dependence on the parameters δ, δ_1, δ_2, γ_1, γ_2, and q.

IV. CONCLUSIONS

In conclusion we have studied the influences of ac Stark shifts on atomic coherent population trapping in two important Raman-coupling systems which may have wide applications. The results show that in the system of a Λ-type three-level atom coupling to a two-mode quantized field, if the ac Stark shifts due to the transitions $|3\rangle\leftrightarrow|1\rangle$ and $|3\rangle\leftrightarrow|2\rangle$ are neglected, the field which traps the atom initially in $|\Psi_A(0)\rangle$ must be in the state $|\Psi_{F1}(0)\rangle$ [Eq. (15)], whose maximum photon number N is arbitrary, and the probability amplitude is only dependent on the atomic initial state. But for the case including the ac Stark shifts, the trapping field must be a two-mode SU(2) coherent field $|\Psi_{F2}(0)\rangle$ [Eq. (17)], whose maximum photon number N is also arbitrary, but the parameter describing the partition of photons is related not only to the atomic initial state but also to the atom-field coupling constants g_1 and g_2. In the Raman-type photoionization system consisting of a two-mode quantized field interacting with a Λ-type atom with two bound levels and a flat continuum of levels, in order to keep the atom initially in $|\Psi_A(0)\rangle$ unionized, the values of ac

Stark shifts due to the transitions $\{|\varepsilon\rangle\}\leftrightarrow|1\rangle$ and $\{|\varepsilon\rangle\}\leftrightarrow|2\rangle$ must satisfy $q\gamma_1+\delta_1=-(q\gamma_2+\delta_2)$ and the field must be in the state $|\Psi_{F2}(0)\rangle$ except that N is fixed by Eq. (37). When the ac Stark shifts are neglected, the trapping field cannot be found even if the atom is initially in its semiclassical trapping state. These results indicate that the ac Stark shifts play an important role in atomic coherent population trapping, and the properties of the detailed atomic structure have a strong effect on the trapping field.

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