Phase control of spatial interference from two duplicated two-level atoms

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(Received 23 January 2010; published 14 February 2011)

We report the phase control of spatial interference of resonance fluorescence from two duplicated two-level atoms driven by two orthogonally polarized fields. We find that in the strongly driven situation, adjusting the relative phase leads to a redistribution of the atoms and a significant change of the atomic coherences so that the pattern could survive. In order to improve the experimental realizability, we therefore propose a scheme to recover the visibility with fixed relative phase by adjusting the relative intensity between the two driving fields or, alternatively, by using a standing-wave field.

DOI: 10.1103/PhysRevA.83.023410 PACS number(s): 32.50.+d, 42.50.St, 42.50.Ct

I. INTRODUCTION

Young's double-slit experiment is important to our understanding of the wave nature of light, which exhibits the first-order coherence properties of light [1]. Recently, there has been considerable interest in the interference of the fluorescence light from two driven atoms which play the role of the slits in Young's experiment [2–6]. Remarkably, Eichmann *et al.* carried out a very nice experiment where the two slits were replaced by two $^{198}Hg^{+}$ ions in a trap and observed the interference pattern in the light scattered from the two ions [2]. However, it was shown that, in the strong-field limit, the two-particle collective dressed states are uniformly populated so the interference vanishes at strong driving [3–6]. This restricts potential applications, e.g., in coherent backscattering from disordered structures of atoms [7], the generation of squeezed coherent light by scattering light off of a regular structure [8], the lithography [9], or precision measurements and optical information processing.

Macovei *et al.* investigated the radiation from a collective of atoms [10] and, very recently, they proposed a scheme to recover first-order interference with almost full visibility in strong fields by tailoring the surrounding electromagnetic bath with a suitable frequency dependence, e.g., with the help of cavities [11]. In the modified reservoirs, the collective many-particle dressed states were repopulated so that the possible scattering pathways were modified and resulted in the recovery of the interference.

In this article, we propose a different scheme to recover the spatial interference of resonance fluorescence from two duplicated two-level atoms via controlling the relative phase of driving fields. The atomic system has been investigated before by Bouchene and coworkers [12]. In these articles, the authors focused on the coherent control of the medium gain for the probe pulse and the effective susceptibility, as well as slow light caused by coherent Zeeman oscillations. The precision of a two-beam interferometer could be doubled by replacing the direct detection of the beat signal with twofold degenerate atomic vapor resonant with the laser [13]. The propagation effect of elliptical polarized short pulses in such kind of atomic medium was investigated too [14]. In our scheme, the atomic system is driven by two orthogonally polarized fields, and thus a closed-loop system is formed. As the same as the common statement, the spatial interference vanishes when atoms are driven by strong fields. In this loop system, the relative phase significantly impacts the populations and the atomic coherences of each atom. Thus with a proper relative phase, even driven by strong fields, the atoms are no longer equally populated. The interference pattern could be recovered accordingly. However, we find later that if the two fields have the same intensity and a relative phase $\pi/2$, which is equivalent to a circularly polarization, no resonance fluorescence could be detected. This restriction can be removed by adjusting the relative intensity or, alternatively, by replacing one driving field with a standing-wave field and then adjusting the distance between the atoms and the observing screen. Based on the technology of phase control [15], this scheme may provide experimental maneuverability.

II. THE MODEL AND EQUATIONS

The atoms used here are modeled as duplicated two-level atoms [see Fig. 1(a)]. We consider the $F = 1/2 \leftrightarrow F = 1/2$ transition (energy $\hbar\omega_0$) excited by orthogonally polarized fields. The system could be realized in the ^{6}Li atom. The two lower (upper) states $\{|1\rangle,|2\rangle\}$ $(\{|3\rangle,|4\rangle\})$ with energies $E_1 = E_2$ $(E_3 = E_4)$ correspond to the degenerate states of the level $^2S_{1/2}$ $F = 1/2$ $(^2P_{1/2}$ $F = 1/2)$ with $m_F = \pm 1/2$.

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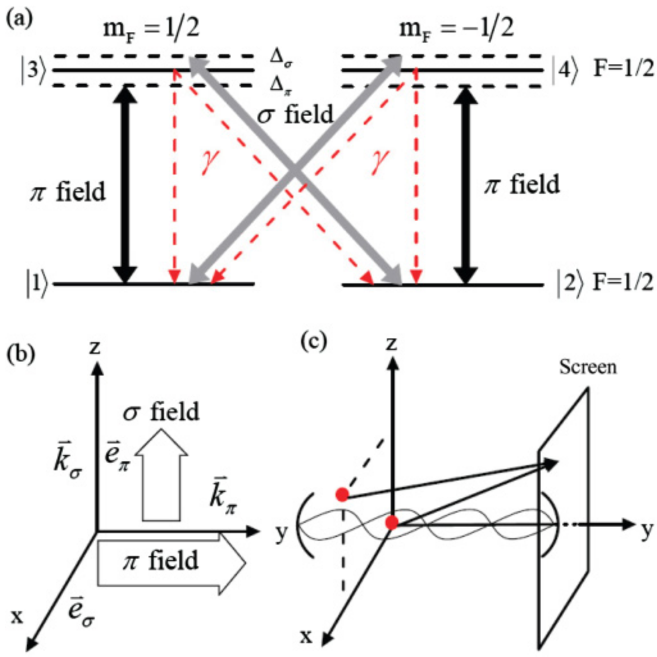

FIG. 1. (Color online) (a) Energy level structure for consideration. The transitions with identical m_F ($|1\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |4\rangle$) are coupled by the π-polarized field, while the transitions with different m_F ($|2\rangle \leftrightarrow |3\rangle$ and $|1\rangle \leftrightarrow |4\rangle$) are coupled by σ-polarized field. (b) Fields configurations. (c) Schematic diagram of the setup.

The transitions with identical m_F (the transitions $|1\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |4\rangle$) are coupled by the π-polarized field, while the transitions with different m_F ($|2\rangle \leftrightarrow |3\rangle$ and $|1\rangle \leftrightarrow |4\rangle$) are coupled by the σ-polarized field. Thus, a closed-loop system is formed, and it allows us to control optical properties of the medium by the phases of the laser fields. The electric fields, with the same frequency ω, are $\vec{E}_\pi(y,t) = \vec{e}_z \epsilon_\pi(y) e^{-i(\omega t - ky)} + \text{c.c.}$ and $\vec{E}_\sigma(z,t) = \vec{e}_x \epsilon_\sigma(z) e^{-i(\omega t - kz)} e^{-i\phi} + \text{c.c.}$, where ϵ_i is the amplitude with $i \in \{\sigma,\pi\}$, ω is the frequency, k is the wave vector, and ϕ is the relative phase between these two driving fields. We assume that both excited states have the same decay rate γ to the each lower level. During these decades, great achievements in trapping neutral atoms have been made [16], among which the magneto-optical trap [17] is a useful tool for producing laser-cooled and trapped neutral atoms. Single-atom experiments were reported [18], too. Based on the experimental research on trapping neutral atoms [16–18], the atoms could be trapped for seconds and be cooled below the Doppler limit. If the atom is initially well trapped, during the detection, the atom will move a few micrometers. Thus the quantized motion and the effect of recoil could be neglected in our calculation.

The atomic dipole operator is the sum of atomic raising $\mu^\uparrow$ and lowering $\mu^\downarrow$ operators whose components are [19]

$$\mu_x^\downarrow = \mu(|1\rangle\langle 4| + |2\rangle\langle 3|)\hat{x}, \tag{1a}$$

$$\mu_y^\downarrow = -i\mu(|1\rangle\langle 4| - |2\rangle\langle 3|)\hat{y}, \tag{1b}$$

$$\mu_z^\downarrow = \mu(|2\rangle\langle 4| - |1\rangle\langle 3|)\hat{z}, \tag{1c}$$

where $\mu_k^\downarrow$ is the k component of the atomic dipole, μ is the dipole matrix element, and $\hat{x}$, $\hat{y}$, and $\hat{z}$ are the usual Cartesian unit vectors. In the interaction picture, the Hamiltonian of the system in an appropriate rotating frame can be written as

$$H = \hbar \begin{pmatrix} 0 & 0 & \Omega_\pi & -\Omega_\sigma e^{-i\phi} \\ 0 & 0 & -\Omega_\sigma^* e^{-i\phi} & -\Omega_\pi \\ \Omega_\pi & -\Omega_\sigma e^{i\phi} & \Delta & 0 \\ -\Omega_\sigma^* e^{i\phi} & -\Omega_\pi & 0 & \Delta \end{pmatrix}, \tag{2}$$

where $\Delta = \omega_0 - \omega$ is the detuning, and the Rabi frequencies are defined as $\Omega_\pi = \mu\epsilon_\pi/2\hbar$ and $\Omega_\sigma = \mu\epsilon_\sigma/2\hbar$. The dynamics of the system can be described using density-matrix approach as

$$\dot{\rho} = -\frac{i}{\hbar}[H,\rho] + L[\rho(t)]. \tag{3}$$

The Liouvillian matrix $L[\rho(t)]$, which describes relaxation by spontaneous decay, is given by

$$L[\rho(t)] = \begin{pmatrix} \gamma(\rho_{33}+\rho_{44}) & 0 & -\gamma\rho_{13} & -\gamma\rho_{14} \\ 0 & \gamma(\rho_{33}+\rho_{44}) & -\gamma\rho_{23} & -\gamma\rho_{24} \\ -\gamma\rho_{31} & -\gamma\rho_{32} & -2\gamma\rho_{33} & -2\gamma\rho_{34} \\ -\gamma\rho_{41} & -\gamma\rho_{42} & -2\gamma\rho_{43} & -2\gamma\rho_{44} \end{pmatrix}. \tag{4}$$

We define $n_g = \rho_{11} + \rho_{22}$ and $n_e = \rho_{33} + \rho_{44}$ as the ground and excited populations, and the coherences $\rho_\pi = \rho_{42} - \rho_{31}$, $\rho_\sigma = \rho_{32} + \rho_{41}$ are responsible for the π- and σ-polarized radiated fields, respectively. We solve the density-matrix equation (3) in steady state while considering the situation that both driving fields are exactly resonant with corresponding transitions ($\Delta = 0$) and all parameters are dimensionless and normalized by γ and we have

$$n_e = \frac{1}{2}\left[1 - \frac{\Omega_\sigma^2 + \Omega_\pi^2}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2 + 4\Omega_\sigma^2\Omega_\pi^2(\cos 2\phi + 1)}\right], \tag{5a}$$

$$\rho_\sigma = \frac{\Omega_\sigma\left[(\Omega_\pi^2 - \Omega_\sigma^2)\sin\phi + i(\Omega_\pi^2 + \Omega_\sigma^2)\cos\phi\right]}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2 + 4\Omega_\sigma^2\Omega_\pi^2(\cos 2\phi + 1)}, \tag{5b}$$

$$\rho_\pi = \frac{\Omega_\pi\left[i(\Omega_\pi^2 + \Omega_\sigma^2\cos 2\phi) - \Omega_\sigma^2\sin 2\phi\right]}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2 + 4\Omega_\sigma^2\Omega_\pi^2(\cos 2\phi + 1)}. \tag{5c}$$

From the steady-state solution, it can be easily seen that there is an additional parameter ϕ to control the atomic dynamics. Note that if the driven fields are strong, the absorption and dispersion for each polarization, as well as the populations, change rapidly around $\phi = \frac{\pi}{2} \pm 2n\pi$ (n is an arbitrary integer).

III. THE INTERFERENCE PATTERN

Our aim is to investigate the far-field interference pattern from two duplicated two-level atoms. For the case of a single atom that interacts with two classical laser light fields linearly polarized along the x axis and z axis, respectively, the steady-state solutions for the atomic coherences and populations have been calculated. Now in our calculation of the far-field interference pattern, we consider that the separation between the atoms is large enough that they may be treated independently. The observing screen is placed in the far-field (large y) and oriented in the xz plane, illustrated in Fig. 1(c). At a point (τ_1,τ_2) on the screen (where τ_i is the light travel time from the ith atom to the observation point, $i = 1,2$), the intensity of the light is

$$I(\tau_1,\tau_2) \propto \langle E_x^{\uparrow} E_x^{\downarrow} + E_z^{\uparrow} E_z^{\downarrow} \rangle, \tag{6}$$

where

$$E_k^{\uparrow}(t;\tau_1,\tau_2) \propto e^{-i\omega(t-\tau_1)} u_k^{\uparrow} + e^{-i\omega(t-\tau_2)} \mathcal{U}_k^{\uparrow}, \tag{7}$$

for $k \in x,z$, u, and $\mathcal{U}$ are the atomic dipoles of the first and second atoms, respectively, and ω is the angular frequency of the laser light. Since the atoms can be considered to be independent and identical, the intensity of the interference pattern when all the light is detected is given by

$$\begin{aligned} I(\tau_1,\tau_2) \propto & \langle u_x^{\uparrow} u_x^{\downarrow} + \mathcal{U}_x^{\uparrow} \mathcal{U}_x^{\downarrow} + u_z^{\uparrow} u_z^{\downarrow} + \mathcal{U}_z^{\uparrow} \mathcal{U}_z^{\downarrow} \rangle \\ & + \langle u_x^{\uparrow} \mathcal{U}_x^{\downarrow} \rangle e^{i\omega(\tau_1-\tau_2)} + \langle u_x^{\downarrow} \mathcal{U}_x^{\uparrow} \rangle e^{-i\omega(\tau_1-\tau_2)} \\ & + \langle u_z^{\uparrow} \mathcal{U}_z^{\downarrow} \rangle e^{i\omega(\tau_1-\tau_2)} + \langle u_z^{\downarrow} \mathcal{U}_z^{\uparrow} \rangle e^{-i\omega(\tau_1-\tau_2)}. \end{aligned} \tag{8}$$

The components in Eq. (8) in the steady state are

$$\begin{aligned} \langle u_x^{\uparrow} u_x^{\downarrow} \rangle_{ss} &= \langle \mathcal{U}_x^{\uparrow} \mathcal{U}_x^{\downarrow} \rangle_{ss} \\ &\propto \mu^2 \langle (|4\rangle\langle 1| + |3\rangle\langle 2|)(|1\rangle\langle 4| + |2\rangle\langle 3|) \rangle_{ss} \\ &= \mu^2 \langle |4\rangle\langle 4| + |3\rangle\langle 3| \rangle_{ss} \\ &= \mu^2 n_e, \end{aligned} \tag{9a}$$

and similarly

$$\langle u_z^{\uparrow} u_z^{\downarrow} \rangle_{ss} = \langle \mathcal{U}_z^{\uparrow} \mathcal{U}_z^{\downarrow} \rangle_{ss} = \mu^2 n_e, \tag{9b}$$

$$\langle u_x^{\uparrow} \mathcal{U}_x^{\downarrow} \rangle_{ss} = \langle u_x^{\uparrow} \mathcal{U}_x^{\uparrow} \rangle_{ss}^* \propto \mu^2 \langle u_x^{\uparrow} \rangle_{ss} \langle \mathcal{U}_x^{\downarrow} \rangle_{ss} = \mu^2 \rho_\sigma \rho_\sigma^*, \tag{9c}$$

$$\langle u_z^{\uparrow} \mathcal{U}_z^{\downarrow} \rangle_{ss} = \langle u_z^{\uparrow} \mathcal{U}_z^{\uparrow} \rangle_{ss}^* \propto \mu^2 \langle u_z^{\uparrow} \rangle_{ss} \langle \mathcal{U}_z^{\downarrow} \rangle_{ss} = \mu^2 \rho_\pi \rho_\pi^*. \tag{9d}$$

Unlike the results in Refs. [5,6], the cross terms $\langle u_x^{\uparrow} \mathcal{U}_x^{\downarrow} \rangle_{ss}$ and $\langle \mathcal{U}_x^{\uparrow} u_x^{\downarrow} \rangle_{ss}$ now contribute to the total intensity due to the driving of the σ-polarized field so $|\rho_\sigma| \neq 0$. Thus the intensity in Eq. (8) is

$$I(\tau_1,\tau_2) \propto 4n_e \left\{ 1 + \frac{1}{2n_e}(\rho_\sigma \rho_\sigma^* + \rho_\pi \rho_\pi^*) \cos[\omega(\tau_1 - \tau_2)] \right\}, \tag{10}$$

The visibility of the interference pattern is defined as $V = (I_{\max} - I_{\min})/(I_{\max} + I_{\min})$. In our duplicated two-level atomic system, the visibility can be calculated by using the steady-state solutions [Eqs. (5b) and (5c)]

$$\begin{aligned} V &= \frac{1}{2n_e}(\rho_\sigma \rho_\sigma^* + \rho_\pi \rho_\pi^*) \\ &= \frac{1}{2} \frac{\Omega_\sigma^2 + \Omega_\pi^2}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2 + 4\Omega_\sigma^2 \Omega_\pi^2 (\cos 2\phi + 1)}. \end{aligned} \tag{11}$$

Compared with Eq. (5a), it is easy to see that

$$V + n_e = 1/2. \tag{12}$$

We note that both of the two components polarized in the x and z axis contribute to the total intensity detected on the screen. As the π- and σ-polarized fields are applied simultaneously, the two components could not be separated. The visibility is always less than one-half, as the σ-polarized scattering light is incoherent [2].

From Eq. (11), we can see that the interference pattern of the resonance fluorescence from two duplicated two-level atoms is related to the Rabi frequencies of the driving fields, and, what is more, to the relative phase ϕ. V as the function of the relative phase between these two driving fields reaches its maximum $V_{\max}$ when $\cos 2\phi = -1$ ($\phi = \frac{\pi}{2} \pm 2n\pi$, n is an arbitrary integer):

$$V_{\max} = \frac{1}{2} \frac{\Omega_\sigma^2 + \Omega_\pi^2}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2}. \tag{13}$$

From Eq. (5a), in this case, the excited population reaches the minimum

$$n_{e\,\min} = \frac{1}{2}\left[1 - \frac{\Omega_\sigma^2 + \Omega_\pi^2}{\Omega_\sigma^2 + \Omega_\pi^2 + 2(\Omega_\sigma^2 - \Omega_\pi^2)^2}\right]. \tag{14}$$

It has been confirmed that in the strong driving situation, the interference pattern vanishes [3–6]. From our main results, Eqs. (11)–(14), we find that in our scheme, the visibility could be realized even in the strong driving situation due to the relative phase ϕ. Figure 2 shows how the visibility evolves under different driving situations. Without the phase difference, the visibility will fall toward zero rapidly while increasing the driving field intensities [see Fig. 2(a)], because when $\phi = 0$ the atoms are equally populated under strong driving [$n_e(\Omega_{\sigma,\pi} \to \infty) \to 1/2$]. When the phase difference between the two driving fields ϕ is nonzero, the equally populated situation will be destroyed in the strong driving fields, therefore phase-dependent interference of resonance fluorescence will show up. When $\phi = \pi/2$, the visibility will reach its maximum, shown in Fig. 2(b), because the atomic coherences $\rho_{\sigma,\pi}$ have been changed significantly around $\phi = \pi/2$. In Fig. 2(c), we show the visibility as a function of the phase ϕ and the driving intensities (we assume that the two fields have the same intensities). As we have analyzed from the expression of V, when $\phi = \pi/2$, the visibility reaches its maximum [also see Fig. 2(d)]. However, with the same driving intensities, from Eqs. (13) and (14) we find that $V_{\max} = 1/2$, and there is no population on the excited state, i.e., no fluorescence could be detected. Physically, if the driving

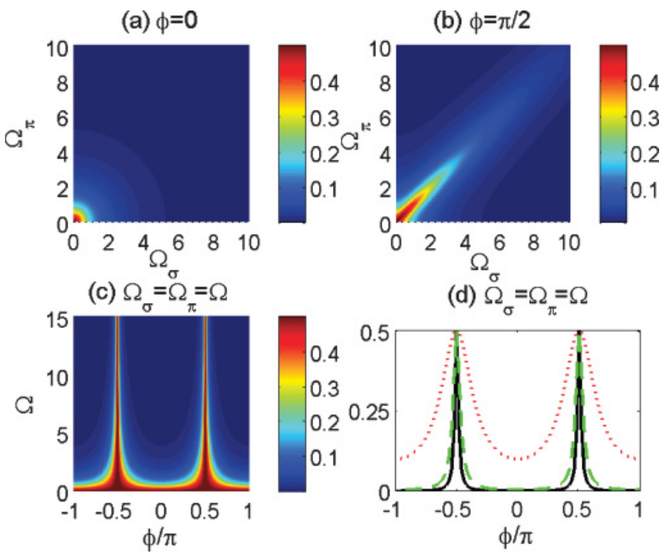

FIG. 2. (Color online) The visibility V in different driving situations: (a) the relative phase is zero; (b) the relative phase is $\pi/2$; and (c) the intensities of the driving fields are equal. (d) Some examples in (c). The dotted (red) curve: $\Omega = 1$, the dashed (green) curve: $\Omega = 5$, and the solid (black) curve: $\Omega = 10$.

fields have equal amplitudes and a mutual phase shift of $\pi/2$, each atom is actually excited by circularly polarized field. In the steady-state regime each atom will be spin oriented along y axis (see Fig. 1) and will not interact with the driving field. The system is transparent for both π- and σ-polarized fields, confirmed by Eqs. (5a)–(5c) that $n_e = \rho_\sigma = \rho_\pi = 0$. We find from these results that the relative phase is the key parameter in the recovery of the interference pattern. Only with a proper relative phase that the visibility could be recovered in strong driven fields. But there is a restriction that when the two orthogonally polarized fields have equal intensities, $\phi \neq \pi/2$ should be satisfied.

In order to remove the above restriction and improve the experimental realizability, we investigate the influence of the strengths of those two driving fields on the interference, exactly when $\phi = \pi/2$. We define the ratio of the Rabi frequencies $r = \Omega_\pi/\Omega_\sigma$. It is shown in Fig. 3(a) that a peak emerges when $r = 1$. While increasing the strengths of the driving fields, the peak becomes narrower. This, however, provides the feasibility to recover the interference under strong driving with $\phi = \pi/2$ by choosing the ratio between the two driving fields properly. We choose $r = 0.9$, 0.95, and 0.99 for examples in Fig. 3(b). It is shown that as r gets closer to 1, the visibility could survive

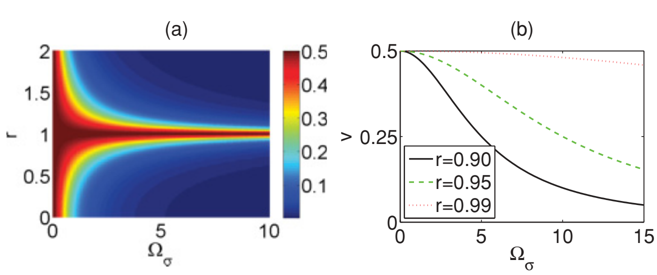

FIG. 3. (Color online) The visibility V by adjusting the relative intensities when $\phi = \pi/2$. (a) V as the functions of r and the intensity Ω_σ. (b) Some examples in (a).

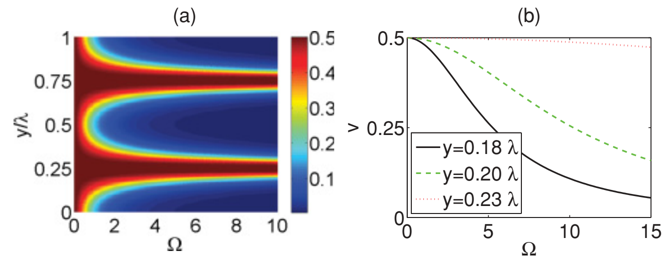

FIG. 4. (Color online) The visibility V by adjusting the distance between the atoms and the screen when $\phi = \pi/2$. (a) V as the functions of the position y and the intensity Ω. (b) Some examples in (a).

even when driven by strong fields. Thus, by adjusting the relative intensity of driving fields, the pattern would reappear under strong driving even when $\phi = \pi/2$.

We note that the adjusting of the relative intensity works only when r is modified around 1. It is known that the intensity of a standing-wave field is periodic in space and oscillates between its minimum and maximum. An idea came into our mind that we can replace one of the driving field with a standing-wave field. As the interference pattern is observed in the xz plane, we then use a π-polarized standing-wave field, which is applied along the y axis and therefore $\Omega_\pi(y) = \Omega\sin(ky)$. The observing screen is fixed at the end of the cavity and the cavity can be moved along the y direction [illustrated in Fig. 1(c)], and the atoms are located in the xz plane so they experience the same driving fields. The intensity of the standing wave is position dependent, therefore the interference pattern in the xz plane is related to the detected distance between the screen and the plane where the atoms are located. In order to compare with the above work, we choose $\Omega_\sigma = \Omega$, i.e., $r = \sin(ky)$. The result is shown in Fig. 4(a). Peaks appear at the antinodes, where $r = 1$. By changing the location of the screen, the visibility could be recovered. In Fig. 4(b), we choose the distance y to correspond with the values of r in Fig. 3(b), and we obtain the same results. In other words, when the relative phase is fixed to $\pi/2$, the interference pattern in the xz plane could be revealed by moving the screen along the y direction. Controlling the distance between the atoms and the screen is an alternative choice as compare to adjust the intensities of the driving fields.

IV. CONCLUSION

In summary, the recovery of interference of resonance fluorescence from two duplicated two-level atoms by relative phase control is investigated. The interference pattern can be recovered in the fluorescence light of strongly driven atoms due to effect of the relative phase between the two driving fields on the populations and the atomic coherences. However, when $\phi = \pi/2$ and $\Omega_\sigma = \Omega_\pi$, the atoms do not interact with the driving fields, and no fluorescence could be detected. By adjusting the relative intensities, this problem can be solved. A scheme of recovering the visibility by using a standing-wave field is proposed, too. By replacing the π-polarized field with a standing wave, the interference pattern in the xz plane could by revealed periodically by moving the screen along the y direction.

ACKNOWLEDGMENTS

L. L. Jin thanks Dr. X. T. Xie for useful discussions and the support from the Science Foundation of Northwest University (No. 10NW14). This work is supported by the 973 Program under Grant No. 2006CB806000, and the National Natural Sciences Foundation of China (Grants No. 60708008, No. 60978013, No. 61008016, No. 11074263, and No. 10874194), the Key Basic Research Foundation of Shanghai (Grant No. 08JC 1409702), and Shanghai Science and Technology Commission (Grant No. 10530704800).

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