Effects of spatial interference on intensity–intensity correlations in collective two-atom emission*

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In this paper, we investigate the effects of the spatial variation of driving-laser phase in a collective two-atom system on the intensity–intensity correlations of the resonant fluorescence. It is shown that the intensity–intensity correlations exhibit quite different characteristics for the different values of the spatial phase of the laser at the position of the two atoms in both cases of the weak and strong driving lasers. Our results suggest that the intensity–intensity correlations can serve as a probe of the spatial interference effect arising from the spatial variation of the laser phase.

Keywords: spatial interference, intensity–intensity correlations, collective two-atomic system
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1. Introduction

Cooperative behaviour can occur in a system of two or more atoms interacting with common vacuum radiation field when the interatomic separation is smaller than the relevant resonant wavelength.[1,2] Since the concept of collective spontaneous emission was first introduced by Dicke,[3] collective effects of the interactive atoms have become a subject of great interest in quantum optics and have undergone extensive investigation.[4−11] Especially, recent studies[12,13] show that the collective spontaneous emission can create entanglement between two cooperative atoms which is a crucial element in quantum information processing.

Most recently, in a collective system of two laser-driven two-level atoms, Das *et al.*[14] investigated quantum interference effects which arise from the spatial variation of the pumping-laser phase. They found that when the atoms are weakly excited by the laser, the fluorescence intensity and the entanglement between the two atoms can be considerably enhanced if the propagation of the laser is arranged to be parallel or antiparallel to the interatomic axis, as compared to the case of laser propagation perpendicular to the axis. Therefore, they proposed that the intensity spectrum of the resonant fluorescence can serve as a probe of the new interference effects. In the case of a strong pumping laser, however, it is shown in Ref. [14] that the spatial interference effects on the fluorescence spectrum and interatomic entanglement are less obvious.

In this paper, we therefore investigate the effects of the spatial phase variation of the driving laser on the intensity–intensity correlations of the resonant fluorescence emitted by the collective atomic system, on the basis of the current theoretical and experimental efforts on intensity–intensity correlations,[15−22] and the fact that nonclassical photon correlations (entangled photons[23]) are considerably significant in quantum information processing. It is found that the photon correlations exhibit quite strong dependence on the spatial phase factor θ of the laser. Especially, when the atoms are excited by the strong laser, quantum interference with the spatial phase $\theta = 0$ leads to noticeable oscillations in the intensity correlations even for large time delay τ, quite different from the case of $\theta = \pi/2$ in which the correlations cease to oscillate in the small time delay. Also, we show that when the laser–atom detuning is equal to the interatomic strength, the photon correlation is insensitive to the spatial variation of the laser phase. Our results suggest that the intensity–intensity correlations can serve as a good probe of the spatial interference effects arising from the spatial variation of the laser phase. This paper is arranged as follows: in Section 2, the model is introduced and the working equation is presented. In Section 3, the intensity–intensity correlation function is calculated, and we discuss in detail the influences of the spatial-phase variation of the driving laser on the

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intensity correlations of the resonant fluorescence. In the last section, we give our main summary.

2. Model

Our system consists of two identical two-level atoms with transition frequencies ω_0 and dipole moments $\boldsymbol{\mu}$. The atoms are assumed to be separated by a distance r_{12} and driven by a monochorionic laser of frequency ω_L described by

$$\boldsymbol{E}(t) = \boldsymbol{\varepsilon}_{\rm L} {\rm e}^{-{\rm i}(\boldsymbol{k}_{\rm L}\cdot\boldsymbol{r}+\omega_{\rm L}t)} + {\rm c.c.}, \tag{1}$$

with wave vector $\boldsymbol{k}_{\rm L}$ (see Fig. 1). Taking into account the damping of the atoms in vacuum, in the rotating-wave approximation, the density operator ρ of the system is governed by the following master equation[1]

$$\frac{{\rm d}\rho}{{\rm d}t} = -{\rm i}[H_{\rm int}, \rho] + \sum_{i,j=1}^{2} \gamma_{ij}(2S_i^- \rho S_j^+ - S_j^+ S_i^- \rho - \rho S_j^+ S_i^-), \tag{2}$$

where S_j^+ and S_j^- are the raising and lowering operators for the jth atom. The Hamiltonian $H_{\rm int}$ in the interaction picture is given by

$$H_{\rm int} = \frac{1}{2}\sum_{j=1}^{2}(\Delta S_j^+ S_j^- + \Omega {\rm e}^{-{\rm i}\,\phi_j} S_j^+) + \Omega_{12} S_1^+ S_2^- + {\rm H.c.}, \tag{3}$$

where $\Delta = \omega_0 - \omega_{\rm L}$, $\Omega = \boldsymbol{\mu}\cdot\boldsymbol{\varepsilon}_{\rm L}$ are the Rabi frequencies of the laser field and are taken to be real by choosing the initial phase of the electromagnetic field, and $\phi_j = \boldsymbol{k}_{\rm L}\cdot\boldsymbol{r}_j$ is the spatial phase factor of the laser at the position $\boldsymbol{r}_j$ of the jth atom. The variation of the spatial phase means that the two atoms experience different driving fields, which leads to interesting interference effects such as the modification of the fluorescence intensity and the interatomic entanglement degree.[14] Here, we focus on the spatial interference due to the effect of spatial variation of the laser on the intensity–intensity correlations of the resonance fluorescence emitted by the atoms. The interatomic interaction Ω_{12}, coming from the interaction between the two atoms mediated by the vacuum radiation field, takes the form of

$$\Omega_{12} = \frac{3\gamma}{2}\Big\{-(1-\cos^2\beta)\frac{\cos(k_0 r_{12})}{k_0 r_{12}} + (1-3\cos^2\beta)\Big[\frac{\sin(k_0 r_{12})}{(k_0 r_{12})^2} + \frac{\cos(k_0 r_{12})}{(k_0 r_{12})^3}\Big]\Big\}, \tag{4}$$

where β is the angle between the dipole moment $\boldsymbol{\mu}$ and the vector $\boldsymbol{r}_{12} = \boldsymbol{r}_2 - \boldsymbol{r}_1$, and we will take $\beta = \pi/2$ in the following discussions; $k_0 = \omega_0/c$, $k_{12} = |\boldsymbol{k}_{12}|$, and $r_{12} = |\boldsymbol{r}_{12}|$. We can see that the interatomic interaction described by Ω_{12} is dependent on the interatomic separation r_{12} and it becomes especially significant at the small separation. The collective decay rates γ_{ij} of the atoms satisfy $\gamma_{12} = \gamma_{21}$ and

$$\gamma_{12} = \frac{3\gamma}{2}\Big\{(1-\cos^2\beta)\frac{\sin(k_0 r_{12})}{k_0 r_{12}} + (1-3\cos^2\beta)\Big[\frac{\cos(k_0 r_{12})}{(k_0 r_{12})^2} - \frac{\sin(k_0 r_{12})}{(k_0 r_{12})^3}\Big]\Big\}, \tag{5}$$

with $\gamma = \gamma_{jj}$ being the single-atom spontaneous emission rate. Here, the interatomic interaction Ω_{12} and the cooperative damping γ_{12} of the atoms determine the collective properties of the two-atom system. From Eqs. (4) and (5), for small interatomic distance $k_0 r_{12} \ll 1$, we have

$$\Omega_{12} = \frac{3\gamma}{2(k_0 r_{12})^3}(1-3\cos^2\beta), \quad \gamma_{12} = \gamma, \tag{6}$$

while for large separation between the two atoms, $k_0 r_{12} \gg 1$, one has $\Omega_{12} = 0$ and $\gamma_{12} = 0$ and thus the collective effect vanishes. Now if performing transformation $S_j \to S_j {\rm e}^{-{\rm i}\,\phi_j}$, the master equation (2) is unchanged, except for

$$\Omega_{12} \to \Omega_{12}{\rm e}^{\pm{\rm i}\,\phi_r}, \quad \gamma_{12} \to \gamma_{12}{\rm e}^{\pm{\rm i}\,\phi_r}, \tag{7}$$

where the relative phase

$$\phi_r = \phi_1 - \phi_2 = k_{\rm L}\cdot\boldsymbol{r}_{12} = \frac{2\pi}{\lambda} r_{12}\cos\theta, \tag{8}$$

where λ is the wavelength of the laser and θ the angle between the propagation direction of the laser and the interatomic axis. According to Eq. (7), we see that the variation of the laser spatial phase will modulate the collective parameters Ω_{12} and γ_{12}. Therefore, the collective effects can be modified by arranging the propagation of the laser.

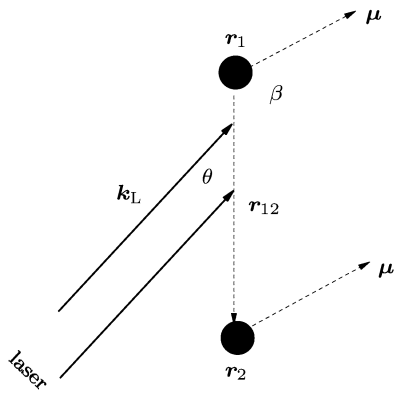


Fig. 1. The geometry of the present system, in which θ is the angle between the propagation vector $\boldsymbol{k}_{\rm L}$ of the laser and the interatomic axis $\boldsymbol{r}_{12}$, and β is the angle between the vector $\boldsymbol{k}_{\rm L}$ and the atomic dipole $\boldsymbol{\mu}$.

By defining the following atomic operators $\hat{X}_j$

$$\begin{aligned}
&\hat{X}_1 = S_1^+, \qquad \hat{X}_2 = S_1^-,\\
&\hat{X}_3 = S_2^+, \qquad \hat{X}_4 = S_2^-,\\
&\hat{X}_5 = S_1^+S_1^-, \qquad \hat{X}_6 = S_2^+S_2^-,\\
&\hat{X}_7 = S_1^+S_2^-, \qquad \hat{X}_8 = S_2^+S_1^-,\\
&\hat{X}_9 = S_1^+S_2^+, \qquad \hat{X}_{10} = S_1^-S_2^-,\\
&\hat{X}_{11} = S_1^+S_1^-S_2^-, \quad \hat{X}_{12} = S_1^+S_1^-S_2^+,\\
&\hat{X}_{13} = S_1^-S_2^+S_2^-, \quad \hat{X}_{14} = S_1^+S_2^+S_2^-,\\
&\hat{X}_{15} = S_1^+S_1^-S_2^+S_2^-,
\end{aligned} \tag{9}$$

according to the master equation (2), the equations of motion of the expectation values of the atomic operators $\bar{X}_j = \langle \hat{X}_j \rangle$ can be derived as

$$\begin{aligned}
\dot{\bar{X}}_1 &= -\bar{\gamma}_1\bar{X}_1 - \bar{\gamma}_-\bar{X}_3 + 2\bar{\gamma}_-\bar{X}_{12} + 2\xi_1^*(\bar{X}_5 - 1/2),\\
\dot{\bar{X}}_3 &= -\bar{\gamma}_1\bar{X}_3 - \bar{\gamma}_-\bar{X}_1 + 2\bar{\gamma}_-\bar{X}_{14} + 2\xi_2^*(\bar{X}_6 - 1/2),\\
\dot{\bar{X}}_5 &= -2\gamma\bar{X}_5 - \xi_1\bar{X}_1 - \xi_1^*\bar{X}_2 - \bar{\gamma}_+\bar{X}_7 - \bar{\gamma}_-\bar{X}_8,\\
\dot{\bar{X}}_7 &= -2\gamma\bar{X}_7 - \xi_2\bar{X}_1 - \xi_1^*\bar{X}_4 - \bar{\gamma}_+\bar{X}_5 - \bar{\gamma}_-\bar{X}_6\\
&\quad + 2\xi_1^*\bar{X}_{11} + 2\xi_2\bar{X}_{14} + 4\gamma_{12}\bar{X}_{15},\\
\dot{\bar{X}}_9 &= -2\bar{\gamma}_1\bar{X}_9 - \xi_2^*\bar{X}_1 - \xi_1^*\bar{X}_3 + 2\xi_1^*\bar{X}_{12} + 2\xi_2^*\bar{X}_{14},\\
\dot{\bar{X}}_{11} &= -2\bar{\gamma}_2\bar{X}_{11} - \bar{\gamma}_-\bar{X}_{14} - \xi_2\bar{X}_5 - \xi_1\bar{X}_7\\
&\quad - \xi_1^*\bar{X}_{10} + 2\xi_2\bar{X}_{15},\\
\dot{\bar{X}}_{13} &= -2\bar{\gamma}_2\bar{X}_{14} - \bar{\gamma}_-\bar{X}_{11} - \xi_1\bar{X}_6 - \xi_2\bar{X}_8\\
&\quad - \xi_2^*\bar{X}_{10} + 2\xi_1\bar{X}_{15},\\
\dot{\bar{X}}_{15} &= -4\gamma\bar{X}_{15} - \xi_2^*\bar{X}_{11} - \xi_2\bar{X}_{12}\\
&\quad - \xi_1^*\bar{X}_{13} - \xi_1\bar{X}_{14},
\end{aligned} \tag{10}$$

where $\bar{\gamma}_j = \gamma_j - \mathrm{i}\,\Delta$ $(j = 1, 2)$, $\bar{\gamma}_+ = \gamma_{12} + \mathrm{i}\,\Omega_{12}$, $\bar{\gamma}_- = \gamma_{12} - \mathrm{i}\,\Omega_{12}$, and $\xi_j = \mathrm{i}\,\Omega \mathrm{e}^{\mathrm{i}\,\phi_j}$. The remaining equations can be obtained by taking complex conjugates of Eq. (10). By writing the Bloch equation as

$$\frac{\mathrm{d}}{\mathrm{d}t}\bar{\boldsymbol{X}}(t) = \boldsymbol{M}\bar{\boldsymbol{X}}(t) + \boldsymbol{N}, \tag{11}$$

where $\boldsymbol{X} = (\bar{X}_1,\ \bar{X}_2, \ldots, \bar{X}_{15})^{\rm T}$, $\boldsymbol{M}$ is a 15×15 matrix and $\boldsymbol{N} = (-\xi_1^*,\ -\xi_1,\ -\xi_2^*,\ -\xi_2,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0,\ 0)^{\rm T}$, the formal solution of the above equation is given by

$$\boldsymbol{X}(t) = \mathrm{e}^{\boldsymbol{M}t}\bar{\boldsymbol{X}}(0) - (I - \mathrm{e}^{\boldsymbol{M}t})\boldsymbol{M}^{-1}\boldsymbol{N}. \tag{12}$$

When $t \to \infty$, the stationary solution is obtained

$$\bar{\boldsymbol{X}}(\infty) = -\boldsymbol{M}^{-1}\boldsymbol{N}. \tag{13}$$

3. Intensity–intensity correlations

The second-order two-time correlation function of the light field is defined as

$$\begin{aligned}
G^{(2)}(\boldsymbol{r}, t; \boldsymbol{r}, t+\tau) = \langle E^{(-)}(\boldsymbol{r}, t)E^{(-)}(\boldsymbol{r}, t+\tau)\\
\times E^{(+)}(\boldsymbol{r}, t+\tau)E^{(+)}(\boldsymbol{r}, t)\rangle,
\end{aligned} \tag{14}$$

where $E^{(\pm)}(\boldsymbol{r}, t)$ are respectively the positive- and negative-frequency parts of the electric field operators at the observation point $\boldsymbol{r} = r\hat{r}$. The two-time correlation function is proportional to the joint probability of detecting one photon at time t and another photon at time $t + \tau$, yielding information about the photon statistics of the emitted field. With the positive-frequency part of the electric field $E^{(+)}$ in the far-field zone[24]

$$E^{(+)}(\boldsymbol{r}, t) = -k_{\rm L}^2 \sum_{j=1}^{2} \frac{\hat{r} \times (\hat{r} \times \boldsymbol{\mu})}{r} S_j^-(t - r/c)\mathrm{e}^{-\mathrm{i}\,k_0\hat{r}\cdot\boldsymbol{r}_j}, \tag{15}$$

the two-time correlation function can be expressed, in terms of the atomic operators, as

$$\begin{aligned}
G^{(2)}(\boldsymbol{r}, t; \boldsymbol{r}, t+\tau) = f^2(r) \sum_{k,l,m,n=1}^{2} \langle S_k^+(t)\\
\times S_l^+(t+\tau)S_m^-(t+\tau)S_n^-(t)\rangle,
\end{aligned} \tag{16}$$

where $f(r) = (\omega_0/c)^4(\mu^2/r^2)$; we have assumed that the observation direction is perpendicular to the atomic dipole moments $\boldsymbol{\mu}$ and the vector $\boldsymbol{r}_{12}$. The normalized intensity–intensity correlation function $g^{(2)}(\boldsymbol{r}, t; \boldsymbol{r}, t+\tau)$ is defined as

$$g^{(2)}(\boldsymbol{r}, t; \boldsymbol{r}, t+\tau) = \frac{G^{(2)}(\boldsymbol{r}, t; \boldsymbol{r}, t+\tau)}{G^{(1)}(\boldsymbol{r}, t)G^{(1)}(\boldsymbol{r}, t+\tau)}, \tag{17}$$

where $G^{(1)}(\boldsymbol{r}, t)$ is the first-order one-time correlation function, defined as

$$G^{(1)}(\boldsymbol{r},t)=\langle E^{(-)}(\boldsymbol{r},t)E^{(+)}(\boldsymbol{r},t)\rangle=f(r)\sum_{m,n=1}^{2}\langle S_m^+(t)S_n^-(t)\rangle, \tag{18}$$

which describes the intensity of the emitted field.

Therefore, the normalized intensity–intensity correlation function in the steady regime can be expressed as

$$g^{(2)}(\tau)=\lim_{t\to\infty}\sum_{k,l,m,n=1}^{2}\frac{\langle S_k^+(t)S_l^+(t+\tau)S_m^-(t+\tau)S_n^-(t)\rangle}{[\sum_{i,j=1}^{2}\langle S_i^+(t)S_j^-(t)\rangle]^2}. \tag{19}$$

Now we evaluate the two-time correlation function appearing in Eq. (19). According to the quantum regression,[25] the two-time correlation functions for certain atomic operators $\hat{X}_j(t)$ satisfy the same equation as the Bloch equation in Eq. (10), i.e.,

$$\frac{\mathrm{d}}{\mathrm{d}\tau}\langle\hat{X}_i^+(t)\hat{X}(t+\tau)\hat{X}_j^-(t)\rangle=\boldsymbol{M}\langle\hat{X}_i^+(t)\hat{X}(t+\tau)\hat{X}_j^-(t)\rangle+\langle\hat{X}_i^+(t)\hat{X}_j^-(t)\rangle\boldsymbol{N}, \tag{20}$$

with its solution

$$\begin{aligned}\langle\hat{X}_i^+(t)\hat{X}(t+\tau)\hat{X}_j^-(t)\rangle=&\,\mathrm{e}^{\boldsymbol{M}\tau}\langle\hat{X}_i^+(t)\hat{X}(t)\hat{X}_j^-(t)\rangle\\&-(I-\mathrm{e}^{\boldsymbol{M}\tau})\boldsymbol{M}^{-1}\boldsymbol{N}\times\langle\hat{X}_i^+(t)\hat{X}_j^-(t)\rangle.\end{aligned} \tag{21}$$

The two-time correlation functions in Eq.(19) can be obtained as

$$\begin{aligned}\langle S_k^+(t)\hat{X}_j(t+\tau)S_n^+(t)\rangle=&\sum_{j'=1}^{15}(\mathrm{e}^{\boldsymbol{M}\tau})_{jj'}\langle S_k^+(t)\hat{X}_{j'}(t)S_n^-(t)\rangle\\&-[(I-\mathrm{e}^{\boldsymbol{M}\tau})\boldsymbol{M}^{-1}\boldsymbol{N}]_{jj'}\times\langle S_k^+(t)S_n^-(t)\rangle.\end{aligned} \tag{22}$$

Therefore, with Eq. (22) the normalized intensity–intensity correlation of the field in the steady-state regime can be calculated numerically.

In Figs. 2–4, the intensity–intensity correlation function $g^{(2)}(\tau)$ is plotted for different Rabi frequencies Ω and atomic distances. From them we can see that the photon correlations are dependent on the angle θ between the propagation of the laser and the interatomic axis, the interatomic separation r_{12} and the Rabi frequency Ω of the pumping laser. It is shown that with the decrease of interatomic distance r_{12} the stronger oscillations are exhibited in the intensity correlation of the fluorescence light [by comparing Fig. 2(a) with Fig. 2(b)]. This is because the photon-mediated interaction between the two atoms becomes stronger, as a consequence of the reduction of interatomic separations. In addition, from Eq. (8), one can find that the variation of the spatial angle factor θ can be visualized effectively as the changes of the interatomic separation r_{12} which in turn determines the indirect-interactive strength Ω_{12} between the atoms. As a result, as shown in Figs. 2–4, we can see that the effects of the spatial interference which results from the variation of the spatial phase factor θ are obviously demonstrated in the intensity correlations of the fluorescent light. In particular, as shown in the insets in Figs. 2 and 3, the nonclassical antibunching feature $[g^{(2)}(0)<1]$ of the emitted field is much stronger for $\theta=\pi/2$ than that in the case of $\theta=0$. More interestingly, from Fig. 4 we can see that when $\theta=\pi/2$ (i.e., the propagation direction of the laser is perpendicular to the interatomic axis), fast stopping of the oscillation of the correlation function is observed for the strong Rabi frequency Ω of the laser. Inversely, for the case $\theta=0$, very obvious quantum beating is exhibited in the photon correlations. For example, as shown in Fig. 4, for $\theta=\pi/2$ the intensity correlation function $g^{(2)}(\tau)\to 1$ and ceases to oscillate when $\tau>2$, while for $\theta=0$ the damped oscillation of the correlation function can be noticeably observed as $\tau>2$ and the antibunching $g^{(2)}(\tau)<1$ can still exhibit at a certain time delay τ. This can be explained in terms of the collective states of the driven two-atom system.

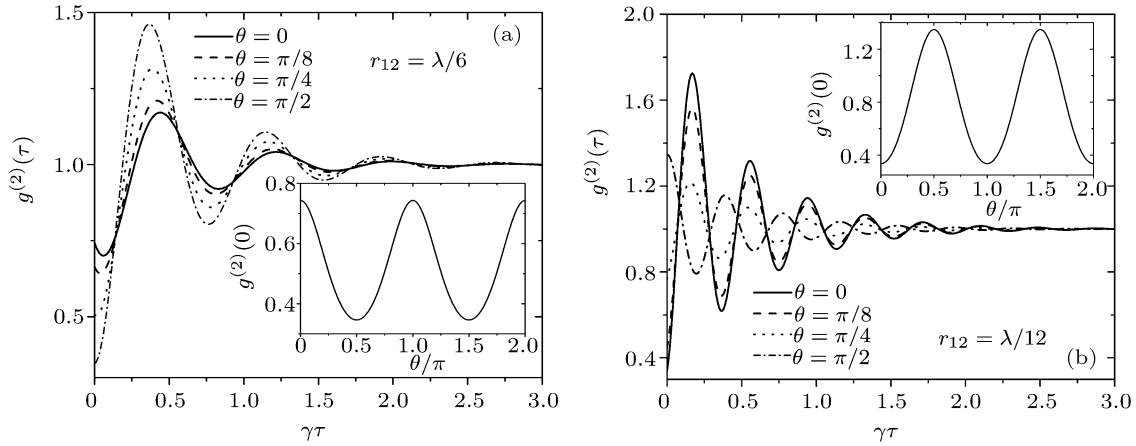


Fig. 2. Dependence of the intensity–intensity correlation function on the spatial angle factor θ for the interatomic separation $r_{12} = \lambda/6$ (a) and $r_{12} = \lambda/12$ (b) with $\Omega/\gamma = 0.1$ and $\Delta/\gamma = 7$. The insets depict the correlation function with $\tau = 0$.

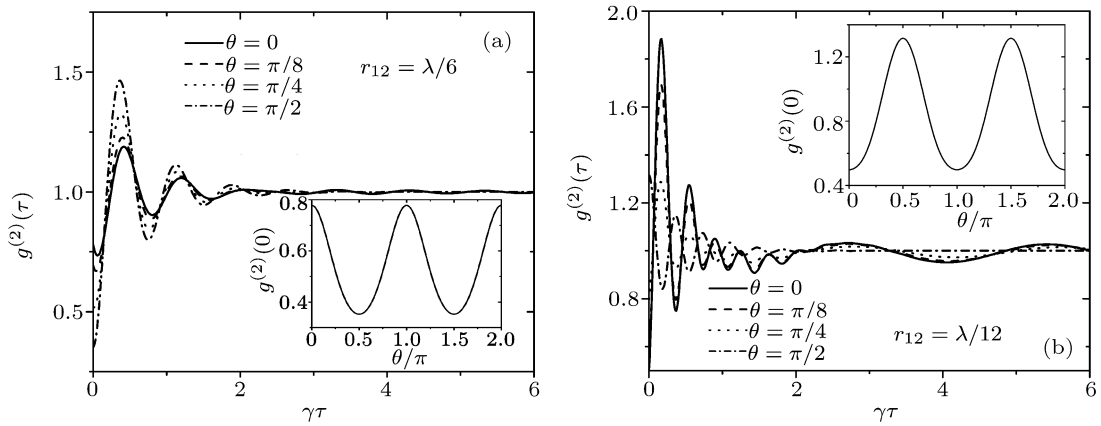


Fig. 3. Dependence of the intensity–intensity correlation function on the spatial angle factor θ for the interatomic separation $r_{12} = \lambda/6$ (a) and $r_{12} = \lambda/12$ (b) with $\Omega/\gamma = 2$ and $\Delta/\gamma = 7$. The insets depict the correlation function with $\tau = 0$.

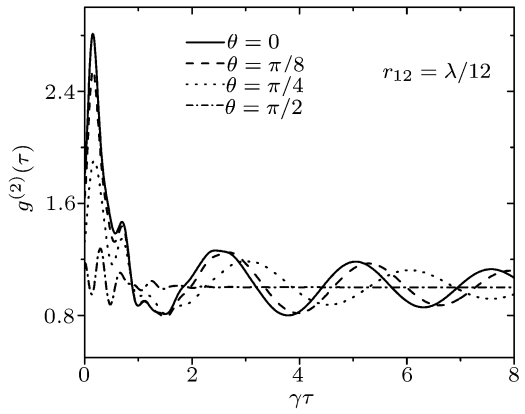


Fig. 4. Dependence of the intensity–intensity correlation function on the spatial angle factor θ for the interatomic separation $r_{12} = \lambda/12$ with $\Omega/\gamma = 5$ and $\Delta/\gamma = 7$.

In the framework of the two-atom collective states, the two-atom system is equivalent to a single four-level atom (shown in Fig. 5). The oscillatory behaviour of the photon correlation originates from the quantum interference between the two possible transition amplitudes $|s(a)\rangle \to |g\rangle$. The amplitude of this modulation becomes maximum when the spatial phase factor $\theta = 0$ (i.e., the propagation direction of the laser is parallel to the interatomic axis). In addition,

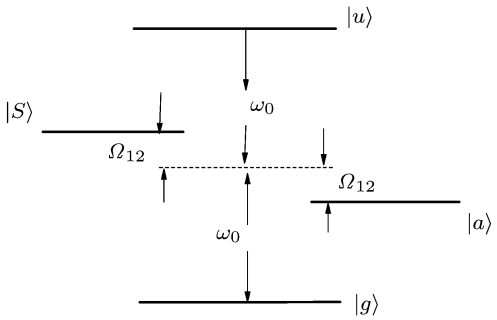


Fig. 5. Collective states of the two-atom system with the upper state $|u\rangle = |e_1e_2\rangle$, the ground state $|g\rangle = |g_1g_2\rangle$, the symmetric state $|s\rangle = (|e_1g_2\rangle + |g_1e_2\rangle)/\sqrt{2}$, and the asymmetric state $|a\rangle = (|e_1g_2\rangle - |g_1e_2\rangle)/\sqrt{2}$. The frequency shift Ω_{12} is due to the dipole–dipole interaction.

for large interatomic separations, Ω_{12} is approximatively zero, and the transitions from $|s\rangle$ and $|a\rangle$ to $|g\rangle$ have the same frequency, leading to the disappearance of the oscillation. As a result, as revealed in Fig. 4, the obvious oscillation of the intensity–intensity correlation function can be observed for $\theta = 0$ and for the small atomic distance r_{12}.

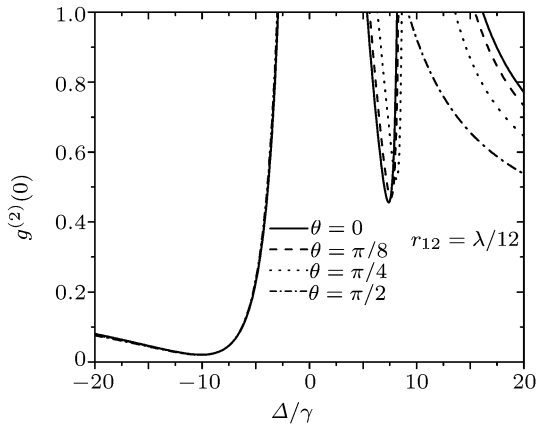


Fig. 6. The intensity–intensity correlation function $g^{(2)}(0)$ versus the laser–atom detuning with $\Omega/\gamma = 2$.

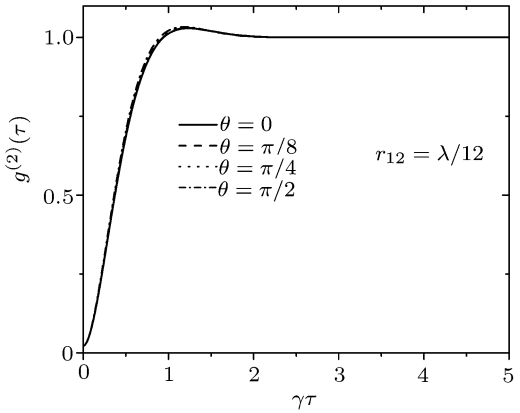


Fig. 7. Dependence of the intensity–intensity correlation function $g^{(2)}(\tau)$ on the spatial angle factor θ for the interatomic separation $r_{12} = \lambda/12$ with $\Omega/\gamma = 2$ and $\Delta/\gamma = -\Omega_{12}$.

Therefore, the interference effects due to the spatial phase variation of the laser apparently manifest in the intensity–intensity correlation function of the resonant fluorescence, especially in the case of the high laser intensity. This suggests that the photon correlations can serve as a probe for such a kind of spatial interference effect, which is obviously advantageous over the detection scheme of studying the fluorescent spectrum since the effect of spatial interference on the fluorescent intensity is almost insignificant in the case of the strong pumping.[14] In Fig. 6 we plot the dependence of the intensity–intensity correlation function $g^{(2)}(0)$ on the laser–atom detuning. From it, one can see that when the detuning $-\Delta$ is in the vicinity of the interatomic strength $\Omega_{12}(\approx 9.3)$, the correlation function $g^{(2)}(0) \approx 0$, indicating the maximum antibunching effect. Furthermore, from Figs. 6 and 7 the spatial variation of the phase angle θ has little effect on the intensity–intensity correlations when the laser–atom detuning $(-\Delta)$ is turned in the vicinity of the interaction strength Ω_{12}. This is because for $\Delta = -\Omega_{12}/\gamma$, the fluorescent light results from the transition from the symmetric state $|s\rangle$ to the ground state $|g\rangle$, which radiates in all directions.[9] Thus, the photon correlations of the resonant fluorescence are little influenced by the variation of the angle θ.

4. Conclusions

In conclusion, we investigate the effects of the spatial interference arising from the spatial phase variation of the laser on the intensity–intensity correlations of the resonant fluorescence emitted by a cooperative two-atom system. We have shown that the photon correlations can exhibit quite different features for different values of the spatial phase factor θ of the laser, especially in the case of the strong pumping laser. This suggests that the intensity–intensity correlations can serve as a probe of such spatial interference effects as are due to the spatial phase variation.

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