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The revival–collapse phenomenon in the higher-order fluctuations of quadrature field components of the multiphoton Jaynes–Cummings model

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Abstract

For a system of multiphoton Jaynes–Cummings model, which represents the interaction of a two-level atom with the radiation field, we study the connection between the atomic inversion, in particular the occurrence of revival–collapse phenomenon (RCP), and the higher-order fluctuation factors. We assume that the atom and the field are initially prepared in the excited and the k-photon coherent states, respectively. We show that there is a class of states for which the higher-order fluctuation factor can provide RCP similar to that involved in the corresponding atomic inversion. Moreover, for initial coherent light we prove that the higher-order fluctuation factor of the three-photon transition can provide RCP similar to that occurring in the atomic inversion of the one-photon transition. As an example we discuss the fourth-order fluctuation in detail.

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In the previous paper [1] for the Jaynes–Cummings model (JCM) we have discussed the possible occurrence of the revival–collapse phenomenon (RCP) of the atomic inversion $\langle\hat{\sigma}_z(T)\rangle$ in the evolution of the quadrature fluctuation (i.e. squeezing), particularly for the normal and the amplitude-squared fluctuations. We have shown that there are two approaches, namely, the natural-phenomenon and numerical-simulation approaches. In the natural-phenomenon approach the fluctuation factors naturally include information on the corresponding atomic inversion. Nevertheless, for the numerical-simulation approach we have deduced the rescaled fluctuation factor for the three-photon JCM, which can exhibit the RCP similar to that involved in the evolution of the atomic inversion of the standard JCM (i.e., single-photon JCM with initial coherent light). We should stress that the term 'natural-phenomenon approach' is given to show that the fluctuation factors of the system provide information on the atomic inversion of the same system; however, the term 'numerical-simulation approach' demonstrates that the fluctuation factors of another system, i.e. three-photon JCM, give information on the atomic inversion of the standard JCM. Additionally, the rescaled fluctuation factor of the latter approach is obtained through very complicated numerical procedures. In this report we complete the results given in [1] by discussing such phenomenon for the higher-order fluctuation and investigating in detail the occurrence of the RCP in the evolution of the fourth-order fluctuation as an example. Like one that exhibits normal squeezing, a field that is squeezed to a higher order is a pure quantum mechanical light and has no classical description. Moreover, the higher-order squeezing for the JCM has been investigated, e.g. in [2, 3]. The results given in [1] and in the present paper draw the attention for the first time to an important fact: the evolution of the $\langle\hat{\sigma}_z(T)\rangle$ can be measured by a homodyne detector [4], in which the signal coming from the microwave cavity is optically mixed with a strong coherent local oscillator via a 50:50 beam splitter and the emerging fields are detected by photodetector [5]. Moreover, the progress in both trapped ions [6] and micromaser [7] is promising to produce such a type of phenomena. Additionally, the higher-order fluctuation is motivated by the development in the higher-order correlation measurement techniques, which can efficiently extract information from a quantum system [8]. Also the motivation is a direct consequence of the interest which has been given to the higher-order fluctuation in the literature, e.g. [9].

The Hamiltonian of the multiphoton JCM in the rotating wave approximation [1, 10, 11] is

$$\frac{\hat{H}}{\hbar} = \omega_0\hat{a}^\dagger\hat{a} + \omega_a\hat{\sigma}_z + \lambda(\hat{a}^m\hat{\sigma}_+ + \hat{a}^{\dagger m}\hat{\sigma}_-), \tag{1}$$

where $\hat{\sigma}_\pm$ and $\hat{\sigma}_z$ are the Pauli spin operators; ω_0 and ω_a are the frequencies of the cavity mode and the atomic transition, respectively; λ is the atom–field coupling constant; m is the transition parameter and $\hbar$ is the Dirac constant. According to the lines given in [12], the use of the Hamiltonian (1) is called an effective Hamiltonian approach. In this approach the probability amplitudes of the dynamical wavefunction of the system include trivial (or intensity independent) overall phase, which has no effect on the evolution of the fluctuation factors. In contrast the modified effective Hamiltonian approach, for the phenomena under consideration, provides similar results to those in [1], i.e. for natural phenomenon the fluctuation factors give typical information on the atomic inversion, while for the numerical-simulation approach they fail. Also as in [1] the values of atomic-relative phases act only in the natural approach. Thus in this report we use the effective Hamiltonian approach and neglect the role of the atomic-relative phases. We assume that the atom is initially prepared in the excited states $|+\rangle$ and the field is in the k-photon coherent states [13, 14] having the form

$$|\psi(0)\rangle = \sum_{n=0}^{\infty} C_n|kn\rangle, \qquad C_n = \frac{\alpha^n}{\sqrt{n!}}\exp\left(-\frac{1}{2}\alpha^2\right), \tag{2}$$

where k is a parameter whose value will be specified in the text. For these initial conditions the interaction dynamical wavefunction can be evaluated [1, 15] as

$$|\Psi(T)\rangle = \sum_{n=0}^{\infty} C_n\left[\cos\left(T\sqrt{\frac{(kn+m)!}{(kn)!}}\right)|+, kn\rangle - \mathrm{i}\sin\left(T\sqrt{\frac{(kn+m)!}{(kn)!}}\right)|-, kn+m\rangle\right], \tag{3}$$

where $T = \lambda t$ and $|-\rangle$ denotes the atomic ground state. In (3) we have considered the resonance case $m\omega_0 = \omega_a$ and α to be real.

The atomic inversion associated with (3) is

$$\langle\hat{\sigma}_z(T)\rangle = \sum_{n}^{\infty} P(n)\cos\left(2T\sqrt{\frac{(kn+m)!}{(kn)!}}\right), \tag{4}$$

where $P(n) = C_n^2$. Additionally, we deduce the asymptotic form for (4) when $m = 1$ in the framework of strong-intensity regime (SR), i.e. $\alpha^2 \gg 1$. It should be borne in mind that $P(n)$

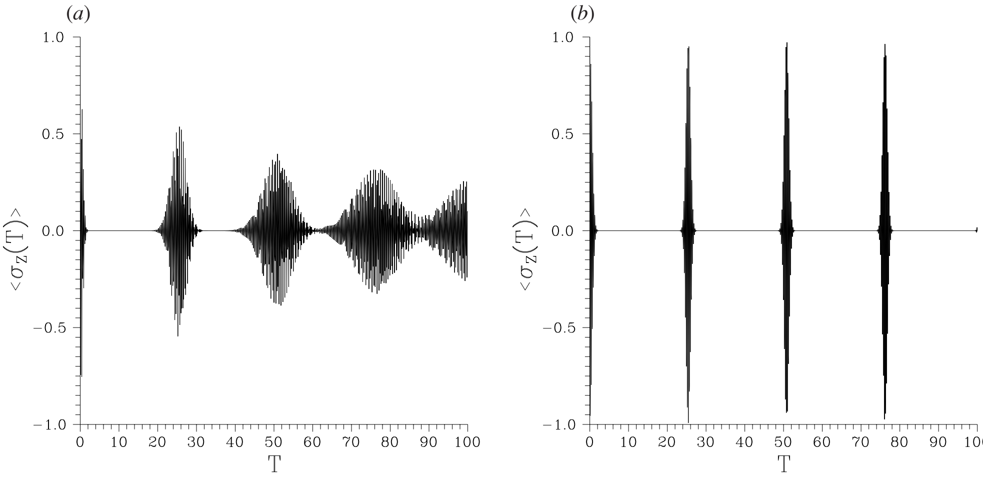

Figure 1. The atomic inversion $\langle\hat{\sigma}_z(T)\rangle$ for the exact (*a*) and the asymptotic (*b*) forms for $m = 1$ against the scaled time T when the optical cavity mode and atom are initially prepared in the three-photon coherent (with $|\alpha| = 7$) and excited states, respectively.

has a Poissonian distribution and consequently in the SR the terms contributing effectively in the summation (4) are those for which $n \simeq \alpha^2 = \bar{n}$, i.e. the harmonic approximation [16]. The argument of the cosine in (4) can be asymptotically expressed as

$$\sqrt{kn+\epsilon_1} = \sqrt{k\bar{n}}\left[1+\frac{n+(\epsilon_1/k)-\bar{n}}{\bar{n}}\right]^{\frac{1}{2}} \simeq \frac{1}{2}(\eta_1+\eta_2 n), \tag{5}$$

where

$$\eta_1 = \sqrt{k\bar{n}}+\frac{\epsilon_1}{\sqrt{k\bar{n}}}, \qquad \eta_2 = \sqrt{\frac{k}{\bar{n}}} \tag{6}$$

and ϵ_1 is a finite c-number, for the present case $\epsilon_1 = 1$. On substituting (5) into (4) and after some minor algebra we obtain the asymptotic form as

$$\langle\hat{\sigma}_z(T)\rangle \simeq \exp\left[-2\alpha^2\sin^2\left(\frac{\eta_2 T}{2}\right)\right]\cos[T\eta_1+\alpha^2\sin(T\eta_2)]. \tag{7}$$

It is evident that the exponential part in (7), which is responsible for the occurrence of the revival patterns, is dependent on the parameter k. More illustratively, the revivals occur around $\eta_2 T = 2m'\pi$, i.e. $T = \frac{2\pi m'\sqrt{\bar{n}}}{\sqrt{k}}$ where m' is a positive integer. This indicates that when k increases the series of the revival patterns increases, too. This is in a good agreement with that occurring for the superposition states. In figures 1(*a*) and (*b*) we have plotted the atomic inversion for the exact (4) and asymptotic (7) forms, respectively, for given values of the interaction parameters. A comparison between these figures shows that there is agreement and disagreement in their behaviour. Qualitatively, they are approximately in agreement in the locations of the revival–collapse patterns in the interaction-time ‘domain‘, but the corresponding revivals cannot provide typical amplitude and shape, i.e. for the exact form the revivals are broader than those for the asymptotic one. This problem can be solved by including higher-order approximation in (6), however, the price that should be paid, we cannot obtain a simple closed form for $\langle\hat{\sigma}_z(T)\rangle$. For convenience we compare figure 1 with those given in the literature for initial coherent light, e.g. [17–19]. Generally, for the exact form one can see that the two types of behaviour (i.e. in the literatures and here) are similar, i.e. an initial collapse of these oscillations is followed by regular revivals that slowly become broader and eventually overlap. The basic difference between figure 1(*a*) and those given for initial coherent light [17–19] is that the revival patterns occur for interaction time much shorter than that of coherent light. Actually, we have found that $T_r^{(k)} = \frac{T_r^{(c)}}{\sqrt{k}}$, where $T_r^{(k)}$ and $T_r^{(c)}$ denote the revival times associated with the initial k-photon coherent state and coherent state, respectively.

Now we draw attention to the higher-order fluctuation using the definition given in [20], which is based on the two quadratures $\hat{X}_N$ and $\hat{Y}_N$ having the forms

$$\hat{X}_N = \frac{1}{2}(\hat{a}^N + \hat{a}^{\dagger N}), \qquad \hat{Y}_N = \frac{1}{2\mathrm{i}}(\hat{a}^N - \hat{a}^{\dagger N}), \tag{8}$$

where N is a positive integer. The fluctuation factors associated with (8) can be expressed as [20]

$$\begin{aligned} F_N(T) &= \langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle + \mathrm{Re}\langle \hat{a}^{2N}(T)\rangle - 2(\mathrm{Re}\langle \hat{a}^N(T)\rangle)^2, \\ S_N(T) &= \langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle - \mathrm{Re}\langle \hat{a}^{2N}(T)\rangle - 2(\mathrm{Im}\langle \hat{a}^N(T)\rangle)^2. \end{aligned} \tag{9}$$

As we mentioned above the attention in this report is focused on the occurrence of the RCP in the evolution of $F_N(T)$ and $S_N(T)$. This will be investigated in the following two parts for natural phenomenon and numerical-simulation approaches.

(i) Natural phenomenon. This approach is based on the fact that $\omega_0\hat{a}^\dagger\hat{a} + \omega_a\hat{\sigma}_z$ is a constant of motion and consequently $\langle \hat{a}^\dagger(T)\hat{a}(T)\rangle$ and $\langle\hat{\sigma}_z(T)\rangle$ provide information on each other. Nevertheless, for the higher-order moments, i.e. $N > 1$, the operator $\hat{a}^{\dagger N}\hat{a}^N + \hat{\sigma}_z$ is not a constant of motion even though, in the SR, $\langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle$ can provide information on the corresponding $\langle\hat{\sigma}_z(T)\rangle$. We show this as follows. From (9), $F_N(T)$ and $S_N(T)$ provide RCP when there are initial field states for which $\langle \hat{a}^{s'}(T)\rangle = 0$ for all values of s'. For instance, state (2) can achieve this when $2N/k = l$ and $N/k = l'$ where l and l' are fractions. More illustratively, for the cubic- ($N = 3$) and fourth-order ($N = 4$) fluctuations this occurs when the parameter $k = 4$ (four-photon coherent state) and 3 (three-photon coherent state), respectively. Actually, three-photon and four-photon coherent states can be generated by the conditional measurement [21] and dispersive cavity QED [22]. We deduce the general form which connects $\langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle$ with the corresponding $\langle\hat{\sigma}_z(T)\rangle$ for the state (3) when $m = 1$. From the above discussion, expressions (9) reduce to

$$\begin{aligned} F_N(T) = S_N(T) &= \langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle \\ &= \sum_{n=0}^{\infty} C_n^2 \left[\frac{(kn)!}{(kn-N)!}\cos^2(T\sqrt{kn+1}) + \frac{(kn+1)!}{(kn-N+1)!}\sin^2(T\sqrt{kn+1})\right] \\ &= \langle \hat{a}^{\dagger N}(0)\hat{a}^N(0)\rangle + \frac{N}{2}\langle \hat{a}^{\dagger N-1}(0)\hat{a}^{N-1}(0)\rangle \\ &\quad - \frac{N}{2}\sum_{n=0}^{\infty} C_n^2 \frac{(kn)!}{(kn-N+1)!}\cos(2T\sqrt{kn+1}). \end{aligned} \tag{10}$$

In the SR and finite N the argument $\sqrt{kn+\epsilon_2}$, where ϵ_2 is a finite and arbitrary c-number, can be replaced by $\sqrt{kn+1}$ and the summation in the third line of (10) can be evaluated in terms of $\langle\hat{\sigma}_z(T)\rangle$, and we obtain

$$\langle\hat{\sigma}_z(T)\rangle = \frac{\langle \hat{a}^{\dagger N}(0)\hat{a}^N(0)\rangle + \frac{N}{2}\langle \hat{a}^{\dagger N-1}(0)\hat{a}^{N-1}(0)\rangle - S_N(T)}{\frac{N}{2}\langle \hat{a}^{\dagger N-1}(0)\hat{a}^{N-1}(0)\rangle}. \tag{11}$$

Formula (11) can be deduced by induction. As an example we evaluate the fourth-order fluctuation when the field is initially in the three-photon coherent state for $m = 1$. The fourth-order squeezing is already studied for the multiphoton JCM in [2]. For $k = 3$, $m = 1$, $N = 4$, (3) and (9) give

$$
\begin{aligned}
F_4(T) = S_4(T) &= \langle \hat{a}^{\dagger 4}(T)\hat{a}^4(T)\rangle \\
&= \langle \hat{a}^{\dagger 4}(0)\hat{a}^4(0)\rangle + 2\langle \hat{a}^{\dagger 3}(0)\hat{a}^3(0)\rangle - 6\alpha^2 \sum_{n=0}^{\infty} P(n)[9\alpha^4 \cos(2T\sqrt{3n+10}) \\
&\quad + 18\alpha^2 \cos(2T\sqrt{3n+7}) + 2\cos(2T\sqrt{3n+4})].
\end{aligned} \tag{12}
$$

Using the discussion given below (10) expression (12) can be expressed as

$$
F_4(T) = S_4(T) = \langle \hat{a}^{\dagger 4}(0)\hat{a}^4(0)\rangle + 2\langle \hat{a}^{\dagger 3}(0)\hat{a}^3(0)\rangle - 2\langle \hat{a}^{\dagger 3}(0)\hat{a}^3(0)\rangle\langle \hat{\sigma}_z(T)\rangle. \tag{13}
$$

Also we have numerically checked (11) for the case expressed by (13) using as values of the interaction parameters those given for figure 1(a). We obtained a typical shape as in figure 1(a) and this verifies the above discussion. We conclude this part by referring to a misprint in (38) in [1] where the term $3\langle \hat{n}(0)\rangle$ should be added to the numerator.

(ii) *Numerical simulation*. In this part we prove that the RCP of the atomic inversion of the standard JCM can be observed in the evolution of the higher-order (i.e. Nth-power) fluctuation of the multiphoton JCM. From (3) one can easily find that $\mathrm{Re}\langle \hat{a}^N(T)\rangle \neq 0$ and $\mathrm{Im}\langle \hat{a}^N(T)\rangle = 0$, where α is real. Therefore, $F_N(T)$ (cf (9)) cannot provide RCP where it includes a squared quantity, which destroys the RCP. This means that RCP almost occurs in the evolution of the $S_N(T)$. Additionally, when $m > 2$ the evolution of $\langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle$ is chaotic and in the SR one can asymptotically prove that $\langle \hat{a}^{\dagger N}(T)\hat{a}^N(T)\rangle \simeq \bar{n}^N$ (see the argument below (14)). Thus if $S_N(T)$ exhibits RCP this would be connected to the evolution of $\mathrm{Re}\langle \hat{a}^{2N}(T)\rangle$, which for $k = 1$ takes the form

$$
\begin{aligned}
\langle \hat{a}^{2N}(T)\rangle = \bar{n}^N \sum_{n=0}^{\infty} P(n) \Bigg\{ &\cos\left(T\sqrt{\frac{(n+m+2N)!}{(n+2N)!}}\right)\cos\left(T\sqrt{\frac{(n+m)!}{n!}}\right) \\
&+ \sqrt{\prod_{j=1}^{2N} \frac{\left(1+\frac{m+j}{n}\right)}{\left(1+\frac{j}{n}\right)}}\sin\left(T\sqrt{\frac{(n+m+2N)!}{(n+2N)!}}\right)\sin\left(T\sqrt{\frac{(n+m)!}{n!}}\right)\Bigg\}.
\end{aligned} \tag{14}
$$

In the framework of both the SR and harmonic approximation the quantity under the square root in the second line of (14) reduces to unity and we obtain

$$
\langle \hat{a}^{2N}(T)\rangle \simeq \bar{n}^N \sum_{n=0}^{\infty} P(n)\cos\left[T\left(\sqrt{\frac{(n+m+2N)!}{(n+2N)!}} - \sqrt{\frac{(n+m)!}{n!}}\right)\right]. \tag{15}
$$

Regardless of the constant prefactor $\bar{n}^N$, the comparison between expression (4) when $(k, m) = (1, 1)$ and (15) leads to the following fact: the fluctuation factor $S_N(T)$ can provide information on $\langle \hat{\sigma}_z(T)\rangle$ of the standard JCM provided that the arguments of the cosines in the two expressions are comparable. Thus the object is to find the proportionality factor μ_N, say, and consequently the value of m, which can achieve this goal. This can be realized by evaluating the asymptotic form for the following quantity:

$$
\mu_N = \frac{\sqrt{\frac{(n+m+2N)!}{(n+2N)!}} - \sqrt{\frac{(n+m)!}{n!}}}{2\sqrt{n+1}}. \tag{16}
$$

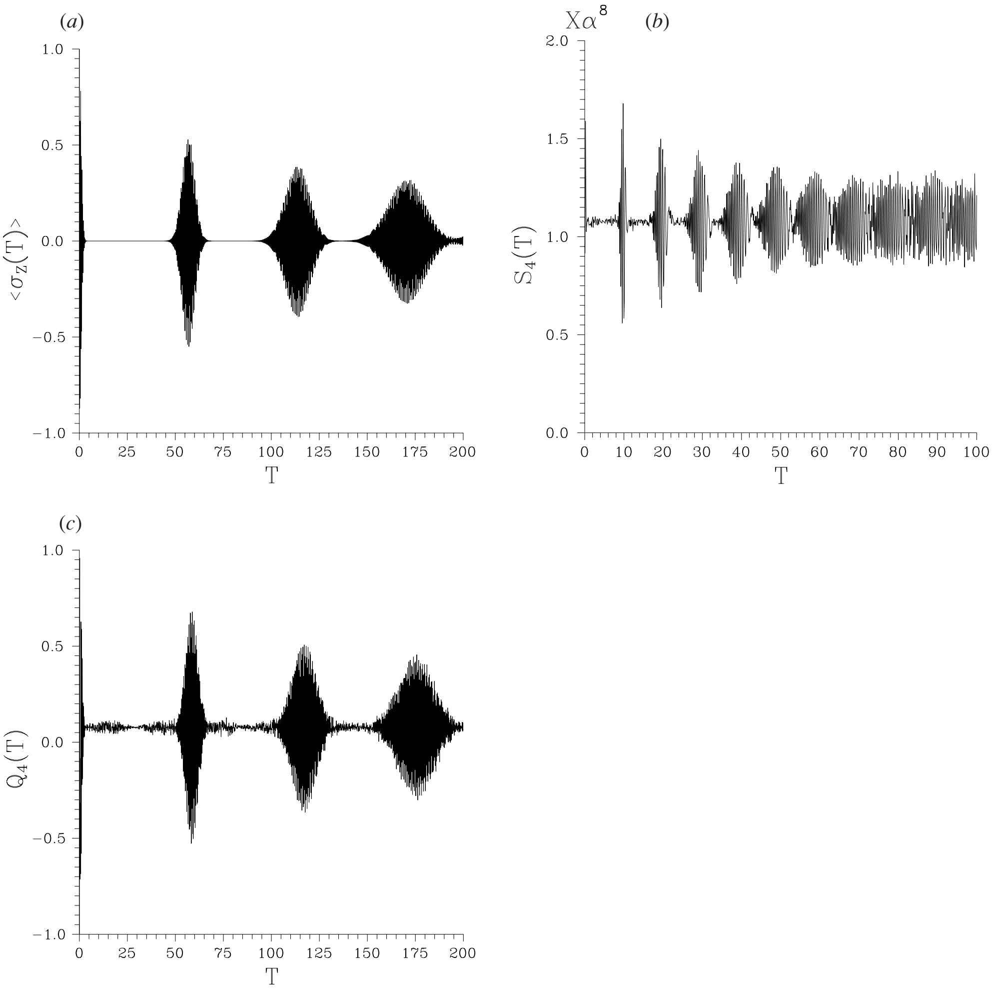

Figure 2. The atomic inversion $\langle\hat{\sigma}_z(T)\rangle$ (*a*), the fluctuation factor $S_4(T)$ (*b*) and the rescaled fluctuation factor $Q_4(T)$ (*c*) plotted against the scaled time T when the optical cavity mode and the atom are initially prepared in the coherent and excited atomic states, respectively, with $(m, |\alpha|) = (3, 9)$.

Expression (16) can be re-expressed as

$$\mu_N = \frac{1}{2\sqrt{(n+1)\prod_{j=1}^{2N}(n+m+j)}}\sqrt{\frac{(n+m)!}{n!}}\left[\frac{\prod_{j=1}^{2N}(n+m+j) - \prod_{j=1}^{2N}(n+j)}{\sqrt{\prod_{j=1}^{2N}(n+m+j)} + \sqrt{\prod_{j=1}^{2N}(n+j)}}\right]$$
$$= \frac{1}{2\sqrt{(n+1)\prod_{j=1}^{2N}(n+m+j)}}\sqrt{\frac{(n+m)!}{n!}}$$
$$\times\left[\frac{2Nmn^{2N-1} + n^{2N-2}(\cdots) + n^{2N-3}(\cdots) + \cdots + n^0(\cdots)}{\sqrt{\prod_{j=1}^{2N}(n+m+j)} + \sqrt{\prod_{j=1}^{2N}(n+j)}}\right]. \qquad (17)$$

It is worth recalling that we are working in the SR. Therefore, for finite values of N we can use $(n+\varepsilon) \to n = \bar{n}$, where ε is a finite c-number. In this case (17) can be modified to

$$\mu_N \simeq \frac{\bar{n}^m}{4\bar{n}^{2N+1}} \sqrt{\prod_{j=0}^{m} \left(1 + \frac{m-j}{\bar{n}}\right)} [2Nm\bar{n}^{2N-1} + \bar{n}^{2N-2}(\cdots) + \bar{n}^{2N-3}(\cdots) + \cdots + \bar{n}^0(\cdots)]$$
$$\simeq \frac{1}{4}\left[2Nm\bar{n}^{\frac{m-3}{2}} + \bar{n}^{\frac{m-5}{2}}(\cdots) + \bar{n}^{\frac{m-7}{2}}(\cdots) + \cdots + \bar{n}^{\frac{m-4N-2}{2}}(\cdots)\right]. \tag{18}$$

In the SR only the first term in (18) can survive provided that $m = 3$, i.e. three-photon JCM. In this case the proportionality factor μ_N reduces to

$$\mu_N = \frac{3N}{2}. \tag{19}$$

From the above discussion and (19) we can write the general form for the rescaled higher-order fluctuation factor, which can give information on $\langle\hat{\sigma}_z(T)\rangle$ of the standard JCM, as

$$Q_N(T) = \frac{S_N\left(\frac{2T}{3N}\right) - \langle\hat{n}(0)\rangle^N}{\langle\hat{n}(0)\rangle^N}. \tag{20}$$

One can easily check that the rescaled factors for the normal and amplitude-squared fluctuation in [1] are particular cases of (20). Finally, in figures 2(a)–(c) we have plotted $\langle\hat{\sigma}_z(T)\rangle$ of the standard JCM, and $S_4(T)$ and $Q_4(T)$ of the three-photon JCM, respectively, for given values of the interaction parameters. Figure 2(b) shows that $S_4(T)$ exhibits a series of revival patterns, which are much greater than those for $\langle\hat{\sigma}_z(T)\rangle$ of the standard JCM (see figure 2(a)). Nevertheless, a comparison between figures 2(a) and (c) leads to that $Q_4(T)$ provides RCP quite similar to that of $\langle\hat{\sigma}_z(T)\rangle$. This concludes our demonstration, i.e. $Q_N(T)|_{m=3} \equiv \langle\hat{\sigma}_z(T)\rangle_{m=1}$ for the same initial conditions. We should stress that when the squeezing-order N increases the initial mean-photon number has to be increased to obtain better information from (20) on $\langle\hat{\sigma}_z(T)\rangle_{m=1}$. Details about this issue will be published elsewhere. The comparison between figures 2(a) and (c) with those given in the literature to the atomic inversion of the standard JCM with initial coherence, e.g. [17–19], shows that they are quite similar. Moreover, in [18] for the JCM it has been shown that the evolution of both the phase variance and the phase distribution can carry certain information on the RCP of the corresponding atomic inversion. These facts and the results given here and in [1] indicate that the evolution of the fluctuation factors for the natural and numerical approaches can partially give information on the phase distribution of the system, and vice versa.

In conclusion we have deduced the general forms of the rescaled fluctuation factors for both natural phenomenon and numerical-simulation approaches. These results with those given in [1] complete the information about the occurrence of the RCP in evolution of the fluctuation factors.

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