Optical Fock-state generation with large number of photons based on atoms coupled to an optical parametric oscillator

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An optical field with a definite number of photons is very important for quantum metrology and quantum information. Some theoretical protocols for creating such a Fock-state have been proposed, but it is still a big challenge to produce it with a large number photons experimentally. We revisit the system of atoms inside an optical parametric oscillator that was proposed in 1990s, and it is found that for the atom ensemble, the optical Fock-state with an arbitrary number of photons can be generated. Compared to the previous proposals, the scheme presented here is simple and seems physically realizable. The system also provides the possibility to demonstrate the strong interaction between nonclassical light and atoms in a confined space.
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1. INTRODUCTION

Fock states play an important role in quantum optics due to their fundamental nonclassical nature, such as antibunching and negativity of the Wigner function, as well as their potential practical applications in quantum metrology [1] and quantum information science [2], such as quantum cryptography [3] and quantum computing [4]. They can also improve the sensitivity of an interferometer to the Heisenberg limit [5].

However, until now it was experimentally challenging to produce such light fields that contain definite number of photons. Several proposals aimed at creating a Fock-state "on demand" have been offered theoretically. Brown *et al.* [6] presented a scheme for the deterministic production of N-photon Fock states from coupling N three-level atoms in a high-finesse optical cavity to an external field, and it is suitable for experimental implementation by using a cloud of cold atoms trapped in a cavity. O'Sullivan *et al.* [7] considered the heralded N-photon states created from the photons produced by an unseeded optical parametric amplifier using time-multiplexed detectors. An experimentally viable approach for preparing arbitrary photon-number states of a cavity mode using continuous measurement and real-time quantum feedback is presented by Geremia [8]. Following this method, Haroche and Raimond have proposed a real-time stabilizing quantum feedback scheme for the preparation and protection of photon-number states of light trapped in a high-Q microwave cavity [9,10] based on their experimental techniques.

For one and two photon states, they have been experimentally demonstrated by using other processes including atomic and molecular fluorescence [11,12], Coulomb blockade for electrons [13], and micromaser [14]. Waks *et al.* [15] have demonstrated Fock-state generation with 1, 2, 3, and 4 photons, by using a visible light photon counter (VLPC) along with the process of parametric down-conversion. Achilles *et al.* [16] got the heralded two-photon-number state by conditionally detecting a two-mode squeezed state generated by parametric down-conversion. Yet producing and detecting a Fock-state containing large number of photons is still a very large challenge. So it is of great importance to explore a feasible and simpler approach to generate Fock states with a large number of photons.

In this paper, we revisit the combined "atom + DOPO" system. DOPO is the degenerate optical parametric oscillator. The system was first provided and discussed by Xiao and Jin *et al.* [17], and was discussed extensively from then on for a large number of atoms [17–19] and also for single atoms [20–23]. We discuss the photon statistics, which is one of the main features of the light field in the case of a large number of atoms. It shows that, when there are many atoms in the cavity, very weak antibunching but strong sub-Poissonian (Mandel Q goes to -1) can be obtained, which shows the attainability of generation of Fock states with arbitrary photon numbers. For comparison, both antibunching and sub-Poissonian are weak for the normal bistability (OAB) system, and it does not show the ability of generating Fock states with a large number of photons. The numerical simulations are explained quite clearly by the consequence of the bistability and concrete parameters of the system. Consider the present state-of-the-art of cold atom and optical parametric process; the system is quite feasible in experiment.

2. BASIC MODEL AND THE CORRELATION FUNCTIONS

Let us consider a combined system consisting of an optical cavity with a pair of nonlinear crystals and the two-level atoms

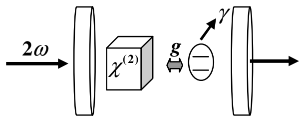

Fig. 1. A schematic of the physical system under consideration.

(see Fig. 1) [17]. The nonlinear medium is prepared to fulfill the phase-matching condition for parametric down-conversion. The system is pumped by a strong external coherent field with the frequency nearly resonant at twice the atomic transition frequency. The total Hamiltonian describing this system is given by [18]

$$
\begin{aligned}
\hat{H} &= \sum_{i=1}^{4} \hat{H}_i, \\
\hat{H}_1 &= \hbar\omega_c \hat{a}^+\hat{a} + \frac{1}{2}\hbar\omega_a \sum_\mu \hat{\sigma}_\mu^z + \hbar\omega_2 \hat{a}_2^+\hat{a}_2, \\
\hat{H}_2 &= i\hbar\sum_\mu g(\hat{a}^+\hat{\sigma}_\mu^- e^{-i\phi} - \hat{a}\hat{\sigma}_\mu^+ e^{i\phi}) + i\hbar\frac{\kappa}{2}(\hat{a}^{+2}\hat{a}_2 - \hat{a}^2\hat{a}_2^+), \\
\hat{H}_3 &= i\hbar(\varepsilon_2\hat{a}_2^+ e^{-i\omega_p t} - \varepsilon_2\hat{a}_2 e^{i\omega_p t}), \\
\hat{H}_4 &= \sum_\mu \hat{\sigma}_\mu^z(\hat{\Gamma}_a^-\hat{\sigma}_\mu^+ + \hat{\Gamma}_a^+\hat{\sigma}_\mu^-) + \hat{a}_2\hat{B}_2^+ + \hat{a}_2^+\hat{B}_2 + \hat{a}\hat{B}_1^+ + \hat{a}^+\hat{B}_1,
\end{aligned} \tag{1}
$$

where $\hat{H}_1$ is the free energy of the system, $\hat{H}_2$ describes the atom-field and parametric interactions in the system, $\hat{H}_3$ is the external driving term, and $\hat{H}_4$ gives decay processes of the atoms and the intracavity field. $(\hat{a}_2^+, \hat{a}_2)$ denote the creation and annihilation operators of the fundamental field mode and $(\hat{a}^+, \hat{a})$ are for the subharmonic field mode. $\hat{\sigma}^\pm$, $\hat{\sigma}^z$ are the Pauli atomic operators. ω_a is the atomic transition frequency, and ω_c is the frequency of the cavity subharmonic field. g is the coupling constant between the atoms and the intracavity subharmonic field. κ is the coupling coefficient of the nonlinear down-conversion process. $\hat{\Gamma}_a^\pm$, $(\hat{B}_2^+, \hat{B}_2)$, $(\hat{B}_1^+, \hat{B}_1)$ are the noise operators of the atoms, the fundamental field, and the down-conversion field, respectively. ε_2 stands for the complex classical amplitude of the external driving field.

The Fokker–Planck equation can be obtained from this Hamiltonian by the standard technique [24]. In deriving the equation, all the operators have been translated into the corresponding c numbers defined by

$$
(\hat{a}, \hat{a}^+, \hat{a}_2, \hat{a}_2^+, \hat{J}^-, \hat{J}^+, \hat{J}^z) \to (\alpha, \alpha^+, \alpha_2, \alpha_2^+, J^-, J^+, J^z),
$$

where

$$
\hat{J}^z \equiv \sum_\mu \hat{\sigma}_\mu^z, \qquad \hat{J}^\pm = \sum_\mu e^{\pm i\phi}\hat{\sigma}_\mu^\pm, \tag{2}
$$

are the collective operators of atoms. From the Fokker–Planck equation, one can get a set of equations of motion for the variables as follows:

$$
\begin{aligned}
\frac{d\alpha}{dt} &= -\gamma_1\alpha + gJ^- + \kappa\alpha^+\alpha_2 + \Gamma_\alpha(t), \\
\frac{d\alpha^+}{dt} &= -\gamma_1\alpha^+ + gJ^+ + \kappa\alpha\alpha_2^+ + \Gamma_{\alpha^+}(t), \\
\frac{d\alpha_2}{dt} &= -\gamma_2\alpha_2 - \frac{\kappa}{2}\alpha^2 + \varepsilon_2, \\
\frac{d\alpha_2^+}{dt} &= -\gamma_2\alpha_2^+ - \frac{\kappa}{2}\alpha^{+2} + \varepsilon_2^*, \\
\frac{dJ^-}{dt} &= -(\gamma_\perp + i\Delta_a)J^- + g\alpha J^z + \Gamma_-(t), \\
\frac{dJ^-}{dt} &= -(\gamma_\perp + i\Delta_a)J^+ + g\alpha^+ J^z + \Gamma_+(t), \\
\frac{dJ^z}{dt} &= -2\gamma_\perp(N + J^z) - 2g(\alpha^+J^- + \alpha J^+) + \Gamma_z(t),
\end{aligned} \tag{3}
$$

and the correlation terms

$$
\begin{aligned}
\langle\Gamma_\alpha(t)\Gamma_\alpha(t')\rangle &= \kappa\alpha_2\delta(t - t'), \\
\langle\Gamma_{\alpha^+}(t)\Gamma_{\alpha^+}(t')\rangle &= \kappa\alpha_2^+\delta(t - t'), \\
\langle\Gamma_-(t)\Gamma_-(t')\rangle &= 2g\alpha J^-\delta(t - t'), \\
\langle\Gamma_+(t)\Gamma_+(t')\rangle &= 2g\alpha^+J^+\delta(t - t'), \\
\langle\Gamma_z(t)\Gamma_z(t')\rangle &= [4\gamma_\perp(N + J^z) - 4g(\alpha^+J^- + \alpha J^+)]\delta(t - t'),
\end{aligned} \tag{4}
$$

where $\Delta_c = \omega_c - \omega_p/2$, $\Delta_2 = \omega_2 - \omega_p$, $\Delta_a = \omega_a - \omega_p/2$. Without loss of generality, we consider a double-resonance cavity for the fundamental and subharmonic fields such that $\Delta_c = 0$, $\Delta_2 = 0$, and $\Delta_a = \omega_c - \omega_a$, and a pure radiative decay process (no collisional decay for atoms); i.e., $\gamma_\perp = \gamma_\parallel/2$. γ_1 and γ_2 are the cavity-decay rate of the subharmonic and pumping field mode, respectively. N is the total number of atoms in the cavity. Several normalized parameters are introduced as follows:

$$
\begin{aligned}
&\Delta \equiv \frac{\Delta_a}{\gamma_\perp}, \qquad C \equiv \frac{Ng^2}{2\gamma_\perp\gamma_1}, \qquad N_s \equiv \frac{\gamma_\perp\gamma_\parallel}{4g^2}, \\
&n_0 \equiv \frac{2\gamma_1\gamma_2}{\kappa^2}, \qquad |\varepsilon_0| \equiv \frac{\gamma_1\gamma_2}{\kappa},
\end{aligned} \tag{5}
$$

where C is the atomic-cooperative parameter, which is defined as the ratio of atomic absorption to the cavity-decay rate of the subharmonic mode. N_s is the single-atom saturation photon number. n_0 is the scaling photon number (saturation photon number) for the DOPO system. ε_0 is the threshold of the pumping amplitude for the DOPO system when the atoms are absent. The normalized field variables are

$$
x = \frac{\alpha}{\sqrt{N_s}}, \qquad y = \frac{\varepsilon_2}{\varepsilon_0}. \tag{6}
$$

Together with the replacements [25]

$$
\tau = \gamma_\perp t, \qquad J^- = \frac{N}{\sqrt{2}}\upsilon, \qquad J^+ = \frac{N}{\sqrt{2}}\upsilon^*, \qquad J^z = \frac{N}{2}m, \tag{7}
$$

we can get the equivalent equations of Eq. (3),

$$\begin{aligned}\frac{dx}{d\tau} &= -\mu x + \mu(y - r\bar{x}^* \cdot \bar{x})x^* + 2\mu C v,\\ \frac{dx^*}{d\tau} &= -\mu x^* + \mu(y - r\bar{x}^* \cdot \bar{x})x + 2\mu C v^*,\\ \frac{dv}{d\tau} &= -(1 + i\Delta)v + \frac{1}{2}xm,\\ \frac{dv^*}{d\tau} &= -(1 + i\Delta)v^* + \frac{1}{2}x^* m,\\ \frac{dm}{d\tau} &= -2m - 4 - 2(x^* v + x v^*).\end{aligned} \tag{8}$$

In order to investigate the photon statistics of the intracavity field, we have adiabatically eliminated the operators of the pumping field. From the steady-state solution, we obtain the following relations:

$$\begin{aligned}&\bar{v} = -\frac{(1 - i\Delta)\bar{x}}{1 + \Delta^2 + \bar{x} \cdot \bar{x}^*}, \qquad \bar{v}^* = -\frac{(1 + i\Delta)\bar{x}^*}{1 + \Delta^2 + \bar{x} \cdot \bar{x}^*},\\ &\bar{m} = \frac{-2(1 + \Delta^2)}{1 + \Delta^2 + \bar{x} \cdot \bar{x}^*},\\ &y = \frac{\bar{x}}{\bar{x}^*}\left[1 + \frac{2C(1 - i\Delta)}{1 + \Delta^2 + \bar{x} \cdot \bar{x}^*}\right] + r\bar{x} \cdot \bar{x}^*.\end{aligned} \tag{9}$$

Let us define the pumping intensity $Y = yy^*$ and the intracavity intensity $X = xx^*$; then we have

$$\left[1 + \frac{2C}{1 + \Delta^2 + X} + rX\right]^2 + \left[\frac{2C\Delta}{1 + \Delta^2 + X}\right]^2 = Y, X = 0. \tag{10}$$

The threshold of the DOPO of this combined system is modified by the atoms [17]

$$|\bar{\varepsilon}^c| = |\varepsilon_0| \left\{\left[1 + \frac{2C}{1 + \Delta^2}\right]^2 + \left[\frac{2C\Delta}{1 + \Delta^2}\right]^2\right\}^{1/2}. \tag{11}$$

Generally, the threshold of the DOPO is increased due to the atoms, and in the case of resonant interaction, it reaches the maximum threshold value $|\bar{\varepsilon}^c_{\max}| = |\varepsilon_0|(1 + 2C)$. A bistable result related to the Eq. (10) with $\Delta = 0$ and $C = 5$, $r = 0.5$, is given in Fig. 2(b), and the corresponding result for the normal optical absorptive bistability is shown in Fig. 2(a) for comparison. The turning points of intracavity intensity X for OAB are $C - 1 \pm \sqrt{C(C - 4)}$, while for atoms in DOPO, it is $\sqrt{2C/r} - 1$, as pointed out in the figures.

For simplicity, let us assume $\bar{x} = \bar{x}^*$ and $\bar{v} = \bar{v}^*$; i.e., the $\bar{x}$ and $\bar{v}$ are real quantities. Generally, the physical quantities u can be described as their steady-state mean value $\bar{u}$ and the noise ξ around $\bar{u}$ [25]; that is, $u = \bar{u} + \xi$. Equation (8) can then be simplified as

$$\frac{d\hat{u}}{d\tau} = \underline{A}\,\hat{u} + \underline{D}^{1/2}\hat{\eta}(\tau), \tag{12}$$

with

$$\begin{aligned}u &= (x, x^*, v, v^*, m)^T,\\ \bar{u} &= \left(\bar{x}, \bar{x}, -\frac{\bar{x}}{1 + \bar{x}^2}, -\frac{\bar{x}}{1 + \bar{x}^2}, -\frac{-2}{1 + \bar{x}^2}\right)^T,\\ \xi &= (\Delta x, \Delta x^*, \Delta v, \Delta v^*, \Delta m)^T,\end{aligned} \tag{13}$$

and

$$\underline{A} = \begin{pmatrix} -\mu & \mu\left(1 + \frac{2C}{1+\bar{x}^2} - 2r\bar{x}^2\right) & 2\mu C & 0 & 0\\ \mu\left(1 + \frac{2C}{1+\bar{x}^2} - 2r\bar{x}^2\right) & -\mu & 0 & 2\mu C & 0\\ \frac{-1}{1+\bar{x}^2} & 0 & -1 & 0 & \bar{x}/2\\ 0 & \frac{-1}{1+\bar{x}^2} & 0 & -1 & \bar{x}/2\\ \frac{2\bar{x}}{1+\bar{x}^2} & \frac{2\bar{x}}{1+\bar{x}^2} & -2\bar{x} & -2\bar{x} & -2 \end{pmatrix}, \tag{14}$$

$$\begin{aligned}\underline{D} = N^{-1}\mathrm{diag}\bigg(&4\mu^2 C\left(1 + \frac{2C}{1 + \bar{x}^2} - 2r\bar{x}^2\right),\\ &4\mu^2 C\left(1 + \frac{2C}{1 + \bar{x}^2} - 2r\bar{x}^2\right), \frac{-2\bar{x}^2}{1 + \bar{x}^2}, \frac{-2\bar{x}^2}{1 + \bar{x}^2}, \frac{32\bar{x}^2}{1 + \bar{x}^2}\bigg).\end{aligned} \tag{15}$$

As we know, the steady-state covariance matrix $\underline{G}$ satisfies [26]

$$\underline{A}\,\underline{G} + \underline{G}\underline{A}^T = -\underline{D}. \tag{16}$$

$\underline{D}$ actually represents a non-positive-definite diffusion. Equation (16) defines a set of 15 linear equations for the elements of the covariance matrix $\underline{G}$. Since the fluctuations ξ are

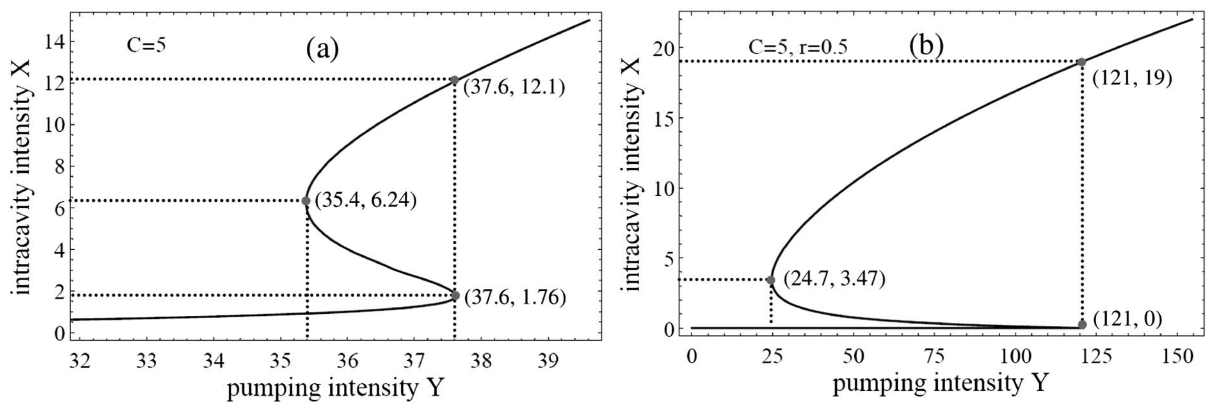

Fig. 2. Typical bistable curve of the intracavity intensity X versus pumping intensity Y; (a) normal optical absorptive bistability with $C = 5$, (b) atoms in DOPO with $\Delta = 0$, $C = 5$, and $r = 0.5$. The turning points have been pointed out.

Gaussian distributed, the moments of all orders can be calculated once these equations have been solved. Since $\underline{A}$ and $\underline{D}$ are real, the total 15 equations decouple into a set of nine effective equations for the real part of $\underline{G}$. By solving these equations, one can in principle get all the elements of the covariance matrix $\underline{G}$.

3. NUMERICAL SIMULATIONS

It is well known that the photon statistics can be characterized by the second order degree of coherence $g^{(2)}$ and the Mandel parameter Q. The former describes the bunching ($g^{(2)} > 1$) or antibunching ($g^{(2)} < 1$) of the photons, and the latter tells how fluctuated the photon number of the field is. $Q > 0(< 0)$ indicates super(sub)-Poissonian distribution of the photon number, which means large(small) fluctuations of the intensity of the field [27]. According to the definitions of these two parameters, we get

$$g^{(2)}(0) = \frac{\langle a^{+2}a^2\rangle}{\langle a^+a\rangle^2} = \frac{\langle(\bar{a}^+ + \Delta a^+)^2(\bar{a}+\Delta a)^2\rangle}{\langle(\bar{a}^+ + \Delta a^+)(\bar{a}+\Delta a)\rangle^2} = 1 + \frac{\bar{a}^{+2}\langle\Delta a^2\rangle + \bar{a}^2\langle\Delta a^{+2}\rangle + 2\bar{a}^+\bar{a}\langle\Delta a^+\Delta a\rangle + \langle\Delta a^{+2}\Delta a^2\rangle - \langle\Delta a^+\Delta a\rangle^2}{(\bar{a}^+\bar{a} + \langle\Delta a^+\Delta a\rangle)^2},$$
$$Q = \frac{N}{4\mu C}(g^{(2)}(0) - 1). \quad (17)$$

For large number of atoms, they are approximated to be

$$g^{(2)}(0) \simeq 1 + \frac{2(\langle\Delta x\Delta x\rangle + \langle\Delta x^*\Delta x\rangle)}{\bar{x}^2},$$
$$Q \simeq \frac{N}{4\mu C}\frac{2(\langle\Delta x\Delta x\rangle + \langle\Delta x^*\Delta x\rangle)}{\bar{x}^2}. \quad (18)$$

Figures 3 and 4 show the second order degree of coherence $g^{(2)}$ at zero delay time and the Mandel Q parameter of the intracavity field versus intracavity intensity X with various cavity-decay rates and atomic-cooperative parameter C, respectively. Here we have chosen the intracavity atom number $N = 10000$. Figures (a) and (c) are for OAB, and Figs. (b) and (d) are for atoms in the DOPO, respectively. We can see that very weak antibunching and sub-Poissonian statistics appear for the usual OAB system when the intracavity field is weak. There exist two sudden changes: one is from super-Poissonian to sub-Poissonian, and the other is just the opposite. These two sharp peaks exactly correspond to the two turning points of the bistability of the system. Between the two turning points, the system is unstable and the result of photon statistics is not valid anymore. In the case of a large number of atoms, the normal OAB system shows weak antibunching and sub-Poissonian behaviors when the intracavity field is very weak, for which the nonclassical properties appear more likely. On the other hand, for atoms in DOPO, the intracavity field shows strong super-Poissonian when the X is approaching zero; this is because the subharmonic field is a squeezed vacuum field below threshold that contains only even photons and is a typical photon-bunched state. When the system operates above threshold, weak antibunching and strong sub-Poissonian statistics appear. The unstable regimes are also shown [see Figs. 3(b) and 4(b), and the two turning points are at 0 and 3.47]. In addition, when increasing the intracavity intensity, the Mandel parameter Q goes to −1, while $g^{(2)}(0)$ is very close to 1, which shows that the scheme can be used for

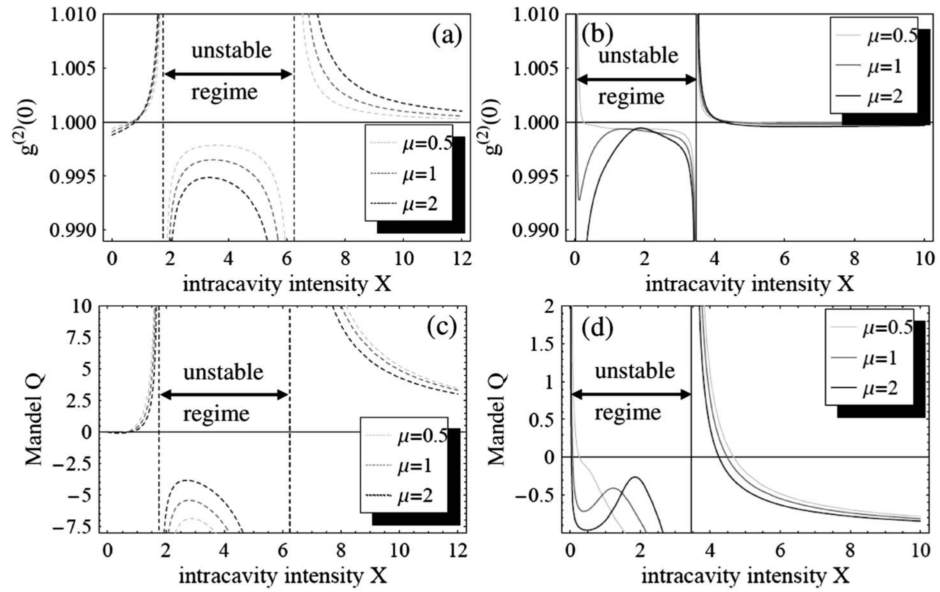

Fig. 3. The second order degree of coherence for zero delay time $g^{(2)}(0)$ and Mandel Q versus intracavity intensity with different cavity-decay rates ($N = 10000$). (a) and (c) are for OAB, $C = 5$; (b) and (d) are for atoms in DOPO, $C = 5$ and $r = 0.5$.

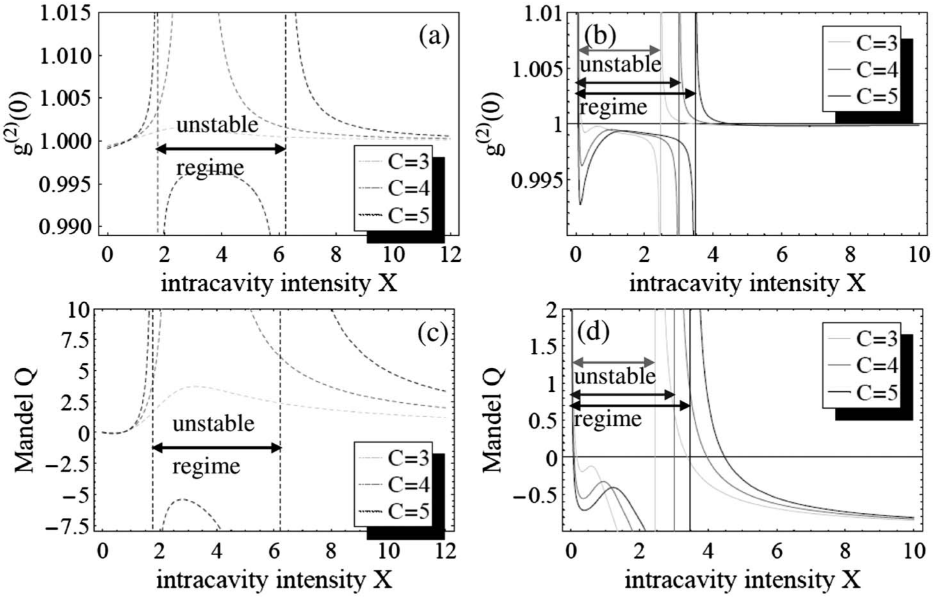

Fig. 4. The second order degree of coherence $g^{(2)}(0)$ and Mandel Q versus intracavity intensity with different atomic-cooperative parameter C ($N = 10000$). (a) and (c) are for OAB, $\mu = 1$; (b) and (d) are for atoms in DOPO, $\mu = 1$ and $r = 0.5$.

generating Fock-state with a definite number of photons. This is reasonable because there is a certain number of atoms in the cavity, which are coupled to photons through the atom-photon interaction. The results also indicate that larger cavity-decay rate μ and atomic-cooperative parameter C lead to stronger antibunching and stronger sub-Poissonian effects. The reason is that cavity decay decreases the intracavity squeezing, which eventually makes the intracavity field show sub-Poissonian for the weak field [28]. On the other hand, atom-field coupling is helpful to increase the nonclassicality for a certain number of atoms [29]. It is worth to mention that, if $C \to 0$, the OAB system goes back to the case of the coherent light field in the optical cavity, and the combined "atom + DOPO" system turns to the situation of a normal DOPO [30]. All the parameters above can be reached based on the present state-of-the-art of cold atom and the optical parametric oscillator, and the system is feasible in experiment.

4. CONCLUSIONS

We have investigated the photon statistics of the combined "atom + NOPO" system. The second order degree of coherence $g^{(2)}(0)$ and the Mandel parameter Q of the intracavity field are calculated. The numerical result of a large number of atoms is obtained for various parameters based on the present experimental situations. The result is also compared to the normal optical absorptive bistability system. It shows that weak antibunching but strong sub-Poissonian statistics can be obtained in this combined system, which means that the system can be used for Fock-state generation with a large number of photons. The scheme does not require a strong nonlinearity at the single quantum level or high-Q cavity in the strong coupling regime. In contrast to the normal bistability system, there is no strong sub-Poissonian effect. This system thus provides an approach to demonstrate the generation of the optical Fock-state with a large number of photons. The present investigation is also significant to understand the photon statistical properties of nonclassical light interacting with atoms. Since the generated nonclassical state is just within the cavity, people do not need to inject the quantum light source from outside, as was demonstrated in the early time experiment [31]. The system is thus more efficient for investigating the interaction between quantum light sources and atoms.

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