Partial entropy change and entanglement in the mixed state for a Jaynes–Cummings model with Kerr medium*

Zhang Yu-Qing(张玉青), Tan Lei(谭 磊)†, Zhu Zhong-Hua(朱中华), Xiong Zu-Zhou(熊祖周), and Liu Li-Wei(刘利伟)

Institute of Theoretical Physics, Lanzhou University, Lanzhou 730000*, China*

(Received 8 June 2009; revised manuscript received 15 June 2009)

By using the algebraic dynamical approach, an atom–field bipartite system in mixed state is employed to investigate the partial entropy change and the entanglement in a cavity filled with Kerr medium. The effects of different nonlinear intensities are studied. One can find that the Kerr nonlinearity can reduce the fluctuation amplitudes of the partial entropy changes and the entanglement of the two subsystems, and also influence their periodic evolution. Meanwhile, increasing the Kerr nonlinear strength can convert the anti-correlated behaviour of the partial entropy change to the positively correlated behaviour. Furthermore, the entanglement greatly depends on the temperature. When the temperature or the nonlinear intensity increases to a certain value, the entanglement can be suppressed greatly.

Keywords: entanglement, algebraic dynamical method, Jaynes–Cummings model, Kerr medium

PACC: 4250, 0365, 0367

1. Introduction

Quantum entropy, which was proposed by von Neumann, is one of the most fundamental concepts of quantum information.[1−3] It can be used to measure the quantum entanglement, quantum optical correlation and quantum decoherence, etc. Entropy and entanglement have been studied widely, especially for bipartite systems. Even the entanglement for multipartite systems is also based on the study of the entanglement of the bipartite system, where the N individual constituents are arbitrarily partitioned into two subgroups.[4] In recent years, the bipartite system without taking into consideration the influence of the environment[5,6] and the bipartite pure open quantum system[7−9] are well investigated. It is well known that for a bipartite system that starts in a pure state, the partial entropies of the two subsystems are equal at all times, and the entanglement can be measured by either of them. However, if one aspires to describe realistic states observed in experiments, it is imperative to allow for a proper quantification also of mixed states entanglement, since there is no system that could be decoupled perfectly from environmental influences, and mixing is thus unavoidable. While for the mixed state, the entanglement is difficult to be measured, because it is not easy to define an analog of the Schmidt decomposition for a general mixed state of a composite system. However, intensive efforts have been done to understand and quantify the entanglement of a mixed state.[10−13] In particular, the relative entropy of entanglement, the entanglement formation and the entanglement distillation are the most commonly used methods to measure the entanglement for a mixed state, which are related to each other through the amount of available classical information about the decomposition of the state. Based on the previous methods, more extensive and further investigations are presented to study the entanglement of mixed states for the dissipative scheme[14] and the multi-dimensional system[15] or infinite dimensional system.[16]

The model of a two-level atom interacting with a single-mode thermal cavity field is widely used to study the entropy for a bipartite system, which can be easily treated mathematically. Based on this simple model, a great number of researches have been

*Project supported by the National Natural Science Foundation of China (Grant No. 10704031), the National Science Foundation for Fostering Talents in Basic Research of the National Natural Science Foundation of China (Grant No. J0630313), the Fundamental Research Fund for Physics and Mathematics of Lanzhou University (Grant No. Lzu05001), and the Natural Science Foundation of Gansu Province, China (Grant No. 3ZS061-A25-035).

†Corresponding author. E-mail: tanlei@lzu.edu.cn

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conducted such as the revivals and the entanglement of a damped Jaynes–Cummings model (JCM),[17] the macroscopic superposition and the entanglement for displaced thermal fields[18] and the regime of entropy exchange between light and matter.[19] The influences of the atomic motion and the field-mode structure on the entropy exchange and the entanglement are also investigated.[20] However, the effect of Kerr nonlinearity process has not been taken into account. Introducing the Kerr nonlinearity in a system will cause various effects, such as nonclassical states,[21] quantum nondemolition,[22] quantum fluctuations,[23] and quantum information.[24] In particular, the Kerr nonlinearity can affect the entanglement in a quantum optical system. It can influence its fluctuation amplitudes, change its evolution period,[25] control the entanglement sudden death[26] and realize the entanglement transfer.[27] But these studies are limited to the situation of pure state. What effects will occur when it is applied to a mixed system is just what we focus on here. At the same time, for a mixed state, the Hamiltonian containing Kerr nonlinearity is not easy to treat through the traditional methods. Here, we propose the algebraic dynamical approach[28] that is an effective tool to deal with the nonlinear system, whose key idea is introducing a canonical transformation that transforms the Hamiltonian into a linear function in terms of a set of Lie algebraic generators. According to algebraic dynamics, linear systems are integrable and solvable, then the time evolution operator and the density operator can be obtained easily. In this paper, we will adopt the algebraic dynamical approach to study the partial entropy change and the entanglement of an atom–field bipartite system in mixed states and in the presence of a nonlinear Kerr medium.

The paper is organized as follows: In Section 2, the basic model, an atom–field bipartite system in the mixed states including the nonlinear effect, is introduced. Then, the density operator of the considered system is derived by means of the algebraic dynamical approach. Section 3 is devoted to the investigation of the entropy correlations and the entanglement between two subsystems. In Section 4, we discuss the results of the partial entropy change and the entanglement for different nonlinear intensities of the Kerr medium. Finally, we present our conclusion in Section 5.

2. Model

In this section, we will derive the analytical expression of the density operator using the algebraic dynamical approach. The interaction between a two-level atom and a quantized cavity field in a nonlinear process can be described by the Hamiltonian ($\hbar = 1$)[29,30]

$$H = H_{\rm a} + H_{\rm f} + H_{\rm af}, \tag{1}$$

where $H_{\rm a}$ is the Hamiltonian of the two-level atom, $H_{\rm f}$ is the Hamiltonian of the field in the presence of a nonlinear Kerr medium and $H_{\rm af}$ is the Hamiltonian of the atom–field interaction, which can be given by

$$\begin{aligned} H_{\rm a} &= \omega S_z, \\ H_{\rm f} &= \omega a^\dagger a + \chi (a^\dagger)^2 a^2, \\ H_{\rm af} &= g(a^\dagger S_- + a S_+)\ . \end{aligned} \tag{2}$$

Here a and $a^\dagger$ are the annihilation and creation operators of the cavity field; S_+ and S_- represent the atomic raising and lowering operators; S_z is the atomic inversion operator; ω is the frequency for both field and the two-level atom in resonant case; g is the atom–field coupling strength, and the coupling constant χ is related to the dispersive part of the nonlinearity of the Kerr-like medium. A straightforward analysis of the system shows that there exists a conserved quantity N for the Hamiltonian (1)

$$N = a^\dagger a + S_z + \frac{1}{2}, \tag{3}$$

and N commutes with the Hamiltonian (1). Based on the algebraic dynamical approach, introduce SU(2) algebra generators $\{J_0,\ J_+,\ J_-\}$, with $J_0 = S_z$, $J_+ = N^{-1/2} a S_+$, $J_- = N^{-1/2} a^\dagger S_-$, which are nonlinear expressions and obey the following commutation relations

$$\begin{aligned} [J_0, J_+] &= J_+, \\ [J_0, J_-] &= -J_-, \\ [J_+, J_-] &= 2J_0. \end{aligned} \tag{4}$$

In terms of the SU(2) algebra generators, we can rewrite the Hamiltonian (1) as

$$H = \Delta J_0 + E + g N^{1/2}(J_+ + J_-), \tag{5}$$

where

$$\Delta = -2\chi(N-1),$$

$$E = \chi(N^2 - 2N + 1) + \omega(N - 1/2). \quad (6)$$

Then, Hamiltonian (1) can be diagonalized by the following canonical transformation operator

$$U_g = \exp(\theta J_+ - \theta J_-) \quad (7)$$

with $\theta = -\arctan[(\sqrt{\Delta^2/4 + g^2N} - \Delta/2)/(gN^{1/2})]$. After the canonical transformation, the dressed Hamiltonian is expressed as

$$H' = U_g^{-1} H U_g = E + \lambda J_0, \quad (8)$$

where $\lambda = \sqrt{\Delta^2 + 4g^2N}$. Then the time evolution operator can be written as

$$\begin{aligned} U(t) &= \mathrm{e}^{-\mathrm{i}Ht} \\ &= U_g \mathrm{e}^{-\mathrm{i}H't} U_g^{-1} \\ &= \mathrm{e}^{-\mathrm{i}Et}\Big[\cos\frac{\lambda t}{2} - 2\mathrm{i}J_0 \sin\frac{\lambda t}{2}\cos 2\theta \\ &\quad + \mathrm{i}(J_+ + J_-)\sin\frac{\lambda t}{2}\sin 2\theta\Big]. \end{aligned} \quad (9)$$

What should be noticed here is that using canonical transformation operator to diagonalize the nonlinear Hamiltonian (1) does not change its intrinsic qualities, which is a convenient way to deal with nonlinear systems.

Under the same initial conditions as those in Ref. [20], the cavity field is in a single-mode thermal state

$$\rho_{\mathrm{f}}(0) = \sum_{n=0}^{\infty} P_n |n\rangle\langle n|, \quad P_n = \frac{m^n}{(m+1)^{n+1}}, \quad (10)$$

where $m = 1/[\exp(\omega/T) - 1]$ is the mean photon number of the cavity field, and the two-level atom is initially in the mixed state

$$\rho_{\mathrm{a}}(0) = C_e|e\rangle\langle e| + C_g|g\rangle\langle g|, \quad (11)$$

where C_e and C_g are the components of the two-state density matrix of the atom with $C_e + C_g = 1$. So, the initial density operator for the atom–field system can be written as

$$\begin{aligned} \rho_{\mathrm{af}}(0) &= \rho_{\mathrm{f}}(0) \otimes \rho_{\mathrm{a}}(0) \\ &= \sum_n P_n C_e |n, e\rangle\langle n, e| \\ &\quad + \sum_n P_n C_g |n, g\rangle\langle n, g|. \end{aligned} \quad (12)$$

By using Eqs. (9) and (12), the density operator with the time evolution for the total system can be derived

$$\begin{aligned} \rho_{\mathrm{af}}(t) &= U(t)\,\rho_{\mathrm{af}}(0)\,U^+(t) \\ &= C_e \sum_n P_n \Bigg\{ \left[\cos^2\left(\frac{\lambda_{n+1}t}{2}\right) + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\cos^2(2\theta_{n+1})\right] |n, e\rangle\langle n, e| \\ &\quad + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1})\,|n+1, g\rangle\langle n+1, g| \\ &\quad + \sin\left(\frac{\lambda_{n+1}t}{2}\right)\sin(2\theta_{n+1})\left[\mathrm{i}\cos\frac{\lambda_{n+1}t}{2} - \sin\frac{\lambda_{n+1}t}{2}\cos(2\theta_{n+1})\right] |n+1, g\rangle\langle n, e| \\ &\quad + \sin\left(\frac{\lambda_{n+1}t}{2}\right)\sin(2\theta_{n+1})\left[-\mathrm{i}\cos\frac{\lambda_{n+1}t}{2} - \sin\frac{\lambda_{n+1}t}{2}\cos(2\theta_{n+1})\right] |n, e\rangle\langle n+1, g| \Bigg\} \\ &\quad + C_g \sum_n P_n \Bigg\{ \left[\cos^2\left(\frac{\lambda_n t}{2}\right) + \sin^2\left(\frac{\lambda_n t}{2}\right)\cos^2(2\theta_n)\right] |n, g\rangle\langle n, g| \\ &\quad + \sin^2\left(\frac{\lambda_n t}{2}\right)\sin^2(2\theta_n)\,|n-1, e\rangle\langle n-1, e| \\ &\quad + \sin\left(\frac{\lambda_n t}{2}\right)\sin(2\theta_n)\left[-\mathrm{i}\cos\frac{\lambda_n t}{2} + \sin\frac{\lambda_n t}{2}\cos(2\theta_n)\right] |n, g\rangle\langle n-1, e| \\ &\quad + \sin\left(\frac{\lambda_n t}{2}\right)\sin(2\theta_n)\left[\mathrm{i}\cos\frac{\lambda_n t}{2} + \sin\frac{\lambda_n t}{2}\cos(2\theta_n)\right] |n-1, e\rangle\langle n, g| \Bigg\}, \end{aligned} \quad (13)$$

where $\lambda_n = \sqrt{\Delta_n^2 + 4g^2 n}$, $\theta_n = -\arctan[(\sqrt{\Delta_n^2/4 + g^2 n} - \Delta_n/2)/(gn^{1/2})]$ with $\Delta_n = -2\chi(n-1)$.

3. Entropy correlations and the dynamical behaviour of entanglement

We now turn to investigate the entropy correlations and the dynamical behaviour of the entanglement between the two subsystems. As is known, for a bipartite system that starts from a pure state, the partial entropies of the two subsystems are equal at all times, because of the equal eigenvalues of the partial density matrices. Here, we focus on the case of the system initially in the mixed state and discuss the effects of nonlinear process on the partial entropy change and the entanglement.

Based on the time-dependent density operator $\rho_{\rm af}(\rm t)$ in Eq. (13), the reduced density operators of the atom and the cavity field can be obtained respectively

$$\begin{aligned}\rho_{\rm a}(t) &= {\rm Tr_f}\rho_{\rm af}(t) = \xi_1(t)|e\rangle\langle e| + \xi_2(t)|g\rangle\langle g|,\\ \rho_{\rm f}(t) &= {\rm Tr_a}\rho_{\rm af}(t) = \xi_3(t)|n\rangle\langle n|,\end{aligned} \tag{14}$$

where

$$\begin{aligned}\xi_1(t) &= C_e\sum_n P_n\left[\cos^2\left(\frac{\lambda_{n+1}t}{2}\right) + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\cos^2(2\theta_{n+1})\right] + C_g\sum_n P_{n+1}\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1}),\\ \xi_2(t) &= C_e\sum_n P_{n-1}\sin^2\left(\frac{\lambda_n t}{2}\right)\sin^2(2\theta_n) + C_g\sum_n P_n\left[\cos^2\left(\frac{\lambda_n t}{2}\right) + \sin^2\left(\frac{\lambda_n t}{2}\right)\cos^2(2\theta_n)\right],\\ \xi_3(t) &= C_e\sum_n P_n\left[\cos^2\left(\frac{\lambda_{n+1}t}{2}\right) + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\cos^2(2\theta_{n+1})\right] + C_e\sum_n P_{n-1}\sin^2\left(\frac{\lambda_n t}{2}\right)\sin^2(2\theta_n)\\ &\quad + C_g\sum_n P_{n+1}\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1}) + C_g\sum_n P_n\left[\cos^2\left(\frac{\lambda_n t}{2}\right) + \sin^2\left(\frac{\lambda_n t}{2}\right)\cos^2(2\theta_n)\right].\end{aligned} \tag{15}$$

To investigate the entropy correlations or the anticorrelated behaviour, we should first calculate the partial entropy changes. For simplicity, we adopt the von Neumann entropy, which is defined by their respective reduced density operators of the components of the system

$$S_i = -{\rm Tr}_i(\rho_i \ln \rho_i), \tag{16}$$

where the subscript i may represent the atom, the field or the atom–field system. With the results of Eqs. (13), (14) and (15), the partial entropy changes for the atom, the field and the sum of them can be written as: $\Delta S_{\rm a}(t) = S_{\rm a}(t) - S_{\rm a}(0)$, $\Delta S_{\rm f}(t) = S_{\rm f}(t) - S_{\rm f}(0)$ and $\Delta S(t) = \Delta S_{\rm a}(t) + \Delta S_{\rm f}(t)$, respectively.

The entanglement formation or, equivalently, concurrence,[12] can quantify the entanglement of 2×2 systems as well as other specific finite systems with special symmetries. But what we will discuss here is a $2\times\infty$ system. However, we can better evaluate the entanglement of $\rho_{\rm af}(t)$ using the lower bound concurrence by projecting $\rho_{\rm af}(t)$ onto a 2×2 system for each photon state n.

The projection of $\rho_{\rm af}(t)$ on the subspace spanned by the basis $|g,n\rangle$, $|e,n\rangle$, $|g,n+1\rangle$ and $|e,n+1\rangle$ can be expressed as

$$\rho_{\rm af}^n(t) = \begin{pmatrix} D_{n-1} & 0 & 0 & 0\\ 0 & E_n & F_n & 0\\ 0 & F_n^* & D_n & 0\\ 0 & 0 & 0 & E_{n+1}\end{pmatrix}, \tag{17}$$

where

$$\begin{aligned}D_n &= C_e P_n\left[\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1})\right] + C_g P_{n+1}\left[\cos^2\left(\frac{\lambda_{n+1}}{2}\right) + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\cos^2(2\theta_{n+1})\right],\\ E_n &= C_e P_n\left[\cos^2\left(\frac{\lambda_{n+1}}{2}\right) + \sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\cos^2(2\theta_{n+1})\right] + C_g P_{n+1}\left[\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1})\right],\\ F_n &= C_e P_n\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1})\left[-{\rm i}\cos\left(\frac{\lambda_{n+1}t}{2}\right) - \sin\left(\frac{\lambda_{n+1}t}{2}\right)\cos(2\theta_{n+1})\right]\\ &\quad + C_g P_{n+1}\sin^2\left(\frac{\lambda_{n+1}t}{2}\right)\sin^2(2\theta_{n+1})\left[{\rm i}\cos\left(\frac{\lambda_{n+1}t}{2}\right) + \sin\left(\frac{\lambda_{n+1}t}{2}\right)\cos(2\theta_{n+1})\right].\end{aligned} \tag{18}$$

With the above results, the concurrence of the $\rho_{\rm af}^n(t)$ can be given by

$$C^n = 2\max[0, |F_n| - \sqrt{D_{n-1} \times E_{n+1}}] \ . \quad (19)$$

Then, we can obtain the lower bound concurrence C by calculating the entanglement of $\rho_{\rm af}^n(t)$ and taking an average over all possible values of n (weighted by the probability of obtaining $\rho_{\rm af}^n(t)$). This holds because $\rho_{\rm af}^n(t)$ is the result of a local projection on $\rho_{\rm af}(t)$, which cannot increase the entanglement on average.[13,17] If this lower bound concurrence is positive, then the system is entangled.

4. Results and discussion

Based on the analytical solutions in Section 3, in this section we will present some numerical results for different parameters in Figs. 1–4 to demonstrate the effects of the Kerr nonlinearity on the partial entropy change and the entanglement.

In order to analyse the influence of Kerr nonlinearity on the partial entropy change, we consider the field is in the weak thermal state with average photon number $m = 0.1$. It clearly shows that when $\chi \ll g$, the behaviours of the partial entropy change are quite different from each other for different atomic initial states, as shown in Figs. 1(a)–4(a). When the atom is initially in the ground state, there is a distinct anti-correlated behaviour for the partial entropy changes of the two subsystems as depicted in Fig. 1(a), and with the increase of the probability of the atom in the excited state, this behaviour is gradually replaced by a positively correlated behaviour, for which the partial entropy changes of the atom and the field are nearly rising and reducing together; and this is similar to the situation that the field and the atom are initially in pure states. At the same time, it is worth while noting that in the situation of $\chi \ll g$ and the atom is initially in a mixed state, the partial entropy change of the field is greater than that of the atom, but opposite behaviour occurs when the atom is initially in a pure state. For the situations of $\chi \simeq g$ or $\chi > g$ (Figs. 1(b), 1(c), 1(d)–Figs. 4(b), 4(c), 4(d)), whether

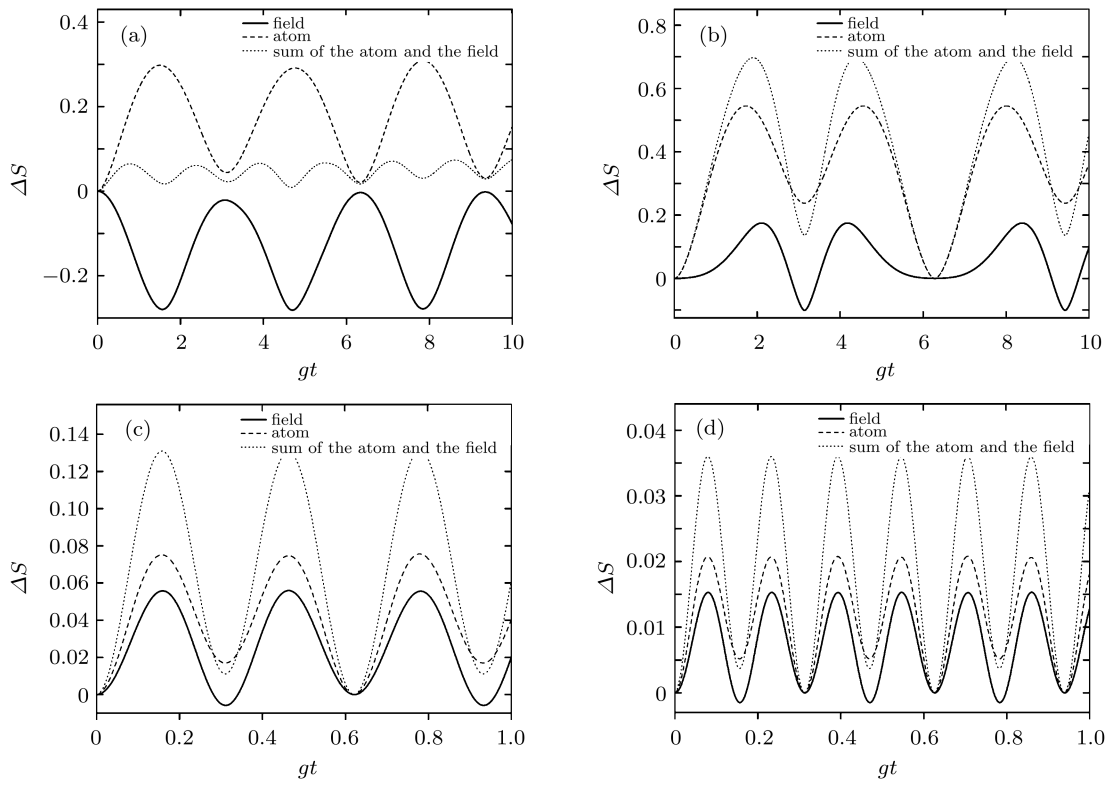


Fig. 1. Partial entropy changes for different Kerr nonlinear strengths. Initially, the atom is in the ground state ($C_g = 1$, $C_e = 0$) and the field is in a weakly excited thermal state ($m = 0.1$). (a) $\chi/g = 0.01$; (b) $\chi/g = 0.5$; (c) $\chi/g = 5$; (d) $\chi/g = 10$.

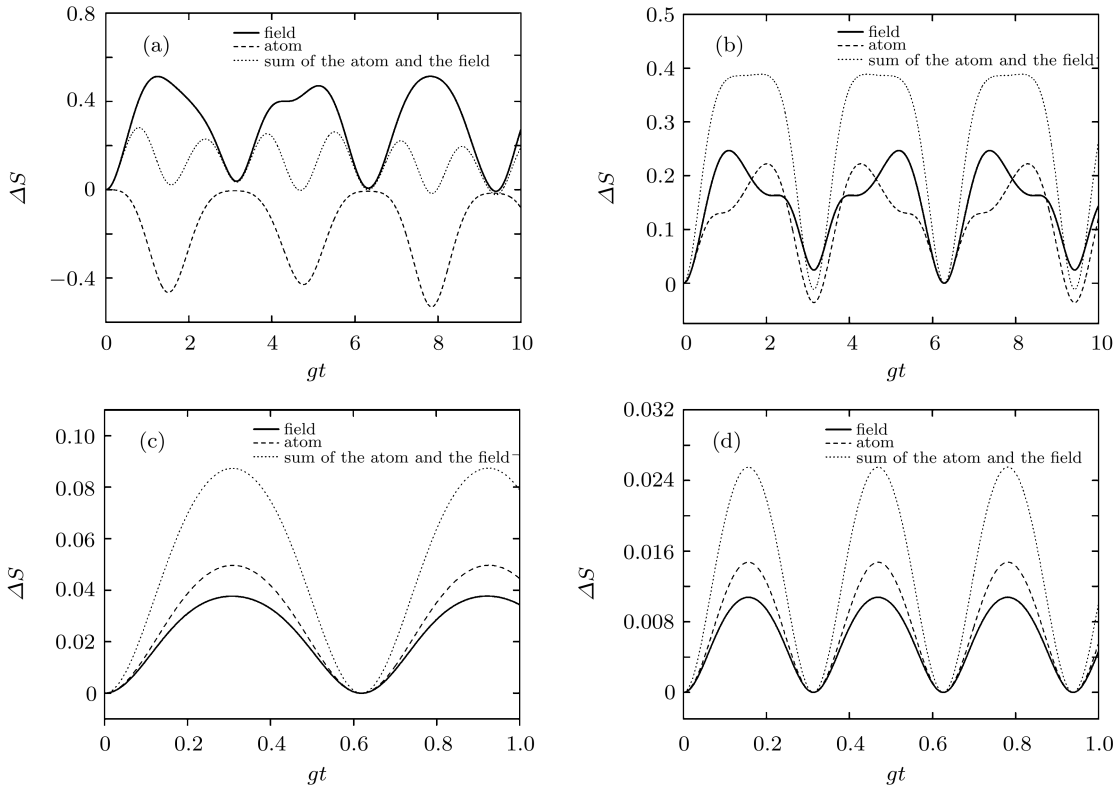


Fig. 2. Partial entropy changes for different Kerr nonlinear strengths. Initially, the atom is in the maximally mixed state ($C_e = C_g = 0.5$) and the field is in a weakly excited thermal state ($m = 0.1$). (a) $\chi/g = 0.01$; (b) $\chi/g = 0.5$; (c) $\chi/g = 5$; (d) $\chi/g = 10$.

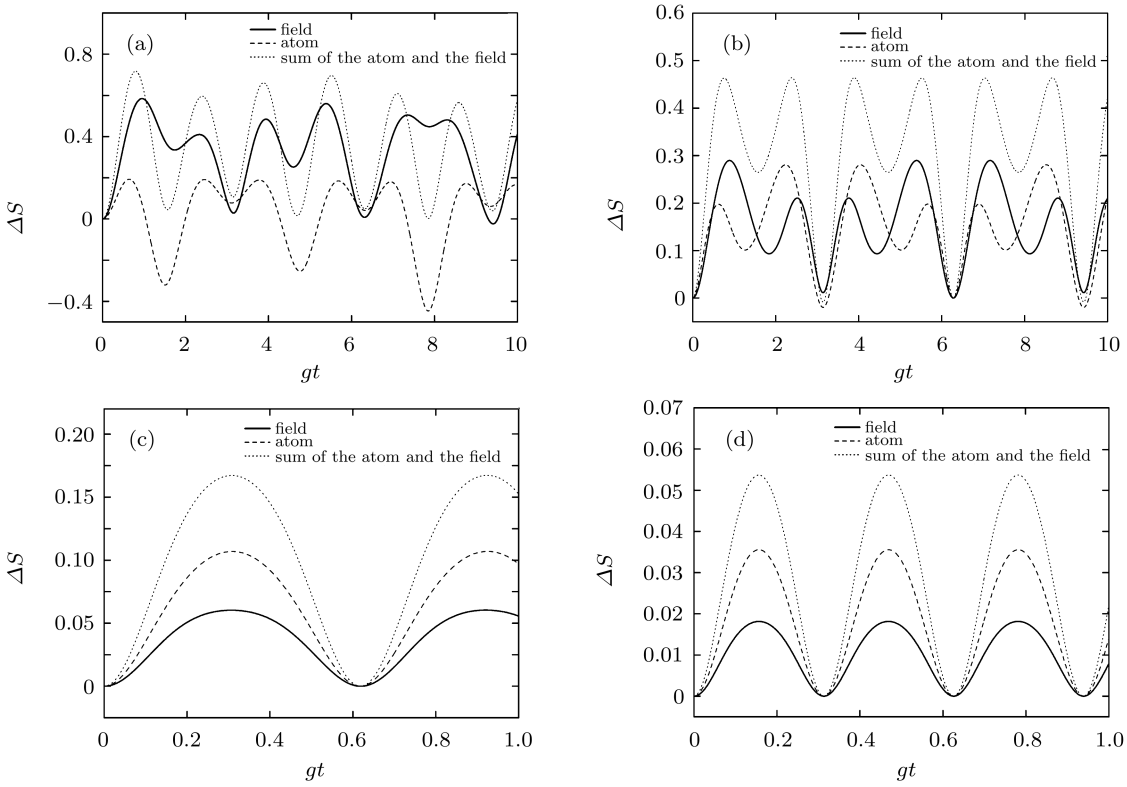


Fig. 3. Partial entropy changes for different Kerr nonlinear strengths. Initially, the atom is in the mixed state ($C_e = 0.8$, $C_g = 0.2$) and the field is in a weakly excited thermal state ($m = 0.1$). (a) $\chi/g = 0.01$; (b) $\chi/g = 0.5$; (c) $\chi/g = 5$; (d) $\chi/g = 10$.

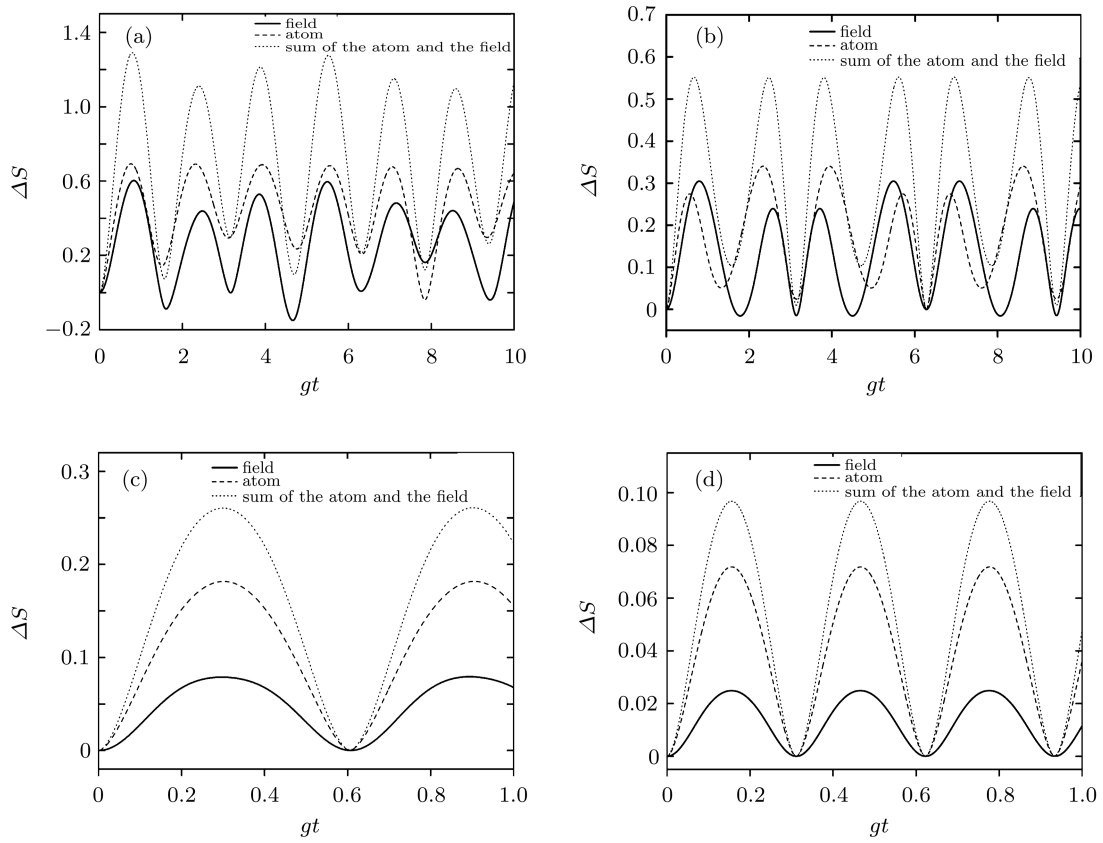


Fig. 4. Partial entropy changes for different Kerr nonlinear strengths. Initially, the atom is in the excited state ($C_e = 1$, $C_g = 0$) and the field is in a weakly excited thermal state ($m = 0.1$). (a) $\chi/g = 0.01$; (b) $\chi/g = 0.5$; (c) $\chi/g = 5$; (d) $\chi/g = 10$.

the atom is in a pure state or in a mixed state, the increase of the nonlinear strength can decrease the fluctuation amplitude of the partial entropy change and shorten its evolution periodicity, and meanwhile the partial entropy change fluctuates with a distinct and regular period as compared to the case of $\chi \ll g$. In addition, we also can find that when the probability of the atom in the ground state is higher, increasing the Kerr nonlinear strength also can convert the anti-correlated behaviour of the partial entropy change to the positively correlated behaviour. This is because for high values of the Kerr nonlinearity there is a competition between the atom–field interaction and nonlinearity effect. As a matter of fact, when $\chi \ll g$, the effect of the Kerr nonlinearity on the entropy evolution is small as compared with the effect of atom–field interaction. As χ is close to g, both of them play the crucial part in the partial entropy change. Once the Kerr nonlinearity dominates over the atom–field coupling, the system goes into a kind of regular evolution, characteristic of the Kerr-like medium.

Figure 5 depicts the effect of Kerr nonlinear process on the entanglement. Our calculation shows that when the atom is initially in the ground state, the concurrence is very small. So, we just discuss the case of the atom is initially in the excited state. When $\chi \ll g$ the atom and the field almost maintain a strong entanglement with a regular periodic time evolution, as shown in Fig. 5(a). The increase of the nonlinear intensity can lead to the decrease of concurrence and the destruction of the regular periodic time evolution, just as shown in Fig. 5(b). Once the Kerr nonlinearity dominates the interaction process, the fluctuation amplitude of concurrence becomes very small and the two subsystems are weakly entangled in this case, see Figs. 5(c) and 5(d). In addition, one can see that the concurrence decreases with the increase of the mean photon number m, which is proportional to the temperature. When the temperature (or the mean photon number) increases to a certain value, the concurrence can be greatly suppressed and the entanglement nearly vanishes.

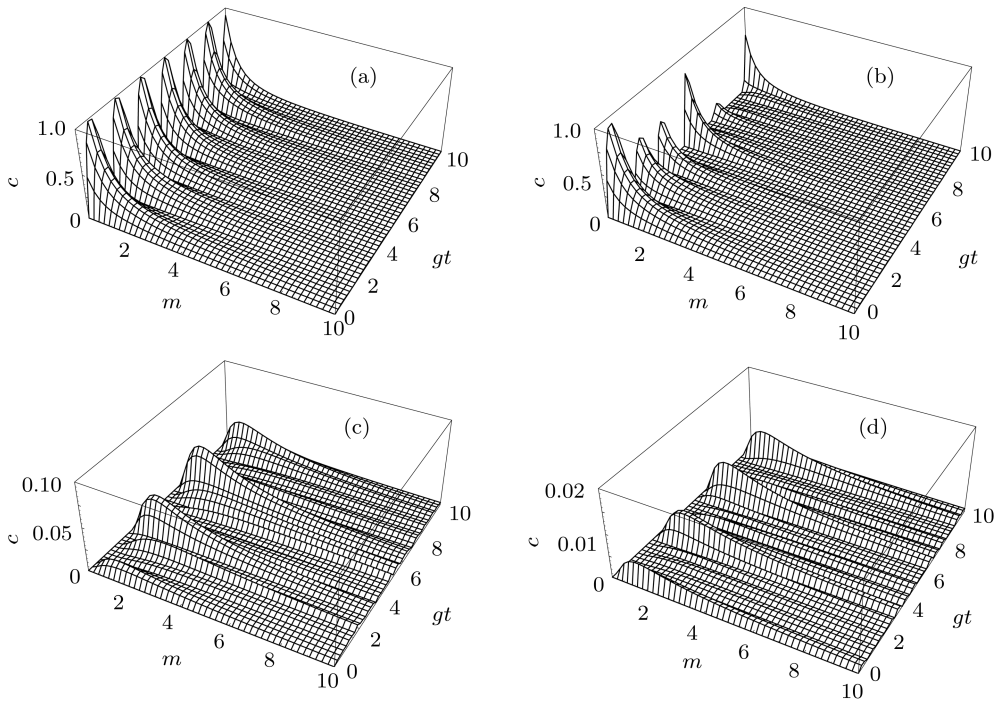


Fig. 5. Lower bound concurrence for different Kerr nonlinear strengths. Initially, the atom is in the excited state. (a) $\chi/g = 0.01$; (b) $\chi/g = 0.5$; (c) $\chi/g = 5$; (d) $\chi/g = 10$.

5. Conclusion

We have studied the partial entropy change and the entanglement for the JCM in the mixed state with the Kerr-like nonlinear process, using the algebraic dynamical approach. It is shown that the partial entropy changes of the two subsystems and their correlated behaviour depend greatly on the Kerr nonlinear intensity and initial states of the atom and the mean photon number of the field. The higher the probability of the atom in the ground state, the easier the anti-correlated behaviour occurs. Meanwhile, the increase of the nonlinear intensity can lead to the decrease of the partial entropy change and convert the anti-correlated behaviour into the positively correlated behaviour. On the other hand, the entanglement is greatly influenced by he tKerr-like nonlinear intensity and the temperature. When the temperature or nonlinear intensity increases to a certain value, the entanglement of the two subsystems can be greatly suppressed.

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