Controlling entanglement sudden death in cavity QED by classical driving fields

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(Dated: July 17, 2008)

We investigate the entanglement dynamics of a quantum system consisting of two-level atoms interacting with vacuum or thermal fields with classical driving fields. We find that the entanglement of the system can be improved by adjusting the classical driving field. The influence of the classical field and the purity of the initial state on the entanglement sudden death is also studied. It is shown that the time of entanglement sudden death can be controlled by the classical driving fields. Particularly, the entanglement sudden death phenomenon will disappear if the classical driving fields are strong enough.

PACS numbers: 03.67.Mn; 03.65.Ud

I. INTRODUCTION

Entanglement, one of the most striking features of quantum mechanics, has been considered as a key resource of quantum information processing [1, 2, 3, 4]. In recent years, the manipulation of quantum entanglement for the system of the cavity quantum electrodynamics (QED) has been extensively investigated [5, 6, 7, 8, 9, 10, 11, 12, 13, 14]. The cavity QED, where atoms interacting with quantized electromagnetic fields inside a cavity, can be used to create the entanglement between atoms in cavities and establish quantum communications between different optical cavities. It has been shown that the entanglement can arise in the interaction of a atom with a cavity field in a thermal state [6, 7] and two atoms with a cavity field [13, 14]. In Ref. [13], Solano et al. have shown that multipartite entanglement can be generated by putting several two-level atoms in a cavity of high quality factor. With the help of a strong classical driving fields, the Schrödinger cat state and other entangled states can be produced.

On the other hand, many efforts have been devoted to the study of the evolution of the joint system formed by two qubits [15, 16, 17, 18, 19, 20, 21, 22, 23]. Particularly, Yu and Eberly [16] pointed out that the single-qubit dynamics and the global dynamics of an entangled two-qubit system subjected to independent environments may be rather different. For example, for a single-qubit system subjected to an environment, the local coherence decays asymptotically. However, the entanglement of an entangled two-qubit system may disappear within a finite time during the dynamics evolution. The nonsmooth finite-time disappearance of entanglement is called “entanglement sudden death” (ESD). Recently, it is reported that the ESD phenomenon has been observed in a quantum optics experiment [24].

In the present paper, we propose a scheme to enhance the entanglement of a quantum system consisting of two-level atoms within cavities by applying and controlling classical driving fields. The influence of the classical field and the purity of the initial state on the ESD is also studied. First, we study the entanglement of the two-level atom and the field by employing logarithmic negativity [25, 26]. Then, we consider a quantum system consisting of two noninteracting atoms each locally interacting with its own vacuum field. The two atoms, which are initially prepared in extended Werner-like (EWL) states [22], are driven by two independent classical fields additionally. It is shown that the ESD may appear in this system and the ESD time can be retarded by increasing the purity of the initial state of the two atoms. In addition, the amount of the entanglement of the two atoms can be significantly increased by applying classical fields. We find that the time of entanglement sudden death can be controlled by the classical driving fields. It is interesting to point out that the entanglement sudden death phenomenon disappear if the classical driving fields are strong enough.

The present paper is organized as follows. In section II, we propose a scheme to improve the entanglement of a two-level atom interacting with vacuum or thermal field by applying a classical driving field. In section III, we study the influence of the classical driving fields on the ESD when the atoms are initially prepared in EWL states. A conclusion is given in section IV.

II. ENTANGLEMENT DYNAMICS OF A TWO-LEVEL ATOM IN A CA VITY WITH A CLASSICAL DRIVING FIELD

A. The model

Now, we consider a system consisting of a two-level atom inside a single mode cavity. The atom is driven
by a classical field additionally. The Hamiltonian of the system can be described by
\begin{equation}
H = \omega a^\dagger a + \frac{\omega_0}{2} \sigma_z + g(\sigma_+ a + \sigma_- a^\dagger) + \lambda(e^{-i\omega t}\sigma_+ + e^{i\omega t}\sigma_-),
\end{equation}
where $\omega$, $\omega_0$ and $\omega_c$ are the frequency of the cavity, atom and classical field, respectively. The operators $\sigma_z$ and $\sigma_\pm$ are defined by $\sigma_z = |e\rangle\langle e| - |g\rangle\langle g|$, $\sigma_+ = |e\rangle\langle g|$, and $\sigma_- = |g\rangle\langle e|$. Here, $a$ and $a^\dagger$ are the annihilation and creation operators of the cavity; $g$ and $\lambda$ are the coupling constants of the interactions of the atom with the cavity and with the classical driving field, respectively. Note that we have set $\hbar = 1$ throughout this paper.

In the rotating reference frame the Hamiltonian of the system is transformed to the Hamiltonian $H_1$ under a unitary transformation $U_1 = \exp(-i\omega t\sigma_z/2)$
\begin{equation}
H_1 = U_1^\dagger H U_1 - iU_1^\dagger \frac{\partial U_1}{\partial t} = H_1^{(1)} + H_1^{(2)},
\end{equation}
with
\begin{align}
H_1^{(1)} &= \omega a^\dagger a + g(e^{i\omega_c t}\sigma_+ a + e^{-i\omega_c t}\sigma_- a^\dagger), \\
H_1^{(2)} &= \Delta_1 \sigma_z + \lambda(\sigma_+ + \sigma_-),
\end{align}
and $\Delta_1 = \omega_0 - \omega_c$. Using the method similar to that used in Ref.\[27\], diagonalizing the Hamiltonian $H_1^{(2)}$, and neglecting the terms which do not conserve energies (rotating wave approximation), we can recast the Hamiltonian $H_1$ as follows
\begin{align}
H_1 &= \omega a^\dagger a + \frac{\Omega_1}{2} \sin \theta (\sigma_+ + \sigma_-) + g \cos \frac{\theta}{2} e^{i\omega_c t} \\
& \quad \times (\sin \frac{\theta}{2} \sigma_z + \cos \frac{\theta}{2} \sigma_+ - \sin \frac{\theta}{2} \sigma_-)a + h.c.,
\end{align}
with $\theta = \arctan(\frac{\Delta_1}{\Omega_1})$. Here $h.c.$ stands for Hermitian conjugation. The Hamiltonian (4) can be diagonalized by a final unitary transformation $U_2$ with $U_2 = \exp \left[ i\omega_c t \sigma_x/2 \right]$. Then, we can rewrite the Hamiltonian of the system
\begin{align}
H_2 &= \omega a^\dagger a + \omega' \sin \theta (\sigma_+ + \sigma_-) + g' \left( -\sin \frac{\theta}{2} \sigma_z \\
& \quad + \cos \frac{\theta}{2} \sigma_+ - \sin \frac{\theta}{2} \sigma_- \right) a + h.c.,
\end{align}
where $\omega' = \sqrt{\Delta_2^2 + 4\lambda^2} + \omega_c$ and $g' = g \cos \frac{\theta}{2}$. It is worth noting that the unitary transformations $U_1$ and $U_2$ are both local unitary transformations. As we known the entanglement of a quantum system does not change under local unitary transformations \[28\]. Thus, the entanglement of the system considered here will not be changed by applying the local unitary transformations $U_1$ and $U_2$.

### B. Entanglement dynamics

We study the entanglement of the system by employing the logarithmic negativity. For a bipartite system described by density matrix $\rho$, the logarithmic negativity is
\begin{equation}
E(\rho) \equiv \log_2 (1 + 2N) = \log_2 ||\rho^T||,
\end{equation}
where $\rho^T$ is the partial transpose of $\rho$ and $||\rho^T||$ stands for the trace norm of $\rho^T$ and $N$ is negativity which is defined by \[22\]
\begin{equation}
N \equiv \frac{||\rho^T|| - 1}{2},
\end{equation}
which is just the absolute value of the sum of the negative eigenvalues of $\rho^T$. The logarithmic negativity has been proven to be an operational good entanglement measure\[29\]. We first assume that the atom is initially prepared in the state $|+\rangle = \cos \frac{\theta}{2} |e\rangle + \sin \frac{\theta}{2} |g\rangle$ and the field is in the Fock state $|n\rangle$. Then, we can find that the state vector at time $t$ is
\begin{align}
|\psi(t)\rangle &= \alpha_n(t)|+\rangle + |\beta_{n+1}(t)|-\rangle, \\
\alpha_n(t) &= e^{i\Delta_2 t/2} \left( \cos (\Omega_n t) - \frac{i\Delta_2}{2\Omega_n} \sin (\Omega_n t) \right), \\
\beta_n(t) &= -ig \cos \frac{\theta}{2} \sqrt{n} e^{-i\Delta_2 t/2} \sin (\Omega_n t) / \Omega_n, \\
\Delta_2 &= \sqrt{(\omega_0 - \omega_c)^2 + 4\lambda^2} + \omega_c - \omega, \\
\Omega_n &= \sqrt{\Delta_2^2 / 4 + (n + 1) (g \cos \theta / 2)^2}, \\
|+\rangle &= \cos \frac{\theta}{2} |e\rangle + \sin \frac{\theta}{2} |g\rangle, \\
|-\rangle &= -\sin \frac{\theta}{2} |e\rangle + \cos \frac{\theta}{2} |g\rangle,
\end{align}
Inserting the state vector into Eq.(7) leads to the logarithmic negativity of the state $|\psi(t)\rangle$
\begin{equation}
E(|\psi\rangle) = \log_2[1 + 2|\alpha_n(t)\beta_{n+1}(t)|].
\end{equation}

Next, we consider the entanglement dynamics of the system when the atom and the field are initially prepared in the state $|+\rangle$ and the thermal state, respectively. The initial density matrix of the system is
\begin{align}
\rho(0) &= |+\rangle \langle +| \otimes \left( \sum_{n=0}^{\infty} p_n |n\rangle \langle n| \right), \\
p_n &= \frac{m^n}{(1 + m)^{n+1}},
\end{align}
where $m = 1/(e^{\beta \omega} - 1)$ is the mean photon number at the inverse temperature $\beta$ and $\omega$ is the frequency of the optical mode. A straightforward calculation shows that
the density matrix at time $t$ is

$$
\rho(t) = \sum_{n=0}^{\infty} \left\{ (p_{n-1}|\beta_n(t)|^2|\rangle \langle n| \right\}^2 + p_n|\alpha_n(t)|^2
\times |+\rangle \langle +| \otimes |n\rangle \langle n| + [p_n|\alpha_n(t)|^2|\beta_{n+1}^*(t)|^2
\times |+\rangle \langle +| \otimes |n+1\rangle \langle n+1| + h.c \right\}
$$

which leads to the logarithmic negativity of the above density matrix

$$
E(\rho) = \log_2 \left[ 1 + \sum_{n=0}^{\infty} (|\xi_n^-| - |\xi_n^+|) \right],
$$

$$
\xi_n^- = \frac{1}{2} \left[ (p_{n-1}|\beta_n|^2 + p_{n+1}|\alpha_{n+1}|^2
- [p_{n-1}|\beta_n|^2 - p_{n+1}|\alpha_{n+1}|^2]^2
+ 4[p_n|\alpha_n|^2|\beta_{n+1}^*|^2]^{1/2} \right].
$$

The logarithmic negativity is plotted as a function of time $t$ in Fig.1 when the field is initially prepared in the vacuum state $|0\rangle$ or in the Fock state $|1\rangle$. For the sake of simplicity, we set $g = 1$ in the numerical calculation. We can see clearly from Fig.1, the atom-field entanglement can be increased by applying the classical driving field.

In Fig.2, the logarithmic negativity $E(\rho)$ is plotted as a function of $t$ with $\omega = 5, \omega_0 = 1$ for $\omega_c = \lambda = 0$ (dotted line) and $\omega_c = \lambda = 2$ (solid line). The mean photon number $\overline{\mu}$ for the upper panel and the lower panel are 0.1 and 0.3, respectively. It is not difficult to find that the entanglement of the two-atom system decreases with the mean photon number $\overline{\mu}$. In other words, the amount of entanglement decreases with the increase of the temperature of the thermal field. However, we can increase the amount of entanglement by controlling the classical field.

III. CONTROLLING ENTANGLEMENT

SUDDEN DEATH

In the previous section, we have investigated the entanglement dynamics of a two-level atom interacting with a cavity field. Recently, Bellomo et al. [22, 23] investigated the influence of the Markovian and Non-Markovian effects on the dynamics of two-qubit entanglement. In this section, we consider a quantum system consisting of two noninteracting atoms each locally interacting with its own vacuum field. The two atoms, which are initially prepared in EWL states, are driven by two independent classical fields additionally. Our purpose here is to investigate the influence of the classical driving field and the purity of the initial states on the entanglement dynamics of the system.

A. Reduced density matrix of the two-atom system

Here, we study the entanglement dynamics of two independent atoms each locally interacting with a vacuum field by using the procedure of Ref.[23]. In addition, each atom is driven by a classical field. In the following, we find the reduced density matrix of the two-atom system. In the basis $|\pm\rangle$, the initial state of the atom interacting with the vacuum field can be cast in the following form

$$
\begin{pmatrix}
\rho_{++}(0) & \rho_{+-}(0) \\
\rho_{-+}(0) & \rho_{--}(0)
\end{pmatrix} \otimes |0\rangle \langle 0|,
$$

where $|+\rangle$ and $|-\rangle$ are defined by Eq.(8). Using the state vector in Eq.(8) and tracing over the freedom of the field, we obtain the single-atom density matrix evolution

$$
\begin{pmatrix}
\rho_{++}(t) & \rho_{+-}(t) \\
\rho_{-+}(t) & \rho_{--}(t)
\end{pmatrix},
$$

FIG. 1: The logarithmic negativity $E(|\psi\rangle)$ is plotted as a function of $t$ with $\omega = 5, \omega_0 = 1$ for $\omega_c = \lambda = 0$ (dotted line) and $\omega_c = \lambda = 2$ (solid line). Upper panel: The field is initially prepared in the vacuum state. Lower panel: The field is initially prepared in the Fock state $|1\rangle$.

FIG. 2: The logarithmic negativity $E(\rho)$ is plotted as a function of $t$ with $\omega = 5, \omega_0 = 1$ for $\omega_c = \lambda = 0$ (dotted line) and $\omega_c = \lambda = 2$ (solid line). Upper panel: The mean photon number of the thermal field is 0.1, i.e., $\overline{\mu} = 0.1$. Lower panel: The mean photon number of the thermal field is 0.3, i.e., $\overline{\mu} = 0.3$. 

where
\[
\rho_{++}(t) = |a_0(t)|^2 \rho_{--}(0),
\]
\[
\rho_{--}(t) = [1 - |a_0(t)|^2] \rho_{++}(0),
\]
\[
\rho_{+-}(t) = a_0(t) \rho_{+}(0) = \rho^*_{--}(t).
\]
By making use of the above equations, we can find the dynamics of the two atoms via a purely algebraic way. Note that this procedure is applicable to arbitrary initial states of the whole system. In the basis \(|1\rangle = |++\rangle, |2\rangle = |+-\rangle, |3\rangle = |-+\rangle, |4\rangle = |--\rangle\), and using Eq.(8) and Eq.(15), the density matrix for the two-atom system is calculated as follows
\[
\rho_{11}(t) = |a_0(t)|^4 \rho_{11}(0),
\]
\[
\rho_{22}(t) = |a_0(t)|^2 [1 - |a_0(t)|^2] \rho_{11}(0) + |a_0(t)|^2 \rho_{22}(0),
\]
\[
\rho_{33}(t) = |a_0(t)|^2 [1 - |a_0(t)|^2] \rho_{11}(0) + |a_0(t)|^2 \rho_{33}(0),
\]
\[
\rho_{44}(t) = [1 - |a_0(t)|^2]^2 \rho_{11}(0) + \rho_{44}(0)
\]
\[
+ [1 - |a_0(t)|^2] [\rho_{22}(0) + \rho_{33}(0)],
\]
\[
\rho_{12}(t) = \rho_{13}(t) = \rho_{23}(t) = \rho_{34}(t) = 0,
\]
\[
\rho_{14}(t) = \alpha_0(t) |\alpha_0(t)|^2 \rho_{11}(0),
\]
\[
\rho_{24}(t) = \alpha_0(t) [1 - |\alpha_0(t)|^2] \rho_{11}(0) + \alpha_0(t) \rho_{24}(0),
\]
\[
\rho_{34}(t) = \alpha_0(t) [1 - |\alpha_0(t)|^2] \rho_{11}(0) + \alpha_0(t) \rho_{34}(0),
\]
with \(\rho_{ij}(t) = \rho^*_{ji}(t)\). We would like to point out that the above procedure allows us to obtain the reduced density matrix of the two-atom system for any initial state.

### B. Extended Werner-like states and logarithmic negativity

We assume the initial states of the two-atom system are the extended Werner-like states defined by
\[
\rho_\Phi(0) = r |\Phi\rangle \langle \Phi| + \frac{1 - r}{4} I,
\]
\[
\rho_\Psi(0) = r |\Psi\rangle \langle \Psi| + \frac{1 - r}{4} I,
\]
\[
|\Phi\rangle = \mu |+-\rangle + \nu |-+\rangle,
\]
\[
|\Psi\rangle = \mu |-+\rangle + \nu |+-\rangle,
\]
where \(r\) is a real number which indicates the purity of initial states, \(I\) is a 4 \times 4 identity matrix, \(\mu\) and \(\nu\) are complex numbers with \(|\mu|^2 + |\nu|^2 = 1\). It is noted that the EWL states belong to the class of the ‘X’ states. Explicitly, if the density matrix of a quantum state is of the form
\[
\begin{pmatrix}
\rho_{11} & 0 & 0 & \rho_{14} \\
0 & \rho_{22} & \rho_{23} & 0 \\
0 & \rho_{23} & \rho_{33} & 0 \\
\rho_{14} & 0 & 0 & \rho_{44}
\end{pmatrix},
\]
then it belongs to the class of the X states.

The EWL states have the following advantages. First, we can easily find that the density matrix at arbitrary time \(t\) is still X structure under the single atom evolution determined by the Hamiltonian \(H_2\) (in the basis \(|1\rangle = |++\rangle, |2\rangle = |+-\rangle, |3\rangle = |-+\rangle, |4\rangle = |--\rangle\) if the initial state is X structure, such as the Bell states and the EWL states. Second, the EWL states allow us to clearly show the influence of the purity and the amount of entanglement of the initial states on the entanglement dynamics simultaneously. The purity of the EWL states are dependent on the parameter \(r\) and the amount of the entanglement of the EWL states are related to \(r\) and \(\mu\). If \(r = 1\) the EWL states reduce to the Bell-like states \(|\Phi\rangle\) and \(|\Psi\rangle\). In the case of \(r = 1, \mu = \nu = 1/\sqrt{2}\) the EWL states become the Bell states while in the case of \(r = 0\) they are the maximally mixed states.

Note that we assume that the initial states of the system are the EWL states. The X structure of the EWL states is maintained during the evolution of the two-atom system. The explicit analytical expression of the logarithmic negativity for the two-atom system can be obtained as
\[
E_X = \log_2(1 + 2N),
\]
\[
N = \max(0, \frac{1}{2} [\sqrt{(\rho_{22} - \rho_{33})^2 + 4|\rho_{14}|^2} - \rho_{22} - \rho_{33}] \\
+ \max(0, \frac{1}{2} [\sqrt{(\rho_{11} - \rho_{44})^2 + 4|\rho_{23}|^2} - \rho_{11} - \rho_{44}]).
\]
We plot the logarithmic negativity as a function of time for different values of the purity of the initial states \(r\) and the strength of the classical field \(\lambda\). In what follows, we denote the entanglement of the two-atom system which is initially prepared in \(\rho_\Phi(0)\) (\(\rho_\Psi(0)\)) by \(E_\Phi\) (\(E_\Psi\)).

The logarithmic negativity \(E_\Phi\) and \(E_\Psi\) are plotted as a function of \(\lambda\) and \(\omega\), with \(\omega = \omega_0 = 1, |\mu|^2 = 1/6, |\nu|^2 = 5/6, r = 2/3, t = 2\) in the upper panel and the lower panel of Fig.3, respectively. From Fig.3, one can clearly see that if the parameters \(\omega_0\) and \(\lambda\) are small, then the logarithmic negativity \(E_\Phi\) and \(E_\Psi\) are zero at time \(t = 2\). The situation is different if the parameters of the classical driving fields increase. When the values of \(\omega_0\) and \(\lambda\) are large enough the two atoms become entangled. The behavior of \(E_\Phi\) and \(E_\Psi\) are different when the classical fields are applied. For example, the area for \(E_\Phi = 0\) is small than that of \(E_\Psi = 0\) which means that the entanglement of the two-atom system is more sensitive with the classical driving fields if they are initially prepared in \(\rho_\Phi\). This feature can be seen more clearly in Fig.4 and Fig.5. It also interesting to note that the entanglement of the two-atom system can be significantly increased by controlling the classical driving fields both for the initial states \(\rho_\Phi\) and \(\rho_\Psi\).

In Fig.4, we plot the logarithmic negativity of the two-atom system as a function of time \(t\) for several values of \(\omega_c\) and \(\lambda\). The upper panel of Fig.4 shows that the ESD occurs without the classical field (see the dotted
line in this panel). After finite dark periods, the entanglement $E_{\Phi}$ revivals completely. The entanglement can be increased by applying the classical field (see the dashed and solid lines of the panel). It is interesting to point out that the ESD disappears when the classical field is strong enough. Comparing the upper panel with the lower panel of Fig. 4, we find that the behavior of entanglement $E_{\Phi}$ and $E_{\Psi}$ are different since the ESD appears in $E_{\Phi}$ with $\omega_c = 0.5, \lambda = 1$ while there is no ESD for $E_{\Psi}$ with the same values of $\omega_c$ and $\lambda$. Besides, the entanglement dark periods of $E_{\Psi}$ is longer than that of $E_{\Phi}$ when all of the other corresponding parameters are the same. In other words, it takes more time for $E_{\Psi}$ to revive the initial entanglement. The parameters $\mu$ and $\nu$ in Fig. 4 are chosen to be $\sqrt{1/6}$ and $\sqrt{5/6}$, respectively.

It is worth noting that different choices of $\mu$ and $\nu$ do not give entanglement dynamics qualitatively different from the case treated here.

Another aspect of interest is how the entanglement dynamics is influenced by the purity of the initial states. In order to show it intuitively, we plot the logarithmic negativity of the two-atom system as a function of $t$ for different values of $r$ in Fig. 5, fixing the parameters $\mu$ and $\nu$. From Fig. 5, one can see that the entanglement dark periods increase with the decrease of the purity of the initial states. The classical driving field again can increase the amount of the entanglement of the system.
IV. CONCLUSIONS

In the present paper, we have considered a quantum system consisting of one two-level atom interacting with a single mode field. The atom is driven by a classical field additionally. The entanglement dynamics is described by the logarithmic negativity. We first investigated the entanglement dynamics of the atom-field system. The field is initially prepared in the Fock state or the thermal state. We find that the atom-field entanglement can be increased by applying the classical driving field.

Then, we consider a quantum system consisting of two noninteracting atoms each locally interacting with its own vacuum field. The two atoms, which are driven by two independent classical fields, are initially prepared in EWL states, i.e., $\rho_\Phi$ and $\rho_\Psi$. We find that classical driving fields can increase the amount of entanglement of the two-atom system. The behavior of entanglement $E_\Phi$ and $E_\Psi$ are different when the classical fields are applied. The entanglement of the two-atom system is more sensitive with the classical driving fields if they are initially prepared in $\rho_\Phi$. It is worth noting that the time of entanglement sudden death can be controlled by the classical driving fields. Specially, the entanglement sudden death phenomenon will disappear if the classical driving fields are strong enough. It is interesting to verify the scheme of our paper in cavity QED experimentally.

ACKNOWLEDGEMENTS

This project was supported by the National Natural Science Foundation of China (Grant no.10774131) and the National Key Project for Fundamental Research of China (Grant no. 2006CB921403).

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