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Decoherence of two-qubit system in a non-Markovian squeezed reservoir*

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The decoherence of two initially entangled qubits coupled with a squeezed vacuum cavity separately is investigated exactly. The results show that, first, in principle, the disentanglement time decreases with the increase of squeeze parameter r, due to the augmenting of average photon number of every mode in the squeezed vacuum cavity. Second, there appear entanglement revivals after the complete disentanglement for the case of even parity initial Bell state, while there occur the entanglement decrease and the entanglement revival before the complete disentanglement for the case of odd parity initial Bell state. The results are quite different from those for the case of qubits in a vacuum cavity.

Keywords: decoherence, squeezed reservoir, non-Markovian, qubit
PACC: 4250, 0365

1. Introduction

In recent years, the entanglement dynamics of qubits coupled with an environment has attracted much attention. Much work has been devoted to the phenomenon termed 'entanglement sudden death' (ESD).[1−15] It is shown that spontaneous disentanglement for a two-level atom model may take only a finite time to be completed, while local decoherence (the normal single-atom transverse and longitudinal decay) takes an infinite time.[3] And, for a non-Markovian reservoir of initially vacuum cavity, there exists a revival phenomenon, which allows the two-qubit entanglement to reappear after a dark period of time, during which the concurrence is zero.[12] In general, the characteristic of the environment plays an important role in the evolving of multi-particle entanglement. Up to now, many typical environments have been investigated, such as, vacuum, squeezed vacuum, multimode vacuum cavity, single mode cavity and so on.

Recently, cavity systems with very strong couplings have been discussed.[16] Generally, in atom-field cavity systems, this ratio of the coupling strength to the atomic transition frequency is typically on the order of $10^{-7} - 10^{-6}$. However, the ratio may become, in the order of magnitude, larger in solid state systems, and the full Hamiltonian, including the virtual processes (counter-rotating terms), must be considered.[17] In this paper, we focus on the decoherence of two qubits strongly coupled with a non-Markovian squeezed vacuum reservoir by using the method given in Ref.[18]. In Section 2, the reduced non-perturbative non-Markovian quantum master equation of atom is derived and its exact solution is obtained. In Section 3, the decoherence of two initially entangled atoms, coupled with two cavities separately, is discussed. The conclusion is given in Section 4.

2. Model and exact solution

2.1. Hamiltonian and non-perturbative master equation

Now we restrict our attention to two noninteracting two-level atoms A and B coupled individually with a cavity field. We first consider the Hamiltonian of the subsystem of a single qubit coupled with its reservoir as

$$H = H_a + H_r + H_{ar}, \tag{1}$$

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where

$$H_a = \omega_0 \frac{\sigma_z}{2}, \tag{2}$$

$$H_r = \sum_k \omega_k a_k^\dagger a_k, \tag{3}$$

$$H_{ar} = (\sigma_+ + \sigma_-) \sum_k g_k \left(a_k^\dagger + a_k \right). \tag{4}$$

Here, ω_0 is the atomic transition frequency between the ground state $|0\rangle$ and excited state $|1\rangle$, $\sigma_z = |1\rangle\langle 1| - |0\rangle\langle 0|$, $\sigma_+ = |1\rangle\langle 0|$ and $\sigma_- = |0\rangle\langle 1|$ are pseudo-spin operators of the atom, the index k labels the field mode of the reservoir with frequency ω_k, $a_k^\dagger$ and a_k are the creation and annihilation operators of mode k respectively, and g_k is the frequency-dependent coupling constant between the transition $|1\rangle - |0\rangle$ and the field mode k.

The reduced non-perturbative non-Markovian quantum master equation of the atom could be obtained through path integral[18]

$$\frac{\partial}{\partial t}\rho_a = -\mathrm{i}\mathcal{L}_a \rho_a - \int_0^t \mathrm{d}s \langle \mathcal{L}_{ar} \mathrm{e}^{-\mathrm{i}\mathcal{L}_0(t-s)} \mathcal{L}_{ar} \mathrm{e}^{-\mathrm{i}\mathcal{L}_0(s-t)} \rangle_r \rho_a, \tag{5}$$

where $\mathcal{L}_0$, $\mathcal{L}_a$ and $\mathcal{L}_{ar}$ are Liouvillian operators defined as

$$\mathcal{L}_0 \rho \equiv [H_a + H_r, \rho],$$
$$\mathcal{L}_a \rho \equiv [H_a, \rho],$$
$$\mathcal{L}_{ar} \rho \equiv [H_{ar}, \rho],$$

and $\langle ... \rangle_r$ denotes partial trace of the reservoir.

Then, we assume that the reservoir is initially in a squeezed vacuum state[19]

$$\rho_r = \prod_k S_k |0_k\rangle\langle 0_k| S_k^\dagger, \tag{6}$$

where

$$S_k = \exp(r\mathrm{e}^{-\mathrm{i}\theta} a_{k_c+k} a_{k_c-k} - r\mathrm{e}^{\mathrm{i}\theta} a_{k_c+k}^\dagger a_{k_c-k}^\dagger) \tag{7}$$

with r and θ being squeeze parameters, and $|0_k\rangle$ is the vacuum state of mode k. The central frequency of squeezing device is $\omega_0 = ck_c$, which corresponds to a multimode squeezed vacuum state with the central frequency equal to the atom transition frequency. And the spectral density of the reservoir is in the Lorentzian form[12,20]

$$J(\omega) = \sum_k g_k^2 \{\delta(\omega - \omega_k) + \delta(\omega + \omega_k)\} = \frac{1}{2\pi} \frac{\lambda\gamma^2}{(\omega - \omega_0)^2 + \gamma^2}, \tag{8}$$

where γ represents the width of the spectral distribution of the reservoir mode, and it is related to the correlation time of the noise induced by the reservoir by $\tau_r = 1/\gamma$. The parameter λ is related to the subsystem-reservoir coupling strength. There are two correlation functions in this model,[21] i.e.

$$\alpha_1 = \int_{-\infty}^{\infty} J(\omega) \mathrm{e}^{-\mathrm{i}(\omega-\omega_0)t} \mathrm{d}\omega = \frac{\gamma\lambda}{2} \mathrm{e}^{-\gamma t}, \tag{9}$$

and

$$\alpha_2 = \int_{-\infty}^{\infty} J(\omega) \mathrm{e}^{\mathrm{i}(\omega+\omega_0)t} \mathrm{d}\omega = \frac{\gamma\lambda}{2} \mathrm{e}^{(-\gamma+\mathrm{i}2\omega_0)t}, \tag{10}$$

where α_1 comes from the rotating-wave interaction and α_2 arises from the counter-rotating wave interaction.

So, the non-perturbative master equation of the subsystem could be derived from Eq.(5) as

$$\frac{\partial}{\partial t}\rho_a = -\Gamma\rho_a + [\varepsilon_0 J_0 + \varepsilon_+ J_+ + \varepsilon_- J_-]\rho_a + [\upsilon_0 K_0 + \upsilon_+ K_+ + \upsilon_- K_-]\rho_a, \tag{11}$$

where J_0, J_+, J_-, K_0, K_+ and K_- are superoperators defined as

$$J_0\rho_a \equiv \left[\frac{\sigma_z}{4}, \rho_a\right],$$
$$J_+\rho_a \equiv \sigma_+\rho_a\sigma_+,$$
$$J_-\rho_a \equiv \sigma_-\rho_a\sigma_-,$$
$$K_0\rho_a \equiv (\sigma_+\sigma_-\rho_a + \rho_a\sigma_+\sigma_- - \rho_a)/2,$$
$$K_+\rho_a \equiv \sigma_+\rho_a\sigma_-,$$
$$K_-\rho_a \equiv \sigma_-\rho_a\sigma_+,$$

and

$$\Gamma = 2M^R f + 2(M^R\alpha^R + M^I\alpha^I) + (2N+1)(f + \alpha^R),$$
$$\varepsilon_0 = -\mathrm{i}2[\omega_0 + 2(M^I\alpha^R - M^I f - M^R\alpha^R) - (2N+1)\alpha^R],$$
$$\varepsilon_+ = 2Mf + 2M^*\alpha + (2N+1)(f+\alpha),$$
$$\varepsilon_- = 2M\alpha^* + 2M^* f + (2N+1)(f+\alpha^*),$$

$$\upsilon_0 = 2(\alpha^R - f),$$
$$\upsilon_+ = 2[M^R f + M^R\alpha^R + M^I\alpha^I + Nf + (N+1)\alpha^R],$$
$$\upsilon_- = 2[M^R f + M^R\alpha^R + M^I\alpha^I + (N+1)f + N\alpha^R],$$

$$\alpha = \frac{\lambda\gamma(1 - \mathrm{e}^{-(\gamma+\mathrm{i}2\omega_0)t})}{2(\gamma + \mathrm{i}2\omega_0)},$$
$$f = \frac{\lambda}{2}[1 - \exp(-\gamma t)],$$

where $N = \sinh^2 r$; $M = -\mathrm{e}^{\mathrm{i}\theta}\sinh r\cosh r$; α^R, α^I, α^* and M^R, M^I, M^* are real part, image part and conjugate of α and M, respectively; N is the average photon number of every mode in the squeezed vacuum cavity. Here, M and M^* represent the phase-dependent correlation between different modes as $\langle b_k b'_k\rangle$ and $\langle b_k^\dagger b_k'^\dagger\rangle$, respectively.[19]

2.2. Exact solution of master equation

The time evolution of density operator in Eq.(11) could be obtained with algebraic approach given in Ref.[22] because the superoperators herein satisfy SU(2) Lie algebraic communication relations, i.e.

$$\begin{aligned}
[J_-,\, J_+]\ \rho_a &= -2J_0\rho_a,\\
[J_0,\, J_\pm]\ \rho_a &= \pm J_\pm\rho_a,\\
[K_-,\, K_+]\ \rho_a &= -2K_0\rho_a,\\
[K_0,\, K_\pm]\ \rho_a &= \pm K_\pm\rho_a,\\
[K_i,\, J_j] &= 0,
\end{aligned} \tag{12}$$

where $i,\, j = 0, \pm$. By directly integrating Eq.(11), the formal solution is obtained as[23]

$$\begin{aligned}
\rho_a(t) = {}& \mathrm{e}^{-\Gamma_k}\hat{T}\mathrm{e}^{\int_0^t \mathrm{d}t(\varepsilon_0 J_0 + + \varepsilon_+ J_+ + \varepsilon_- J_-)}\\
& \times \hat{T}\mathrm{e}^{\int_0^t \mathrm{d}t(v_0 K_0 + + v_+ K_+ + v_- K_-)}\rho_a(0),
\end{aligned} \tag{13}$$

where $\hat{T}$ is time ordering operator and

$$\begin{aligned}
\Gamma_k &= (2M^R + 2N + 1)(F + \tilde{\alpha}^R) + 2M^I\tilde{\alpha}^I),\\
\tilde{\alpha} &= \int_0^t \alpha\mathrm{d}t = \frac{\lambda\gamma\left(t - \dfrac{1 - \mathrm{e}^{-(\gamma + \mathrm{i}2\omega_0)t}}{\gamma + \mathrm{i}2\omega_0}\right)}{2(\gamma + \mathrm{i}2\omega_0)}\\
&\equiv \tilde{\lambda}^R + \mathrm{i}\alpha^I,\\
\tilde{\alpha}^* &= \tilde{\alpha}^R - \mathrm{i}\tilde{\alpha}^I,\\
F(t) &= \lambda\{t - [1 - \exp(-\gamma t)]/\gamma\}/2
\end{aligned}$$

with $\tilde{\alpha}^R$, $\tilde{\alpha}^I$ and $\tilde{\alpha}^*$ being real part, image part and conjugate of $\tilde{\alpha}$, respectively.

The exponential functions of superoperators in Eq.(13) could be disentangled into the following forms:[23]

$$\hat{T}\mathrm{e}^{\int_0^t \mathrm{d}t(\varepsilon_0 J_0 + + \varepsilon_+ J_+ + \varepsilon_- J_-)} = \mathrm{e}^{j_+ J_+}e^{j_0 J_0}\mathrm{e}^{j_- J_-}, \tag{14}$$

and

$$\hat{T}\mathrm{e}^{\int_0^t \mathrm{d}t(v_0 K_0 + + v_+ K_+ + v_- K_-)} = \mathrm{e}^{k_+ K_+}e^{k_0 K_0}\mathrm{e}^{k_- K_-}, \tag{15}$$

where j_+, j_0, j_- and k_+, k_0, k_- satisfy the following differential equations:

$$\dot{X}_+ = \mu_+ - \mu_- X_+^2 + \mu_0 X_+, \tag{16}$$

$$\dot{X}_0 = \mu_0 - 2\mu_- X_+, \tag{17}$$

and

$$\dot{X}_- = \mu_- \exp(X_0), \tag{18}$$

where $\mu = \varepsilon$ for $X = j$; $\mu = \nu$ for $X = k$.

Using the results above, the exact solution of the master equation Eq.(11) is obtained as

$$\rho_a(t) = \mathrm{e}^{-\Gamma_k}\tilde{\rho}(t), \tag{19}$$

where

$$\tilde{\rho}(t) = \begin{pmatrix} l\rho_a^{11}(0) + m\rho_a^{00}(0) & x\rho_a^{10}(0) + y\rho_a^{01}(0)\\ q\rho_a^{01}(0) + r\rho_a^{10}(0) & n\rho_a^{00}(0) + p\rho_a^{11}(0)\end{pmatrix}, \tag{20}$$

$$\begin{aligned}
l &= \mathrm{e}^{k_0/2} + \mathrm{e}^{-k_0/2}k_+k_-,\\
m &= \mathrm{e}^{-k_0/2}k_+, \qquad (21)\\
n &= \mathrm{e}^{-k_0/2}, \quad p = \mathrm{e}^{-k_0/2}k_-, \qquad (22)\\
q &= \mathrm{e}^{-j_0/2}, \quad r = \mathrm{e}^{-j_0/2}j_-, \qquad (23)\\
x &= \mathrm{e}^{j_0/2} + \mathrm{e}^{-j_0/2}j_+j_-, \ \text{and}\\
y &= \mathrm{e}^{-j_0/2}j_+. \qquad (24)
\end{aligned}$$

2.3. Concurrence

In order to investigate the entanglement dynamics of the bipartite system, we use Wootters concurrence.[24] For simplicity, we assume that the two subsystems have the same parameters. The concurrence of the whole system could be obtained as[6]

$$\begin{aligned}
C_\xi &= \max\{0,\, c_1,\, c_2\}, \quad (\xi = \Phi,\, \Psi),\\
c_1 &= 2\mathrm{e}^{-2\Gamma_k}(\sqrt{\rho_{23}\rho_{32}} - \sqrt{\rho_{11}\rho_{44}}),\\
c_2 &= 2\mathrm{e}^{-2\Gamma_k}(\sqrt{\rho_{14}\rho_{41}} - \sqrt{\rho_{22}\rho_{33}}),
\end{aligned} \tag{25}$$

corresponding to the initial states of $|\Phi\rangle = \beta|01\rangle + \eta|10\rangle$ and $|\Psi\rangle = \beta|00\rangle + \eta|11\rangle$, respectively, where β is real and $0 < \beta < 1$, $\eta = |\eta|\mathrm{e}^{\mathrm{i}\varphi}$ and $\beta^2 + |\eta|^2 = 1$. For maximum entanglement Bell state, β is equal to $\sqrt{2}/2$. The reduced joint density matrix of the two atoms, in the standard product basis $\mathcal{B} = \{|1\rangle \equiv |11\rangle, |2\rangle \equiv |10\rangle, |3\rangle \equiv |01\rangle, |4\rangle \equiv |00\rangle\}$, could be obtained by using the method presented in Ref.[12]

$$\rho^{AB} = \mathrm{e}^{-2\Gamma_k}\begin{pmatrix} \rho_{11} & 0 & 0 & \rho_{14}\\ 0 & \rho_{22} & \rho_{23} & 0\\ 0 & \rho_{32} & \rho_{33} & 0\\ \rho_{41} & 0 & 0 & \rho_{44}\end{pmatrix}, \tag{26}$$

where the diagonal elements are

$$\rho_{11} = l^2\rho_{11}(0) + lm\rho_{22}(0) + ml\rho_{33}(0) + m^2\rho_{44}(0),$$
$$\rho_{22} = lp\rho_{11}(0) + lm\rho_{22}(0) + mp\rho_{33}(0) + mn\rho_{44}(0),$$
$$\rho_{33} = lp\rho_{11}(0) + pm\rho_{22}(0) + nl\rho_{33}(0) + nm\rho_{44}(0),$$
$$\rho_{44} = p^2\rho_{11}(0) + pn\rho_{22}(0) + np\rho_{33}(0) + n^2\rho_{44}(0),$$

and the nondiagonal elements are

$$\rho_{14} = x^2\rho_{14}(0) + xy\rho_{23}(0) + yx\rho_{32}(0) + y^2\rho_{41}(0),$$
$$\rho_{23} = xr\rho_{14}(0) + xq\rho_{23}(0) + yr\rho_{32}(0) + yq\rho_{41}(0),$$
$$\rho_{32} = rx\rho_{14}(0) + ry\rho_{23}(0) + qx\rho_{32}(0) + qy\rho_{41}(0),$$
$$\rho_{41} = r^2\rho_{14}(0) + rq\rho_{23}(0) + qr\rho_{32}(0) + q^2\rho_{41}(0).$$

3. Numerical results and discussion

In order to study the effects of non-Markovian squeezed reservoir on the decoherence, we assume that λ is equal to 10γ in Eq.(8), which could be realized in a high-Q cavity.[12]

First, we focus on the effects of squeeze parameters on the decoherence of two qubits with initial maximum Bell entanglement states $|\Phi\rangle$ and $|\Psi\rangle$, respectively.

(A) For $\omega_0 = 10\gamma$ and $\beta = \sqrt{2}/2$, figures 1 and 2 show that for the case of initial maximum Bell state

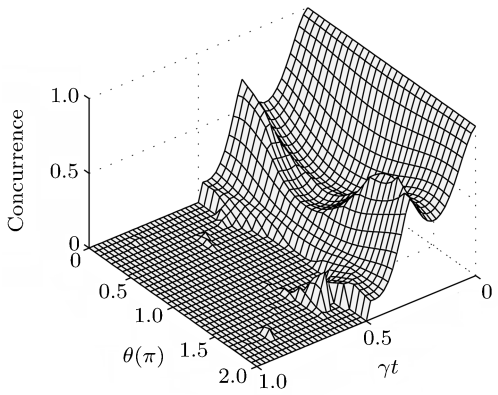


Fig.1. Concurrence C_Φ as a function of γt and θ with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\beta = \sqrt{2}/2$ and $r = 0.2$.

$|\Phi\rangle$, the concurrence first decreases to a certain value and then revives before it vanishes, while that periodically vanishes and revives with its revival amplitude damping for the case of initial maximum Bell state $|\Psi\rangle$. It also reveals that the amplitude and the duration time of entanglement revival are different for the case of $\pi/2 \leq \theta \leq 3\pi/2$ from those for the cases of $0 \leq \theta \leq \pi/2$ and $3\pi/2 \leq \theta \leq 2\pi$.

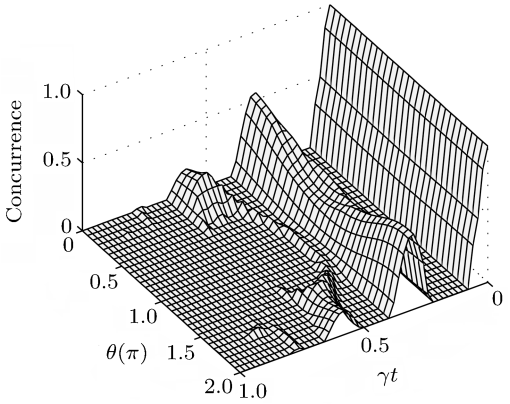


Fig.2. Concurrence C_Ψ as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\beta = \sqrt{2}/2$ and $r = 0.2$.

(B) For $\omega_0 = 10\gamma$ and $\beta = \sqrt{2}/2$, figures 3 and 4 show that for $r \leq 1$, the concurrence first decreases to a certain value and then revives before it vanishes for the case of initial maximum Bell state $|\Phi\rangle$, while that periodically vanishes and revives with amplitude damping for the case of initial maximum Bell state $|\Psi\rangle$. And the concurrence decreases monotonically and vanishes permanently after a short time for $r > 1$. The results exhibit that the disentanglement time decreases with the increase of squeeze parameter r for a fixed value of θ, due to the increasing of average photon number of every mode in the squeezed vacuum cavity.

The above results also reveal that the decoherence behaviour of concurrence is sensitive to squeeze parameters but insensitive to initial state.

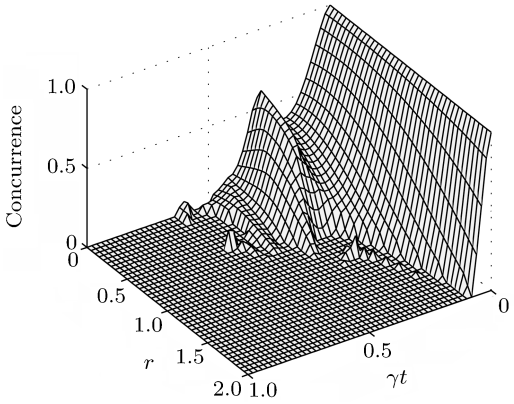


Fig.3. Concurrence C_Φ as a function of γt and r with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\beta = \sqrt{2}/2$ and $\theta = \pi/4$.

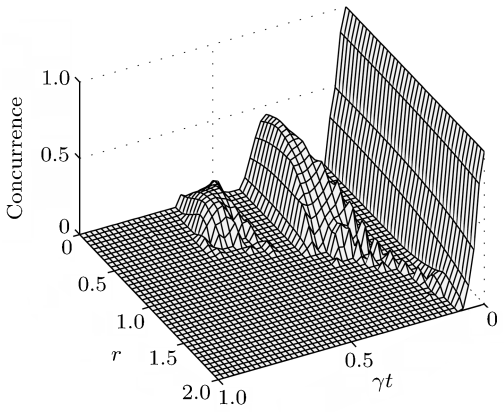


Fig.4. Concurrence C_Ψ as a function of γt and r with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\beta = \sqrt{2}/2$ and $\theta = \pi/4$.

Then, for the fixed squeeze parameters $\theta = \pi/4$ and $r = 0.2$, the decoherences for different initial states are discussed.

(A) The decoherence behaviour of C_Φ is discussed as follows.

(i) From Figs.5–7, we could find that the concurrence C_Φ decreases and revives before it vanishes permanently. With the ratio of the coupling strength to the atomic frequency increasing, the time of entanglement revival decreases because of the increasing of the average value of correlation function α_2, corresponding to the counter-rotating wave term.

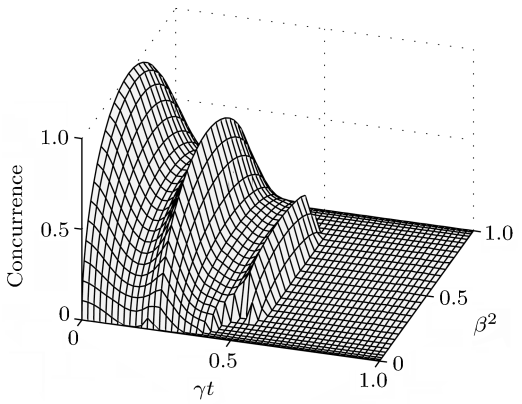


Fig.5. Concurrence C_Φ as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 50\gamma$, $\theta = \pi/4$ and $r = 0.5$.

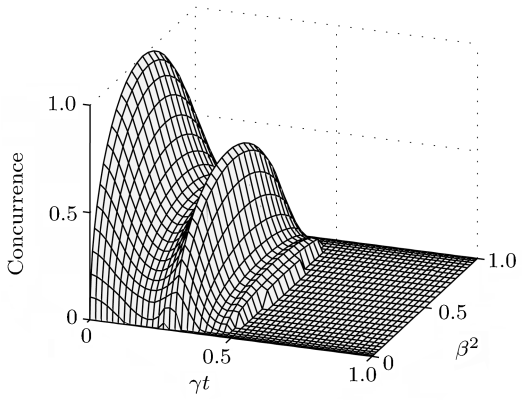


Fig.6. Concurrence C_Φ as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\theta = \pi/4$ and $r = 0.5$.

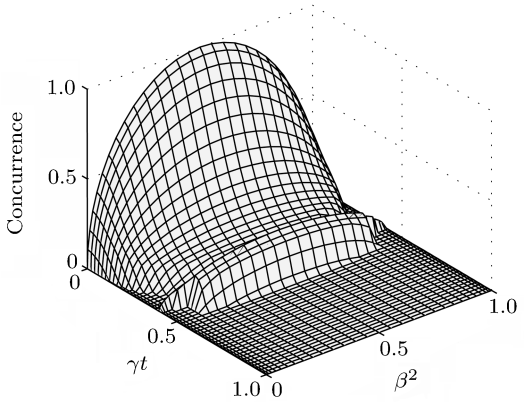


Fig.7. Concurrence C_Φ as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 2\gamma$, $\theta = \pi/4$ and $r = 0.5$.

(ii) For $\omega_0 = 3\gamma$ and $\lambda = 20\gamma$, figure 8 exhibits that the concurrence C_Φ decreases monotonically and vanishes permanently in a short time even in a non-Markovian squeezed reservoir, which arises from the strong interaction between atom and the non-Markovian squeezed reservoir.

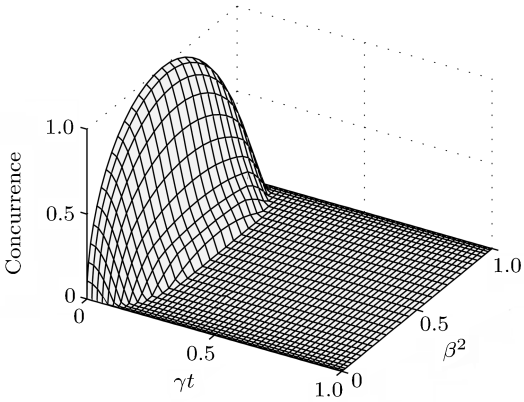


Fig.8. Concurrence C_Φ as a function of γt and β^2 with $\lambda = 20\gamma$, $\omega_0 = 2\gamma$, $\theta = \pi/4$ and $r = 0.5$.

(B) The decoherence behaviour of C_Ψ is discussed as follows.

(i) From Figs.9–11, we find that the concurrence periodically vanishes and revives with its revival amplitude damping. With the ratio of the coupling strength to the atomic frequency increasing, the time of entanglement revival decreases. There appears an entanglement revival after a period of time of disentanglement, which is different from that of the case of

C_{Φ}.

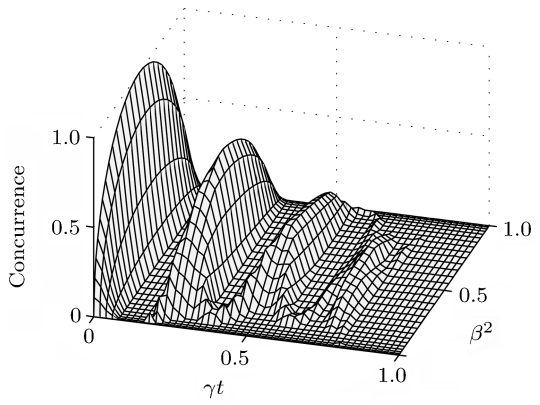


Fig.9. Concurrence C_{Ψ} as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 50\gamma$, $\theta = \pi/4$ and $r = 0.5$.

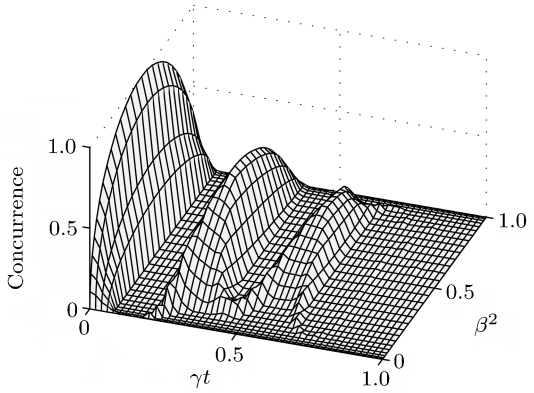


Fig.10. Concurrence C_{Ψ} as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 5\gamma$, $\theta = \pi/4$ and r=0.5.

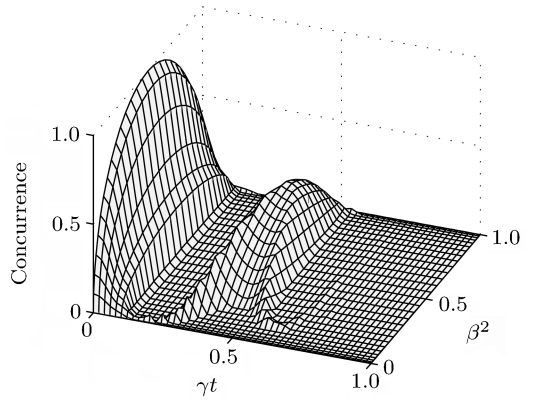


Fig.11. Concurrence C_{Ψ} as a function of γt and β^2 with $\lambda = 10\gamma$, $\omega_0 = 2\gamma$, $\theta = \pi/4$ and r=0.5.

(ii) For $\omega_0 = 3\gamma$ and $\lambda = 20\gamma$, the evolution dynamics of concurrence C_{Ψ} is almost the same as that in Fig.8. Unlike the two cases above, the evolution behaviour of concurrence C_{Ψ} becomes symmetric because of the strong coupling of atom with non-Markovian reservoir and the effect of counter-rotating wave interaction.

The above results reveal that in principle, the decoherence of C_{Φ} is symmetrical with respect to β^2 because of the symmetry of initial state $|\Phi\rangle$, while that is unsymmetrical with respect to β^2 because the initial state $|\Psi\rangle$ is unsymmetrical with respect to β^2. However, the strong coupling and the counter-rotating wave interaction could make the decoherence of C_{Ψ} symmetrical with respect to β^2. With coupling strength enhanced, the memory effect of the counter-rotating wave terms becomes dominant because the average value of correlation function α_2 increases. The results are quite different from those for the case of qubits in a vacuum cavity.[12]

4. Conclusion

The reduced non-perturbative quantum master equation of atom in an initially squeezed vacuum cavity is derived and its exact solution is obtained. The decoherence behaviours of two qubits with squeeze parameters and other parameters are discussed.

The results show that the decoherence behaviour of two qubits in the squeezed reservoir is dependent on the squeeze parameter, the ratio of the coupling strength to the atomic transition frequency and the ratio of the width of reservoir spectral density to the atomic transition frequency. First, in principle, the disentanglement time decreases with the increase of squeeze parameter r, due to the increasing of average photon number of every mode in the squeezed vacuum cavity. Second, there appears the entanglement revival after the complete disentanglement for the case of even parity initial Bell state, while there occur the entanglement decrease and the entanglement revival before the complete disentanglement for the case of odd parity initial Bell state. Third, with the coupling strength enhanced, the time of entanglement revival decreases due to the effect of counter-rotating wave term. Fourth, there exists the entanglement revival due to the memory effect of the non-Markovian squeezed environment.

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