Measurement-induced disturbance and nonequilibrium thermal entanglement in a qutrit–qubit mixed spin XXZ model*

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We investigate the dynamics of nonequilibrium thermal quantum correlation of a qutrit–qubit mixed spin system coupled to two bosonic reservoirs at different temperatures using measurement-induced disturbance. The effects of initial states of the spins and temperatures of the reservoirs on measurement-induced disturbance and entanglement are discussed. The results demonstrate that measurement-induced disturbance is more robust than entanglement against the influence of both these factors and there is no sudden death phenomenon for measurement-induced disturbance. The dependences of steady-state measurement-induced disturbance and entanglement on coupling constant and anisotropy parameter are also studied. Steady-state entanglement vanishes for a ferromagnetic qutrit–qubit model, while steady-state measurement-induced disturbance exists for both the antiferromagnetic and ferromagnetic cases. Appropriately modulating the coupling constant and anisotropy parameter can strengthen quantum correlation.

Keywords: measurement-induced disturbance, entanglement

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1. Introduction

Quantum entanglement plays a valuable role as a physical resource in quantum computation and quantum information processing.[1−4] However, implementation of all quantum information tasks cannot be attributed to entanglement. There are several instances where quantum advantage is reflected in the absence or near absence of entanglement, such as quantum nonlocality without entanglement,[5,6] Grover search,[7] and quantum speedup with separable states.[8,9] It has been proved experimentally that entanglement fails to explain exponential speedups in deterministic quantum computation with one pure qubit model.[10,11] These quantum tasks cannot be simulated by classical methods; they are surely correlations of a quantum nature, rather than entanglement. In the past, entanglement was often identified with quantum correlation. In fact, this identification is appropriate for pure states, but problematic for mixed states, hence the need to investigate the difference and relation between quantum correlation and quantum entanglement.

Ollivier and Zurek first proposed a measure of quantumness of correlations termed quantum discord.[12] Quantum discord is defined as the difference of two classically identical expressions for mutual information in a quantum case. Although the state of a system is separated, this does not guarantee vanishing of the discord. This shows that the absence of entanglement does not imply classicality. Similar methods have been made independently by Henderson and Vedral in their endeavor to separate classical and quantum correlations.[13,14] Oppenheim *et al.*[15] proposed to characterize quantum correlations in terms of entropy production and work extraction. Luo[16] proposed a kind of quantum correlation that is understood in a broad sense as statistical correlation in the quantum scenario, and is something more general than entanglement.[16] These proposals have attracted much attention to the dynamics of quantum correlation.[17−29] Recently, several papers were presented to study equilibrium thermal quantum correlation.[26−29] Two coupled qubits captured

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in different heat baths comprise a system which is not only of theoretical interest but also commonplace in nanophysics. However, there is no report on quantum correlation under nonequilibrium thermal conditions.

Quantum spin systems as a kind of solid-state qubit have played a central role in some applications for quantum communication[30] and quantum computation.[31] In semiconductor quantum dots the transfer of quantum information between nuclear and electronic spins has recently been considered.[32] The nuclear spin is generally weakly coupled to its environment while the electronic spin is strongly coupled to a great variety of degrees of freedom within the solid. This motivates us to investigate the dynamics of quantum correlation of spin chains. Besides the general spin-1/2 system, the properties of the quantum mixed-spin system are also attractive. For an inorganic compound, quasi-one-dimensional bimetallic chains $ACu(pbaOH)(H_2O_3)\cdot 2H_2O$, with each unit cell containing two different spins $(S, 1/2)$, have been synthesized.[33] The compounds of ACu may be considered as a Heisenberg mixed spin $(S, 1/2)$ with antiferromagnetic interactions. As an example, the bimetallic compound chain of NiCu[33] can be considered as a mixed spin $(1, 1/2)$ model. In this paper, we study quantum correlation and entanglement of a system composed of two mixed spins $(1, 1/2)$ coupled to two bosonic baths at different temperatures.

This paper is organized as follows. We briefly introduce quantum correlation via measurement-induced disturbance (MID) in Section 2. In Section 3 we put forward the model of a qutrit–qubit mixed spin coupled to two different bosonic reservoirs, and obtain a solution for the systemic dynamics. We discuss the dynamics of nonequilibrium thermal MID and entanglement and make a comparison in Section 4. We conclude our results in Section 5.

2. Quantum correlation via measurement-induced disturbance

In this section, a brief introduction is given to MID as a measurement of quantum correlations. Given a bipartite state ρ which in general is not classical, any complete local projective measurement $\{\Pi_i^A \otimes \Pi_j^B\}$ will induce a classical state $\Pi(\rho)$,

$$\Pi(\rho) = \sum_{ij} (\Pi_i^A \otimes \Pi_j^B)\rho(\Pi_i^A \otimes \Pi_j^B), \tag{1}$$

where Π_i^A and Π_j^B are complete projective measurements consisting of one-dimensional orthogonal projections for parties A and B. This measurement Π is induced by the spectral resolutions of the reduced states $\rho^A = \sum_i p_i^A \Pi_i^A$ and $\rho^B = \sum_j p_j^B \Pi_j^B$. So $\Pi(\rho)$ is a classical state that is closest to the original state ρ since this kind of measurement can leave the reduced states invariant. Hence any reasonable distance between ρ and $\Pi(\rho)$ can be used to measure the quantum correlation in ρ.

In this paper, we will use the measurement method proposed by Luo, i.e., quantum mutual information difference between ρ and $\Pi(\rho)$ to measure quantum correlation in ρ. The total correlation in any bipartite state ρ can be quantified by the quantum information

$$I(\rho) = H(\rho^A) + H(\rho^B) - H(\rho), \tag{2}$$

with the reduced states $\rho^{A(B)} = \mathrm{Tr}_{B(A)}\rho$ and von Neumann entropy $H(\rho) = -\mathrm{Tr}\rho\log\rho$ for arbitrary density matrix ρ. $\Pi(\rho)$ is a classical state and thus $I(\Pi(\rho))$ quantifies the classical correlations in $\Pi(\rho)$. Hence quantum correlations can be captured by the measurement-induced disturbance, i.e., the difference between $I(\rho)$ and $I(\Pi(\rho))$,

$$Q(\rho) = I(\rho) - I(\Pi(\rho)). \tag{3}$$

3. Basic model

We consider a mixed-spin $(1, 1/2)$ Heisenberg XXZ chain. The Hamiltonian of this system reads

$$\begin{aligned} H_s = {} & \frac{\varepsilon_1}{2}S_A^z + \frac{\varepsilon_2}{2}s_B^z \\ & + J\big(S_A^x s_B^x + S_A^y s_B^y + kS_A^z s_B^z\big), \end{aligned} \tag{4}$$

where ε_1 and ε_2 denote the energy of spins A and B, respectively; J is the coupling constant and k stands for the anisotropy parameter. S_A^η $(\eta = x, y, z)$ is spin-1 operator, s_B^ς $(\varsigma = x, y, z)$ is a Pauli matrix

$$S_A^x = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \quad S_A^y = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & -\mathrm{i} & 0 \\ \mathrm{i} & 0 & -\mathrm{i} \\ 0 & \mathrm{i} & 0 \end{pmatrix},$$

$$S_A^z = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}. \tag{5}$$

Each spin is coupled to a separate bosonic bath with different temperatures and the Hamiltonians of the reservoirs $j = 1, 2$ are as follows:

$$H_{Fj} = \sum_n \omega_{n,j} b_{n,j}^\dagger b_{n,j}. \tag{6}$$

For the sake of simplicity, we suppose the coupling between the levels of the spin-1 to be identical. The interaction between each spin and the corresponding bath is described by

$$\begin{aligned} H_{\rm SF1} &= S_A^+ \sum_n g_n^{(1)} b_{n,1} + S_A^- \sum_n g_n^{(1)*} b_{n,1}^\dagger \\ &= \sum_\nu V_{1,\nu} f_{1,\nu}, \\ H_{\rm SF2} &= S_B^+ \sum_n g_n^{(2)} b_{n,2} + S_B^- \sum_n g_n^{(2)*} b_{n,2}^\dagger \\ &= \sum_\nu V_{2,\nu} f_{2,\nu}, \end{aligned} \tag{7}$$

with $S_A^+ = |2\rangle_{AA}\langle 1| + |1\rangle_{AA}\langle 0|$, $S_B^+ = |1\rangle_{BB}\langle 0|$. Here $b_{n,j}^\dagger$ creates an excitation in mode n of reservoir j with a coupling strength $g_n^{(j)}$, $V_{j,\nu}$ are operators satisfying $[H_s, V_{j,\nu}] = -\omega_\nu V_{j,\nu}$ and the operator $f_{j,\nu}$ are the random operators of reservoirs and act on degrees of freedom for baths. The index ν corresponds to transitions between eigenstates of the system induced by the bath. Such a decomposition is always possible. As an example, we choose[34]

$$V_{j,\nu} = |\psi_i\rangle\langle\psi_i|H_{SFj}|\psi_{i'}\rangle\langle\psi_{i'}|, \tag{8}$$

where $|\psi_i\rangle$ and $|\psi_{i'}\rangle$ are eigenstates of H_s. Then

$$\omega_\nu = \langle\psi_{i'}|H_s|\psi_{i'}\rangle - \langle\psi_i|H_s|\psi_i\rangle. \tag{9}$$

Note that this is the case where the frequency spectrum $\{\omega_\nu\}$ is not degenerate. If there are different transition operators belonging to the same frequency, a sum is needed over all transitions with a fixed energy difference of ω_ν.

The whole system of spin chain with reservoirs is described by the Liouville equation

$$\begin{aligned} \frac{{\rm d}\alpha}{{\rm d}t} &= -{\rm i}[H, \alpha], \\ H &= H_s + \sum_{j=1,2} H_{Fj} + H_{SFj}. \end{aligned} \tag{10}$$

We assume that the bosonic reservoirs recover quickly after being disturbed so that the states of environments remain constant during the evolution. Hence the density operator of the whole system is $\alpha(t) = \rho(t)\rho_1(0)\rho_2(0)$ where each reservoir is described by its own canonical equilibrium density operator $\rho_j = {\rm e}^{-\beta_j H_{Fj}}/{\rm tr}({\rm e}^{-\beta_j H_{Fj}})$ with $\beta_j = 1/K_{\rm B}T_j$ ($K_{\rm B} = 1$) and $\rho(t)$ is the reduced density operator for the quantum system of interest. So in the Born–Markov approximation[34] the master equation for the reduced density matrix for the spin chain is as follows:

$$\frac{{\rm d}\rho}{{\rm d}t} = -{\rm i}[H_{\rm s}, \rho] + \mathscr{L}_1(\rho) + \mathscr{L}_2(\rho). \tag{11}$$

In order to solve this master equation, we go to the basis consisting of the eigenstates $|\psi_i\rangle$ of the system Hamiltonian $H_{\rm s}$ with the corresponding eigenvalues λ_i,

$$\begin{aligned} &|\psi_1\rangle = |00\rangle,\ \lambda_1 = -\frac{1}{2}(\varepsilon_1 + \varepsilon_2 - 2Jk), \\ &|\psi_6\rangle = |21\rangle,\ \lambda_6 = \frac{1}{2}(\varepsilon_1 + \varepsilon_2 + 2Jk), \\ &|\psi_2\rangle = -\sin\theta_1|10\rangle + \cos\theta_1|01\rangle, \\ &\lambda_2 = -\frac{1}{4}(\varepsilon_1 + 2Jk + \Delta^+), \\ &|\psi_3\rangle = \cos\theta_1|10\rangle + \sin\theta_1|01\rangle, \\ &\lambda_3 = -\frac{1}{4}(\varepsilon_1 + 2Jk - \Delta^+), \\ &|\psi_4\rangle = -\sin\theta_2|20\rangle + \cos\theta_2|11\rangle, \\ &\lambda_4 = \frac{1}{4}(\varepsilon_1 - 2Jk - \Delta^-), \\ &|\psi_5\rangle = \cos\theta_2|20\rangle + \sin\theta_2|11\rangle, \\ &\lambda_5 = \frac{1}{4}(\varepsilon_1 - 2Jk + \Delta^-), \end{aligned} \tag{12}$$

with

$$\begin{aligned} \Delta^\pm &= \sqrt{32J^2 + (\varepsilon_1 - 2\varepsilon_2 \pm 2Jk)^2}, \\ \sin\theta_1 &= \frac{4\sqrt{2}J}{\sqrt{32J^2 + (\varepsilon_1 - 2\varepsilon_2 + 2Jk + \Delta^+)^2}}, \\ \sin\theta_2 &= \frac{4\sqrt{2}J}{\sqrt{32J^2 + (\varepsilon_1 - 2\varepsilon_2 - 2Jk + \Delta^-)^2}}. \end{aligned} \tag{13}$$

In this representation the dissipative operators are written as

$$\begin{aligned} &\mathscr{L}_j(\rho) \\ &= \sum_{\nu=1}^{8} J^{(j)}(-\omega_\nu)\big(2V_{j,\nu}\rho V_{j,\nu}^\dagger - \{\rho, V_{j,\nu}^\dagger V_{j,\nu}\}_+\big) \\ &\quad + J^{(j)}(\omega_\nu)\big(2V_{j,\nu}^\dagger \rho V_{j,\nu} - \{\rho, V_{j,\nu}V_{j,\nu}^\dagger\}_+\big), \end{aligned} \tag{14}$$

where the corresponding transition frequencies and transition operators are

$$\begin{aligned}
\omega_1 &= \lambda_6-\lambda_5, \quad V_{j,1}=[\delta_{j,1}\sin\theta_2+\delta_{j,2}\cos\theta_2]|\psi_5\rangle\langle\psi_6|,\\
\omega_2 &= \lambda_6-\lambda_4, \quad V_{j,2}=[\delta_{j,1}\cos\theta_2-\delta_{j,2}\sin\theta_2]|\psi_4\rangle\langle\psi_6|,\\
\omega_3 &= \lambda_5-\lambda_3, \quad V_{j,3}=[\delta_{j,1}\cos(\theta_1+\theta_2)+\delta_{j,2}\sin\theta_2\cos\theta_1]|\psi_3\rangle\langle\psi_5|,\\
\omega_4 &= \lambda_5-\lambda_2, \quad V_{j,4}=[\delta_{j,1}\sin(\theta_2-\theta_1)-\delta_{j,2}\sin\theta_1\sin\theta_2]|\psi_2\rangle\langle\psi_5|,\\
\omega_5 &= \lambda_4-\lambda_3, \quad V_{j,5}=[\delta_{j,1}\sin(\theta_1-\theta_2)+\delta_{j,2}\cos\theta_1\cos\theta_2]|\psi_3\rangle\langle\psi_4|,\\
\omega_6 &= \lambda_4-\lambda_2, \quad V_{j,6}=[\delta_{j,1}\cos(\theta_1+\theta_2)-\delta_{j,2}\sin\theta_1\cos\theta_2]|\psi_2\rangle\langle\psi_4|,\\
\omega_7 &= \lambda_3-\lambda_1, \quad V_{j,7}=[\delta_{j,1}\cos\theta_1+\delta_{j,2}\sin\theta_1]|\psi_1\rangle\langle\psi_3|,\\
\omega_8 &= \lambda_2-\lambda_1, \quad V_{j,8}=[-\delta_{j,1}\sin\theta_1+\delta_{j,2}\cos\theta_1]|\psi_1\rangle\langle\psi_2|. \qquad (15)
\end{aligned}$$

The spectral density in Eq. (14) is

$$J^{(j)}(\omega_\nu) = \int_0^\infty \mathrm{d}\tau\, \mathrm{e}^{\mathrm{i}\omega_{j,\nu}}\langle \mathrm{e}^{-\mathrm{i}H_{Bj}\tau} f_{j,\nu}^\dagger \mathrm{e}^{\mathrm{i}H_{Bj}\tau} f_{j,\nu}\rangle. \quad (16)$$

The bath is considered as an infinite set of harmonic oscillators in this paper, so the spectral density becomes $J^{(j)}(\omega_\nu)=\gamma_j(\omega_\nu)n_j(\omega_\nu)$, with $n_j(\omega_\nu)=(\mathrm{e}^{\beta_j\omega_\nu}-1)^{-1}$, and $J^{(j)}(-\omega_\nu)=\mathrm{e}^{\beta_j\omega_\nu}J^{(j)}(\omega_\nu)$. For the sake of simplicity we regard the coupling constant as independent of frequency, i.e., $\gamma_j(\omega_\nu)=\gamma_j$.

In the eigenstates basis $|\psi_i\rangle$, the differential equation set for the diagonal elements of the density matrix of spin subsystem can be expressed as

$$\begin{aligned}
&\Big(\dot\varrho_{11}(t)\ \dot\varrho_{22}(t)\ \dot\varrho_{33}(t)\ \dot\varrho_{44}(t)\ \dot\varrho_{55}(t)\ \dot\varrho_{66}(t)\Big)^{-1}\\
&= M\Big(\varrho_{11}(t)\ \varrho_{22}(t)\ \varrho_{33}(t)\ \varrho_{44}(t)\ \varrho_{55}(t)\ \varrho_{66}(t)\Big)^{-1}, \qquad (17)
\end{aligned}$$

where M is a 6×6 matrix and can be obtained by substituting matrix forms of Eqs. (4), (14), and (15) in Eq. (11) in the representation of eigenstates. This differential equation set is complex so it is solved numerically. The differential equations for the nondiagonal elements are

$$\dot\varrho_{i,j}(t)=[-\mathrm{i}(\lambda_i-\lambda_j)-c_{i,j}]\varrho_{i,j}(t), \quad (18)$$

where $c_{i,j}$ can be written as the similar progress of finding M. The analytical solution for the nondiagonal elements is easy to obtain:

$$\varrho_{i,j}(t)=\mathrm{e}^{-\mathrm{i}(\lambda_i-\lambda_j)t-c_{i,j}t}\varrho_{i,j}(0). \quad (19)$$

In the basis $\{|21\rangle,|20\rangle,|11\rangle,|10\rangle,|01\rangle,|00\rangle\}$, we consider that the initial state of the system is

$$\begin{aligned}
\rho_0 &= \rho_{11}(0)|21\rangle\langle21|+\rho_{22}(0)|20\rangle\langle20|\\
&\quad+\rho_{33}(0)|11\rangle\langle11|+\rho_{44}(0)|10\rangle\langle10|\\
&\quad+\rho_{55}(0)|01\rangle\langle01|+\rho_{66}(0)|00\rangle\langle00|\\
&\quad+\rho_{23}(0)|20\rangle\langle11|+\rho_{32}(0)|11\rangle\langle20|\\
&\quad+\rho_{45}(0)|10\rangle\langle01|+\rho_{54}(0)|01\rangle\langle10|. \quad (20)
\end{aligned}$$

We can easily obtain the expressions of ρ_t^A,ρ_t^B, and $\Pi(\rho_t)$ by simple calculation, and find $\rho_t^A=[\Pi(\rho_t)]^A,\rho_t^B=[\Pi(\rho_t)]^B$.

According to the MID measurement method in Section 2, we can obtain the quantum correlation in ρ_t

$$\begin{aligned}
Q(\rho_t) &= I(\rho_t)-I(\Pi(\rho_t))\\
&= H(\Pi(\rho_t))-H(\rho_t)\\
&= \lambda_1\log\lambda_1+\lambda_2\log\lambda_2+\lambda_3\log\lambda_3+\lambda_4\log\lambda_4\\
&\quad-\rho_{22}(t)\log\rho_{22}(t)-\rho_{33}(t)\log\rho_{33}(t)\\
&\quad-\rho_{44}(t)\log\rho_{44}(t)-\rho_{55}(t)\log\rho_{55}(t), \quad (21)
\end{aligned}$$

with

$$\begin{aligned}
\lambda_{1,2} &= \frac{1}{2}\Big\{\rho_{22}(t)+\rho_{33}(t)\\
&\quad\pm\sqrt{\big[\rho_{22}(t)-\rho_{33}(t)\big]^2+4|\rho_{23}(t)|^2}\Big\},\\
\lambda_{3,4} &= \frac{1}{2}\Big\{\rho_{44}(t)+\rho_{55}(t)\\
&\quad\pm\sqrt{\big[\rho_{44}(t)-\rho_{55}(t)\big]^2+4|\rho_{45}(t)|^2}\Big\}. \quad (22)
\end{aligned}$$

In order to compare MID with entanglement of two spins in the same scale, we choose the logarithmic negativity.[35] The logarithmic negativity of the two spins is

$$LN(\rho)=\log\|\rho^{\Gamma_A}\|_1=\log[2N(\rho)+1], \quad (23)$$

with the negativity[36]

$$\begin{aligned}
N(\rho) &= \max\{0,-\mu_1\}+\max\{0,-\mu_2\},\\
\mu_1 &= \frac{1}{2}\Big\{\rho_{11}(t)+\rho_{44}(t)
\end{aligned}$$

$$- \sqrt{\left[\rho_{11}(t) - \rho_{44}(t)\right]^2 + 4|\rho_{23}(t)|^2}\Big\},$$

$$\mu_2 = \frac{1}{2}\Big\{\rho_{33}(t) + \rho_{66}(t) - \sqrt{\left[\rho_{33}(t) - \rho_{66}(t)\right]^2 + 4|\rho_{45}(t)|^2}\Big\}. \quad (24)$$

4. Results and discussion

We studied the dynamics of entanglement and MID for two spins in two baths at different temperatures. From Fig. 1, we can observe that the entanglement and MID damply oscillate in a short time. This indicates that the environments cause the entanglement and MID of two spins to decay and can also entangle the separate states at the same time. The state of the system then converges to a steady state for a long time. The steady-state entanglements quickly disappear with increasing temperature. However, the steady-state MID is not greatly affected by temperature and only declines slightly with increasing temperature. MID is more robust to resisting the influence of temperature and there is no sudden death phenomenon.

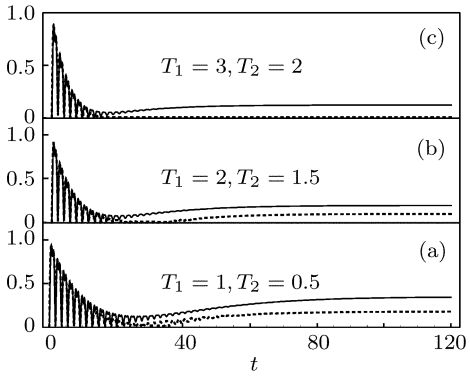


Fig. 1. Logarithmic negativity and measurement-induced disturbance of the mixed spin chain (1,1/2) for the initial state $\rho_0 = |20\rangle\langle 20|$. Here $\gamma_1 = \gamma_2 = 0.02$, $J = 1$, $k = 1.4$, $\varepsilon_1 = 7$, $\varepsilon_2 = 4$, t in units of $(50\gamma)^{-1}$. Solid line: measurement-induced disturbance. Dashed line: logarithmic negativity.

The dynamic evolutions of entanglement and MID for the different initial states are shown in Fig. 2. First we can see that there is no sudden death of MID for these initial states. When the two spins are initially in $|21\rangle$, there is no oscillation and then the state of the system stabilizes directly. However, entanglement and MID can cause oscillation for some initial states as shown in Figs. 1 and 2. This phenomenon shows the competition between unitary evolution and dissipative dynamics.[37] Although the initial states of the system are different, the entanglement and MID reach the same value with increasing time, that is, the steady-state entanglement and the steady-state MID are independent of the initial states of the system. The steady-state MID is larger than the steady value of entanglement.

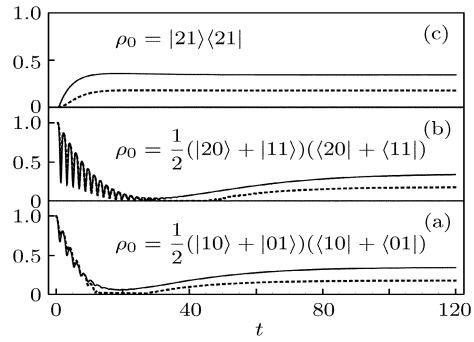


Fig. 2. Logarithmic negativity and measurement-induced disturbance of the mixed spin chain (1,1/2) for the different initial states. $T_1 = 1$, $T_2 = 0.5$, $\gamma_1 = \gamma_2 = 0.02$, $J = 1$, $k = 1.4$, $\varepsilon_1 = 7$, $\varepsilon_2 = 4$. Solid line: measurement-induced disturbance. Dashed line: logarithmic negativity.

From Figs. 1 and 2, we can easily see that the states of the system will reach a steady state with increasing time. We can obtain the stationary solution by setting $\dot{\varrho}_t = 0$ and the normalization condition $\mathrm{Tr}\varrho_t = 1$.

Evolutions of the steady-state entanglement and the steady-state MID are shown in Fig. 3. With increasing bath temperature the steady-state entanglement abruptly vanishes, while the MID decreases only asymptotically. This indicates the robustness of steady-state MID against steady-state entanglement with temperature increase. Steady-state entanglement and MID can reach their maxima when the temperatures of the two baths take a certain value. The maximal value of steady-state MID is larger than that of entanglement.

Figure 4 shows that there is no entanglement for a ferromagnetic ($J < 0$) qutrit–qubit model, while MID exists for both the antiferromagnetic ($J > 0$) and ferromagnetic ($J < 0$) cases. At the same time, this figure denotes that MID is more robust than entanglement against high temperature. After the entanglement reaches its minimum, it vanishes at all times. However, MID decreases with temperature and starts increasing again after reaching a minimum. In Fig. 4(b) MID vanishes at the critical point $J = 0$. This signals the presence of the quantum phase transition at finite temperature T but steady-state entanglement cannot signal this. Figure 4(b) also shows that MID of the steady state increases with the

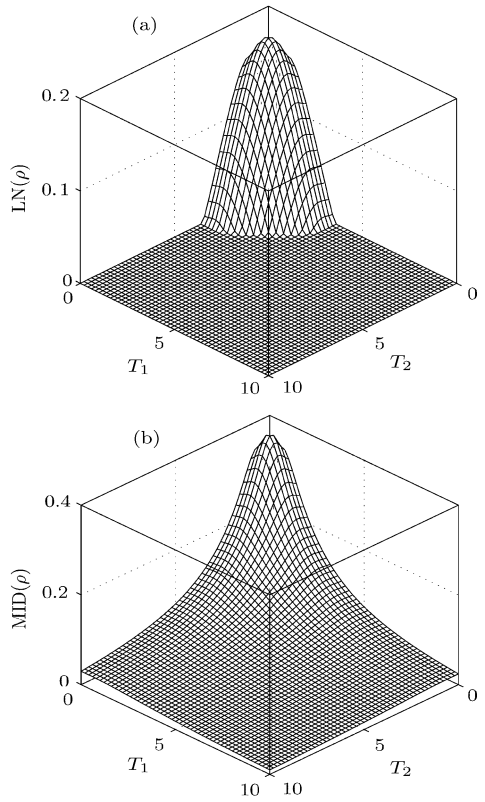


Fig. 3. Variation of steady-state logarithmic negativity and measurement-induced disturbance for the mixed spin chain (1,1/2) with respect to two bath temperatures T_1 and T_2, $\gamma_1 = \gamma_2 = 0.02$, $J = 1$, $k = 1.4$, $\varepsilon_1 = 7$, $\varepsilon_2 = 4$.

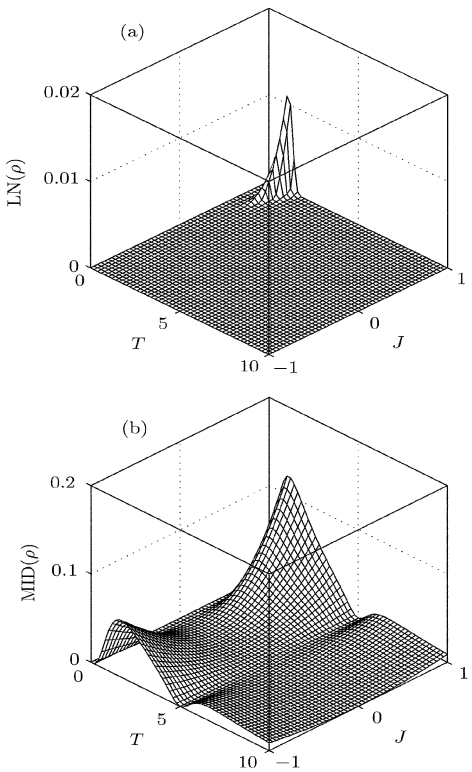


Fig. 4. Steady-state logarithmic negativity and measurement-induced disturbance of the mixed spin chain (1,1/2) as functions of T and J. $T_1 = T_2 = T$, $\gamma_1 = \gamma_2 = 0.02$, $\varepsilon_1 = 7$, $\varepsilon_2 = 4$, $k = 0.5$.

increase of the absolute value of J: the larger this is, the larger the steady-state MID. When the coupling constant J between two spins is certain, the steady-state MID reaches its maximum at an appropriate temperature and then decreases with increasing temperature. After it drops to its minimum, MID recovers.

In Fig. 5, when the value of k changes from negative to positive, steady-state entanglement vanishes in the region of $k < 0$, while MID increases continuously. The larger the values of k that can be obtained in experiments, the larger the stationary value for MID. A large spin–spin stationary MID is very useful for applications to quantum information protocols.

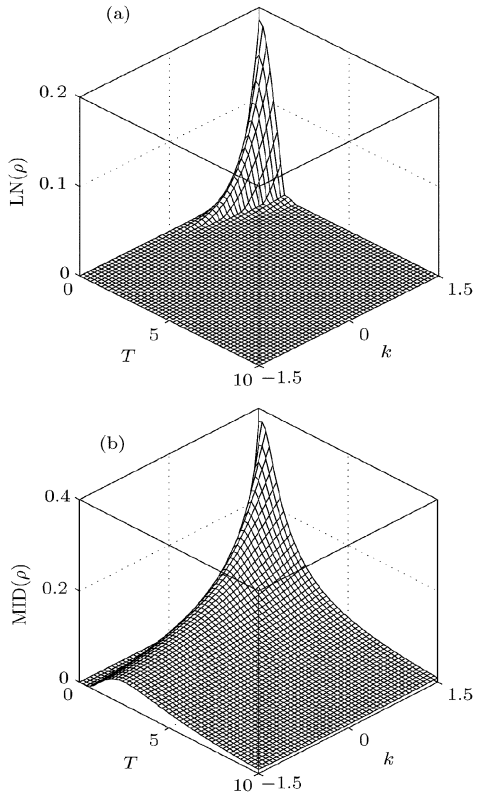


Fig. 5. Steady-state logarithmic negativity and measurement-induced disturbance of the mixed spin chain (1,1/2) as functions of T and k. $T_1 = T_2 = T$, $\gamma_1 = \gamma_2 = 0.02$, $\varepsilon_1 = 7$, $\varepsilon_2 = 4$, $J = 1$.

5. Conclusion

We have investigated the quantum correlations of a qutrit–qubit mixed spin chain coupled to two bosonic reservoirs at different temperatures using measurement-induced disturbance. The influences of the temperatures of the two baths and different initial states of two spins on the entanglement and measurement-induced disturbance are considered. In

this model the system reaches a steady state in the long term. We also studied the effects of the coupling constant J and anisotropy parameter k on steady-state measurement-induced disturbance and entanglement. There is no entanglement for a ferromagnetic qutrit–qubit model, while measurement-induced disturbance exists for both the antiferromagnetic and ferromagnetic cases. Steady-state measurement-induced disturbance and entanglement can be strengthened by tuning J and k. A large spin–spin stationary quantum correlation should be very useful for implementation in quantum information tasks. We also compare steady-state measurement-induced disturbance with steady-state entanglement under the same conditions. The results demonstrate that measurement-induced disturbance is more robust than entanglement as the temperature increases. When entanglement vanishes, measurement-induced disturbance is still present. This shows that separable states can possess quantum correlation. In a word, quantum entanglement is not the same as quantum correlation, only a special kind of it.

References

[1] Nielsen M A and Chuang I L 2000 *Quantum Computation and Quantum Information* (Cambridge: Cambridge University Press)

[2] Tao Y J, Tian D P, Hu M L and Qin M 2008 *Chin. Phys.* B **17** 624

[3] Liang M L and Yuan B 2010 *Chin. Phys.* B **19** 094206

[4] Zheng S B 2010 *Chin. Phys.* B **19** 064204

[5] Horodecki M, Horodecki P, Horodecki R, Oppenheim J, Sen A, Sen U and Synak-Radtke B 2005 *Phys. Rev.* A **71** 062307

[6] Niset J and Cerf N J 2006 *Phys. Rev.* A **74** 052103

[7] Ahn J, Weinacht T C and Bucksbaum P H 2000 *Science* **287** 463

[8] Braunstein S L, Caves C M, Jozsa R, Linden N, Popescu S and Schack R 1999 *Phys. Rev. Lett.* **83** 1054

[9] Meyer D A 2000 *Phys. Rev. Lett.* **85** 2014

[10] Knill E and Laflamme R 1998 *Phys. Rev. Lett.* **81** 5672

[11] Lanyon B P, Barbicri M, Almeida M P and White A G 2008 *Phys. Rev. Lett.* **101** 200501

[12] Ollivier H and Zurek W H 2002 *Phys. Rev. Lett.* **88** 017901

[13] Henderson L and Vedral V 2001 *J. Phys.* A: *Math. Gen.* **34** 6899

[14] Vedral V 2003 *Phys. Rev. Lett.* **90** 050401

[15] Oppenheim J, Horodecki M, Horodecki P and Horodecki R 2002 *Phys. Rev. Lett.* **89** 180402

[16] Luo S L 2008 *Phys. Rev.* A **77** 022301

[17] Datta A, Shaji A and Caves C M 2008 *Phys. Rev. Lett.* **100** 050502

[18] Luo S L 2008 *Phys. Rev.* A **77** 042303

[19] Ali M, Rau A R P and Alber G 2010 *Phys. Rev.* A **81** 042105

[20] Wang B, Xu Z Y, Chen Z Q and Feng M 2010 *Phys. Rev.* A **81** 014101

[21] Lu X M, Ma J, Xi Z and Wang X 2011 *Phys. Rev.* A **83** 012327

[22] Werlang T, Souza S, Fanchini F F and Villas Boas C J 2009 *Phys. Rev.* A **80** 024103

[23] Ali M 2010 *J. Phys.* A: *Math. Theor.* **43** 495303

[24] Wang L C, Yan J Y and Yi X X 2011 *Chin. Phys.* B **20** 040305

[25] Chen Q Y, Fang M F, Xiao X and Zhou X F 2011 *Chin. Phys.* B **20** 050302

[26] Zhang G F, Jiang Z T and Abliz A 2011 *Ann. Phys.* **326** 867

[27] Zhang G F, Hou Y C and Ji A L 2011 *Solid State Commun.* **151** 790

[28] Werlang T and Rigolin G 2010 *Phys. Rev.* A **81** 044101

[29] Wang Q, Liao J Q and Zeng H S 2010 *Chin. Phys.* B **19** 100311

[30] Bose S 2003 *Phys. Rev. Lett.* **91** 207901

[31] Kane B E 1998 *Nature* **393** 133

[32] Imamoglu A, Knill E, Tian L and Zoller P 2003 *Phys. Rev. Lett.* **91** 017402

[33] van Koningsbruggen P J, Kahn O, Nakatani K, Pei Y, Renard J P, Drillon M and Legoll P 1990 *Inorg. Chem.* **29** 3325

[34] Goldman M 2001 *J. Magn. Reson.* **149** 160

[35] Plenio M B 2005 *Phys. Rev. Lett.* **95** 090503

[36] Życzkowski K, Horodecki P, Sanpera A and Lewenstein M 1998 *Phys. Rev.* A **58** 883

[37] Sinaysky I, Petruccione F and Burgarth D 2008 *Phys. Rev.* A **78** 062301