Evolution of Bell-nonlocality of two cavity fields in the double Jaynes-Cummings model

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The Bell-nonlocality of two initially entangled macroscopic fields in the double Jaynes-Cummings model is investigated. Moreover, the process by which detuning between the atomic transition frequency and the field frequency affects the evolution of the Bell-nonlocality of two macroscopic fields is studied. The effect of the disparity between the two coupling strengths is discussed.

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Quantum entanglement is one of the most striking features of quantum systems. It is an indispensable element of quantum computation and quantum information processing[1−4]. Over the past few years, there has been considerable interest in the dynamics of entanglement based on the different quantum systems[5]. Yu *et al.* demonstrated how two entangled qubits which individually interact with vacuum noise can become completely disentangled in a finite time. This interesting phenomenon is termed as entanglement sudden death (ESD)[6] and studied further by the authors in other cases[7−9]. Moreover, ESD has been theoretically studied in a variety of systems[10−14] and demonstrated in quantum optics experiments[15,16].

Nonlocality of quantum state is the other basic characteristic of quantum mechanics. In recent work, Jaeger *et al.* pointed out that the Bell-nonlocality has new features which are similar to ESD[17]. They showed that multipartite Bell-inequality violations could disappear suddenly in a finite time, a phenomenon currently referred to as the Bell-nonlocality sudden death. Yang *et al.* investigated tripartite nonlocality evolution in two-atom Tavis–Cummings model with consideration to cavity decay[18]. Previous studies on the evolution of entanglement and nonlocality have been limited to microscopic systems because of the uncertainty of proper quantification of entanglement in relation to macroscopic systems. More recently, Luo *et al.* extended the analysis of the evolution of nonlocality to the case of two macroscopic fields that interact with a resonant atom[19]. They referred to the study of Chen *et al.* on the formalism of Bell CHSH's inequality[20], which is based on pseudospin operators, to study the evolution of nonlocality of continuous-variable states. They showed that the collapse and revival of the Bell-nonlocality are similar to the collapse and revival of the atomic population inversion of the Jaynes-Cummings model (JCM). Liao *et al.* investigated quantum nonlocality dynamics of the interaction between a three-level atom and a class of two-mode non-classical states (including the entangled coherent state, the pair coherent state, and the two-mode squeezed state) in the non-degenerate two-photon JCM[21].

We consider the most general case in the double JCM, which is different from the situation in Ref. [19]. We study how the mean photon number of the macroscopic field, as well as the detuning between the atomic transition frequency and the field frequency, affects the evolution of two initial entangled macroscopic fields. Furthermore, we discuss the effects of asymmetric couplings on the evolution of Bell-nonlocality between the two fields.

We consider two spatially separated JCM cavity fields that individually interact with a two-level atom (Fig. 1). It should be emphasized that there is no interaction and communication between the two cavities or between the two atoms. Moreover, we suppose that the two cavities are treated as lossless. In a frame rotating with the photon frequency and assuming $\hbar = 1$, the Hamiltonian equation for the whole system is given by

$$\begin{aligned} H_{\rm tot} =& \frac{1}{2}(\Delta\sigma_z^A + \Delta\sigma_z^B) + g_1(a_1^+\sigma_-^A + a_1^-\sigma_+^A) \\ &+ g_2(a_2^+\sigma_-^B + a_2^-\sigma_+^B), \end{aligned} \tag{1}$$

where Δ is the detuning between the atomic transition frequency ω and the field frequency; σ_+^j, σ_-^j, σ_z^j $(j = A, B)$ denote the rising, lowering, and population inversion operators for the jth atom; a_i^+ and a_i^- $(i = 1, 2)$ are the creation and annihilation operators for the ith cavity mode, respectively; and g_1 and g_2 are the atom-cavity coupling strengths.

For the macroscopic fields, we adhere to the Chen *et al.* version of Bell CHSH's inequality[20], which is based on pseudospin operators. It is defined as

$$s_z = \sum_{n=0}^{\infty} [|2n+1\rangle\langle 2n+1| - |2n\rangle\langle 2n|], \tag{2}$$

$$s_- = \sum_{n=0}^{\infty} |2n\rangle\langle 2n+1| = (s_+)^+, \tag{3}$$

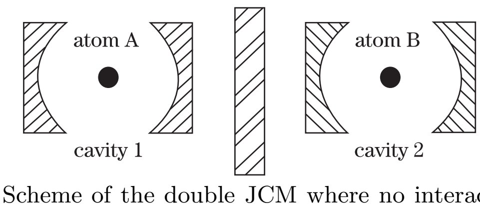

Fig. 1. Scheme of the double JCM where no interaction between the two cavities exists.

$$s_x \pm \mathrm{i}s_y = 2s_\pm, \tag{4}$$

$$\hat{a}\cdot\hat{s} = s_z\cos\theta_a + \sin\theta_a(\mathrm{e}^{\mathrm{i}\varphi_a}s_- + \mathrm{e}^{-\mathrm{i}\varphi_a}s_+), \tag{5}$$

where θ_a and φ_a are the polar and azimuthal angles of the unit vector $\hat{a}$, respectively. The Bell-CHSH operator based on the pseudospin operators is defined as[20]

$$B_{\mathrm{CHSH}} = (\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2) + (\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2) + (\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2) - (\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2), \tag{6}$$

where $\hat{a}$ and $\hat{b}$ are the unit vectors and the subscripts 1 and 2 are the labels of the two continuous-variable systems. In the continuous-variable system, $\langle B_{\mathrm{CHSH}}\rangle$ is the expectation value of B_{CHSH} for a given quantum state, which is bounded by $2\sqrt{2}$. If $|\langle B_{\mathrm{CHSH}}\rangle| = 2\sqrt{2}$ for a particular state, the Bell-CHSH inequality is indicated to be maximally violated by the given quantum state.

Supposing that the atoms are prepared in the ground state $|g\rangle$ and the two cavities are initially in the entangled coherent state, then

$$|\psi_{\mathrm{field}}\rangle = \frac{1}{\sqrt{N_a}}(\cos\,\theta|\alpha\rangle_1|-\alpha\rangle_2 - \sin\,\theta|\alpha\rangle_1|-\alpha\rangle_2),$$
$$N_\alpha = 1 - \sin\,2\theta\mathrm{e}^{-4|\alpha|^2}. \tag{7}$$

After an interaction time t, the evolution of the system is given by

$$\begin{aligned}|\Psi(t)\rangle = \frac{1}{\sqrt{N_\alpha}}\Big\{\cos\theta\sum_{n=0}^{\infty}\Big[C_n\Big(\cos\frac{\Omega_{n-1}}{2}t + \mathrm{i}\frac{\Delta}{\Omega_{n-1}}\\ \sin\frac{\Omega_{n-1}}{2}t\Big)|g\rangle_{\mathrm{A}}|n\rangle_1 - \mathrm{i}C_{n+1}\frac{2g\sqrt{n+1}}{\Omega_n}\\ \sin\frac{\Omega_n}{2}t|e\rangle_{\mathrm{A}}|n\rangle_1\Big]\sum_{m=0}^{\infty}\Big[C'_m\Big(\cos\frac{\Omega'_{m-1}}{2}t + \mathrm{i}\frac{\Delta}{\Omega'_{m-1}}\\ \sin\frac{\Omega'_{m-1}}{2}t\Big)|g\rangle_{\mathrm{B}}|m\rangle_2 - \mathrm{i}C'_{m+1}\frac{2g\sqrt{m+1}}{\Omega'_m}\sin\frac{\Omega'_m}{2}t|e\rangle_{\mathrm{B}}|\\ m\rangle_2\Big] - \sin\theta\sum_{n=0}^{\infty}\Big[C'_n\Big(\cos\frac{\Omega_{n-1}}{2}t + \mathrm{i}\frac{\Delta}{\Omega_{n-1}}\sin\frac{\Omega_{n-1}}{2}t\Big)|\\ g\rangle_{\mathrm{A}}|n\rangle_1 - \mathrm{i}C'_{n+1}\frac{2g\sqrt{n+1}}{\Omega_n}\sin\frac{\Omega_n}{2}t|e\rangle_{\mathrm{A}}|n\rangle_1\Big]\quad\sum_{m=0}^{\infty}\\ \Big[C_m\Big(\cos\frac{\Omega'_{m-1}}{2}t + \mathrm{i}\frac{\Delta}{\Omega'_{m-1}}\sin\frac{\Omega'_{m-1}}{2}t\Big)|g\rangle_{\mathrm{B}}|m\rangle_2\\ - \mathrm{i}C_{m+1}\frac{2g\sqrt{m+1}}{\Omega'_m}\sin\frac{\Omega'_m}{2}t|e\rangle_{\mathrm{B}}|m\rangle_2\Big]\Big\},\end{aligned} \tag{8}$$

where $C_n = \frac{\alpha^n}{\sqrt{n!}}\mathrm{e}^{-|\alpha|^2/2}$, $C'_m = \frac{(-\alpha)^m}{\sqrt{m!}}\mathrm{e}^{-|\alpha|^2/2}$, $\Omega_n = \sqrt{\Delta^2 + 4g_1^2(n+1)}$, $\Omega'_m = \sqrt{\Delta^2 + 4g_2^2(m+1)}$.

In the present paper, the expectation value of B_{CHSH} for the state described by Eq. (8) is given by

$$\begin{aligned}\langle B_{\mathrm{CHSH}}\rangle &= \langle\Psi(t)|B_{\mathrm{CHSH}}|\Psi(t)\rangle\\ &= \langle(\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2)\rangle + \langle(\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2)\rangle\\ &+ \langle(\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2)\rangle - \langle(\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2)\rangle.\end{aligned} \tag{9}$$

The evolution of $\langle B_{\mathrm{CHSH}}\rangle$ of the two fields that individually interact with a resonant atom has been discussed in Ref. [19]. In this section, we study how the atom-field detuning affects the evolution of the Bell-nonlocality of two initially entangled coherent fields. Thus, assuming that the two cavities are in the maximally entangled state, we set $g_1 = g_2 = g, \theta = \frac{\pi}{4}$ and $\Delta = ag$ in Eq. (9). Subsequently, we choose the azimuthal angles $\varphi_a = \varphi_b = \varphi_{a'} = \varphi_{b'} = 0$ and the polar angles $\theta_a = 0, \theta_{a'} = \frac{\pi}{2}, \theta_b = -\theta_{b'}$. From Eqs. (2)–(5), (8), and (9), the result can be obtained as

$$\langle(\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2)\rangle = \langle(\hat{a}\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2)\rangle = \frac{4\cos\theta_b\mathrm{e}^{-2|\alpha|^2}}{N_\alpha}(N_1 - N_4)(N_3 - N_2), \tag{10}$$

$$\langle(\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}\cdot\hat{s}_2)\rangle = -\langle(\hat{a}'\cdot\hat{s}_1)\otimes(\hat{b}'\cdot\hat{s}_2)\rangle = \frac{-4\sin\theta_b\mathrm{e}^{-2|\alpha|^2}}{N_\alpha}[(N_5 + N_7)^2 + N_6^2], \tag{11}$$

where
$$N_\alpha = 1 - \mathrm{e}^{-4|\alpha|^2}, \tag{12}$$

$$N_1 = \sum_{n=0}^{\infty}\frac{\alpha^{4n+2}}{(2n+1)!}\Big\{\cos^2\Big(\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big) + \Big[\frac{a}{\sqrt{a^2+4(2n+1)}}\sin\Big(\frac{g\sqrt{a^2+4(2n+1)}}{2}\Big)\Big]^2\Big\}, \tag{13}$$

$$N_2 = \sum_{n=0}^{\infty}\frac{\alpha^{4n}}{(2n)!}\Big\{\cos^2\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big) + \Big[\frac{a}{\sqrt{a^2+8n}}\sin\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big)\Big]^2\Big\}, \tag{14}$$

$$N_3 = \sum_{n=0}^{\infty}\frac{\alpha^{4n+4}}{(2n+2)!}\Big[\frac{2\sqrt{2n+2}}{\sqrt{a^2+4(2n+2)}}\Big]^2 \sin^2\Big[\frac{g\sqrt{a^2+4(2n+2)}}{2}t\Big], \tag{15}$$

$$N_4 = \sum_{n=0}^{\infty}\frac{\alpha^{4n+2}}{(2n+1)!}\Big[\frac{2\sqrt{2n+1}}{\sqrt{a^2+4(2n+1)}}\Big]^2 \sin^2\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big], \tag{16}$$

$$N_5 = \sum_{n=0}^{\infty}\frac{\alpha^{4n+1}}{\sqrt{(2n)!(2n+1)!}}\Big\{\frac{a^2}{\sqrt{a^2+8n}\sqrt{a^2+4(2n+1)}} \sin\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big]\sin\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big) + \cos\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big]\cos\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big)\Big\}, \tag{17}$$

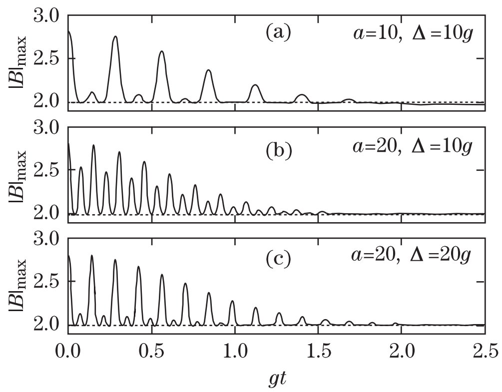

Fig. 2. Evolution of the $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ in a short period of time.

$$N_6=\sum_{n=0}^{\infty}\frac{\alpha^{4n+1}}{\sqrt{(2n)!(2n+1)!}}\Big\{\frac{a^2}{\sqrt{a^2+4(2n+1)}}\cos\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big)\sin\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big]-\frac{a^2}{\sqrt{a^2+8n}}\cos\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big]\sin\Big(\frac{g\sqrt{a^2+8n}}{2}t\Big)\Big\},\qquad(18)$$

$$N_7=\sum_{n=0}^{\infty}\frac{\alpha^{4n+3}}{\sqrt{(2n+1)!(2n+2)!}}\frac{4\sqrt{(2n+1)(2n+2)}}{\sqrt{[a^2+4(2n+1)][a^2+4(2n+2)]}}\cdot\sin\Big[\frac{g\sqrt{a^2+4(2n+1)}}{2}t\Big]\sin\Big[\frac{g\sqrt{a^2+4(2n+2)}}{2}t\Big].\qquad(19)$$

Consequently, $\langle B_{\rm CHSH}\rangle$ can be given by

$$\langle B_{\rm CHSH}\rangle=\frac{8{\rm e}^{-2|\alpha|^2}}{N_\alpha}(N_1-N_4)(N_3-N_2)(\cos\theta_b+U\sin\theta_b),\qquad(20)$$

where $U=\frac{N_6^2-(N_5+N_7)^2}{(N_1-N_4)(N_3-N_2)}$. From Eq. (20), the maximum value of $\langle B_{\rm CHSH}\rangle$ can be easily derived as

$$|\langle B_{\rm CHSH}\rangle|_{\rm max}=\frac{8{\rm e}^{-2|\alpha|^2}}{N_\alpha}\,|(N_1-N_4)(N_3-N_2)|\,\sqrt{1+U^2}.\qquad(21)$$

The maximum value of $|\langle B_{\rm CHSH}\rangle|$ for the state with different atom-field detuning in a short period of time is plotted in Fig. 2. The function of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ facilitates rapid oscillations in a short period of time. The Bell-nonlocality violations of the two fields will disappear in a finite time and will be recovered soon because the atoms interact with the cavity fields. Afterwards, the atomic dipole stores the phase information of the field, eventually destroying the coherence of the fields. If the atoms and the cavities are maximally entangled, then the two fields are in mixed states after tracing over the atoms. This will cause the disappearance of the nonlocality between these two different field phase elements. In Fig. 2, the oscillation frequency of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ decreases when the atom-field detuning increases. Moreover, the oscillation frequency of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ increases when the number of the mean photon increases.

The oscillation period of the entanglement between the two fields is $T=\pi/g\sqrt{a^2+4|\alpha|^2}$, which is equal to the oscillation period of entanglement between the atomic system and field system[22]. Furthermore, the oscillation period is half of the corresponding period of the atomic population inversion because we used squares of summations rather than summations for the original inversion calculations in Eq. (21)[23]. Thus, we set $\alpha=20,\Delta=10g,\alpha=20,\Delta=20g$, and $\alpha=10,\Delta=10g$, and substitute them into $T=\pi/g\sqrt{a^2+4|\alpha|^2}$ resulting in oscillation periods of $0.0762/g$, $0.0702/g$, and $0.1404/g$, respectively. These numerical results are in good agreement with those shown in Fig. 2.

The peak value of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ decreases gradually in relation to time because the atom and field cannot be completely disentangled at time τ. With the increase in τ, the remaining entanglement between the atoms and cavities increases, while the entanglement between the two fields decreases. Moreover, as the mean photon number of the fields and the atom-field detuning increase, the remaining entanglement between the atoms and fields decreases. Thus, the decrease of the oscillation peaks becomes slower.

Furthermore, an interesting phenomenon in the non-resonant case exists (Fig. 2). The function of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ has two group peaks and the odd peaks are higher than the even ones because the coefficients of the summations in Eq. (21) involve $\cos^2 t\sqrt{\Delta^2+8g^2n}/2$, $\cos^2 t\sqrt{\Delta^2+4g^2(2n+1)}/2$, $\sin^2 t\sqrt{\Delta^2+8g^2n}/2$, and $\sin^2 t\sqrt{\Delta^2+4g^2(2n+1)}/2$. The values of the odd and even peaks are determined by the coefficients of the cosine and sine summations, respectively. When the detuning is not excessively large, a slight difference between the cosine and sine summations coefficients is produced. The difference increases when the atom-field detuning increases. The even peaks will become lower as the atom-field detuning increases until it disappears.

Compared with the theoretical analysis in Ref. [23], the summations in Eq. (21) are similar to those for the atomic population inversion in the original discussion of quantum revivals. Thus, we expect $|\langle B_{\rm CHSH}\rangle|_{\rm max}$ to have a similar revival behavior. The evolution of $|\langle B_{\rm CHSH}\rangle|_{\rm max}$

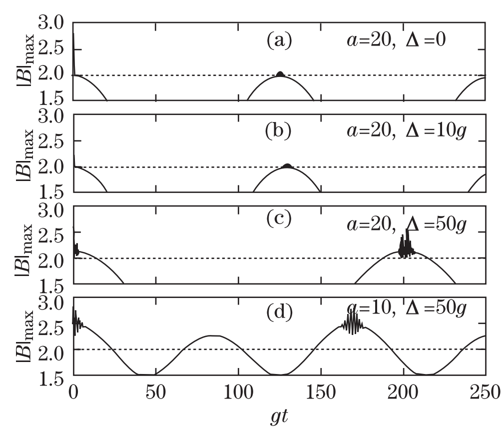

Fig. 3. The maximum value of $|\langle B_{\rm CHSH}\rangle|$ versus time for maximally entangled coherent states.

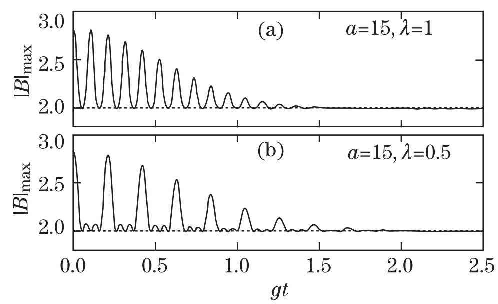

Fig. 4. Maxima of $|\langle B_{\rm CHSH}\rangle|$ versus time for entangled coherent states, where $g_2 = \lambda g_1$.

under different conditions for a long period of time is shown in Fig. 3 where the Bell-nonlocality violations revival after an interval is clarified. The interval time is equal to that of the atomic population inversion, which is given by $t_R = \pi\left(\Delta^2 + 4g^2|\alpha|^2\right)^{1/2}/g^2$[23]. Substituting $\alpha = 20, \Delta = 0, 10g, 50g$; $\alpha = 10$, and $\Delta = 50g$ into the above equation, the results are in good agreement with those shown in Fig. 3. The revival interval will increase with the mean photon number of the fields and the increase of the detuning between the atoms and fields. Figure 3 shows that the function of $|\langle B_{\rm CHSH}\rangle|_{\max}$ creates rapid oscillations with period $T = \pi/g\sqrt{a^2 + 4|\alpha|^2}$ in the same regions that are close to the revival time.

In Figs. 3(b), (c), and (d), the value of $|\langle B_{\rm CHSH}\rangle|_{\max}$ which is close to the revival regions increases when the atom-field detuning increases. An increasing atom-field detuning means that the energy exchange between the atomic system and the cavity decreases. This indicates the decreasing influence of the atoms on the fields. Therefore, the entangled information of the cavity fields leaking into the atoms is reduced when the atom-field detuning increases after one or several full periods of the energy exchange[24].

Guo *et al.*[25] investigated the effects of asymmetric couplings between atoms and a cavity field on entanglement dynamics. In some initial states, the inhomogeneous couplings not only induced but also enhanced the entanglement. In the current study, we discuss the effects of asymmetric couplings on the evolution of Bell-nonlocality of the fields that individually interact with a resonant atom. By setting $g_1 = g$, $g_2 = \lambda g_1$, and $\Delta = 0$ the evolution of Bell-nonlocality of the fields in this case can be derived from Eqs. (8) and (9). In Fig. 4, $|\langle B_{\rm CHSH}\rangle|_{\max}$ is plotted as a function of gt. Comparing $\theta = \pi/4$ with the equal couplings case[19], the function of $|\langle B_{\rm CHSH}\rangle|_{\max}$ creates rapid oscillations in a short period of time for $g_2 = 0.5g_1$. In Fig. 4(b), two low oscillation peaks between the two adjacent high peaks are shown.

In conclusion, we investigate the Bell-nonlocality of two initially entangled macroscopic fields in the double JCM. It should be emphasized that there is no interaction between the two cavities and that the cavities in our double JCM are lossless. In our model, the two spatially separated cavities individually interact with a non-resonant atom. Similar to the resonant case[19], collapse and revival of Bell-nonlocality of the two macroscopic fields are observed. The revival time ($t = \pi(\Delta^2 + 4g^2|\alpha|^2)^{1/2}/g^2$) and the oscillation period ($T = \pi/g\sqrt{a^2 + 4|\alpha|^2}$) increase when the detuning between the atomic transition frequency and the field frequency increase. The collapse of the Bell-nonlocality is caused by the atom and field that cannot be completely disentangled. The disappearance of the field phase information stored in the two atoms leads to the recovery of Bell-nonlocality. In addition, we discuss how the disparity between the two coupling strengths affects the evolution of Bell-nonlocality of the fields.

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