Deterministic generation of Greenberger–Horne–Zeilinger and W states for three distant atoms via adiabatic passage*

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We propose two schemes for generating Greenberger–Horne–Zeilinger and W states of three distant atoms. In the present schemes, the atoms are individually trapped in three spatially separated optical cavities coupled by two optical fibres. Performing an adiabatic passage along dark states, the population of cavities and fibres excited is negligible under certain conditions. In addition, the spontaneous decay of atoms is also efficiently suppressed based on our proposals. Furthermore, the discussion about the entanglement fidelity is given and we point out that our schemes work robustly with small fluctuations of experimental parameters.

Keywords: Greenberger–Horne–Zeilinger state, W state, adiabatic passage

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1. Introduction

Entanglement, as one of the most unique and fundamental features of quantum mechanics, plays a significant role in the quantum information field. It is an essential ingredient for testing fundamental quantum-mechanical principles, but it also provides promising and wide practical applications in quantum information processing.[1−11] To date, the entanglement of bipartite systems has been well understood and generated in various quantum systems,[12−16] while entanglement in multi-qubits is more complicated and fascinating than in two qubits. For the three-particle case, there are two classes of tripartite entanglement states: the Greenberger–Horne–Zeilinger (GHZ) state[17,18] and the W state.[19] These two classes are inequivalent in the sense that they cannot be converted to each other under stochastic local operations and classical communication. It is known that the GHZ state gives the maximal violation of correlation function Bell inequalities[20,21] and the W state has the maximal bipartite entanglement among all three qubit states. These two classes have found applications not only in testing quantum nonlocality without inequality, but also in realizing quantum information processing, including multiqubit quantum teleportation,[22−24] quantum dense coding,[25,26] quantum cloning machine,[27−30] quantum secret sharing[31−33] and so on.

For decades, various quantum systems have been suggested as being possible candidates for performing quantum state preparation and quantum information processing and a large number of theoretical and experimental schemes were proposed for generating GHZ state and W state of photons,[34−38] atoms,[39−49] ions,[50,51] superconducting qubits,[52,53] semiconductor quantum dots,[54] etc. Among them, the cavity quantum electrodynamics (QED) is one of the most popular systems to engineer quantum entanglement.[55−60] It provides a promising avenue for achieving varieties of novel ideas, by utilizing the strong coupling between a high-fitness optical cavity and atoms. A quantum electrodynamic model involving a large-detuned single-mode cavity field and N identical two-level atoms was investigated by Guo *et al.*,[42] and its application to the generation of mul-

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tiparticle entanglement states was analysed. Subsequently, Zou *et al.*[43,44] proposed several schemes to generate GHZ state and W state for three distant atoms trapped separately in leaky cavities, utilizing linear optical elements and high efficient multiphoton coincidence detection. More recently, Gonta *et al.*[45] proposed two schemes for deterministically realizing four-partite entangled GHZ and W states by using Rydberg atoms within the framework of cavity QED. Their schemes are based on the resonant interaction of a chain of atoms with a bimodal cavity that supports two independent modes of the photon field.

On the other hand, the adiabatic passage[61−64] has become a useful tool for realizing quantum information processing. It has several advantages in quantum state engineering since the decoherence due to spontaneous emission from the excited states can be compressed and it is robust to some experimental parameter errors. Lately, based on the techniques of adiabatic passage, a large number of schemes have been proposed for the entanglement generation.[65−75] For example, the proposals for atom–photon, atom–atom and photon–photon entanglement preparation were investigated in Ref. [65] by fractional adiabatic passage. Moreover, several robust schemes for deterministically generating or transferring three-dimensional entanglement of distant atoms were investigated via adiabatic passage, in which the atomic spontaneous emission and the decays of fibre and cavities were successfully avoided.[70,76]

In view of all the above mentioned, in the present paper, we propose two schemes for deterministically generating three-particle GHZ state and W state for three atoms trapped in distant cavities connected by optical fibres. The present schemes are different from the previous schemes of generating GHZ state and W state. In Ref. [71], schemes for generating GHZ state of three distant atoms was proposed. The evolution through adiabatic passage was first performed between the first and second atoms, then the second atom and the third atom underwent transition through another adiabatic passage to obtain the GHZ state. While in our scheme, we just utilize one dark state for preparing the GHZ state. In Ref. [69], schemes for generating W states of spatially separated atoms were proposed based on the stimulated Raman adiabatic passage (STIRAP) or fractional STIRAP techniques. In their schemes, the atoms fall through the cavity with a certain velocity and then meet the laser beams. While in our schemes, the atoms are trapped in distant cavities for the following preparation of entangled states. Based on the adiabatic passage along dark states, our schemes have the following characteristics. (i) The atomic excited states are unpopulated during the whole interaction process, hence the proposals are robust to the atomic spontaneous emission. (ii) Under certain conditions, the probabilities with which cavity modes and fibre modes are excited can be negligible, thus the decays of fibre and cavities are effectively eliminated. (iii) Taking advantage of adiabatic passage, the schemes are insensitive to small fluctuations of experimental parameters.

2. Model and Hamiltonian

The framework of our scheme is shown in Fig. 1. Three atoms are trapped in three distant optical cavities (cavities 1–3), which are connected by two fibres (fibres 1 and 2). Figures 1(a) and 1(b) display the configurations of atoms for preparing GHZ state and W state, respectively. In what follows we will present our models in detail and obtain the effective Hamiltonian.

2.1. GHZ state

For implementing our proposal to achieve three-partite GHZ state, we choose cavity 1 and cavity 3 as single mode optical cavities whose cavity modes are denoted as $a_{1\mathrm{l}}$ and $a_{3\mathrm{r}}$, respectively and cavity 2 are selected as double-mode optical resonant cavity with cavity modes $a_{2\mathrm{l}}$ and $a_{2\mathrm{r}}$. Two Λ-type atoms are individually trapped in cavities 1 and 3. The states $|0\rangle$ and $|1\rangle$ are two ground levels of the atom and $|e\rangle$ is an excited level. In cavities 1 and 3, the atomic transitions $|1\rangle_1 \leftrightarrow |e\rangle_1$ and $|0\rangle_3 \leftrightarrow |e\rangle_3$ are dispersively driven by two classical fields with time-dependent Rabi frequencies $\Omega_1(t)$ and $\Omega_3(t)$, respectively. The cavity modes with coupling strengths $g_{1\mathrm{l}}$ and $g_{3\mathrm{r}}$ interact with the atomic transitions $|0\rangle_1 \leftrightarrow |e\rangle_1$ and $|1\rangle_3 \leftrightarrow |e\rangle_3$, respectively. A tripod-type four-level atom is trapped in cavity 2 and the atomic transition $|a\rangle_2 \leftrightarrow |e\rangle_2$ is dispersively driven by a classical field with the time-dependent Rabi frequency $\Omega_2(t)$. The transition $|e\rangle_2 \leftrightarrow |1\rangle_2(|0\rangle_2)$ is coupled with cavity mode $a_{2\mathrm{l}}$ ($a_{2\mathrm{r}}$) with left-circular (right-circular) polarization whose coupling strength is $g_{2\mathrm{l}}$ ($g_{2\mathrm{r}}$). The energy differences between the atomic transitions $|0\rangle_i \leftrightarrow |e\rangle_i$ ($i = 1, 2$), $|1\rangle_j \leftrightarrow |e\rangle_j$ ($j = 2, 3$) and the relevant cavity mode, which are denoted as Δ_1, Δ_2 and Δ_3, are of the frequency detuning and they satisfy the corresponding two-photon resonance conditions as shown in Fig. 1(a). Without loss of generality, we assume that all the Rabi frequencies and coupling rates are real in this paper.

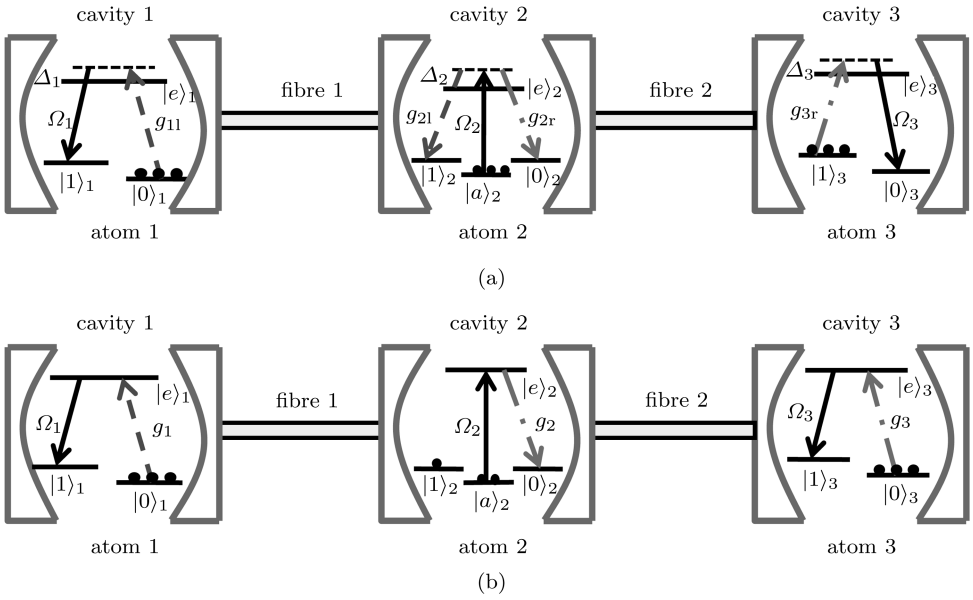


Fig. 1. Schematics of our model system and the level configuration of each atom for realizing (a) GHZ state, (b) W state.

Then, in the interaction picture, in the dipole and rotating wave approximation, the interaction Hamiltonian of the atom–cavity system can be written as ($\hbar = 1$)[77,78]

$$\begin{aligned} H_{\rm GHZ}^{\rm ac} = & \sum_{n=1,2,3} (-\Delta_n |e\rangle_n\langle e|) + (\Omega_1(t)|e\rangle_1\langle 1| \\ & + \Omega_2(t)|e\rangle_2\langle a| + \Omega_3(t)|e\rangle_3\langle 0| + g_{1l}a_{1l}|e\rangle_1\langle 0| \\ & + g_{2l}a_{2l}|e\rangle_2\langle 1| + g_{2r}a_{2r}|e\rangle_2\langle 0| \\ & + g_{3r}a_{3r}|e\rangle_3\langle 1| + \text{H.c.}), \end{aligned} \tag{1}$$

where H.c. denotes the Hermitian conjugate. Next, we assume that the large-detuning condition is satisfied, i.e., $|\Delta_n| \gg |\Omega_n|, |g_{1l}|, |g_{2l}|, |g_{2r}|, |g_{3r}|$ $(n = 1, 2, 3)$. In such a case, the excited state $|e\rangle_n$ of an atom is only virtually excited in the process of atom–field interaction and can be eliminated adiabatically. By applying the standard quantum optical technique,[79] the effective Hamiltonian describing the atom–cavity system reads

$$\begin{aligned} & H_{\rm GHZ}^{\rm eff1} \\ = & \frac{\Omega_1(t)^2}{\Delta_1}|1\rangle_1\langle 1| + \frac{\Omega_2(t)^2}{\Delta_2}|a\rangle_2\langle a| + \frac{\Omega_3(t)^2}{\Delta_3}|0\rangle_3\langle 0| \\ & + \frac{g_{1l}^2}{\Delta_1}a_{1l}^\dagger a_{1l}|0\rangle_1\langle 0| + \frac{g_{2l}^2}{\Delta_2}a_{2l}^\dagger a_{2l}|1\rangle_2\langle 1| \\ & + \frac{g_{2r}^2}{\Delta_2}a_{2r}^\dagger a_{2r}|0\rangle_2\langle 0| + \frac{g_{3r}^2}{\Delta_3}a_{3r}^\dagger a_{3r}|1\rangle_3\langle 1| \\ & + \left(\frac{g_{1l}\Omega_1(t)}{\Delta_1}a_{1l}|1\rangle_1\langle 0| + \frac{g_{2l}\Omega_2(t)}{\Delta_2}a_{2l}^\dagger|1\rangle_2\langle a| \right. \\ & \left. + \frac{g_{2r}\Omega_2(t)}{\Delta_2}a_{2r}^\dagger|0\rangle_2\langle a| + \frac{g_{3r}\Omega_3(t)}{\Delta_3}a_{3r}|0\rangle_3\langle 1| \right. \\ & \left. + \text{H.c.}\right), \end{aligned} \tag{2}$$

where the first seven terms on the right-hand side in Hamiltonian (2) represent cavity and classical field induced atom level shifts, respectively. By using other two lasers coupling the corresponding atom levels , these energy level shifts can be straightforwardly compensated.[80] The Hamiltonian (2) can be further simplified into

$$\begin{aligned} & H_{\rm GHZ}^{\rm ac} \\ = & \Omega_{\rm e1}(t)a_{1l}|1\rangle_1\langle 0| + \Omega_{\rm e2}^{\rm l}(t)a_{2l}^\dagger|1\rangle_2\langle a| \\ & + \Omega_{\rm e2}^{\rm r}(t)a_{2r}^\dagger|0\rangle_2\langle a| + \Omega_{\rm e3}(t)a_{3r}|0\rangle_3\langle 1| + \text{H.c.}, \end{aligned} \tag{3}$$

where $\Omega_{\rm e1}(t) = g_{1l}\Omega_1(t)/\Delta_1$, $\Omega_{\rm e2}^{\rm l}(t) = g_{2l}\Omega_2(t)/\Delta_2$, $\Omega_{\rm e2}^{\rm r}(t) = g_{2r}\Omega_2(t)/\Delta_2$ and $\Omega_{\rm e3}(t) = g_{3r}\Omega_3(t)/\Delta_3$ are the effective Rabi frequencies for the corresponding Raman transitions $|0\rangle_1 \to |e\rangle_1 \to |1\rangle_1$, $|a\rangle_2 \to |e\rangle_2 \to |1\rangle_2$, $|a\rangle_2 \to |e\rangle_2 \to |0\rangle_2$ and $|1\rangle_3 \to |e\rangle_3 \to |0\rangle_3$, respectively.

In our scheme, three cavities are connected via two optical fibres, in the short fibre limit $(2L\bar{\eta})/(2\pi c) \ll 1$, only one (resonant) mode of a fibre will interact with the cavity mode, where L is the length of the fibre, $\bar{\eta}$ is the decay rate of the cavities fields into a continuum of fibre modes. The interaction Hamiltonian $H_{\rm GHZ}^{\rm cf}$ can be approximated to[81]

$$\begin{aligned} & H_{\rm GHZ}^{\rm cf} \\ = & \eta_{\rm A}[b_{\rm A}(a_{1l}^\dagger + a_{2l}^\dagger)] + \eta_{\rm B}[b_{\rm B}(a_{2r}^\dagger + a_{3r}^\dagger)] + \text{H.c.} \end{aligned} \tag{4}$$

Here, we should point out that the polarizations of fibre modes $b_{\rm A}^k$ and $b_{\rm B}^k$ have been chosen as the left and the right circulares, respectively. $\eta_{\rm A(B)}^k$ is the coupling strength between fibre mode k and the corresponding cavity mode, and the relative phase φ between two cavity modes, which originates from the propagation of the field through the fibre of length, has been used in the cavity modes.

Lastly, in the interaction picture, the total Hamiltonian of this cavity–fibre–cavity system can be written as

$$\begin{aligned} & H_{\rm GHZ} \\ = \; & H_{\rm GHZ}^{\rm ac} + H_{\rm GHZ}^{\rm cf} \\ = \; & \Omega_{\rm e1}(t)a_{1\rm l}|1\rangle_1\langle 0| + \Omega_{\rm e2}^{\rm l}(t)a_{2\rm l}^\dagger|1\rangle_2\langle a| \\ & + \Omega_{\rm e2}^{\rm r}(t)a_{2\rm r}^\dagger|0\rangle_2\langle a| + \Omega_{\rm e3}(t)a_{3\rm r}|0\rangle_3\langle 1| \\ & + \eta_{\rm A}[b_{\rm A}(a_{1\rm l}^\dagger + a_{2\rm l}^\dagger)] + \eta_{\rm B}[b_{\rm B}(a_{2\rm r}^\dagger + a_{3\rm r}^\dagger)] + {\rm H.c.} \quad (5) \end{aligned}$$

2.2. W state

For implementing our proposal to achieve a three-partite W state, three atoms are individually trapped in three single-mode optical cavities separately, which are connected by two optical fibres. The ith atomic transition $|0\rangle_i \rightarrow |e\rangle_i$ (i = 1–3) is resonantly coupled with the ith cavity mode a_i. Three classical fields drive the transitions $|1\rangle_1 \rightarrow |e\rangle_1$, $|a\rangle_2 \rightarrow |e\rangle_2$ and $|1\rangle_3 \rightarrow |e\rangle_3$ resonantly with coupling coefficients $\Omega_1(t)$, $\Omega_2(t)$ and $\Omega_3(t)$ respectively. For the coupling between the cavity fields and the fibre modes, we have taken the short fibre limit $(2L\bar{\eta})/(2\pi c) \ll 1$, which is similar to the previous discussion. Subsequently, under the dipole and rotating wave approximation, it is straightforward to write the total effective Hamiltonian for our cavity–fibre–cavity system as[77,78,81]

$$H_{\rm W} = H_{\rm W}^{\rm ac} + H_{\rm W}^{\rm cf}, \quad (6a)$$

$$\begin{aligned} H_{\rm W}^{\rm ac} = & \sum_{i=1}^{3}(g_i a_i|e\rangle_i\langle 0|) + \Omega_1(t)|e\rangle_1\langle 1| \\ & + \Omega_2(t)|e\rangle_2\langle a| + \Omega_3(t)|e\rangle_3\langle 1| + {\rm H.c.}, \quad (6b) \end{aligned}$$

$$H_{\rm W}^{\rm cf} = \eta_{\rm A}[b_{\rm A}(a_1^\dagger + a_2^\dagger)] + \eta_{\rm B}[b_{\rm B}(a_2^\dagger + a_3^\dagger)] + {\rm H.c.}, \quad (6c)$$

where g_i is the coupling coefficient between cavity mode i and atom i and cavity mode a_2 interacts synchronously with fibre modes $b_{\rm A}$ and $b_{\rm B}$.

3. Generation of GHZ state

In this section we show how to deterministically prepare the three-partite GHZ state, $|\Psi_{\rm GHZ}\rangle = (1/\sqrt{2})(|0\rangle_1|0\rangle_2|0\rangle_3 + |1\rangle_1|1\rangle_2|1\rangle_3)$. Consider at the initial time that the system is in the state $|\Psi(0)_{\rm GHZ}\rangle = |0\rangle_1|a\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$, where $|0\rangle_1|a\rangle_2|1\rangle_3$ means that atoms 1, 2, 3 are in the states $|0\rangle_1$, $|a\rangle_2$, $|1\rangle_3$, respectively, $|n_1n_2n_3\rangle_{\rm c}$ denotes the photon numbers in the mode of cavities 1-3, $|m_1m_2\rangle_{\rm f}$ represents m_1 photons in the mode of fibre 1, m_2 photons in the mode of fibre 2. Dominated by Hamiltonian (5), the system state evolution remains in the subspace with one excitation number spanned by the basis vectors

$$\begin{aligned} |\psi_1\rangle &= |0\rangle_1|a\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_2\rangle &= |0\rangle_1|1\rangle_2|1\rangle_3|01^{\rm l}0\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_3\rangle &= |0\rangle_1|1\rangle_2|1\rangle_3|000\rangle_{\rm c}|1^{\rm l}0\rangle_{\rm f}, \\ |\psi_4\rangle &= |0\rangle_1|1\rangle_2|1\rangle_3|1^{\rm l}00\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_5\rangle &= |1\rangle_1|1\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_6\rangle &= |0\rangle_1|0\rangle_2|1\rangle_3|01^{\rm r}0\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_7\rangle &= |0\rangle_1|0\rangle_2|1\rangle_3|000\rangle_{\rm c}|01^{\rm r}\rangle_{\rm f}, \\ |\psi_8\rangle &= |0\rangle_1|0\rangle_2|1\rangle_3|001^{\rm r}\rangle_{\rm c}|00\rangle_{\rm f}, \\ |\psi_9\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}. \quad (7) \end{aligned}$$

In this subspace, there exists a dark state with a null eigenvalue for the Hamiltonian $H_{\rm GHZ}$,

$$\begin{aligned} & |\Psi(t)\rangle_{\rm dark}^{\rm GHZ} \\ = \; & \frac{1}{D_0}(\eta_{\rm A}\eta_{\rm B}\Omega_{\rm e1}\Omega_{\rm e3}|\psi_1\rangle - \eta_{\rm B}\Omega_{\rm e1}\Omega_{\rm e3}\Omega_{\rm e2}^{\rm l}|\psi_3\rangle \\ & + \eta_{\rm A}\eta_{\rm B}\Omega_{\rm e3}\Omega_{\rm e2}^{\rm l}|\psi_5\rangle - \eta_{\rm A}\Omega_{\rm e1}\Omega_{\rm e3}\Omega_{\rm e2}^{\rm r}|\psi_7\rangle \\ & + \eta_{\rm A}\eta_{\rm B}\Omega_{\rm e1}\Omega_{\rm e2}^{\rm r}|\psi_9\rangle), \quad (8) \end{aligned}$$

where

$$D_0 = \sqrt{\eta_{\rm A}^2\eta_{\rm B}^2\Omega_{\rm e1}^2\Omega_{\rm e3}^2 + \eta_{\rm B}^2\Omega_{\rm e1}^2\Omega_{\rm e3}^2\Omega_{\rm e2}^{\rm l2} + \eta_{\rm A}^2\eta_{\rm B}^2\Omega_{\rm e3}^2\Omega_{\rm e2}^{\rm l2} + \eta_{\rm A}^2\Omega_{\rm e1}^2\Omega_{\rm e3}^2\Omega_{\rm e2}^{\rm r2} + \eta_{\rm A}^2\eta_{\rm B}^2\Omega_{\rm e1}^2\Omega_{\rm e2}^{\rm r2}}. \quad (9)$$

In our calculation, we assume that $|\eta_{\rm A}|$, $|\eta_{\rm B}| \gg |\Omega_{\rm e1}|$, $|\Omega_{\rm e2}^{\rm l(r)}|$, $|\Omega_{\rm e3}|$, thus the dark state $|\Psi(t)\rangle_{\rm dark}^{\rm GHZ}$ is approximately given as

$$|\Psi(t)\rangle_{\rm dark}^{\rm GHZ} = \frac{1}{D}(\Omega_{\rm e1}\Omega_{\rm e3}|\psi_1\rangle + \Omega_{\rm e3}\Omega_{\rm e2}^{\rm l}|\psi_5\rangle + \Omega_{\rm e1}\Omega_{\rm e2}^{\rm r}|\psi_9\rangle), \quad (10)$$

where $D=\sqrt{\Omega_{\rm e1}^2\Omega_{\rm e3}^2+\Omega_{\rm e3}^2\Omega_{\rm e2}^{\rm l2}+\Omega_{\rm e1}^2\Omega_{\rm e2}^{\rm r2}}$. Under the adiabatic condition, the initial state $|\Psi(0)\rangle_{\rm GHZ}=|\psi_1\rangle$ undergoes the evolution determined by Eq. (10). For realizing the generation of GHZ state, we can appropriately design the pulse shapes such that

$$\lim_{t\to-\infty}\frac{\Omega_{\rm e2}^{\rm l}(t)}{\Omega_{\rm e1}(t)}=0,\quad \lim_{t\to-\infty}\frac{\Omega_{\rm e2}^{\rm r}(t)}{\Omega_{\rm e3}(t)}=0, \tag{11a}$$

$$\lim_{t\to+\infty}\frac{\Omega_{\rm e1}(t)}{\Omega_{\rm e2}^{\rm l}(t)}=0,\quad \lim_{t\to+\infty}\frac{\Omega_{\rm e3}(t)}{\Omega_{\rm e2}^{\rm r}(t)}=0, \tag{11b}$$

and adiabatically transfer the initial state $|\psi_1\rangle$ into equal probability amplitude superposition states of $|\psi_5\rangle$ and $|\psi_9\rangle$, with $|\psi(t)\rangle_{t\to+\infty}=(1/\sqrt{2})(|\psi_5\rangle+|\psi_9\rangle)$, which is a product state of the three-atom entangled state, the cavity mode state and the fibre mode state. It is a maximally entangled state, i.e., the GHZ state. In order to accomplish the generation of this entangled state, here the pulse shapes of the laser fields are assumed to be Gaussian,[82,83]

$$\Omega_{1,2,3}(t)=\Omega_0{\rm exp}\left[-\frac{(t-T/2-t_{1,2,3})^2}{2\tau^2}\right], \tag{12}$$

where Ω_0 is the amplitude of $\Omega_{1,2,3}$, T is the total adiabatic time and τ is the laser beam waist. In our following simulation, we set $g_{1{\rm l}}=g_{2{\rm l(r)}}=g_{3{\rm r}}=g$, $\Delta_1=\Delta_2=\Delta_3=\Delta$. The system parameters are chosen as $g=1$ GHz, $\Omega_0=0.2g$, $T=100/g$, $\tau=14/g$, $t_1=t_3=-10/g$, $t_2=10/g$ and $\Delta=15g$. In order to show clearly expression (12) of pulse shapes, we plot the time dependences of Rabi frequencies of the laser fields and the time evolutions of the population in Fig. 2. It is obviously seen from Fig. 2(b) when $t\to+\infty$, the populations of $|\psi_5\rangle$ and $|\psi_9\rangle$ are both equal to 50%, which means that the GHZ state $|\Psi_{\rm GHZ}\rangle$ can be deterministically achieved via adiabatic evolution. In the present scheme, the fidelity for generating the GHZ state is higher than 99.99% ($F_{\rm GHZ}=\langle\Psi_{\rm G}|\psi(t)\rangle$) when the dimensionless time $gt>100$ and $t=10^{-7}$ s can be considered to be $+\infty$ here. The plots of pulse shapes are shown in Fig. 2(a), where at the beginning the lasers which act on atoms 1 and 3 are switched on first simultaneously. During the evolution, the laser intensities are slowly changed (increasing Ω_2 and decreasing $\Omega_{1(3)}$ adiabatically) so that the condition $\Omega_2\gg\Omega_{1(3)}$ is fulfilled at the end for realizing the GHZ state.

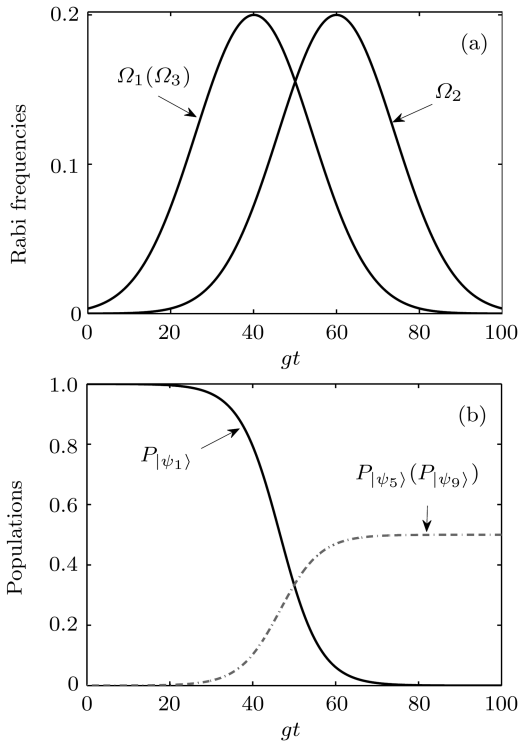


Fig. 2. (a) Time dependences of the Rabi frequencies of the laser fields for generation of GHZ state, (b) time evolutions of the populations of corresponding system states. The system parameters are set to be $\Omega_0=0.2g$, $T=100/g$, $\tau=14/g$, $t_1=t_3=-10/g$, $t_2=10/g$, $g=1$ GHz and $\Delta=15g$.

4. Generation of W state

In this section we show how to deterministically prepare a three-partite W state, $|\Psi_{\rm W}\rangle=(1/\sqrt{3})(|0\rangle_1|0\rangle_2|1\rangle_3+|0\rangle_1|1\rangle_2|0\rangle_3+|1\rangle_1|0\rangle_2|0\rangle_3)$. Suppose at the initial time, atom 1 is in the state $|0\rangle_1$, atom 2 is in the superposition state $(\sqrt{2}|a\rangle+|1\rangle)/\sqrt{3}$, atom 3 in the state $|0\rangle_3$, all the cavity modes and fibre modes are in the vacuum state $|000\rangle_{\rm c}|00\rangle_{\rm f}$, i.e., $|\Psi_{\rm W}(0)\rangle=(1/\sqrt{3})|0\rangle_1|1\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}+(2/\sqrt{3})|0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$. If we consider that atom 2 is in the state $|a\rangle_2$ initially, then dominated by Hamiltonian (6), the system state evolution will be restricted to the subspace with one excitation number spanned by the basis vectors

$$\begin{aligned}
|\psi_1\rangle &= |1\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_2\rangle &= |e\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_3\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|100\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_4\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|10\rangle_{\rm f},\\
|\psi_5\rangle &= |0\rangle_1|e\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_6\rangle &= |0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_7\rangle &= |0\rangle_1|0\rangle_2|e\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},
\end{aligned}$$

$$
\begin{aligned}
|\psi_8\rangle &= |0\rangle_1|0\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_9\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|010\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_{10}\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|001\rangle_{\rm c}|00\rangle_{\rm f},\\
|\psi_{11}\rangle &= |0\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|01\rangle_{\rm f}. \qquad (13)
\end{aligned}
$$

In this subspace, the corresponding dark state of Hamiltonian (6) is

$$
\begin{aligned}
|\Psi(t)\rangle_{\rm dark}^{\rm W} = \frac{1}{K_0}\bigg(&\frac{g_1}{\Omega_1}|\psi_1\rangle - |\psi_3\rangle - \frac{g_2}{\Omega_2}|\psi_6\rangle\\
&+ \frac{g_3}{\Omega_3}|\psi_8\rangle + |\psi_9\rangle - |\psi_{10}\rangle\bigg), \qquad (14)
\end{aligned}
$$

where

$$
K_0 = \sqrt{\frac{g_1^2}{\Omega_1^2} + \frac{g_2^2}{\Omega_2^2} + \frac{g_3^2}{\Omega_3^2} + 3}. \qquad (15)
$$

In our following calculation, we assume that the condition $g_i \gg \Omega_i(t)$ $(i = 1, 2, 3)$ is always satisfied in the whole adiabatic progress, which means that the population of cavity modes in excited states can be dropped safely. Thus the dark state $|\Psi(t)\rangle_{\rm dark}^{\rm W}$ is approximately given as

$$
\begin{aligned}
&|\Psi(t)\rangle_{\rm dark}^{\rm W}\\
&= \frac{1}{K}\left(\frac{g_1}{\Omega_1}|\psi_1\rangle - \frac{g_2}{\Omega_2}|\psi_6\rangle + \frac{g_3}{\Omega_3}|\psi_8\rangle\right), \qquad (16)
\end{aligned}
$$

where $K = \sqrt{g_1^2/\Omega_1^2 + g_2^2/\Omega_2^2 + g_3^2/\Omega_3^2}$. The initial state $|\psi_6\rangle$ will undergo the evolution determined by Eq. (10). For realizing the generation of W state, we can appropriately design the pulse shapes such that

$$
\lim_{t\to-\infty}\frac{\Omega_1(t)}{\Omega_2(t)} = 0, \quad \lim_{t\to-\infty}\frac{\Omega_3(t)}{\Omega_2(t)} = 0, \qquad (17a)
$$

$$
\lim_{t\to+\infty}\frac{\Omega_2(t)}{\Omega_1(t)} = 0, \quad \lim_{t\to+\infty}\frac{\Omega_2(t)}{\Omega_3(t)} = 0, \qquad (17b)
$$

where we have set $g_1 = g_2 = g_3 = g$. The initial state $|\psi_6\rangle$ can evolve into equal probability amplitude superposition states of $|\psi_1\rangle$ and $|\psi_8\rangle$. On the other hand, when atom 2 is initially in the state $|1\rangle_2$, the system state $|0\rangle_1|1\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$ will remain unchanged with the evolvement of time. Consequently, based on our above discussion, the W state $|\psi(t)\rangle_{t\to+\infty} = (1/\sqrt{3})(|0\rangle_1|0\rangle_2|1\rangle_3 + |0\rangle_1|1\rangle_2|0\rangle_3 + |1\rangle_1|0\rangle_2|0\rangle_3)$ can be achieved via adiabatic passage. It is seen that the pulse shapes of the laser fields are the key points for accomplishing the generation of this entangled state. Here they are also assumed to be Gaussian,[82,83]

$$
\Omega_{1,2,3}(t) = \Omega_0 \exp\left[-\frac{(t - T/2 - t_{1,2,3})^2}{2\tau^2}\right]. \qquad (18)
$$

We have chosen the system parameter as follows: $\Omega_0 = 0.2g$, $T = 100/g$, $\tau = 12/g$, $t_1 = t_3 = -10/g$, $t_2 = 10/g$ and $g = 1$ GHz for generating W state. In Fig. 3, we plot the time dependences of Rabi frequencies of the laser fields and the time evolutions of the population for the initial condition $|\Psi(0)\rangle = |\psi_6\rangle = |0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$. It is obviously seen from Fig. 3(b) when $t \to +\infty$, the system state will evolve into $|\Psi_a\rangle = (1/\sqrt{2})(|\psi_1\rangle + |\psi_8\rangle)$. In addition, we have already pointed out for the initial state $|0\rangle_1|1\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$, the system state does not change under the time evolution dominated by Hamiltonian (6) since it is decoupled from the laser interaction. Therefore, the W state $|\Psi_{\rm W}\rangle = (1/\sqrt{3})(|0\rangle_1|0\rangle_2|1\rangle_3 + |0\rangle_1|1\rangle_2|0\rangle_3 + |1\rangle_1|0\rangle_2|0\rangle_3)$ can be deterministically achieved via adiabatic evolution. In the present scheme, the fidelity for generating the W state is higher than 99.99% ($F_{\rm W} = \langle\Psi_{\rm W}|\psi(t)\rangle$), when the dimensionless time $gt > 100$ and $t = 10^{-7}$ s can be considered to be $+\infty$ here. In Fig. 3(a), we show the pulse shapes of the laser fields. In order to realize the W state, we turn on the lasers which act on atoms 1 and 3 simultaneously at the beginning, and in the evolution process, slowly change the laser intensities (increasing Ω_2 and decreasing $\Omega_{1(3)}$ adiabatically) so that at the end Ω_1, Ω_2 and Ω_3 can meet the conditions for realizing W state, which is similar to the previous scheme proposed for the generation of the GHZ state.

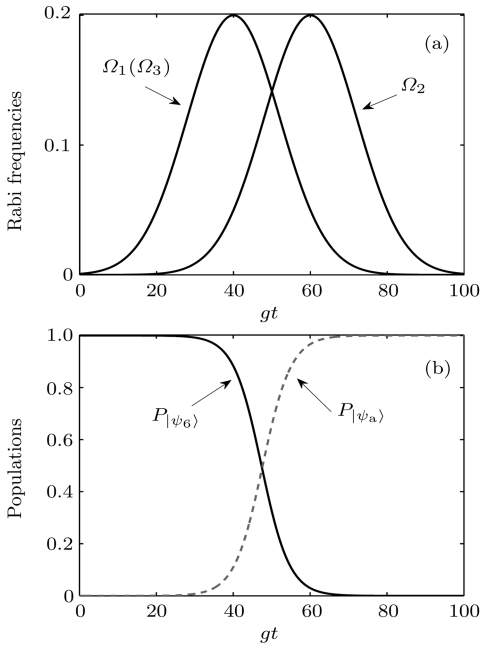


Fig. 3. (a) Time dependences of the Rabi frequencies of the laser fields for the generation of W state, (b) time evolutions of the populations of corresponding system states. The system parameters are set to be $\Omega_0 = 0.2g$, $T = 100/g$, $\tau = 12/g$, $t_1 = t_3 = -10/g$, $t_2 = 10/g$ and $g = 1$ GHz.

5. Discussion and conclusion

In the previous section, we have put forward two schemes to generate GHZ state and W state of three distant atoms, respectively. It is noticed that the influences of atomic spontaneous emission from the excited states, photon leakage out of the fibre and cavity are not taken into account in our discussion. In the next content, we will consider more practical situations and study the effects of these factors on the fidelity of realizing entanglement states. To accomplish the GHZ state, we have used the large-detuning condition, adiabatically eliminated the excited state[84] and set the spontaneous emission rate to be zero. Meanwhile, the effect of photon loss out of the cavity can also be ignored based on our scheme. This is because utilizing the adiabatic passage along dark state, the cavity modes are never appreciably populated and always kept in vacuum states in the whole process. However, it is noted that the terms related to $|\psi_3\rangle$ and $|\psi_7\rangle$ in expression (8) have been dropped for our above simulations since we assume that $|\eta_{\rm A}|$, $|\eta_{\rm B}| \gg |\Omega_{\rm e1}|$, $|\Omega_{\rm e2}^{\rm l(r)}|$, $|\Omega_{\rm e1}|$. In fact, the fibre modes may be excited, that is to say, there is a probability with which the states $(|0\rangle_1|1\rangle_2|1\rangle_3|000\rangle_{\rm c}|1^{\rm l}0\rangle_{\rm f} + |0\rangle_1|1\rangle_2|1\rangle_3|000\rangle_{\rm c}|01^{\rm r}\rangle_{\rm f})/\sqrt{2}$ are populated. Then, due to the fibre decay they may evolve into $(|0\rangle_1|1\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f} + |0\rangle_1|0\rangle_2|1\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f})/\sqrt{2}$ after the photon emits out of fibre, which will cause error for achieving GHZ state. Considering the effect of photon leakage out of fibre, the fidelity can be written as

$$F_{\rm GHZ} = 1 - \gamma_{\rm f}\int_0^T \frac{(\eta_{\rm B}\Omega_{\rm e1}\Omega_{\rm e3}\Omega_{\rm e2}^{\rm l})^2 + (\eta_{\rm A}\Omega_{\rm e1}\Omega_{\rm e3}\Omega_{\rm e2}^{\rm r})^2}{D_0^2}\,{\rm d}t, \tag{19}$$

where D_0 is given in expression (9); $\gamma_{\rm f}$ denotes the decay rate of fibre mode; we assume $\gamma_{\rm fA} = \gamma_{\rm fB} = \gamma_{\rm f}$; T is the total adiabatic time on the order of 100 ns, i.e., $T = 100$ ns in our following simulations. Figure 4(a) exhibits the effects of fibre loss on the fidelity $F_{\rm GHZ}$ for different values of Ω_0, $\Omega_0 = 0.2g$, $\Omega_0 = 0.5g$ and $\Omega_0 = 1.0g$, where we set $\eta_{\rm A} = \eta_{\rm B} = \eta = 0.6g$. In Fig. 4(b) we show the curves for fidelity $F_{\rm GHZ}$ versus η/g for different decay rates of fibre modes, $\gamma_{\rm f} = 0.05g$, $\gamma_{\rm f} = 0.1g$ and $\gamma_{\rm f} = 0.2g$, where $\Omega_0 = 0.5g$. The other parameters are the same as those in Fig. 2. It is seen from Fig. 4(a) that when $\Omega_0 \geq 0.5g$, even if there exists a relatively large fibre decay rate $\gamma_{\rm f} = 0.3g$, we can still obtain GHZ state with high fidelity ($F_{\rm GHZ} > 99\%$). Figure 4(b) indicates that with the increase of fibre–cavity coupling strength, the fidelity increases and when $\eta/g \geq 0.5$, $F_{\rm GHZ}$ will be more stable with its value larger than 99%. It is noticed that based on the present scheme, we do not demand the condition $g \gg \Omega_0$, which will facilitate the experimental achievement of the generation of GHZ state. Next we will discuss the influence of various decoherence process on the fidelity of realizing W state. Based on our scheme, the effect of atomic spontaneous emission from the excited states and photon leakage out of fibre can be greatly eliminated and even ignored because the atomic excited states and the fibre modes are not populated in the whole process via adiabatic passage along dark state. While the photon losses in cavities may have a greater influence on the fidelity of achieving W state. In the present scheme, we set $|\Psi_{\rm W}(0)\rangle = (1/\sqrt{3})|0\rangle_1|1\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f} + (2/\sqrt{3})|0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$. There is a probability with which at the initial time, the system is in the state $|0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$. In our above simulations, it is noted that we have dropped the terms related to $|\psi_3\rangle$, $|\psi_9\rangle$ and $|\psi_{10}\rangle$ in expression (14) since we assume that the atom–cavity coupling strengths g_i ($i = 1, 2, 3$) are much larger than the Rabi frequencies Ω_i ($i = 1, 2, 3$). In fact, it is an ideal situation. In the practical experimental circumstance, if $g_i \gg \Omega_{0i}(t)$ ($i = 1, 2, 3$) cannot be satisfied perfectly, the cavity modes may be excited and the effect of photon losses out of cavities is not negligible any more. That is to say, there is a probability with which the states $(|0\rangle_1|0\rangle_2|0\rangle_3|100\rangle_{\rm c}|00\rangle_{\rm f} + |0\rangle_1|0\rangle_2|0\rangle_3|010\rangle_{\rm c}|00\rangle_{\rm f} + |0\rangle_1|0\rangle_2|0\rangle_3|001\rangle_{\rm c}|00\rangle_{\rm f})/\sqrt{3}$ are populated. Then, due to the cavity decay they may evolve into $|0\rangle_1|0\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$ after the photons are emitted out of cavity, which will reduce the fidelity of achieving W state. We assume $\kappa_1 = \kappa_2 = \kappa_3 = \kappa$, where κ_i ($i = 1, 2, 3$) are the decay rates of the cavity fields, which leads to an error of $\kappa\int_0^T (3/K_0^2)\,{\rm d}t$, where K_0 is given by Eq. (15). Hence, for $|\Psi_{\rm W}(0)\rangle = (1/\sqrt{3})|0\rangle_1|1\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f} + (2/\sqrt{3})|0\rangle_1|a\rangle_2|0\rangle_3|000\rangle_{\rm c}|00\rangle_{\rm f}$, the fidelity of accomplishing W state can be written as

$$F_{\rm W} = \frac{2}{3}\left(1 - \kappa\int_0^T \frac{3}{K_0^2}\,{\rm d}t\right) + \frac{1}{3}, \tag{20}$$

where T is the total adiabatic time on the order of 100 ns, i.e., $T = 100$ ns in our following simulations. Figure 4(c) exhibits the effects of cavity loss on the fidelity $F_{\rm W}$ for different values of Ω_0: $\Omega_0 = 0.2g$, $\Omega_0 = 0.3g$ and $\Omega_0 = 0.4g$. In Fig. 4(b) we plot curves

for fidelity $F_{\rm W}$ versus the ratio of maximal values of Rabi field to the atom–cavity coupling strength Ω_0/g for different decay rates of cavity modes, $\kappa = 0.01g$, $\kappa = 0.02g$ and $\kappa = 0.04g$. The other parameters are the same as those in Fig. 3. It is seen from Fig. 4(d) that if κ reaches $0.04g$, the fidelity reduces a little bit quickly when $\Omega_0/g > 0.4$. Therefore, in order to obtain a high fidelity W state, the condition of $g_i \gg \Omega_{0i}(t)$ $(i = 1, 2, 3)$ are required (in experiment, g_i more than four times larger than Ω_{0i} has been realized[85−87] and the other parameters should also been chosen appropriately.

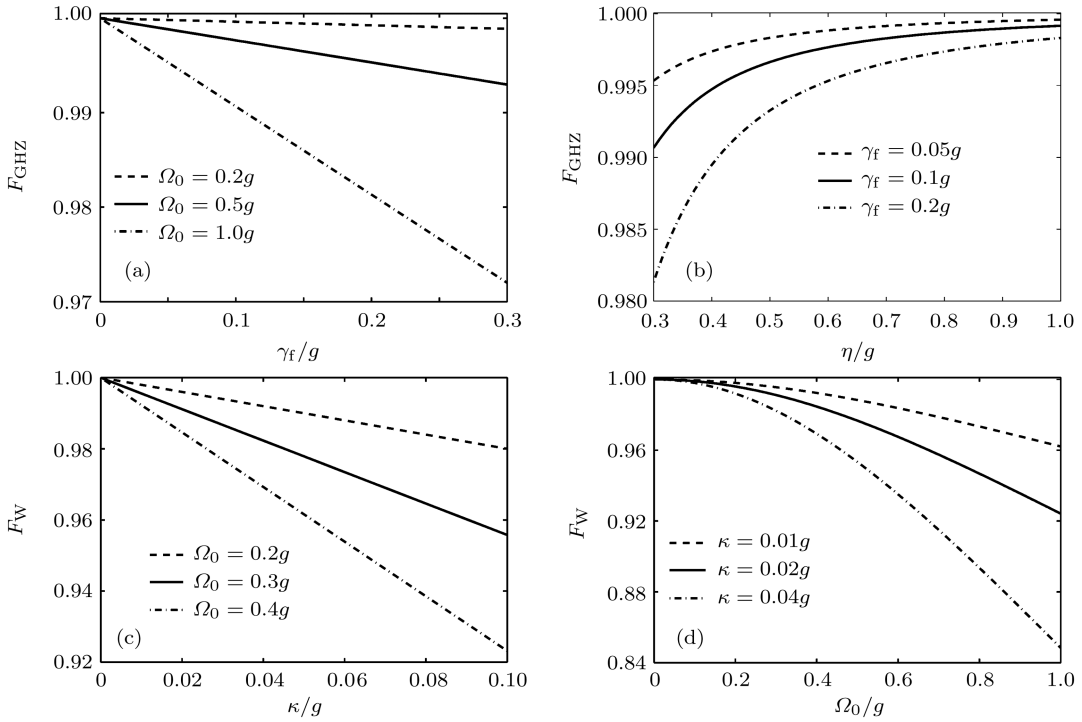


Fig. 4. (a) Fidelities of realizing GHZ state each as a function of $\gamma_{\rm f}/g$ for $\Omega_0 = 0.2g$, $\Omega_0 = 0.5g$ and $\Omega_0 = 1.0g$, $\eta = 0.6g$; (b) fidelities of realizing GHZ state each as a function of η/g for $\gamma_{\rm f} = 0.05g$, $\gamma_{\rm f} = 0.1g$ and $\gamma_{\rm f} = 0.2g$, $\Omega_0 = 0.5g$; (c) fidelities of realizing W state each as a function of κ/g for $\Omega_0 = 0.2g$, $\Omega_0 = 0.3g$ and $\Omega_0 = 0.4g$; (d) fidelities of realizing W state as a function of Ω_0/g for $\kappa = 0.01g$, $\kappa = 0.02g$ and $\kappa = 0.04g$. The other parameters are the same as those in Fig. 2 (for (a) and (b)) and Fig. 3 (for (c) and (d)).

Besides the various decoherence processes discussed above, another challenge we may meet is that based on our scheme for generating GHZ state and W state, at the beginning, the lasers which act on atoms 1 and 3 are switched on first at the same time, i.e., $t_1 = t_3$ set in our simulations. In practical experimental situations, the requirement that they are turned on simultaneously can not be precisely fulfilled. In Figs. 5 (a) and 5 (b), we plot the curves for fidelity $F_{\rm GHZ}$ and $F_{\rm W}$ versus corresponding Δt, where we assume $t_1 = t_3 + \Delta t$, the other parameters are the same as those in Figs. 2 and 3, respectively. It is shown in Fig. 5 that when -1 ns$\leq \Delta t \leq 1$ ns, the fidelity can be larger than 98.5% for generating both GHZ state and W state based on the present schemes. As a result, even though the ideal parameter conditions $t_1 = t_3$ could not be satisfied exactly in practical situations, the entangled states $\Psi_{\rm GHZ}$ and $\Psi_{\rm W}$ can still be realized with high fidelity.

In the following part, we will briefly discuss the experimental feasibility of our scheme. The level configurations of corresponding atoms can be realized by choosing the hyperfine-split levels for the D lines of cold alkali-metal atoms.[88,89] For instance, we can choose the atom of ^{87}Rb. The atomic levels and field polarizations are shown in Fig. 6. On the other hand, based on recent experiments in realizing high-Q cavity and strong atom–cavity coupling,[90−93] we can set the system parameters as follows: for GHZ state, $g = 2\pi\times160$ MHz, $\Omega_0 = 2\pi\times35$ MHz, $\Delta = -2\pi\times2.4$ GHz, $T = 100$ ns, $\tau = 14$ ns, $t_1 = t_3 = 10$ ns, $t_2 = -10$ ns, $\eta = 2\pi \times 80$ MHz and $\gamma_{\rm f} = 2\pi \times 16$ MHz, while for W state, $g = 2\pi \times 160$ MHz, $\Omega_0 = 2\pi \times 35$ MHz, $T = 100$ ns, $\tau = 12$ ns, $t_1 = t_3 = 10$ ns, $t_2 = -10$ ns, $\kappa = 2\pi \times 8.4$ MHz.

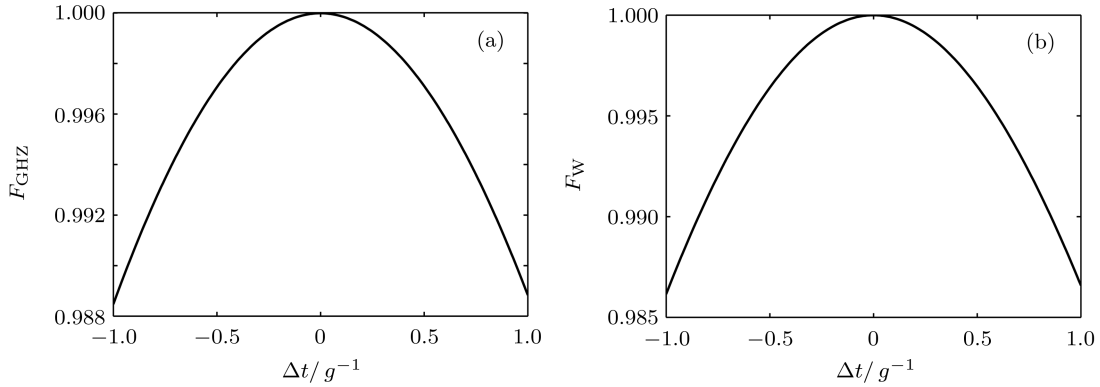


Fig. 5. (a) The fidelity for the generation of GHZ state versus Δt, where the pulse parameters are the same as those in Fig. 2; (b) the fidelity for the generation of W state versus Δt, where the pulse parameters are the same as those in Fig. 3.

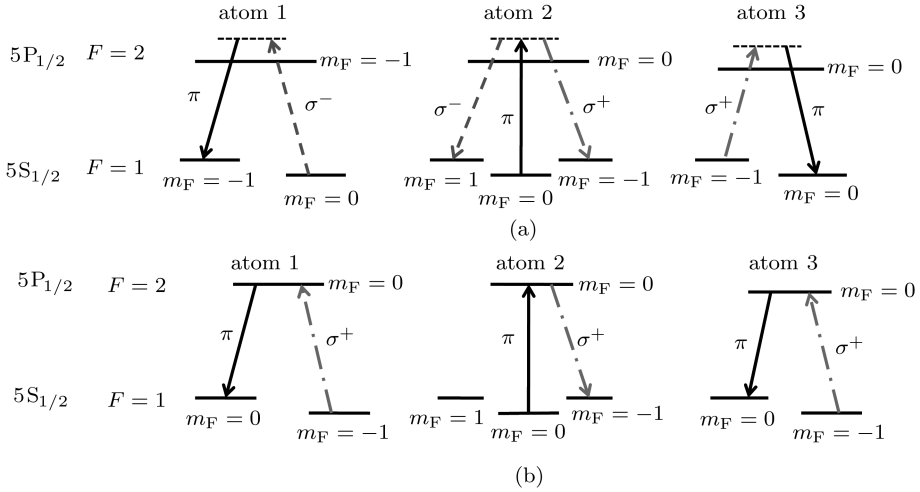


Fig. 6. The level configurations of atoms for the generation of GHZ state (a) and W state (b), σ^-, σ^+ and π denote left circular, right circular and line polarizations (light), respectively.

In conclusion, we have proposed two schemes for deterministically generating three-particle GHZ state and W state for three atoms, which are separately trapped in three distant cavities connected by two optical fibres. By using a sequence of pulsed laser field, our atom–cavity–fibre systems are always in the dark states. In the whole process, the influences of atomic spontaneous emission, photons leakage out of fibres and cavities are effectively compressed via adiabatic passages. Furthermore, we make an estimation on the fidelities of realizing the GHZ state and W state by considering different parameters and point out that our schemes are insensitive to small fluctuations of experimental parameters. As a result, the present schemes provide efficient approaches to realizing multi-particle entanglement and it is expected that our present proposal would be a theoretic supplement to the progress of quantum network science.

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