

Generation of multi-atom entangled states in coupled cavities via transitionless quantum driving

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Received: 25 May 2015 / Accepted: 26 September 2015 / Published online: 8 October 2015
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Abstract In this paper, we propose an efficient scheme to generate three-atom W states in spatially separated cavities connected by optical fibers. In the scheme, we combine the “transitionless quantum driving” with “quantum Zeno dynamics” to construct a shortcut to fast generate W states. Comparing with the traditional adiabatic passage, the significant advantage is that the interaction time required for the creation of the W state is much shorter, which is very important in view of decoherence. Furthermore, the harmful effects of various decoherence such as atomic spontaneous emission, cavity losses and the fiber photon leakages are considered. Numerical simulations illustrate that the shortcut scheme is much faster than the schemes using adiabatic passage and robust against the decoherence. Moreover, this scheme can also be generalized to generation of N -atom W states.

Keywords W states · Adiabatic passage · Shortcuts to adiabatic passage · Transitionless quantum driving

1 Introduction

Quantum entanglement is the soul of quantum mechanics. The manipulation of entangled states is not only the basic task for quantum information processing (QIP) [1,2], but also fundamental for demonstrating quantum nonlocality [3,4]. In general, entanglement of multi-qubit is more useful for QIP and shows more nonclassical effects. There are two main kinds of entangled states for three-qubit entanglement, the W state [5] and the Greenberger–Horne–Zeilinger (GHZ) state [4]. Among these two

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kinds of tripartite entanglement, the W state attracts much attention since it can retain bipartite entanglement when any one of the qubits is traced out, and it has the advantages in quantum teleportation [6]. Thus, many theoretical works have been devoted to the generation of W states via different techniques [7–26]. Among these techniques, the resonant interaction [16, 17] and the stimulated Raman scattering involving adiabatic passage (STIRAP) [18–21] have attracted attention in recent years. However, the schemes based on resonant interaction are very fast, but they depend on the exact knowledge of all parameters and require controlling the interaction time accurately. Although the schemes based on adiabatic passage are robust against experimental parameters and do not require controlling the interaction time accurately, the operation time required to get the goal is relatively long, which may lead to larger decoherence effects. In order to combine advantages of the resonant interaction and the adiabatic passage, a famous technique named “shortcuts to adiabatic passage” (STAP) [27–31], which can fast and robustly generate entangled states, has been extended in recent years.

The “transitionless quantum driving” (TQD) [31] is a well-known method to construct shortcuts to speed up adiabatic processes effectively. Generally speaking, the laser pulses are not strongly limited by using TQD to construct shortcuts, but a non-existent Hamiltonian in real experiment is necessary by using this method. Apparently, if we take no account of that whether the constructed Hamiltonian in experiment is existent or not, TQD is an effective method to fast generate entangled states and implement quantum processing. Because for any time-dependent original Hamiltonian $H_0(t)$, TQD provides a very effective method to construct a Hamiltonian $H(t)$ which accurately derives the instantaneous eigenstates of $H_0(t)$. Moreover, the transitions between them do not occur at all during the evolution of the whole system regardless of the rate of change. In other words, the instantaneous eigenstates of $H_0(t)$ can be regarded as truly moving eigenstates of $H(t)$. As we mentioned above, the outcome as the same as the adiabatic process can be obtained by using TQD in a shorter time. In view of the advantages of TQD, it is worth to finding ways to overcome the problem that the Hamiltonian $H(t)$ designed by using TQD does not exist in experiment. So far, lots of shortcut schemes have been proposed in theory and implemented in experiment [30–53].

It is worth noticing that Chen et al. [33] first constructed shortcuts to perform fast and noise-resistant populations transfer in multi-particle systems, by combining “Lewis–Riesenfeld invariants” with “quantum Zeno dynamics” (QZD). Soon after, a lot of schemes are rapidly proposed to perform fast and noise-resistant QIP [35–37, 54–56], by using similar ideas with slightly breaking the Zeno condition down. However, most of the above shortcut schemes are based on a single cavity, which is still a challenge to manipulate a large number of qubits. This is due to the fact that the spatial separation between neighboring qubits decreases as the number of qubits increases, and thus, individual addressing becomes increasingly difficult. The coupled-cavity systems [57–64] are considered as a suitable candidate for the solution of the above deficiency. In view of that we wonder if it is possible to use TQD to construct shortcuts for generation of multi-atom entanglement, i.e., W states, in coupled cavities.

In this paper, we construct STAP to fast generate W states in spatially separated cavities by combining “TQD” with “QZD.” The scheme has the following advantages:

- (1) The interaction time required for the creation of the W state is much shorter than the adiabatic passage.
- (2) Numerical simulations show that the decoherence such as atomic spontaneous emission, cavity losses and the fiber photon leakages has little influence on this scheme.
- (3) Individual addressing becomes relatively easy in coupled cavities.
- (4) This scheme can be generalized to generation of N -atom W states as well.

This paper is structured as follows. In Sect. 2, we construct a theoretical model by using QZD. In Sect. 3, we construct a shortcut passage based on TQD and show how to use the constructed shortcut to fast generate W states. In Sect. 4, N -atom W states are generated in one step by the same principle. A discussion on experimental feasibility and a summary are given in Sect. 5.

2 Theoretical model

As shown in Fig. 1 a, b, we consider that three identical Λ -type atoms are trapped in three spatially separated cavities C_1 , C_2 and C_3 , respectively. The cavities are linked by optical fibers f_1 and f_2 . Each atom has an excited state $|e\rangle$ and two ground states $|f\rangle$ and $|g\rangle$. The atomic transition $|f\rangle \leftrightarrow |e\rangle$ is resonantly driven by classical field

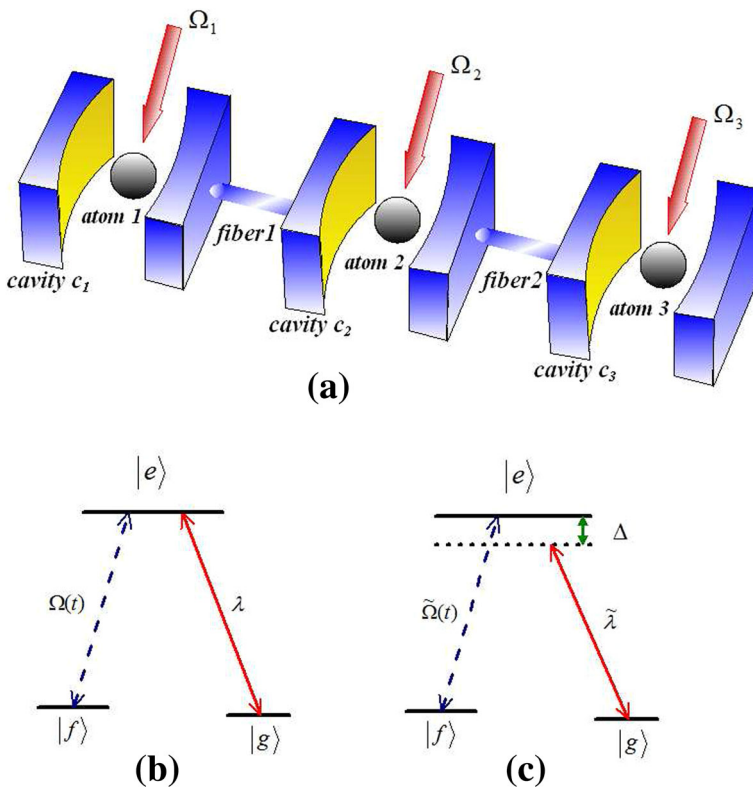


Fig. 1 Experimental setup and level configuration for each atom

$\Omega(t)$, and the transition $|g\rangle \leftrightarrow |e\rangle$ resonantly coupled to the cavity with coupling coefficient λ . In the short-fiber limit, i.e., $2Lv/(2\pi c) \ll 1$ (where L denotes the fiber length, c denotes the speed of light, and ν denotes the decay of the cavity field into a continuum of fiber mode), only one resonant fiber mode interacts with the cavity mode [65]. Under the rotating-wave approximation (RWA), the interaction Hamiltonian for this system can be written as ($\hbar = 1$)

$$\begin{aligned}
 H_o &= H_{al} + H_{ac}, \\
 H_{al} &= \sum_{k=1}^3 \Omega_k(t) |e\rangle_k \langle f| + H.c., \\
 H_{ac} &= \sum_{k=1}^3 \lambda_k a_k |e\rangle_k \langle g| + \nu_1 b_1^\dagger (a_1 + a_2) + \nu_2 b_2^\dagger (a_2 + a_3) + H.c., \tag{1}
 \end{aligned}$$

where a_k^\dagger and a_k denote creation and annihilation operators of the cavities C_k ($k = 1, 2, 3$), respectively; b_j^\dagger and b_j denote the creation and annihilation operators associated with the resonant mode of fiber f_j ($j = 1, 2$), respectively. For the sake of simplicity, we assume $\Omega_2 = \Omega_x$, $\Omega_1 = \Omega_3 = \Omega_y$, $\lambda_k = \lambda$ and $\nu_j = \nu$. If we assume the initial state of the system is in $|g, f, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}$, the whole system evolves in the subspace \mathbb{V} spanned by:

$$\begin{aligned}
 |\psi_1\rangle &= |g, f, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_2\rangle &= |g, e, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_3\rangle &= |g, g, g\rangle_{1,2,3}|1, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_4\rangle &= |g, g, g\rangle_{1,2,3}|0, 1, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_5\rangle &= |g, g, g\rangle_{1,2,3}|0, 0, 1\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_6\rangle &= |g, g, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|1, 0\rangle_{f1,f2}, \\
 |\psi_7\rangle &= |g, g, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 1\rangle_{f1,f2}, \\
 |\psi_8\rangle &= |g, g, e\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_9\rangle &= |g, g, f\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_{10}\rangle &= |e, g, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}, \\
 |\psi_{11}\rangle &= |f, g, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}. \tag{2}
 \end{aligned}$$

Under the Zeno condition $\Omega_x, \Omega_y \ll \sqrt{3}\lambda$, the Hilbert subspace \mathbb{V} is split into seven Zeno subspaces according to the degeneracy of eigenvalues of H_{ac} ,

$$\begin{aligned}
 Z_0 &= \{|\phi_1\rangle, |\phi_2\rangle, |\phi_3\rangle, |\phi_4\rangle, |\phi_5\rangle\}, \\
 Z_1 &= \{|\phi_6\rangle\}, \quad Z_2 = \{|\phi_7\rangle\}, \quad Z_3 = \{|\phi_8\rangle\}, \\
 Z_4 &= \{|\phi_9\rangle\}, \quad Z_5 = \{|\phi_{10}\rangle\}, \quad Z_6 = \{|\phi_{11}\rangle\}, \tag{3}
 \end{aligned}$$

where the eigenstates of H_{ac} are

$$\begin{aligned}
 |\phi_1\rangle &= |\psi_1\rangle, \\
 |\phi_2\rangle &= \frac{v}{\sqrt{2v^2 + \lambda^2}} \left[|\psi_2\rangle + \frac{\lambda}{v} |\psi_7\rangle + |\psi_8\rangle \right], \\
 |\phi_3\rangle &= |\psi_9\rangle, \\
 |\phi_4\rangle &= \frac{v}{\sqrt{2v^2 + \lambda^2}} \left[|\psi_2\rangle - \frac{\lambda}{v} |\psi_6\rangle + |\psi_{10}\rangle \right], \\
 |\phi_5\rangle &= |\psi_{11}\rangle, \\
 |\phi_6\rangle &= \frac{1}{\sqrt{6}} [-|\psi_2\rangle + |\psi_3\rangle - |\psi_4\rangle - |\psi_5\rangle + |\psi_8\rangle + |\psi_{10}\rangle], \\
 |\phi_7\rangle &= \frac{1}{\sqrt{6}} [-|\psi_2\rangle - |\psi_3\rangle + |\psi_4\rangle + |\psi_5\rangle + |\psi_8\rangle + |\psi_{10}\rangle], \\
 |\phi_8\rangle &= \frac{\lambda}{2\sqrt{v^2 + \lambda^2}} \left[-\frac{\sqrt{v^2 + \lambda^2}}{\lambda} |\psi_4\rangle + \frac{\sqrt{v^2 + \lambda^2}}{\lambda} |\psi_5\rangle + \frac{v}{\lambda} |\psi_6\rangle \right. \\
 &\quad \left. - \frac{v}{\lambda} |\psi_7\rangle - |\psi_8\rangle + |\psi_{10}\rangle \right], \\
 |\phi_9\rangle &= \frac{\lambda}{2\sqrt{v^2 + \lambda^2}} \left[\frac{\sqrt{v^2 - \lambda^2}}{\lambda} |\psi_4\rangle - \frac{\sqrt{v^2 + \lambda^2}}{\lambda} |\psi_5\rangle + \frac{v}{\lambda} |\psi_6\rangle \right. \\
 &\quad \left. - \frac{v}{\lambda} |\psi_7\rangle - |\psi_8\rangle + |\psi_{10}\rangle \right], \\
 |\phi_{10}\rangle &= \frac{\lambda}{2\sqrt{3(3v^2 + \lambda^2)}} \left[2|\psi_2\rangle - \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_3\rangle - \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_4\rangle \right. \\
 &\quad \left. - \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_5\rangle + \frac{3v}{\lambda} |\psi_6\rangle + \frac{3v}{\lambda} |\psi_7\rangle + |\psi_8\rangle + |\psi_{10}\rangle \right], \\
 |\phi_{11}\rangle &= \frac{\lambda}{2\sqrt{3(3v^2 + \lambda^2)}} \left[2|\psi_2\rangle + \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_3\rangle + \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_4\rangle \right. \\
 &\quad \left. + \frac{2\sqrt{3v^2 + \lambda^2}}{\lambda} |\psi_5\rangle + \frac{3v}{\lambda} |\psi_6\rangle + \frac{3v}{\lambda} |\psi_7\rangle + |\psi_8\rangle + |\psi_{10}\rangle \right], \tag{4}
 \end{aligned}$$

and the corresponding eigenvalues are

$$\begin{aligned}
 \varsigma_0 &= 0, \quad \varsigma_1 = -\lambda, \quad \varsigma_2 = \lambda, \quad \varsigma_3 = -\sqrt{v^2 + \lambda^2}, \quad \varsigma_4 = \sqrt{v^2 + \lambda^2}, \\
 \varsigma_5 &= -\sqrt{3v^2 + \lambda^2}, \quad \varsigma_6 = \sqrt{3v^2 + \lambda^2}. \tag{5}
 \end{aligned}$$

Under the Zeno condition, we obtain the effective Hamiltonian governing the evolution

$$H_{\text{eff}} = \delta\Omega_x (|\phi_2\rangle + |\phi_4\rangle)\langle\psi_1| + \delta\Omega_y (|\phi_2\rangle\langle\psi_9| + |\phi_4\rangle\langle\psi_{11}|) + H.c., \tag{6}$$

where $\delta = \frac{v}{\sqrt{2v^2 + \lambda^2}}$ is the normalization factors of the eigenstates $|\phi_2\rangle$ and $|\phi_4\rangle$. Then, we use four orthogonal vectors $|\mu_+\rangle = \frac{1}{\sqrt{2}}(|\phi_2\rangle + |\phi_4\rangle)$, $|\mu_-\rangle = \frac{1}{\sqrt{2}}(|\phi_2\rangle - |\phi_4\rangle)$, $|\vartheta_+\rangle = \frac{1}{\sqrt{2}}(|\psi_9\rangle + |\psi_{11}\rangle)$ and $|\vartheta_-\rangle = \frac{1}{\sqrt{2}}(|\psi_9\rangle - |\psi_{11}\rangle)$ to rewrite the Hamiltonian in Eq. (6) as

$$H_{\text{eff}}^1 = \sqrt{2}\delta\Omega_x|\mu_+\rangle\langle\psi_1| + \delta\Omega_y|\mu_+\rangle\langle\vartheta_+| + \delta\Omega_y|\mu_-\rangle\langle\vartheta_-| + H.c.. \tag{7}$$

It is obvious that when the initial state is $|\psi_1\rangle$, the terms containing $|\mu_-\rangle$ and $|\vartheta_-\rangle$ are negligible because they are decoupled to the time evolution of initial state. Then, we can obtain the final effective Hamiltonian

$$H_{\text{fe}} = \sqrt{2}\delta\Omega_x|\mu_+\rangle\langle\psi_1| + \delta\Omega_y|\mu_+\rangle\langle\vartheta_+| + H.c., \tag{8}$$

which can be treated as a simple three-level system with an excited state $|\mu_+\rangle$ and two ground states $|\psi_1\rangle$ and $|\vartheta_+\rangle$. Then, we obtain the instantaneous eigenstates of the final effective Hamiltonian H_{fe}

$$\begin{aligned} |n_0(t)\rangle &= \cos\theta|\psi_1\rangle + \sin\theta|\vartheta_+\rangle, \\ |n_{\pm}(t)\rangle &= \frac{1}{\sqrt{2}}(\sin\theta|\psi_1\rangle \pm \sqrt{2}\beta|\mu_+\rangle - \cos\theta|\vartheta_+\rangle), \end{aligned} \tag{9}$$

and corresponding eigenvalues $\iota_0 = 0$, $\iota_{\pm} = \pm\delta\beta$, respectively. Here, $\beta = \sqrt{2\Omega_x^2 + \Omega_y^2}$ and $\theta = \arctan -\frac{\sqrt{2}\Omega_x}{\Omega_y}$. It is obvious that the state $|\psi_1\rangle = |n_0(0)\rangle$ would follow $|n_0(t)\rangle$ closely, if the adiabatic condition $|\langle n_0(t)|\partial_t n_{\pm}(t)\rangle| \ll |\iota_{\pm}|$ is fulfilled. However, it is undesirable to obtain the target state, because this process would take quite a long time.

3 The shortcut scheme is proposed to generate W states based on TQD

3.1 Construct STAP by using TQD

As far as we know, the instantaneous states $|n_k(t)\rangle$ ($k = 0, \pm$) we mentioned in Eq. (9) do not meet the Schrödinger equation $i\partial_t|n_k(t)\rangle = H_{\text{fe}}(t)|n_k(t)\rangle$. The key to solving the problem is to construct shortcuts for a system governed by $H_0(t)$ with finding out a Hamiltonian $H(t)$ which drives the instantaneous eigenstates $|n_k(t)\rangle$ of $H_0(t)$ exactly. According to the Berry’s transitionless tracking algorithm [31], the Hamiltonian $H(t)$ can be reverse engineered from $H_0(t)$. And disregarding the effect of phases, we obtain the simplest Hamiltonian $H(t)$ in the form of

$$H(t) = i \sum_{k=0,\pm} |\partial_t n_k(t)\rangle\langle n_k(t)|. \tag{10}$$

Then, we can use transitionless tracking algorithm to construct the Hamiltonian that exactly drives the eigenstates of $H_{fe}(t)$,

$$H_{CCD} = i\dot{\theta}|\vartheta_+\rangle\langle\psi_1| + H.c., \tag{11}$$

where $\dot{\theta} = \sqrt{2}[\dot{\Omega}_y(t)\Omega_x(t) - \dot{\Omega}_x(t)\Omega_y(t)]/\beta^2$. However, it is necessary to find out an alternative physically feasible (APF) system whose effective Hamiltonian is equivalent to H_{CCD} , because obtaining the CDD Hamiltonian $H(t)$ is an outstanding challenge in present experimental condition. The model we used with three atoms trapped in coupled cavity is as the same as Fig. 1a. As shown in Fig. 1c, we make all the resonant atomic transitions into nonresonant atomic transitions with detuning Δ . The interaction Hamiltonian for the present system reads

$$\begin{aligned} \tilde{H}_o &= \tilde{H}_{al} + \tilde{H}_{ac} + H_e, \\ \tilde{H}_{al} &= \sum_{k=1}^3 \tilde{\Omega}_k(t)|e\rangle_k\langle f| + H.c., \\ \tilde{H}_{ac} &= \sum_{k=1}^3 \lambda a_k|e\rangle_k\langle g| + vb_1^\dagger(a_1 + a_2) + vb_2^\dagger(a_2 + a_3) + H.c., \\ H_e &= \sum_{k=1}^3 |e\rangle_k\langle e|. \end{aligned} \tag{12}$$

For the sake of simplicity, we set $\tilde{\Omega}_2 = \tilde{\Omega}_x$, $\tilde{\Omega}_1 = \tilde{\Omega}_3 = \tilde{\Omega}_y$, $\tilde{\lambda}_k = \lambda$ and $\tilde{v}_k = v$. Then, for the Hamiltonian in Eq. (12), performing similar processes from Eqs. (1)–(8), we can approximately neglect the terms containing the high oscillating frequencies. Meanwhile, the terms that are decoupled to the time evolution of initial state can also be neglected. We obtain an effective Hamiltonian for present system

$$\tilde{H}_{\text{eff}} = (\sqrt{2}\delta\tilde{\Omega}_x|\mu_+\rangle\langle\psi_1| + \delta\tilde{\Omega}_y|\mu_+\rangle\langle\vartheta_+| + H.c.) + 3\Delta\delta^2|\mu_+\rangle\langle\mu_+|, \tag{13}$$

where $\delta = \frac{v}{\sqrt{2v^2 + \lambda^2}}$. Then, we adiabatically eliminate the state $|\mu_+\rangle$ by using second-order perturbation approximation. The effective Hamiltonian becomes

$$\tilde{H}_{\text{eff}}^1 = \frac{2|\tilde{\Omega}_x|^2}{3\Delta}|\psi_1\rangle\langle\psi_1| + \frac{|\tilde{\Omega}_y|^2}{3\Delta}|\vartheta_+\rangle\langle\vartheta_+| + \left(\frac{\sqrt{2}\tilde{\Omega}_x\tilde{\Omega}_y^*}{3\Delta}|\vartheta_+\rangle\langle\psi_1| + H.c. \right). \tag{14}$$

When we choose $\tilde{\Omega}_x = \tilde{\Omega}_y^* = \Omega_s$, and $\tilde{\Omega}_x = \frac{i\Omega_s}{\sqrt{2}}$, the front two terms caused by Stark shift being removed, the above Hamiltonian becomes

$$\tilde{H}_{fe} = i\Omega|\vartheta_+\rangle\langle\psi_1| + H.c., \tag{15}$$

where $\Omega = \frac{\Omega_s^2}{3\Delta}$. If we set $\Omega = \dot{\theta}$, we can obtain $\tilde{H}_{fe} = H_{CCD}$. That means, under condition $\dot{\Omega}_x, \dot{\Omega}_y \ll \sqrt{3}\lambda, \sqrt{3}\Delta$, the Hamiltonian for speeding up the adiabatic dark-state evolution governed by H_0 has been constructed.

3.2 Generation of W states based on STAP

We will show that the generation of the W state based on STAP is much faster than the adiabatic passage. To generate a three-atom W state, we choose the Rabi frequencies as

$$\begin{aligned} \Omega_x(t) &= \Omega_0 \sin \alpha \exp \left[\frac{-(t - t_0 - t_f/2)^2}{t_c^2} \right], \\ \Omega_y(t) &= \Omega_0 \exp \left[\frac{-(t + t_0 - t_f/2)^2}{t_c^2} \right] + \Omega_0 \cos \alpha \exp \left[\frac{-(t - t_0 - t_f/2)^2}{t_c^2} \right], \end{aligned} \tag{16}$$

where Ω_0 is the amplitude and $\{t_0, t_c\}$ are related parameters. In order to meet the limited conditions, the time-dependent $\Omega_x(t)$ and $\Omega_y(t)$ are gotten as shown in Fig. 2 with parameters $\tan \alpha = 1, t_0 = 0.15t_f$ and $t_c = 0.2t_f$. As we mentioned above, the Ω_s is given

$$\Omega_s = \sqrt{3\Delta\dot{\theta}} = \sqrt{\frac{3\sqrt{2}\Delta(\dot{\Omega}_y\Omega_x - \dot{\Omega}_x\Omega_y)}{\beta^2}}. \tag{17}$$

If we set two dimensionless parameters

$$\eta_1 = \frac{t't_f - t_0 - 0.5t_f}{t_c},$$

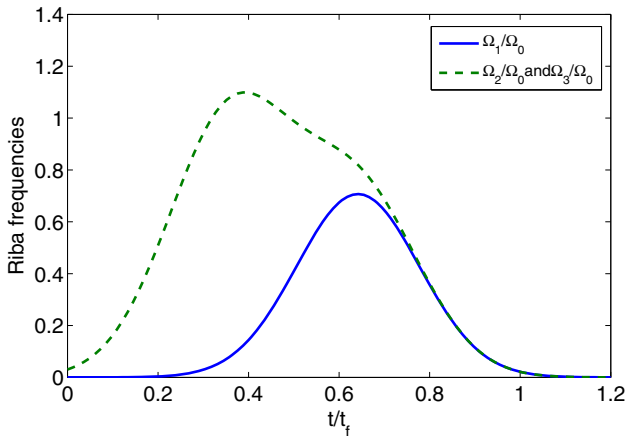


Fig. 2 Dependence on t/t_f of Ω_x/Ω_0 and Ω_y/Ω_0

$$\eta_2 = \frac{t't_f + t_0 - 0.5t_f}{t_c}, \tag{18}$$

where $t' = \frac{t}{t_f}$. Then, putting Eqs. (16) and (18) into Eq. (17), we obtain

$$\Omega_s = \sqrt{\frac{6\sqrt{2}\Delta R^2}{t_f}}, \tag{19}$$

where

$$R = \sqrt{\frac{\eta_1\Omega_x\Omega_y - \Omega_x\Omega_0(\eta_2e^{-\eta_2^2} + \cos\alpha\eta_1e^{-\eta_1^2})}{2\Omega_x^2 + \Omega_y^2}}. \tag{20}$$

We find out the amplitude of Ω_s is mainly dominated by $\chi = \sqrt{\frac{6\sqrt{2}\Delta}{t_f}}$, since the amplitude of R is close to 1. Therefore, the limited conditions become

$$\begin{aligned} \sqrt{\frac{6\sqrt{2}\Delta}{t_f}} &\ll \sqrt{3}\lambda \Rightarrow t_f \gg \frac{2\sqrt{2}\Delta}{\lambda^2}, \\ \sqrt{\frac{6\sqrt{2}\Delta}{t_f}} &\ll \sqrt{3}\Delta \Rightarrow t_f \gg \frac{2\sqrt{2}}{\Delta}. \end{aligned} \tag{21}$$

In order to meet the Zeno condition, it is better to choose a smaller Δ if the interaction time t_f is short, when λ is a constant value. However, a larger Δ is still required due to need to meet the large detuning condition. This can be demonstrated in Fig. 3

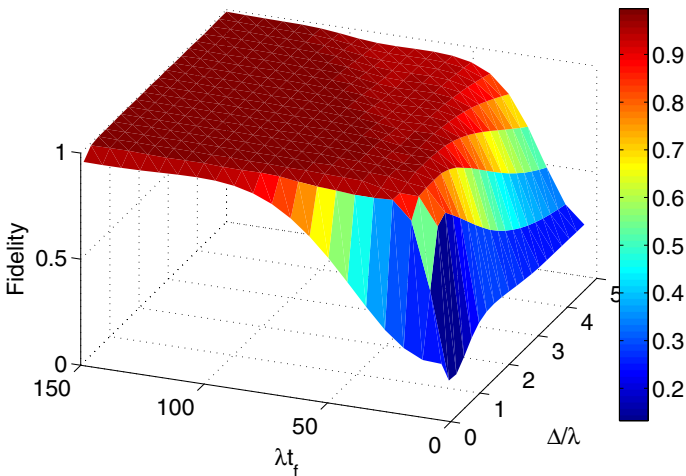


Fig. 3 Fidelity of the W state versus the interaction time λt_f and the detuning Δ/λ in the shortcut scheme

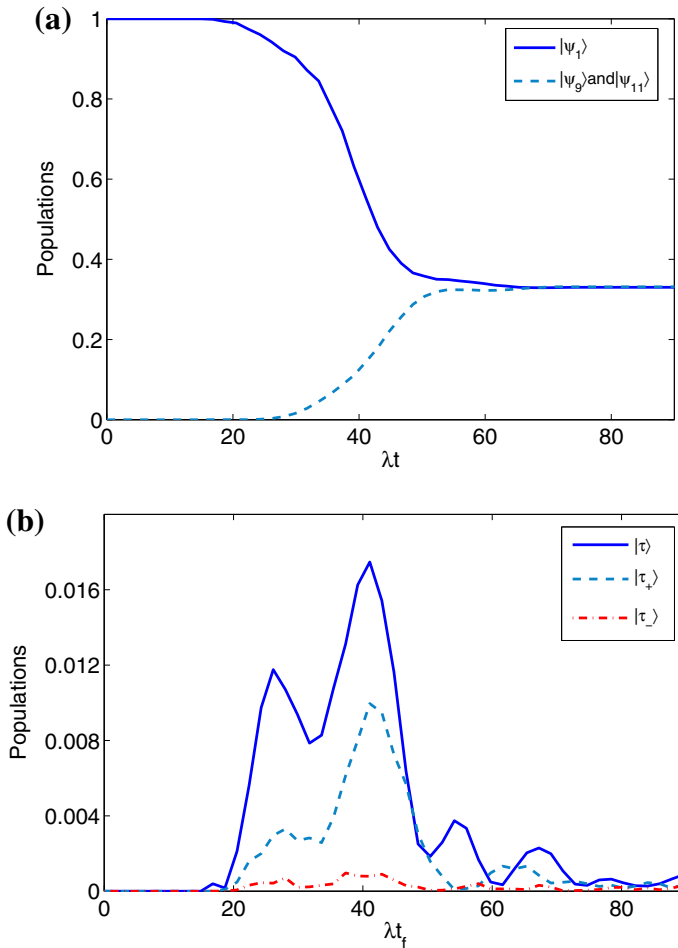


Fig. 4 **a** Time evolution of the populations for the states $|\psi_1\rangle$, $|\psi_9\rangle$ and $|\psi_{11}\rangle$ with $\{t_f = 75/\lambda$ and $\Delta = 3\lambda\}$. **b** Time evolution of the populations for the intermediate states $|\tau\rangle$, $|\tau_+\rangle$ and $|\tau_-\rangle$ with $\{t_f = 75/\lambda$ and $\Delta = 3\lambda\}$

which shows the fidelity of the W state versus parameters λt_f and Δ/λ . As shown in Fig. 3, a long operation time t_f is still required if Δ is too small or too large. The fidelity F of the state is defined as $F = |\langle \psi | \rho(t_f) | \psi \rangle|$, where $\rho(t_f)$ is the density operator of the given system when $t = t_f$. Then, in order to verify the shortcut scheme effectively, we contrast the performances of population transfer from the initial state $|\psi_1\rangle$. The time-dependent population for any state $|\psi\rangle$ is given through relationship $P = |\langle \psi | \rho(t) | \psi \rangle|$, where $\rho(t)$ is the corresponding time-dependent density operator. Figure 4a shows that a perfect populations transfer can be achieved with suitable parameters $\{\Delta = 3\lambda, t_f = 75/\lambda\}$. Unfortunately, slightly breaking the Zeno condition leads to a slight failure of the approximation during the stepped-up evolution. Wherefore, the parameter should be slightly corrected, i.e., a related

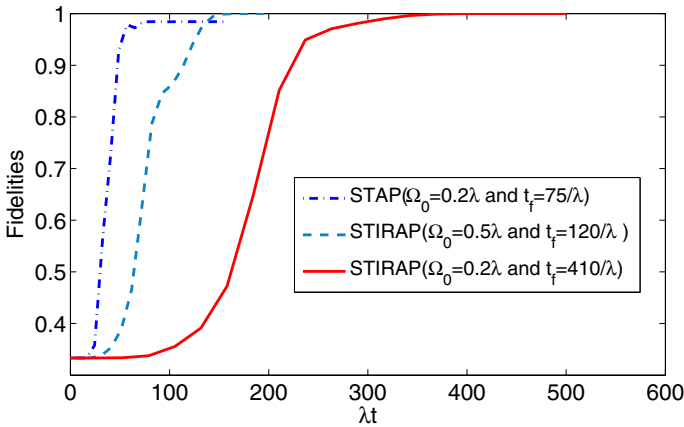


Fig. 5 Comparison between the fidelity of the W state via STIP and STIRAP

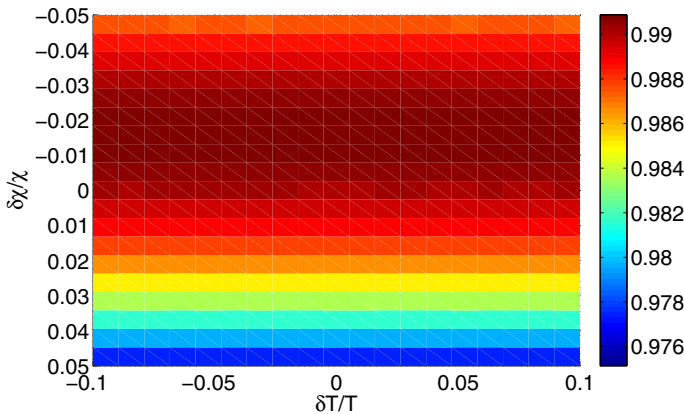


Fig. 6 Fidelity of the W state via STAP versus the variations of T and χ

parameter ($\tilde{\Omega}_x \rightarrow 1.05\tilde{\Omega}_x$). We plot the time evolution of the populations in intermediate states $|\tau\rangle = \frac{1}{\sqrt{3}}(|\psi_2\rangle + |\psi_8\rangle + |\psi_{10}\rangle)$, $|\tau_+\rangle = \frac{1}{\sqrt{3}}(|\psi_3\rangle + |\psi_4\rangle + |\psi_5\rangle)$ and $|\tau_-\rangle = \frac{1}{\sqrt{2}}(|\psi_6\rangle + |\psi_7\rangle)$ in Fig. 4b to prove this. Figure 4b shows that the intermediate states can be effectively neglected, because the populations of the intermediate states are almost zero. In Fig. 5, we compare the fidelity between our scheme and adiabatic scheme. Generally speaking, the interaction time required for the creation of the W state via adiabatic passage becomes more and more shorter with increasing of Ω_0 , but the RWA is not longer effective for the system if Ω_0 is relatively large. Figure 5 shows that the interaction time via adiabatic passage is still much longer than that via STAP, even if we choose $\Omega_0 = 0.5\lambda$. That means our scheme is much faster than the adiabatic passage. However, it is known to all that the fidelity of the W state is influenced with the mutative parameters. The fidelity of the W state versus the variation $\delta\chi$ and δT is plotted in Fig. 6. Here, we define $\delta x = x' - x$ as the deviation of any

parameter x , where x' is the actual value and x is the ideal value. As shown in Fig. 6, the fidelity almost keeps unchanging with the variation $\delta T (T = 85/\lambda)$. Figure 6 also shows a deviation $|\delta\chi/\chi| = 5\%$ which means the variation in the amplitude of Ω_s only causes a reduction about 1.5% in the fidelity. That is to say, our scheme is robust against variations in the experimental parameters.

Now, we will investigate the influence of various decoherence caused by the atomic spontaneous emission, cavity losses and the fiber photon leakages. The master equation of the whole system reads

$$\begin{aligned} \dot{\rho} = & -i[H_0, \rho] + \sum_{k=1}^3 \frac{\gamma_k}{2} (2\sigma_k^- \rho \sigma_k^+ - \sigma_k^+ \sigma_k^- \rho - \rho \sigma_k^+ \sigma_k^-) \\ & + \sum_{k=1}^3 \frac{\kappa_k}{2} (2a_k^- \rho a_k^+ - a_k^+ a_k^- \rho - \rho a_k^+ a_k^-) \\ & + \sum_{k=1}^2 \frac{\beta_k}{2} (2b_k^- \rho b_k^+ - b_k^+ b_k^- \rho - \rho b_k^+ b_k^-), \end{aligned} \tag{22}$$

where γ_k denotes the atomic spontaneous decay rate of the k th atom; κ_k and β_k denote the decay rates of the k th cavity and fiber, respectively. For simplicity, we set $\gamma_1 = \gamma_2 = \gamma_3 = \gamma$, $\kappa_1 = \kappa_2 = \kappa_3 = \kappa$ and $\beta_1 = \beta_2 = \beta$. The master equation can be numerically solved in the subspace $\tilde{\mathcal{V}} \in \{|g, g, g\rangle_{1,2,3}|0, 0, 0\rangle_{c1,c2,c3}|0, 0\rangle_{f1,f2}\}$. Then, we plot the fidelity versus the dimensionless parameters γ/λ , κ/λ and β/λ in Fig. 7. We can see the fidelity is almost unaffected by the fiber loss rate β , because the fiber length is very short. Thus, the cavity decay and the atomic spontaneous emission decay become the main sources of decoherence. It is necessary to focus on the influences of the atomic spontaneous emission rate γ and the cavity decay rate κ . As shown in Fig. 8, the scheme is more sensitive to atomic spontaneous emissions than the cavity decay, since the Zeno condition is not satisfied faultlessly. Therefore, the populations of intermediate states ($|\tau\rangle$, $|\tau_+\rangle$) and $|\tau_-\rangle$) cannot be completely suppressed. Fortunately, we still can create a three-atom W state with a high fidelity 96.27%, even if we set $\gamma = \kappa = 0.01\lambda$. This means the scheme is robust since the fidelity decreases slowly.

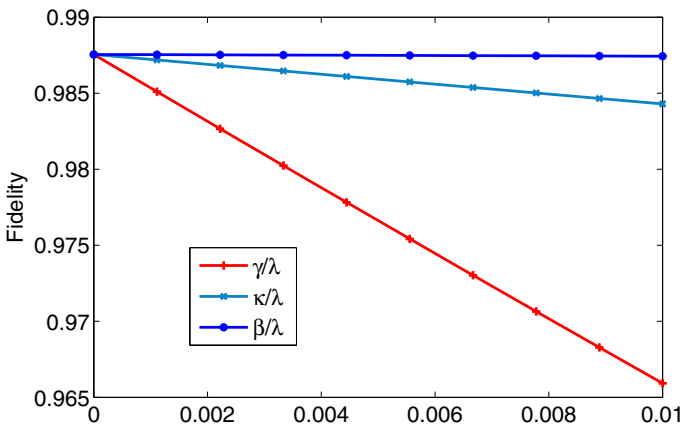


Fig. 7 Fidelity of the W state via STAP versus the decoherence parameters γ/λ , κ/λ and β/λ with $t_f = 75/\lambda$ and $\Delta = 3\lambda$

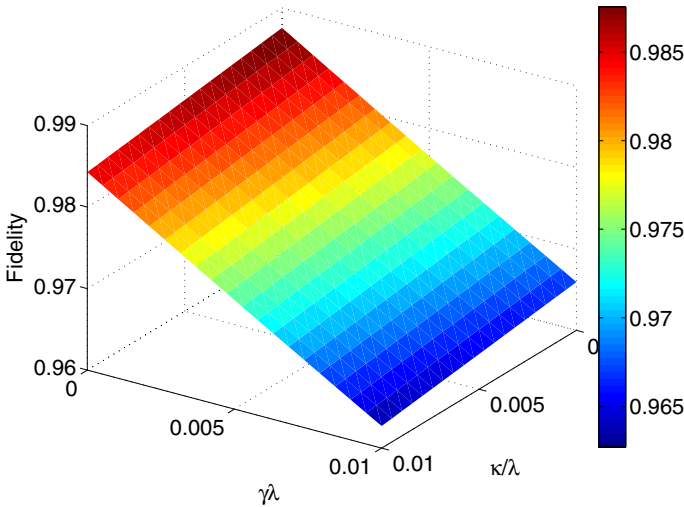


Fig. 8 Dependences on the decoherence parameters γ/λ and κ/λ of the fidelity of the W state via STAP with $t_f = 75/\lambda$ and $\Delta = 3\lambda$

4 Generation of N -atom W states by STAP

In this section, we give a brief description about generation of N -atom ($N > 3$) W states in a similar system by the same principle. As shown in Fig. 9, we assume that N identical Λ -type atoms are trapped in N spatially separated cavities C_1, C_2, \dots, C_N , respectively. The cavities are linked by optical fibers f_1, f_2, \dots, f_{N-1} . The level configurations of the atoms are the same as we mentioned in Fig. 1b. Then, the Hamiltonian for the system in the interaction picture reads

$$\begin{aligned}
 \hat{H}_o &= \hat{H}_{al} + \hat{H}_{ac} + \hat{H}_e, \\
 \hat{H}_{al} &= \sum_{k=1}^N \Omega_k(t) |e\rangle_k \langle f| + H.c., \\
 \hat{H}_{ac} &= \sum_{k=1}^N \lambda_k a_k |e\rangle_k \langle g| + \sum_{k=1}^{N-1} v_k b_k^\dagger (a_1 + a_{k+1}) + H.c., \\
 \hat{H}_e &= \sum_{k=1}^N |e\rangle_k \langle e|.
 \end{aligned} \tag{23}$$

We assume the initial state is $|f, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}$. The whole system evolves in the subspace spanned \check{V} by

$$\begin{aligned}
 |\check{\psi}_1\rangle &= |f, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\check{\psi}_2\rangle &= |e, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}},
 \end{aligned}$$

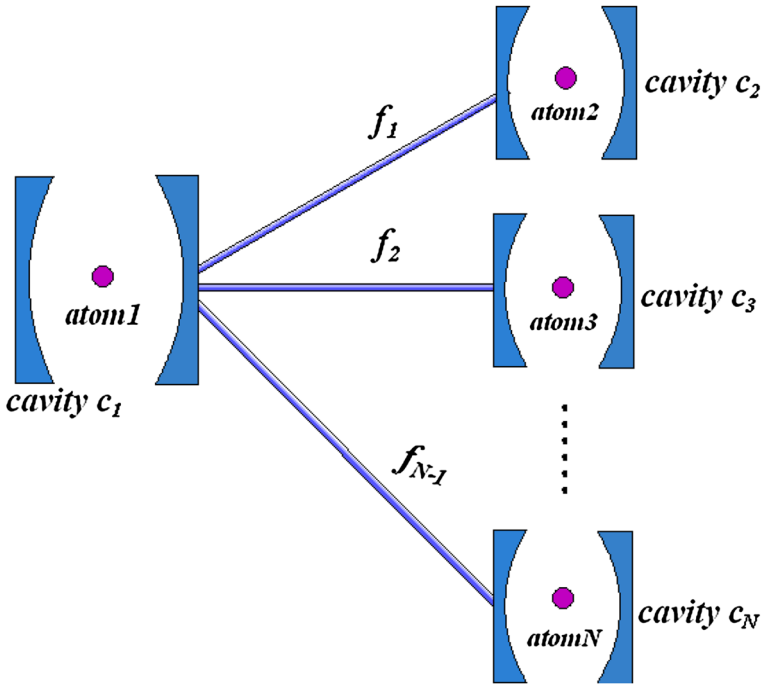


Fig. 9 Experimental setup diagram for generation of N -atom W states

$$\begin{aligned}
 |\psi_3\rangle &= |g, g, g, \dots, g\rangle_{1,2,\dots,N} |1, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_4\rangle &= |g, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |1, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_5\rangle &= |g, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 1, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_6\rangle &= |g, e, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_7\rangle &= |g, f, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_8\rangle &= |g, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 1, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_9\rangle &= |g, g, g, \dots, g\rangle_{1,2,\dots,N} |0, 0, 1, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_{10}\rangle &= |g, g, e, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 |\psi_{11}\rangle &= |g, g, f, \dots, g\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}, \\
 &\dots\dots\dots \\
 |\psi_{4N-1}\rangle &= |g, g, g, \dots, f\rangle_{1,2,\dots,N} |0, 0, 0, \dots, 0\rangle_{c_1, c_2, \dots, c_N} |0, 0, \dots, 0\rangle_{f_1, f_2, \dots, f_{N-1}}.
 \end{aligned}
 \tag{24}$$

The procedure is the same as we mentioned in Sect. 3, and we get an effective Hamiltonian (we set $\hat{\Omega}_2 = \hat{\Omega}_3 = \dots = \hat{\Omega}_N = \hat{\Omega}_y$)

$$\begin{aligned} \hat{H}_{\text{eff}} = & \frac{(N-1)|\hat{\Omega}_1|^2}{N\Delta} |\hat{\psi}_1\rangle\langle\hat{\psi}_1| + \frac{|\hat{\Omega}_y|^2}{N\Delta} |\hat{\psi}_+\rangle\langle\hat{\psi}_+| \\ & + \left(\frac{\sqrt{N-1}\hat{\Omega}_1\hat{\Omega}_y^*}{N\Delta} |\hat{\psi}_+\rangle\langle\hat{\psi}_1| + H.c. \right). \end{aligned} \tag{25}$$

When we choose $\hat{\Omega}_y = \hat{\Omega}_s$ and $\hat{\Omega}_1 = \frac{i\hat{\Omega}_s}{\sqrt{N-1}}$ (here $\hat{\Omega}_s$ is a real number), the above Hamiltonian becomes

$$\hat{H}_{\text{fe}} = i\hat{\Omega} |\hat{\psi}_+\rangle\langle\hat{\psi}_1| + H.c., \tag{26}$$

where $\hat{\Omega} = \frac{|\hat{\Omega}_s|^2}{N\Delta}$, $|\hat{\psi}_+\rangle = \frac{1}{\sqrt{N-1}}(|\hat{\psi}_7\rangle + |\hat{\psi}_{11}\rangle + \dots + |\hat{\psi}_{4N-1}\rangle)$ ($N = 4, 5, \dots, N$). As long as $\hat{\Omega} = \hat{\Omega}$, $\hat{H}_{\text{fe}} = H_{\text{CCD}}$, the Hamiltonian for speeding up the adiabatic dark-state evolution governed by \hat{H}_0 has been constructed. Wherefore, we can obtain N -atom W states, i.e., $|\hat{\psi}(t)\rangle = \frac{1}{\sqrt{2}}(|\hat{\psi}_1\rangle + |\hat{\psi}_+\rangle) = \frac{1}{\sqrt{N}}(|\hat{\psi}_1\rangle + |\hat{\psi}_7\rangle + |\hat{\psi}_{11}\rangle + \dots + |\hat{\psi}_{4N-1}\rangle)$.

5 Experimental feasibility and conclusions

Finally, let us consider the experimental feasibility of the proposed scheme. In experiment, this scheme can be realized in the strong-coupling regime [31,66], if the cesium atoms can be cooled and trapped in coupled cavity. Furthermore, a set of cavity quantum electrodynamics (QED) parameters $(\lambda, \gamma, \kappa)/2\pi = (750, 2.62, 3.5)$ MHz is predicted to be available [67], with the cavity mode wavelength about 850 nm. In this condition, we can obtain a relatively high fidelity 98.42 % in the shortcut scheme.

In summary, we have proposed an efficient theoretical scheme to generate W states for three atoms trapped in coupled cavities linked by optical fibers based on STAP by combining ‘‘TQD’’ with ‘‘QZD.’’ The influences of the decoherence such as atomic spontaneous emission, cavity losses and the fiber photon leakages are numerically studied. Numerical simulations demonstrate that the shortcut scheme is faster and robust against the decoherence. Additionally, N -atom W states can also be generated when we set $\hat{\Omega}_1 = \frac{i\hat{\Omega}_s}{\sqrt{N-1}}$ and $\hat{\Omega}_2 = \hat{\Omega}_3 = \dots = \hat{\Omega}_N = \hat{\Omega}_s$. In fact, the shortcut is possible to be constructed with the same method presented in this scheme, so long as the quantum system whose Hamiltonian is possible to be simplified into a similar form in Eq. (8). That means the shortcut method is useful to realize fast and noise-resistant quantum information processing for multi-particle systems.

Acknowledgments This work was supported by the National Natural Science Foundation of China under Grant No. 11404061.

References

1. Ekert, A.K.: Quantum cryptography based on Bell’s theorem. *Phys. Rev. Lett.* **67**, 661–663 (1991)
2. Gisin, N., Massar, S.: Optimal quantum cloning machines. *Phys. Rev. Lett.* **79**, 2153–2156 (1997)

3. Bell, J.S.: On the Einstein–Podolsky–Rosen paradox. *Physics* (Long Island City, NY) **1**, 195–200 (1965)
4. Greenberger, D.M., Horne, M.A., Shimony, A., Zeilinger, A.: Bell’s theorem without inequalities. *Am. J. Phys.* **58**, 1131–1143 (1990)
5. Dür, W., Vidal, G., Cirac, J.I.: Three qubits can be entangled in two inequivalent ways. *Phys. Rev. A* **62**, 062314(12) (2000)
6. Bennett, C.H., Brassard, G., Crépeau, C., Jozsa, R., Peres, A., Wootters, W.K.: Teleporting an unknown quantum state via dual classical and Einstein–Podolsky–Rosen channels. *Phys. Rev. Lett.* **70**, 1895–1899 (1993)
7. An, N.B.: Cavity-catalyzed deterministic generation of maximal entanglement between nonidentical atoms. *Phys. Lett. A* **344**, 77–83 (2005)
8. Zhang, C.L., Li, W.Z., Chen, M.F.: Preparation of W state among spatially separated atomic ensembles collectively controlled by a single atom via adiabatic passage. *Opt. Commun.* **311**, 301–306 (2013)
9. Song, J., Xia, Y., Song, H.S.: Entangled state generation via adiabatic passage in two distant cavities. *J. Phys. B* **40**, 4503–4511 (2007)
10. Bastin, T., Thiel, C., von Zanthier, J., Lamata, L., Solano, E., Agarwal, G.S.: Operational determination of multiqubit entanglement classes via tuning of local operations. *Phys. Rev. Lett.* **102**, 053601(4) (2009)
11. Wang, X.W., Yang, G.J., Su, Y.H., Xie, M.: Simple schemes for quantum information processing with W-type entanglement. *Quantum Inf. Process.* **8**, 431–442 (2009)
12. Chen, R.X., Shen, L.T.: Tripartite entanglement of atoms trapped in coupled cavities via quantum Zeno dynamics. *Phys. Lett. A* **375**, 3840–3844 (2011)
13. Zhong, Z.R.: Generation of entanglement of multiple atoms via distant cavities. *Opt. Commun.* **283**, 1972–1974 (2010)
14. Lu, M., Xia, Y., Song, J., An, N.B.: Generation of N-atom W-class states in spatially separated cavities. *J. Opt. Soc. Am. B* **30**, 2142–2147 (2013)
15. Chen, Y.H., Xia, Y., Song, J.: Effective protocol for generation of multiple atoms entangled states in two coupled cavities via adiabatic passage. *Quantum Inf. Process.* **12**, 3771–3783 (2013)
16. Castro, A.O., Johnson, N.F., Quiroga, L.: Scheme for on-resonance generation of entanglement in time-dependent asymmetric two-qubit-cavity systems. *Phys. Rev. A* **70**, 020301(R)(4) (2004)
17. Zheng, S.B.: Scalable generation of multi-atom W states with a single resonant interaction. *J. Opt. B Quantum Semiclass. Opt.* **7**, 10–13 (2005)
18. Fewell, M.P., Shore, B.W., Bergmann, K.: Coherent population transfer among three states: full algebraic solutions and the relevance of non adiabatic processes to transfer by delayed pulses. *Aust. J. Phys.* **50**, 281–308 (1997)
19. Bergmann, K., Theuer, H., Shore, B.W.: Coherent population transfer among quantum states of atoms and molecules. *Rev. Mod. Phys.* **70**, 1003–1025 (1998)
20. Vitinov, N.V., Halfmann, T., Shore, B.W., Bergmann, K.: Laser-induced population transfer by adiabatic passage techniques. *Annu. Rev. Phys. Chem.* **52**, 763–809 (2001)
21. Král, P., Thanopoulos, L., Shapiro, M.: Coherently controlled adiabatic passage. *Rev. Mod. Phys.* **79**, 53–77 (2007)
22. Misra, B., Sudarshan, E.C.G.: The Zenos paradox in quantum theory. *J. Math. Phys.* **18**, 756 (1977)
23. Itano, W.M., Heinzen, D.J., Bollinger, J.J., Wineland, D.J.: Quantum Zeno effect. *Phys. Rev. A* **41**, 2295–2300 (1990)
24. Kwiat, P., Weinfurter, H., Herzog, T., Zeilinger, A., Kasevich, M.A.: Interaction-free measurement. *Phys. Rev. Lett.* **74**, 4763–4766 (1995)
25. Facchi, P., Gorini, V., Marmo, G., Pascazio, S., Sudarshan, E.C.G.: Quantum Zeno dynamics. *Phys. Lett. A* **275**, 12–19 (2000)
26. Facchi, P., Pascazio, S.: Quantum Zeno subspaces. *Phys. Rev. Lett.* **89**, 080401(4) (2002)
27. Chen, X., Lizuain, I., Ruschhaupt, A., Guéry-Odelin, D., Muga, J.G.: Shortcut to adiabatic passage in two- and three-level atoms. *Phys. Rev. Lett.* **105**, 123003(4) (2010)
28. Torrontegui, E., Ibáñez, S., Modugno, S., Campo, A.D., Gué-Odelin, D., Ruschhaupt, A., Chen, X., Muga, J.G.: Shortcuts to adiabaticity. *Adv. Atomic Mol. Opt. Phys.* **62**, 117–169 (2013)
29. Campo, A.D.: Shortcuts to adiabaticity by counterdiabatic driving. *Phys. Rev. Lett.* **111**, 100502(5) (2013)
30. Masuda, S., Nakamura, K.: Acceleration of adiabatic quantum dynamics in electromagnetic fields. *Phys. Rev. A* **84**, 043434(11) (2011)

31. Berry, M.V.: Transitionless quantum driving. *J. Phys. A* **42**, 365303(9) (2009)
32. Lu, M., Xia, Y., Shen, L.T., Song, J., An, N.B.: Shortcuts to adiabatic passage for population transfer and maximum entanglement creation between two atoms in a cavity. *Phys. Rev. A* **89**, 012326(7) (2014)
33. Chen, Y.H., Xia, Y., Chen, Q.Q., Song, J.: Efficient shortcuts to adiabatic passage for fast population transfer in multiparticle systems. *Phys. Rev. A* **89**, 033856(12) (2014)
34. Lu, M., Xia, Y., Shen, L.T., Song, J.: Using shortcut to adiabatic passage for the ultrafast quantum state transfer in cavity QED system. *Laser Phys.* **24**, 105201(7) (2014)
35. Chen, Y.H., Xia, Y., Chen, Q.Q., Song, J.: Shortcuts to adiabatic passage for multiparticles in distant cavities: applications to fast and noise-resistant quantum population transfer, entangled states preparation and transition. *Laser Phys. Lett.* **11**, 115201(14) (2014)
36. Chen, Y.H., Xia, Y., Chen, Q.Q., Song, J.: Fast and noise-resistant implementation of quantum phase gates and creation of quantum entangled states. *Phys. Rev. A* **91**, 012325(15) (2015)
37. Chen, Y.H., Xia, Y., Chen, Q.Q., Song, J.: Shortcuts to adiabatic passage for fast generation of Greenberger–Horne–Zeilinger states by transitionless quantum driving (2014). [arXiv:1411.6747](https://arxiv.org/abs/1411.6747)
38. Muga, J.G., Chen, X., Ruschhaupt, A., Guéry-Odelin, D.: Frictionless dynamics of Bose–Einstein condensates under fast trap variations. *J. Phys. B* **42**, 241001(4) (2009)
39. Chen, X., Ruschhaupt, A., Schmidt, S., Campo, A.D., Guéry-Odelin, D., Muga, J.G.: Fast optimal frictionless atom cooling in harmonic traps: shortcut to adiabaticity. *Phys. Rev. Lett.* **104**, 063002(4) (2010)
40. Chen, X., Muga, J.G.: Transient energy excitation in shortcuts to adiabaticity for the time-dependent harmonic oscillator. *Phys. Rev. A* **82**, 053403(7) (2010)
41. Schaff, J.F., Capuzzi, P., Labeyrie, G., Vignolo, P.: Shortcuts to adiabaticity for trapped ultracold gases. *New J. Phys.* **13**, 113017(31) (2011)
42. Torrontegui, E., Ibáñez, S., Chen, X., Ruschhaupt, A., Guéry-Odelin, D., Muga, J.G.: Fast atomic transport without vibrational heating. *Phys. Rev. A* **83**, 013415(9) (2011)
43. Chen, X., Torrontegui, E., Stefanatos, D., Li, J.S., Muga, J.G.: Optimal trajectories for efficient atomic transport without final excitation. *Phys. Rev. A* **84**, 043415(9) (2011)
44. Torrontegui, E., Chen, X., Modugno, M., Schmidt, S., Ruschhaupt, A., Muga, J.G.: Fast transport of Bose–Einstein condensates. *New J. Phys.* **14**, 013031(11) (2012)
45. Li, Y., Wu, L.A., Wang, Z.D.: Fast ground-state cooling of mechanical resonators with time-dependent optical cavities. *Phys. Rev. A* **83**, 043804(5) (2011)
46. Campo, A.D.: Frictionless quantum quenches in ultracold gases: a quantum-dynamical microscope. *Phys. Rev. A* **84**, 031606(R)(4) (2011)
47. Campo, A.D.: Frictionless quantum quenches in ultracold gases: a quantum-dynamical microscope. *Phys. Rev. A* **96**, 60005 (2011)
48. Ruschhaupt, A., Chen, X., Alonso, D., Muga, J.G.: Optimally robust shortcuts to population inversion in two-level quantum systems. *New J. Phys.* **14**, 093040(19) (2012)
49. Schaff, J.F., Song, X.L., Vignolo, P., Labeyrie, G.: Fast optimal transition between two equilibrium states. *Phys. Rev. A* **82**, 033430(5) (2010)
50. Schaff, J.F., Song, X.L., Capuzzi, P., Vignolo, P., Labeyrie, G.: Shortcut to adiabaticity for an interacting Bose–Einstein condensate. *Eur. Phys. Lett.* **93**, 23001(5) (2011)
51. Godsilal, C., Kirkland, S., Severini, S., Smith, J.: Number-theoretic nature of communication in quantum spin systems. *Phys. Rev. Lett.* **109**, 050502(4) (2012)
52. Tseng, S.Y., Chen, X.: Engineering of fast mode conversion in multimode waveguides. *Opt. Lett.* **37**, 5118 (2012)
53. Chen, X., Torrontegui, E., Muga, J.G.: Lewis–Riesenfeld invariants and transitionless quantum driving. *Phys. Rev. A* **83**, 062116(8) (2011)
54. Song, L.C., Xia, Y., Song, J.: Noise resistance of Toffoli gate in an array of coupled cavities. *J. Mod. Opt.* **61**, 1290–1297 (2014)
55. Liang, Y., Wu, Q.C., Su, S.L., Ji, X., Zhang, S.: Shortcuts to adiabatic passage for multiqubit controlled-phase gate. *Phys. Rev. A* **91**, 032304(8) (2015)
56. Liang, Y., Ji, X.: Shortcuts to adiabatic passage for multiqubit controlled phase gate (2014). [arXiv:1411.7434](https://arxiv.org/abs/1411.7434)
57. Pellizzari, T.: Quantum networking with optical fibres. *Phys. Rev. Lett.* **79**, 5242–5245 (1997)
58. Enk, S.J.V., Kimble, H.J., Cirac, J.L., Zoller, P.: Quantum communication with dark photons. *Phys. Rev. A* **59**, 2659–2664 (1999)

59. Zheng, S.B.: Generation of Greenberger–Horne–Zeilinger states for multiple atoms trapped in separated cavities. *Eur. Phys. J. D* **54**, 719–722 (2005)
60. Ye, S.Y., Zhong, Z.R., Zheng, S.B.: Deterministic generation of three-dimensional entanglement for two atoms separately trapped in two optical cavities. *Phys. Rev. A* **77**, 014303(4) (2008)
61. Giampaolo, S.M., Illuminati, F.: Long-distance entanglement and quantum teleportation in coupled-cavity arrays. *Phys. Rev. A* **80**, 050301(R)(4) (2009)
62. Zhang, K., Li, Z.Y.: Transfer behavior of quantum states between atoms in photonic crystal coupled cavities. *Phys. Rev. A* **81**, 033843(10) (2010)
63. Raussendorf, R., Briegel, H.J.: A one-way quantum computer. *Phys. Rev. Lett.* **86**, 5188–5191 (2001)
64. Zhong, Z.R., Huang, X.B.: An unconventional geometric phase gate between two arbitrary qubits in a two-dimensional network. *J. Opt. Soc. Am. B* **32**, 258–264 (2015)
65. Serafini, A., Mancini, S., Bose, S.: Distributed quantum computation via optical fibers. *Phys. Rev. Lett.* **96**, 010503(4) (2006)
66. Ye, J., Vernooy, D.W., Kimble, H.J.: Trapping of single atoms in cavity QED. *Phys. Rev. Lett.* **83**, 4987–4990 (1999)
67. Spillane, S.M., Kippenberg, T.J., Vahala, K.J., Goh, K.W., Wilcut, E., Kimble, H.J.: Ultrahigh-Q toroidal microresonators for cavity quantum electrodynamics. *Phys. Rev. A* **71**, 013817(10) (2005)