Generation of four-atom Greenberger–Horn–Zeilinger state via adiabatic passage*

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We propose a scheme to generate a Greenberger–Horn–Zeilinger (GHZ) state of four atoms trapped in a two-mode optical cavity via an adiabatic passage. The scheme is robust against moderate fluctuations of the experimental parameters. Numerical calculations show that the excited probabilities of both the cavity modes and the atoms are tiny and depend on the pulse peaks of the classical laser fields. For certain decoherence due to the atomic spontaneous emission and the cavity decay, there exits a range of pulse peaks to get a high fidelity.

Keywords: Greenberger–Horn–Zeilinger (GHZ) state, adiabatic passage, cavity quantum electrodynamics

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1. Introduction

Entanglement, as a fundamental feature of quantum mechanics, has practical applications in quantum information processing (QIP), such as quantum cryptography,[1] quantum dense coding,[2] and quantum teleportation.[3] The emblematic bipartite entangled states are the Bell states,[4] which have stronger correlations than those allowed by any local hidden variable theory. In contrast to the biparticle entanglement, the multipartite entanglement is much more complex. It has been illustrated that the multipartite entanglement has advantages over the biparticle entanglement in many ways.[5] There exist various kinds of multipartite entanglements, such as W state,[6] Greenberger–Horn–Zeilinger (GHZ) state,[6,7] and χ state.[8] In recent years, many schemes have been proposed to generate the multiparticle entangled states. In Ref. [9], the authors showed how to generate atomic entangled states with a selective resonant interaction in cavity quantum electrodynamic. Cho *et al.*[10] proposed a scheme of optical pumping, by which a system of atoms coupled to harmonic oscillators is driven to an entangled steady state through the atomic spontaneous emission. In Ref. [11], a new kind of four-mode continuous variable coherent-entangled state was proposed in the Fock space by using the technique of integration within an ordered product. Song *et al.*[12] demonstrated a feedback strategy to create a multi-body Dicke state in a coupled cavity system recently. The GHZ state is maximally entangled in several senses.[13] For instance, it maximally violates the Bell inequalities; the mutual information of measurement outcomes is maximum; it is maximally stable against (white) noise; and one can locally obtain an Einstei–Podolsky–Rosen (EPR) state shared between any two of the particles from a GHZ state with unit probability. Thus, different proposals in experiment[14–16] and theory[17–21] have been investigated for generating the GHZ states.

The physical systems utilized to generate entanglement include superconducting circuits,[22,23] linear optical systems,[24] cavity quantum electrodynamics (QED),[25] trapped ions,[26] quantum dots,[27] etc. It is well known that the cavity QED is an ideal candidate for quantum communication and quantum state engineering. The adiabatic techniques have attracted much attention due to their properties. The outstanding merits of using an adiabatic passage are that it does not need the precise control of the Rabi frequency and the pulse duration, and the population in excited states is negligible due to their absence in the dark state used for evolution. Thus, the techniques of stimulated Raman adiabatic passage (STIRAP)[28] and fractional stimulated Raman adiabatic passage (f-STIRAP)[29] have been extensively used for QIP.[30–36]

In this paper, we propose a scheme for generating a GHZ state of four atoms trapped in a two-mode optical cavity via STIRAP based on the resonant atom–cavity interaction. Our scheme does not need to accurately control the evolution time, which reduces the complexity in experiment. Numerical calculations show that there exits a suitable range of pulse parameters. Thus, the scheme is insensitive to small fluctuations of the parameters. On the other hand, the excited probabilities of the cavity modes and the atoms are extremely small during the evolution, and closely related to the pulse peaks of the classical laser fields. In an actual experiment, we can select reasonable pulse peaks according to the decoherence parameters to get a high fidelity.

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2. Generating four-atom GHZ state via adiabatic passage

We consider four identical six-level atoms trapped in a two-mode cavity. As shown in Fig. 1, the atoms have two degenerate excited states $|e_{\rm L}\rangle_i$ and $|e_{\rm R}\rangle_i$ $(i=1,2,3,4)$, two degenerate ground states $|g_{\rm L}\rangle_i$ and $|g_{\rm R}\rangle_i$, and two intermediate states $|f_{\rm L}\rangle_i$ and $|f_{\rm R}\rangle_i$, where subscript i denotes the i-th atom in the cavity. The atomic transitions $|e_{\rm L}\rangle_i \leftrightarrow |g_{\rm L}\rangle_i$ and $|e_{\rm R}\rangle_i \leftrightarrow |g_{\rm R}\rangle_i$ are coupled resonantly to the left-circularly and the right-circularly polarized modes of the cavity with coupling coefficients $g_{i{\rm L}}$ and $g_{i{\rm R}}$, respectively. The l-th $(l=1,3)$ atomic transition $|f_{\rm L}\rangle_l \leftrightarrow |e_{\rm L}\rangle_l$ is driven resonantly by a classical laser field with coupling coefficient $\Omega_{l{\rm R}}$, while the r-th $(r=2,4)$ atomic transition $|f_{\rm R}\rangle_r \leftrightarrow |e_{\rm R}\rangle_r$ is driven resonantly by a classical laser field with coupling coefficient $\Omega_{r{\rm R}}$. In the interaction picture, under the dipole and the rotating wave approximations, the Hamiltonian for the whole system can be written as $(\hbar=1)$

$$\begin{aligned} H_{\rm I} = & \Bigg[\sum_{i=1}^{4}(g_{i{\rm L}}a_{\rm L}|e_{\rm L}\rangle_{ii}\langle g_{\rm L}| + g_{i{\rm R}}a_{\rm R}|e_{\rm R}\rangle_{ii}\langle g_{\rm R}|) \\ & + \sum_{l=1,3}(\Omega_{l{\rm R}}|e_{\rm L}\rangle_{ll}\langle f_{\rm L}|) + \sum_{r=2,4}(\Omega_{r{\rm R}}|e_{\rm R}\rangle_{rr}\langle f_{\rm R}|)\Bigg] \\ & + {\rm H.c.}, \end{aligned} \tag{1}$$

where $a_{\rm L}$ and $a_{\rm R}$ are the annihilation operators for the left-circularly and the right-circularly polarized cavity modes, respectively. For simplicity, we assume $g_{i{\rm L}} = g_{i{\rm R}} = g$, $\Omega_{1{\rm L}} = \Omega_{2{\rm R}} = \Omega_{\rm A}$, and $\Omega_{3{\rm L}} = \Omega_{4{\rm R}} = \Omega_{\rm B}$. If the system is initially prepared in state $|f_{\rm L}\rangle_1|f_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}$, the evolution subspace $\forall_1$ is spanned by the state vectors

$$\begin{aligned} |\varphi_1\rangle &= |f_{\rm L}\rangle_1|f_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_2\rangle &= |e_{\rm L}\rangle_1|f_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_3\rangle &= |g_{\rm L}\rangle_1|f_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|1\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_4\rangle &= |g_{\rm L}\rangle_1|f_{\rm R}\rangle_2|e_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_5\rangle &= |g_{\rm L}\rangle_1|f_{\rm R}\rangle_2|f_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_6\rangle &= |f_{\rm L}\rangle_1|e_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_7\rangle &= |f_{\rm L}\rangle_1|g_{\rm R}\rangle_2|g_{\rm L}\rangle_3|g_{\rm R}\rangle_4|0\rangle_{\rm L}|1\rangle_{\rm R}, \\ |\varphi_8\rangle &= |f_{\rm L}\rangle_1|g_{\rm R}\rangle_2|g_{\rm L}\rangle_3|e_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}, \\ |\varphi_9\rangle &= |f_{\rm L}\rangle_1|g_{\rm R}\rangle_2|g_{\rm L}\rangle_3|f_{\rm R}\rangle_4|0\rangle_{\rm L}|0\rangle_{\rm R}. \end{aligned} \tag{2}$$

Here, we adopt the notation $|n\rangle_p$ $(n=0,1;\ p={\rm L},{\rm R})$, where n denotes the photon number in the left-circularly $(p={\rm L})$ or right-circularly $(p={\rm R})$ polarized cavity mode. The dark state, which is the eigenstate of zero eigenvalue, is

$$\begin{aligned} |\psi_0\rangle = & N(\Omega_{\rm B}g|\varphi_1\rangle - \Omega_{\rm A}\Omega_{\rm B}|\varphi_3\rangle + \Omega_{\rm A}g|\varphi_5\rangle \\ & - \Omega_{\rm A}\Omega_{\rm B}|\varphi_7\rangle + \Omega_{\rm A}g|\varphi_9\rangle), \end{aligned} \tag{3}$$

where N is the normalization coefficient given by

$$N = (\Omega_{\rm B}^2g^2 + 2\Omega_{\rm A}^2\Omega_{\rm B}^2 + 2\Omega_{\rm A}^2g^2)^{-\frac{1}{2}}. \tag{4}$$

If we design the pulse shapes to satisfy the relations

$$\lim_{t\to-\infty}\frac{\Omega_{\rm A}(t)}{\Omega_{\rm B}(t)} = 0, \quad \lim_{t\to\infty}\frac{\Omega_{\rm B}(t)}{\Omega_{\rm A}(t)} = 0, \tag{5}$$

then based on Eq. (3), we can adiabatically transfer the initial state to

$$\begin{aligned} |\psi\rangle = & \frac{1}{\sqrt{2}}(|g_{\rm L}\rangle_1|f_{\rm R}\rangle_2|f_{\rm L}\rangle_3|g_{\rm R}\rangle_4 \\ & + |f_{\rm L}\rangle_1|g_{\rm R}\rangle_2|g_{\rm L}\rangle_3|f_{\rm R}\rangle_4), \end{aligned} \tag{6}$$

with the two-mode cavity left in the vacuum state.

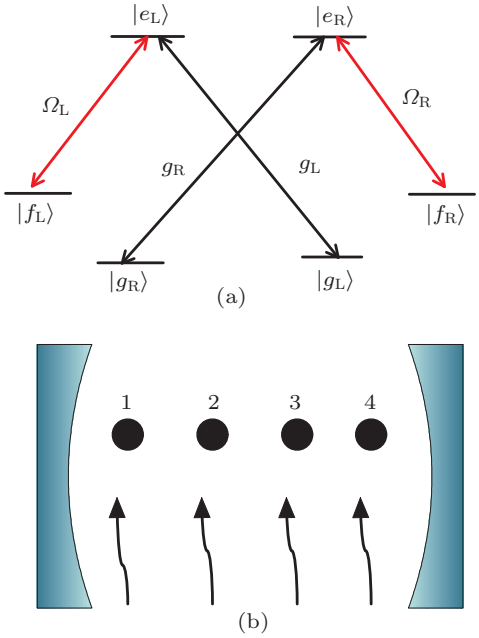


Fig. 1. (color online) (a) Energy levels of the atoms, (b) sketch of the experimental setup. Four atoms are trapped in a two-mode cavity.

The time-dependent Rabi frequencies of the laser field introduced in our scheme have Gaussian envelopes, which were widely used in experiments[37,38] and theoretical schemes for the generation of entangled state.[39,40] They are given by

$$\begin{aligned} \Omega_{\rm A}(t) &= \Omega_{\rm A,max}\,{\rm e}^{-(t-\tau_1)^2/t_{\rm p}^2}, \\ \Omega_{\rm B}(t) &= \Omega_{\rm B,max}\,{\rm e}^{-(t-\tau_2)^2/t_{\rm p}^2}, \end{aligned} \tag{7}$$

where $\Omega_{\rm A,max}$ $(\Omega_{\rm B,max})$, $t_{\rm p}$, and τ_1 (τ_2) are the peak, the width, and the time delay of the laser pulse, respectively. Similar to Refs. [39] and [40], t, τ_1, τ_2, and $t_{\rm p}$ are in units of Γ^{-1}, where Γ is the atomic spontaneous emission rate. In this paper, we assume $\Omega_{\rm A,max} = \Omega_{\rm B,max} = \Omega_{\rm max}$ and set $\tau_1 = 30\Gamma^{-1}$, $\tau_2 = 20\Gamma^{-1}$, and $t_{\rm p} = 8\Gamma^{-1}$ to satisfy Eq. (5). Figure 2(a) shows the time dependences of the Rabi frequencies of the laser fields with $\Omega_{\rm max} = 0.2$, and figure 2(b) shows the time evolutions of the populations. We can see that the population is split at 50% between $|g_{\rm L}\rangle_1|f_{\rm R}\rangle_2|f_{\rm L}\rangle_3|g_{\rm R}\rangle_4$ and $|f_{\rm L}\rangle_1|g_{\rm R}\rangle_2|g_{\rm L}\rangle_3|f_{\rm R}\rangle_4$, while the population of the initial state

is almost zero when $t > 35\Gamma^{-1}$. By using the asymmetric coding[41] for the atomic qubits, i.e., qubits 1 and 2 are encoded in the subspaces spanned by states $|g_L\rangle_1|f_L\rangle_1$ and $|f_R\rangle_2|g_R\rangle_2$, while qubits 3 and 4 are encoded in the subspaces spanned by states $|f_L\rangle_3|g_L\rangle_3$ and $|g_R\rangle_4|f_R\rangle_4$, respectively, we finally generate the four-atom GHZ state

$$|\mathrm{GHZ}\rangle = \frac{1}{\sqrt{2}}(|0\rangle_1|0\rangle_2|0\rangle_3|0\rangle_4 + |1\rangle_1|1\rangle_2|1\rangle_3|1\rangle_4). \quad (8)$$

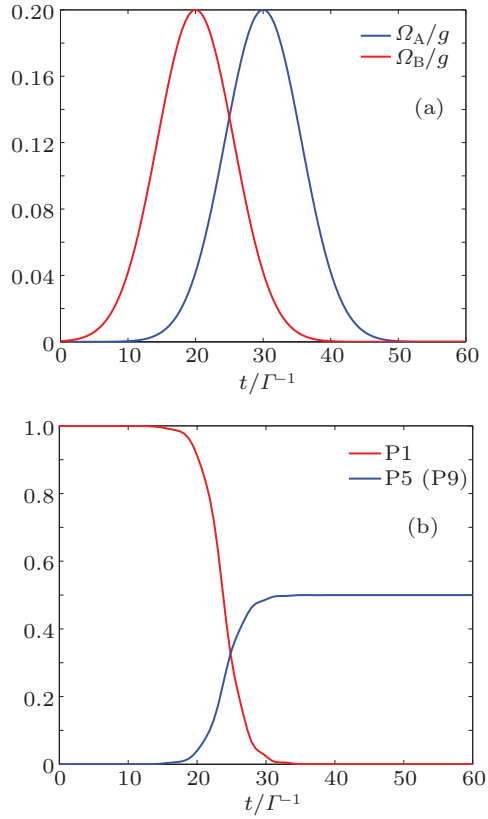


Fig. 2. (color online) (a) Time dependences of the Rabi frequencies of the laser fields for atoms, where we have chosen $\Omega_{\max}/g = 0.2$, $\tau_1 = 30\Gamma^{-1}$, $\tau_2 = 20\Gamma^{-1}$, $t_p = 8\Gamma^{-1}$. (b) The corresponding populations of states $|\varphi_1\rangle$ (P1) and $|\varphi_5\rangle$ ($|\varphi_9\rangle$) (P5 (P9)).

3. Analysis and discussion

Our scheme produces a four-atom GHZ state with a high fidelity via an adiabatic passage. For a pure state, the fidelity is defined by $F = \langle\psi(t)|\psi_{\mathrm{ideal}}\rangle^2$, where $|\psi(t)\rangle$ is the final state which is governed by the Schrödinger equation in the system, and $|\psi_{\mathrm{ideal}}\rangle$ is the ideal four-atom GHZ state in Eq. (6). We plot the fidelity of the four-atom GHZ state versus time t in Fig. 3 with the same pulse parameters given in Fig. 2(a). It shows that within the pulse duration, the fidelity increases with the time going on and reaches nearly 1 when $t > 35\Gamma^{-1}$. So, it is unnecessary to accurately control the evolution time. In the following calculation, we choose $35\Gamma^{-1}$ as the whole evolution time. Figure 4 depicts the relationships between the fidelity and the pulse parameters (a) $\Omega_{\max}/g$, (b) t_p, (c) τ_1 and τ_2. We can see that the fidelity is affected by the pulse parameters. However, as long as the corresponding parameters are set in reasonable ranges, it is unnecessary to exactly adjust the parameters to get a high fidelity. Thus, our scheme is insensitive to moderate fluctuations of the parameters.

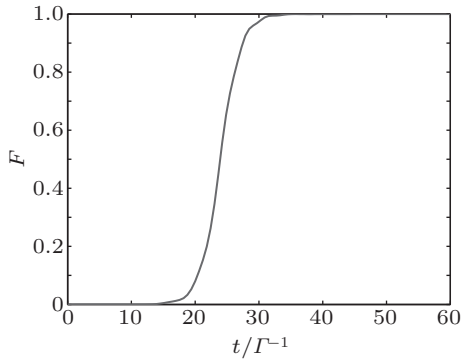


Fig. 3. Fidelity versus time t with the same parameters given in Fig. 2.

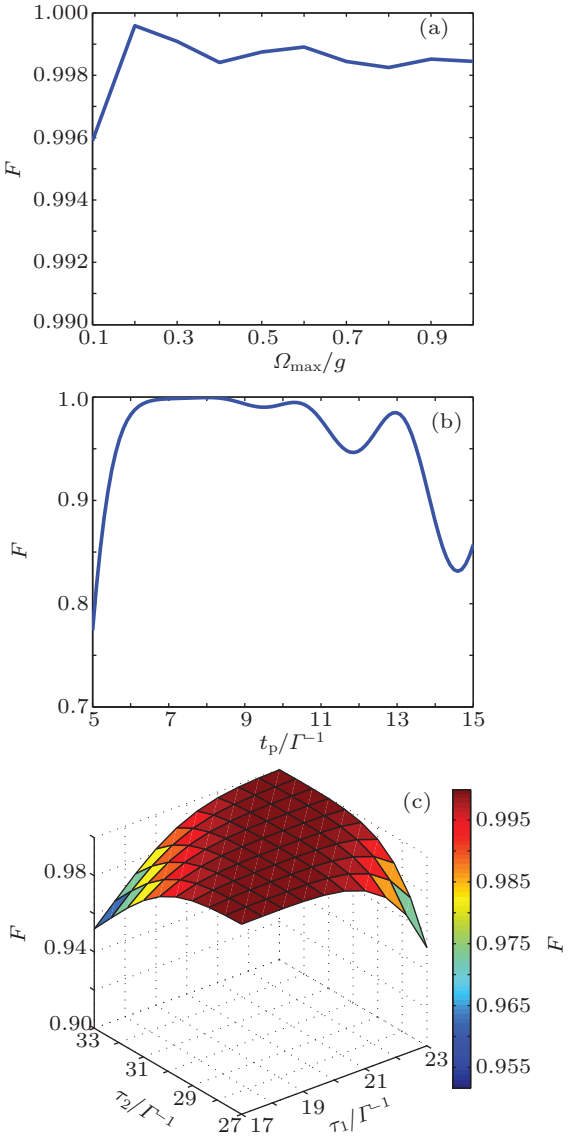


Fig. 4. (color online) (a) Fidelity versus $\Omega_{\max}/g$ with $\tau_1 = 30\Gamma^{-1}$, $\tau_2 = 20\Gamma^{-1}$, and $t_p = 8\Gamma^{-1}$; (b) fidelity versus t_p with $\Omega_{\max}/g = 0.2$, $\tau_1 = 30\Gamma^{-1}$, and $\tau_2 = 20\Gamma^{-1}$; (c) fidelity versus τ_1 and τ_2 with $\Omega_{\max}/g = 0.2$ and $t_p = 8\Gamma^{-1}$.

In the above analysis, we have assumed that the whole system is ideally isolated from the environment, and we have omitted the decoherence in our system. Now, we will introduce the decoherence mainly arising from the atomic spontaneous emission and the cavity decay. When the decoherence is taken into account, the master equation of the whole system can be written as

$$\dot{\rho} = -\mathrm{i}[H,\rho] + \frac{\Gamma}{2}\sum_{i=1}^{4}\sum_{j=f,g}\sum_{p=\mathrm{L,R}}(2S_{ijp}^{-}\rho S_{ijp}^{+} - S_{ijp}^{+}S_{ijp}^{-}\rho - \rho S_{ijp}^{+}S_{ijp}^{-}) + \frac{\kappa}{2}\sum_{p=\mathrm{L,R}}(2a_p\rho a_p^{+} - a_p^{+}a_p\rho - \rho a_p^{+}a_p), \quad (9)$$

where $S_{ijp}^{+} = |e_p\rangle_{ii}\langle j_p|$ and $S_{ijp}^{-} = |j_p\rangle_{ii}\langle e_p|$. Here κ and Γ denote the cavity decay rate and the atomic spontaneous emission rate, respectively. For simplicity, we have assumed that the atomic spontaneous emission rate of $|e_p\rangle \to |g_p\rangle$ $(p = \mathrm{L,R})$ is equal to that of $|e_p\rangle \to |f_p\rangle$, and labeled as Γ. We suppose that the mean number of thermal photons in the cavity is zero. For a mixed state, the fidelity is defined as $F = \left(\mathrm{Tr}\sqrt{(\rho_1)^{\frac{1}{2}}\rho_2(\rho_1)^{\frac{1}{2}}}\right)^2$, where ρ_1 is the density matrix of the ideal state, and ρ_2 is the density matrix governed by master equation (9). Figure 5 plots the fidelity versus both κ/g and Γ/g with $\Omega_{\max}/g = 0.2$. We can see that the fidelity decreases with the increases of κ/g and Γ/g. The fidelity can remain above 0.91 even if κ/g and Γ/g are both up to 0.05, which means that our scheme is robust against the cavity decay and the atomic spontaneous emission. To further support our analysis, we plot the fidelity versus the cooperativity C in Fig. 6, where $C = g^2/(\kappa\Gamma)$, and we have set $\Omega_{\max} = 0.2g$ and $\kappa = \Gamma$. It shows that the fidelity of the four-atom GHZ state is over 0.9 when $C = 300$. We should point out that a dissipative scheme for preparing an entangled state of two Λ atoms in an optical cavity in Ref. [42] can generate an entangled state with a higher fidelity under the same cooperativity C. However, that scheme requires the parameters to satisfy a certain relation to make the system evolve into the target state. Furthermore, the dissipative scheme for preparing multipartite entanglement needs more classical fields with more relations to be satisfied,[43] which makes the experiment difficult. Our scheme does not need to exactly adjust the parameters, which is more achievable in experiment.

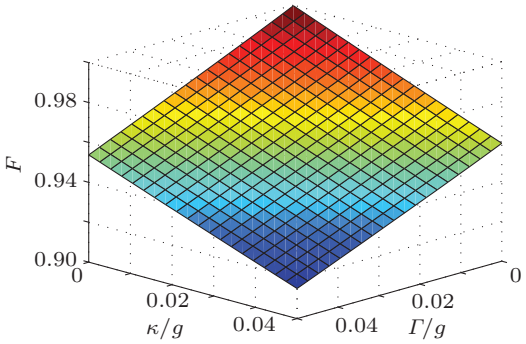


Fig. 5. (color online) Fidelity versus κ/g and Γ/g with $\Omega_{\max}/g = 0.2$.

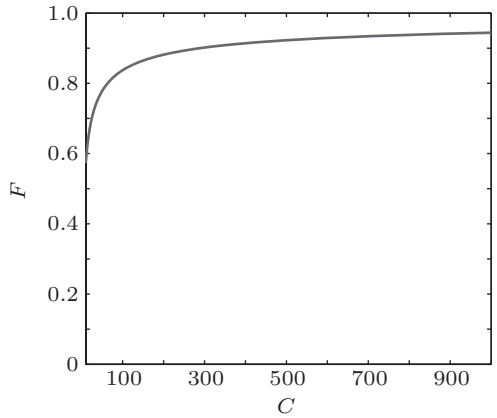


Fig. 6. Fidelity versus cooperativity C with $\Omega_{\max} = 0.2g$ and $\kappa = \Gamma$.

During the operation, we should make sure that the system remains in the dark state $|\psi_0\rangle$. According to Ref. [44], the adiabatic criterion is $\Delta\omega T \gg 1$, where $\Delta\omega$ is the minimum energy gap between the dark and the bright states, and T is the pulse duration. In this paper, we have $\Delta\omega = 2[g^2 - \frac{1}{2}(4g^4 - 2\Omega_{\mathrm{A}}^2\Omega_{\mathrm{B}}^2 + \Omega_{\mathrm{A}}^4 + \Omega_{\mathrm{B}}^4)^{1/2} + \frac{1}{2}\Omega_{\mathrm{A}}^2 + \frac{1}{2}\Omega_{\mathrm{B}}^2]^{1/2}$. For a given T, the criterion can be met by increasing $\Omega_{\max}$. But we can see from Fig. 7 that as $\Omega_{\max}$ increases, the population of the cavity modes increases, while that of the atomic excited states decreases. Therefore, it is important to select suitable pulse peaks. Figure 8(a) (8(b)) shows the fidelity versus κ/g (Γ/g) and $\Omega_{\max}/g$. As Fig. 8(a) shown, the fidelity is almost unaffected by $\Omega_{\max}/g$ when κ/g is small, while it decreases markedly with the increase of $\Omega_{\max}/g$ when κ/g is large. On the other hand, figure 8(b) shows that Γ/g affects the fidelity slightly as long as $\Omega_{\max}/g$ is large enough. When $\Omega_{\max}/g > 0.35$, the fidelity can be 0.91 even if $\Gamma/g = 0.2$. Figure 8 implies that $\Omega_{\max}/g$ should be small if κ/g is the determining decoherence factor, while $\Omega_{\max}/g$ should correspondingly increase if Γ/g dominates the decoherence. For certain κ/g and Γ/g, there exits a reasonable range of $\Omega_{\max}/g$ to reach a high fidelity. As an example, we plot the fidelity versus $\Omega_{\max}/g$ with $\kappa/g = 0.05$ and $\Gamma/g = 0.05$ in Fig. 9. We can see that the fidelity is above 0.9 when $\Omega_{\max}/g$ is between 0.17 and 0.27.

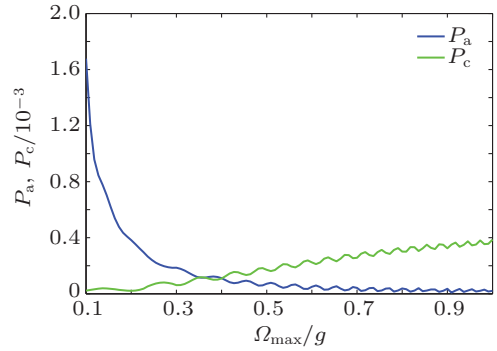


Fig. 7. (color online) Excited probabilities versus Ω_{max}/g with $\kappa/g = 0.05$ and $\Gamma/g = 0.05$, where P_a and P_c denote the excited probabilities of the atoms and the cavity modes, respectively.

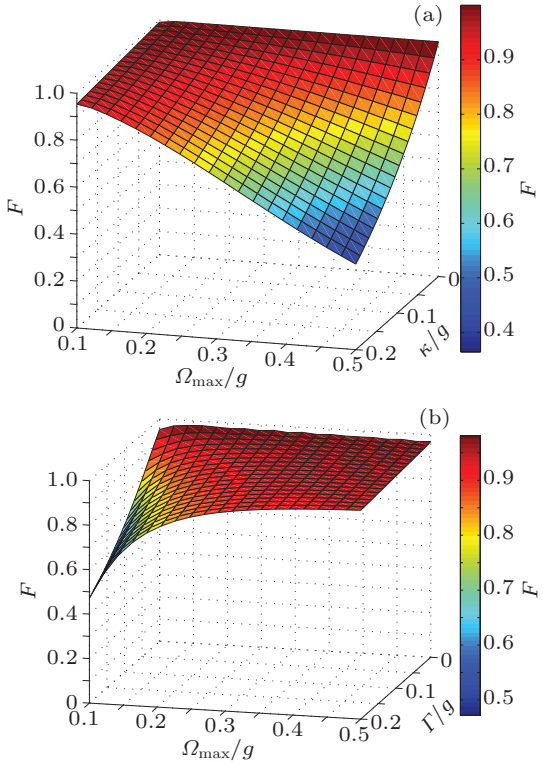


Fig. 8. (color online) (a) Fidelity versus Ω_{max}/g and κ/g with $\Gamma/g = 0$; (b) fidelity versus Ω_{max}/g and Γ/g with $\kappa/g = 0$.

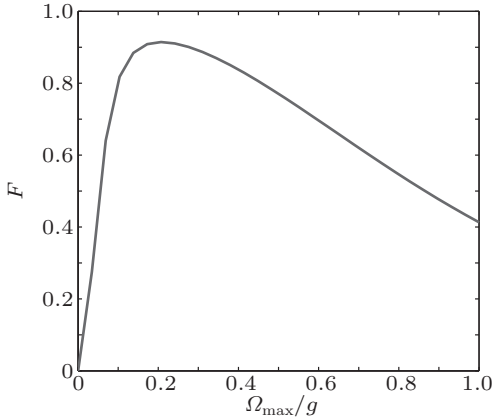


Fig. 9. Fidelity versus Ω_{max}/g with $\kappa/g = 0.05$ and $\Gamma/g = 0.05$.

4. Experimental feasibility and conclusion

Now let us briefly discuss the experimental feasibility of the proposed scheme. The required atomic level configuration can be achieved in $^{198}Hg^{+}$[45] or $^{40}Ca^{+}$.[46,47] For $^{40}Ca^{+}$, the degenerate excited states $|e_L\rangle$ and $|e_R\rangle$ correspond to the hyperfine levels $|F = 1/2, m = -1/2\rangle$ and $|F = 1/2, m = 1/2\rangle$ of $4^2P_{1/2}$, respectively, $|g_L\rangle$ and $|g_R\rangle$ correspond to the the hyperfine levels $|F = 1/2, m = 1/2\rangle$ and $|F = 1/2, m = -1/2\rangle$ of $4^2S_{1/2}$, respectively, while $|f_L\rangle$ and $|f_R\rangle$ correspond to the hyperfine levels $|F = 3/2, m = -3/2\rangle$ and $|F = 3/2, m = 3/2\rangle$ of $3^2D_{3/2}$, respectively. In recent experiments,[47–49] the parameters $(g, \Gamma, \kappa) = 2\pi \times (110, 22, 0.05)$ MHz are realizable. We can obtain from the calculation that the fidelity of the four-atom GHZ state is above 0.91 with the achievable experimental parameters. Here the atoms are separated by at least one optical wavelength so that a single laser pulse can address each atom separately. Additionally, $^{40}Ca^{+}$ ions must be localized at a fixed position in the cavity with a high precision. This can be achieved by integrating the techniques of laser cooling and ion trapping. Reference [50] has reported that a single $^{40}Ca^{+}$ ion can be localized at a fixed position in a cavity with a high precision for a long time. But, it is still difficult to fix four atoms in a cavity simultaneously. We hope the relevant techniques will be improved in the future to make our scheme possible.

In conclusion, we have proposed a scheme for generating a GHZ state of four atoms trapped in a two-mode optical cavity via an adiabatic passage of dark state. We have verified that the parameters of the classical laser fields are allowed to fluctuate moderately in our scheme. Taking advantage of the adiabatic passage, the excited states of the cavity modes and the atoms are rarely populated. By numerical calculations, we find that the excited probabilities of the cavity modes and the atoms rely on the pulse peaks of the classical laser fields. In actual experiments, we can choose appropriate pulse peaks to get a high fidelity according to the experimental decoherence parameters. We have also briefly discussed the experimental feasibility of the proposed scheme. Based on the cavity QED techniques soon to be available, the scheme might be realizable.

References

[1] Ekert A K 1991 *Phys. Rev. Lett.* **67** 661
[2] Bennett C H and Wiesner S J 1992 *Phys. Rev. Lett.* **69** 2881
[3] Bennett C H, Brassard G, Crepeau C, Jozsa R, Peres A and Wootters W 1993 *Phys. Rev. Lett.* **70** 1895
[4] Bell J S 1965 *Physics* **1** 195
[5] Cabello A 2002 *Phys. Rev. Lett.* **89** 100402
[6] Dür W, Vidal G and Cirac J I 2000 *Phys. Rev. A* **62** 062314
[7] Greenberger D M, Horne M A, Shimony A and Zeilinger A 1990 *Am. J. Phys.* **58** 1131
[8] Deng F G and Long G L 2004 *Phys. Rev. A* **70** 012311
[9] Xiao Y F, Zou X B and Guo G C 2007 *Phys. Rev. A* **75** 012310
[10] Cho J, Bose S and Kim M S 2011 *Phys. Rev. Lett.* **106** 020504
[11] Li H Q, Ren T Q, Zhang Y H and Xu X L 2012 *Chin. Phys. B* **21** 010305

[12] Song J, Xia Y and Song H S 2012 *Phys. Rev. A* **86** 034303
[13] Gisin N and Bechmann-Pasquinucci H 1998 arXiv: quant-ph /9804045
[14] Giovannetti V, Lloyd S and Maccone L 2004 *Science* **306** 1330
[15] Gao W B, Lu C Y, Yao X C, Xu P, Guhne O, Goebel A, Chen Y A, Peng C Z, Chen Z B and Pan J W 2010 *Nature Phys.* **6** 331
[16] Monz T, Schindler P, Barreiro J T, Chwalla M, Nigg D, Coish W A, Harlander M, Hänsel W, Hennrich M and Blatt R 2011 *Phys. Rev. Lett.* **106** 130506
[17] Zheng X J, Fang M F, Liao X P, Cai J W and Cao S 2007 *Chin. Phys.* **16** 906
[18] Wang J, Yu L B and Ye L 2007 *Chin. Phys.* **16** 2211
[19] Song J, Xia Y and Song H S 2008 *J. Phys. B* **41** 065507
[20] Zheng S B 2008 *Phys. Lett. A* **372** 591
[21] Xiong W and Ye L 2011 *Opt. Commu.* **284** 3659
[22] Wei L F, Liu Y X and Nori F 2006 *Phys. Rev. Lett.* **96** 246803
[23] Neeley M, Bialczak R C, Lenander M, Lucero E, Mariantoni M, O'Connell A D, Sank D, Wang H, Weides M, Wenner J, Yin Y, Yamamoto T, Cleland A N and Martinis J M 2010 *Nature* **467** 570
[24] Zeilinger A 1999 *Rev. Mod. Phys.* **71** 288
[25] Raimond J M, Brtule M and Haroehe S 2001 *Rev. Mod. Phys.* **73** 565
[26] Leibfried D, Knill E, Seidelin S, Britton J, Blakestad R B, Chiaverini J, Hume D B, Itano W M, Jost J D, Langer C, Ozeri R, Reichle R and Wineland D J 2005 *Nature* **438** 639
[27] Chen G, Bonadeo N H, Steel D G, Gammon D, Katzer D S, Park D and Sham L J 2000 *Science* **289** 1906
[28] Bergmann K, Theuer H and Shore B W 1998 *Rev. Mod. Phys.* **70** 1003
[29] Vitanov N V, Suominen K A and Shore B W 1999 *J. Phys. B* **32** 4535
[30] Talab M A, Guérin S, Sangouard N and Jauslin H R 2005 *Phys. Rev. A* **71** 023805
[31] Talab M A, Guérin S and Jauslin H R 2005 *Phys. Rev. A* **72** 012339
[32] Lacour X, Sangouard N, Guérin S and Jauslin H R 2006 *Phys. Rev. A* **73** 042321
[33] Zheng A S, Liu L B and Chen H Y 2011 *Chin. Phys. Lett.* **28** 080303
[34] Song J, Xia Y and Song H S 2010 *Appl. Phys. Lett.* **96** 071102
[35] Song K H, Xiang S H, Liu Q and Lu D H 2006 *Phys. Rev. A* **74** 064303
[36] Deng Z J, Gao K L and Feng M 2007 *Phys. Rev. A* **75** 032347
[37] Gaubatz U, Rudecki P, Schiemann S and Bergmann K 1990 *J. Chem. Phys.* **92** 5363
[38] Dada A C, Leach J, Buller G S, Padgett M J and Andersson E 2011 *Nature Phys.* **7** 677
[39] Yang Z B, Wu H Z and Zheng S B 2010 *Chin. Phys. B* **19** 094205
[40] Song J, Xia Y, Song H S, Guo J L and Nie J 2007 *Europhys. Lett.* **80** 60001
[41] Yang Z B, Wu H Z, Su W J and Zheng S B 2009 *Phys. Rev. A* **80** 012305
[42] Kastoryano M J, Reiter F and Sfrensen A S 2011 *Phys. Rev. Lett.* **106** 090502
[43] Chen X Y, Shen L T, Yang Z B, Wu H Z and Chen M F 2012 *J. Opt. Soc. Am. B* **29** 001535
[44] Pellizzari T, Gardiner S A, Cirac J I and Zoller P 1995 *Phys. Rev. Lett.* **75** 3788
[45] Kiffner M, Evers J and Keitel C H 2006 *Phys. Rev. Lett.* **96** 100403
[46] Kiffner M, Lange B, Hayasaka K, Lange W and Walther H 2004 *Nature* **431** 1075
[47] Keller M, Lange B, Hayasaka K, Lange W and Walther H 2004 *New J. Phys.* **6** 95
[48] Stute A, Casabnone B, Schindler P, Monz T, Schmidt P O, Brandsätter B, Northup T E and Blatt R 2012 *Nature* **485** 482
[49] Buck J R and Kimble H J 2003 *Phys. Rev. A* **67** 033806
[50] Mundt A B, Kreuter A, Becher C, Leibfried D, Eschner J, Schmidt-Kaler F and Blatt R 2002 *Phys. Rev. Lett.* **89** 103001