Quantum state transfer between distant nodes of a quantum network via adiabatic passage

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A scheme is presented to realize transfer of a three-dimensional quantum state between two atoms trapped in distant cavities connected by an optical fiber, whose mode is resonant with the polarization cavity modes. Performing an adiabatic passage along dark states, the fiber mode remains in the vacuum state and the population of the cavities being excited can be negligible under certain conditions. In addition, atomic transitions in our scheme are largely detuned from cavity modes. As a consequence, the atomic spontaneous emission and the fiber decay can be effectively eliminated. Furthermore, we give some discussions about the fidelity of our scheme.

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In the new era of quantum information science, quantum network [1] has played an important role for it offers a variety of novel applications ranging over distributed quantum computation [2], teleportation [3], and metrology [4]. The realization of quantum networks, which consists of spatially separated nodes connected by quantum communication channels, requires new scientific capabilities for operating quantum coherent and quantum entanglement [5]. Cavity quantum electrodynamics (cavity QED) provides a promising tool for quantum networking [6] because of its important paradigm for coherent coupling of optical and atomic qubits [7]. From a practical perspective, photon, as a good flying qubit, is perfect for the transmission of quantum information while atom is particularly well suited for storing qubit in its long-lived internal states. So in the implementation of quantum networks [5,8–11], atom represents the quantum memory while photon propagating in an optical fiber [10] or free space serves as the data bus to realize the quantum channel. Through a coherent coupling mediated by optical fiber, which is a good method to entangle two separate atomic qubit, Mancini *et al.* [5] obtained an effective direct interaction between two atoms placed in distant cavities. Specially, a deterministic gates between two-level atoms in separate optical cavities has been realized in Ref. [9]. About the quantum transfer, Cirac *et al.* [11] outlined the first scheme to realize the quantum transmission between distant atoms with unit efficiency by tailoring time-symmetric photon wave packets, so precise control of the laser-pulse shape is required. This problem has also been addressed by Pellizzari [10], whose approach is based on an adiabatic passage through two cavities which remain in their respective vacuum states during the whole operation. In the above two schemes, it is easy to dephase during the transmission process because the qubits of cavities and fiber are represented by vacuum one-photon states, respectively [12]. Recently, based on the cavity input-output process, Refs. [13,14]. have constructed the schemes of quantum state exchange between atoms or between atom and photon, by dint of single-pulse injection. However, they are probabilistic schemes as they depend on the detection of the photon decaying from the leaking cavities and, thus, high efficient photon detectors are required. In particular, initial single-photon injection is very difficult to meet.

In this Brief Report, we present a scheme to transfer a three-dimensional quantum state between two atoms located at spatially separated cavities connected by an optical fiber. Based on the adiabatic passage along dark states, our scheme has the following characteristics: (i) The fiber modes remain in the vacuum state and the population of the cavities being exited can be negligible under certain conditions, so the decays of fiber and cavities can be effectively eliminated. (ii) Atomic spontaneous emission can be efficiently reduced since the cavity frequency is largely detuned from the atomic transition frequency. (iii) The information of cavities and fiber are encoded in the photonic polarization states. (iv) Our scheme works robustly beyond Lamb-Dicke limit, which is important for current experimental technique.

The framework of our proposal is shown in Fig. 1. The atoms are modeled by five-level systems with three ground states, $|0\rangle, |g\rangle, |1\rangle$, and two excited states, $|e_1\rangle, |e_2\rangle$, as shown in Fig. 2. In the model, two classical laser pulses incident from one mirror of the cavity and collinear with the cavity axis, couple to the atom transition $|e_0\rangle \leftrightarrow |0\rangle (|e_1\rangle \leftrightarrow |1\rangle)$ with the time-dependent Rabi frequency $2\Omega_{L(R)}(t)$. The transition $|e_{0(1)}\rangle \leftrightarrow |g\rangle$ is coupled with cavity mode $a_{L(R)}$ with left-circular (right-circular) polarization with coupling strength $g_{L(R)}$ [15]. Without loss of generality, all the Rabi frequencies and coupling rates are assumed to be real. In the rotation frame, the Hamiltonian has the following form ($\hbar = 1$):

$$H_i = \Delta_L |e_0\rangle_i\langle e_0| + \Delta_R |e_1\rangle_i\langle e_1| + \Omega_L^{(i)}(|e_0\rangle_i\langle 0| + \text{H.c.}) + \Omega_R^{(i)}(|e_1\rangle_i\langle 1| + \text{H.c.}) + g_L(|e_0\rangle_i\langle g| a_L^{(i)} + \text{H.c.}) + g_R(|e_1\rangle_i\langle g| a_R^{(i)} + \text{H.c.}) (i = A, B), \quad (1)$$

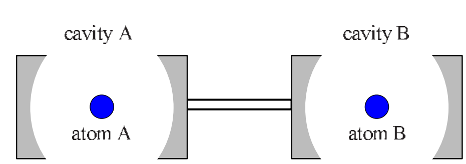

FIG. 1. (Color online) Schematic representation of quantum state transmission between two atoms in spatially separate optical cavities connected by an optical fiber.

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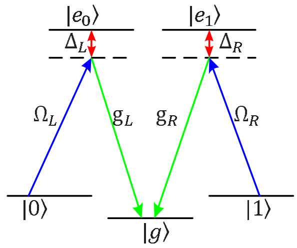

FIG. 2. (Color online) The relevant level structure and transition of the atom.

where the detuning $\Delta_{L(R)}$ is defined as energy difference between the atomic transition $|e_{0(1)}\rangle \leftrightarrow |g\rangle$ and the relevant cavity mode, and H.c. stands for the Hermitian conjugate. If we assume $\Delta_{L(R)} \gg \Omega_{L(R)}^{(i)}$, $g_{L(R)}$, the excited state will then essentially remain unpopulated and, thus, we can adiabatically eliminate the excited states from temporal evolution state space. For simplicity and feasibility, we assume that $\Omega_L^{(i)}(t)=\Omega_R^{(i)}(t)=\Omega_i(t)$, $g_L=g_R=g$, and $\Delta_L=\Delta_R=\Delta$. Under these conditions, it is natural to describe the atom-cavity system by the effective Hamiltonian

$$H_i^{\text{eff}} = -\frac{\Omega_i^2(t)}{\Delta}(|0\rangle_i\langle 0| + |1\rangle_i\langle 1|) - \frac{g^2}{\Delta}|g\rangle_i\langle g|(a_L^{(i)\dagger}a_L^{(i)} + a_R^{(i)\dagger}a_R^{(i)}) - \frac{\Omega_i(t)g}{\Delta}[(|0\rangle_i\langle g|a_L^{(i)} + |1\rangle_i\langle g|a_R^{(i)}) + \text{H.c.}]\ (i=A,B). \quad (2)$$

Next step we will consider two atom-cavity systems connected by an optical fiber. The interaction Hamiltonian for the fiber coupled to the modes of two cavities are assumed to take the following form [9,10]:

$$H_{\text{fib}} = \sum_{k=1}^{\infty} \sum_{j=L,R} \nu_{j,k}[b_{j,k}^{\dagger}(a_j^A + (-1)^k e^{-i\varphi}a_j^B)] + \text{H.c.}, \quad (3)$$

where $b_{j,k}^{\dagger}$ is the creation operator of the fiber mode k with polarization j (=L,R, corresponding to left-circular and right-circular polarizations), $\nu_{j,k}$ is the coupling strength between the cavity mode j and the fiber mode k, and the phase φ is due to the propagation of the field through the fiber of length l [9]. The subscripts A and B distinguish the two atom-cavity subsystems. In the short fiber limit $2l\Gamma/(2\pi c) \leq 1$, where Γ is the decay rate of the cavity fields into a continuum of fiber modes, only resonant modes of the fiber interact with the cavity modes [9]. For this case, the Hamiltonian H_{fib} may be approximated to

$$H_{\text{fib}} = \sum_{j=L,R} \nu[b_j^{\dagger}(a_j^A + a_j^B) + \text{H.c.}], \quad (4)$$

where the phase $(-1)^k e^{-i\varphi}$ in Eq. (3) has been absorbed into the annihilation and creation operators of the modes of the second cavity field [16,17]. Here $\nu_R=\nu_L$ has been assumed. The Hamiltonian for the total system of atom-cavity fiber is

$$H_{\text{tot}} = H_A^{\text{eff}} + H_B^{\text{eff}} + H_{\text{fib}}, \quad (5)$$

specified in Eqs. (2) and (4). For the initial state $|0\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$, the state evolution remains in the subspace spanned by the basis-state vectors $\{|0\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|L0\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|0L\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|00\rangle|L\rangle_f$, $|g\rangle_A|0\rangle_B|00\rangle|\text{vac}\rangle_f\}$. Here, the first and second kets refer to the quantum states of atoms A and B. The third ket represents the polarization state of cavity modes of cavities A and B, respectively, where $|L(R)\rangle$ represents there is a photon with left-circular (right-circular) polarization in the cavity. The last ket $|\text{vac}\rangle_f$ denotes the subsystem of the fiber modes in the vacuum state. Then the corresponding dark state of the Hamiltonian H_{tot} is

$$|D_0\rangle \propto g\Omega_B|0\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f - \Omega_A\Omega_B|g\rangle_A|g\rangle_B(|L0\rangle - |0L\rangle)|\text{vac}\rangle_f - g\Omega_A|g\rangle_A|0\rangle_B|00\rangle|\text{vac}\rangle_f. \quad (6)$$

For the initial state $|1\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$, the state evolution remains in the subspace spanned by the basis-state vectors $\{|1\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|R0\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|0R\rangle|\text{vac}\rangle_f$, $|g\rangle_A|g\rangle_B|00\rangle|R\rangle_f$, $|g\rangle_A|1\rangle_B|00\rangle|\text{vac}\rangle_f\}$, and we can write the other dark state

$$|D_1\rangle \propto g\Omega_B|1\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f - \Omega_A\Omega_B|g\rangle_A|g\rangle_B(|R0\rangle - |0R\rangle)|\text{vac}\rangle_f - g\Omega_A|g\rangle_A|1\rangle_B|00\rangle|\text{vac}\rangle_f. \quad (7)$$

For the initial state $|g\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$, it does not change under the time evolution generated by H_{tot} because it is decoupled from the laser interaction [10]. So it is the third dark state of the system,

$$|D_g\rangle = |g\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f. \quad (8)$$

Note that if the system remains in the dark states, the two atoms are always in ground states and the fiber mode is in vacuum state. In the dark states $|D_0\rangle$ and $|D_1\rangle$ the fiber modes are not populated due to destructive quantum interference. Taking $|D_0\rangle$ for example, since the transition $|g\rangle_A|g\rangle_B|L0\rangle|\text{vac}\rangle_f \leftrightarrow |g\rangle_A|g\rangle_B|0L\rangle|\text{vac}\rangle_f$ is mediated by the intermediate state $|g\rangle_A|g\rangle_B|00\rangle|L\rangle_f$ and the transition paths $|g\rangle_A|g\rangle_B|L0\rangle|\text{vac}\rangle_f \rightarrow |g\rangle_A|g\rangle_B|00\rangle|L\rangle_f$ and $|g\rangle_A|g\rangle_B|0L\rangle|\text{vac}\rangle_f \rightarrow |g\rangle_A|g\rangle_B|00\rangle|L\rangle_f$ interfere destructively, the fiber mode remains in vacuum state. Moreover, we assume that the condition $g \gg \Omega_i(t)(i=A,B)$ is always satisfied throughout the whole process. Therefore, the population of the cavity modes in excited states can be negligible as shown in Fig. 3. Here the pulse shape of the laser fields is assumed to be Gaussian [18]: $\Omega_{A,B}(t)=\Omega_0 \exp[-(t-2/T-t_{A,B})^2/(2\tau^2)]$, where Ω_0 is the maximal values of the $\Omega_{A,B}$, T is the total adiabatic time, and τ is the laser beam waist. We have chosen the pulse parameters as $\Omega_0=0.2g$, $T=60/g$, $t_A=-t_B=8/g$, $\tau=12/g$ and $g=1$ GHz. It is obvious that all population is completely transferred from the state $|n\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f$ to the state $-|g\rangle_A|n\rangle_B|00\rangle|\text{vac}\rangle_f(n=0,1)$. Especially, the population of the cavity modes in excited states $|g\rangle_A|g\rangle_B(|j0\rangle - |0j\rangle)|\text{vac}\rangle_f/\sqrt{2}(j=L,R)$ is close to zero. So we could write the Eqs. (6) and (7) into

$$|D_0\rangle \propto g\Omega_B(t)|0\rangle_A|g\rangle_B|00\rangle|\text{vac}\rangle_f - g\Omega_A(t)|g\rangle_A|0\rangle_B|00\rangle|\text{vac}\rangle_f, \quad (9a)$$

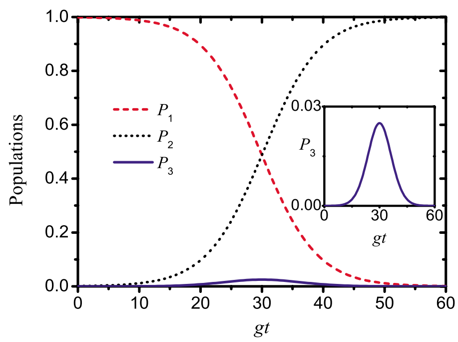

FIG. 3. (Color online) The time dependence of the populations. P_1 represents the population of the state $|n\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f$, P_2 represents the population of the state $-|g\rangle_A|n\rangle_B|00\rangle|\mathrm{vac}\rangle_f(n=0,1)$, and P_3 represents the population of the state $|g\rangle_A|g\rangle_B(|j0\rangle-|0j\rangle)|\mathrm{vac}\rangle_f/\sqrt{2}(j=L,R)$. Here we assume the pulse shape of the laser fields is Gaussian [18], $\Omega_{A,B}(t)=\Omega_0\exp[-(t-2/T-t_{A,B})^2/(2\tau^2)]$, with $g=1$ GHz, $\Omega_0=0.2g,T=60/g,t_A=-t_B=8/g,\tau=12/g$.

$$|D_1\rangle \propto g\Omega_B(t)|1\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f - g\Omega_A(t)|g\rangle_A|1\rangle_B|00\rangle|\mathrm{vac}\rangle_f. \tag{9b}$$

Our goal is to accomplish the quantum state transfer

$$(\alpha|0\rangle_A+\beta|1\rangle_A+\gamma|g\rangle_A)\otimes|g\rangle_B\otimes|00\rangle|\mathrm{vac}\rangle_f \rightarrow |g\rangle_A\otimes(\alpha|0\rangle_B+\beta|1\rangle_B+\gamma|g\rangle_B)\otimes|00\rangle|\mathrm{vac}\rangle_f, \tag{10}$$

where $\alpha,\beta,$ and γ are complex numbers and $|\alpha|^2+|\beta|^2+|\gamma|^2=1$. Using the dark states of Eqs. (4)–(6) for adiabatic passage is our central idea. If the system is prepared in a superposition of the three dark states in Eqs. (4)–(6) and the laser intensities are changed slowly (i.e., adiabatically) no other eigenstates of H_{tot} will be populated. Initially, subsystem A is prepared in an unknown quantum superposition as given in Eq. (10). In practice, this initial state can be written as a superposition of the three dark states $|D_{0,1,g}\rangle$ provided that $\Omega_B(t)\gg\Omega_A(t)$,

$$(\alpha|0\rangle_A+\beta|1\rangle_A+\gamma|g\rangle_A)\otimes|g\rangle_B\otimes|00\rangle|\mathrm{vac}\rangle_f = \alpha|D_0\rangle+\beta|D_1\rangle+\gamma|D_g\rangle_f. \tag{11}$$

In fact, this means that at the beginning of the transfer the laser which acts on the atom B is switched on first. In this case, the dark states $|D_0\rangle$ and $|D_1\rangle$ are approximately $|0\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f$ and $|1\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f$, respectively. During the transfer process, we slowly change the laser intensities, i.e., increase $\Omega_A(t)$ and decrease $\Omega_B(t)$ so that at the end the inequality $\Omega_A\gg\Omega_B$ holds and the state transfer $|n\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f\rightarrow-|g\rangle_A|n\rangle_B|00\rangle|\mathrm{vac}\rangle_f$ $(n=0,1)$ is realized adiabatically. If the change is carried out slowly enough, the finial state of the system is still the above superposition of $|D_0\rangle$, $|D_1\rangle$, and $|D_g\rangle$ and as mentioned above we can write it as

$$\alpha|D_0\rangle+\beta|D_1\rangle+\gamma|D_g\rangle = |g\rangle_A\otimes(-\alpha|0\rangle_B-\beta|1\rangle_B+\gamma|g\rangle_B)\otimes|00\rangle|\mathrm{vac}\rangle. \tag{12}$$

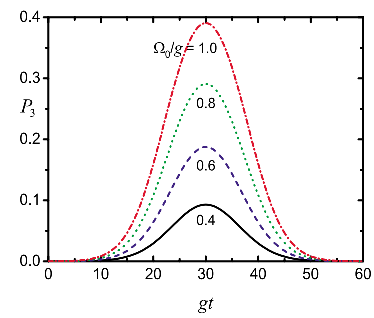

FIG. 4. (Color online) The time dependence of the population (P_3) of the state $|g\rangle_A|g\rangle_B(|j0\rangle-|0j\rangle)|\mathrm{vac}\rangle_f/\sqrt{2}(j=L,R)$ for $\Omega_0/g=0.4$, 0.6, 0.8, and 1.0 (solid, dashed, dotted, and dot-dashed lines, respectively). Other parameters are as in Fig. 3.

This quantum state corresponds to the desired final state of the transfer, apart from phase factors that can be eliminated by the appropriate subsequent logic operations.

In what follows, we will consider more practical situation. Considering the experimental feasibility of the present context, various decoherence processes such as atomic spontaneous emission from the atomic excited states and photon loss may influence the fidelity of our scheme. We know that under the large-detuning condition, the effect of spontaneous emission can be well suppressed [19], so we ignore it and set the spontaneous emission rate to zero in our scheme. Photon loss mainly comes from the leakage of cavities and fiber. We find that the effect of photon leakage out of the fiber can be greatly eliminated and even ignored. This is because the fiber mode is really not excited and is always kept in the vacuum state in the whole process. However, compared to the fiber, the photon losses in cavities have greater impact on the fidelity of our scheme. We note that, as shown in Fig. 4, if the condition $g\gg\Omega_{A,B}(t)$ is not ideally satisfied, the population of the cavities being excited cannot be negligible any more. That is to say cavity modes may be excited. In this case, during the adiabatically transferring procedure $|n\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f\rightarrow-|g\rangle_A|n\rangle_B|00\rangle|\mathrm{vac}\rangle_f$ $(n=0,1)$, according to the Eqs. (6) and (7), there is a probability of the states $|g\rangle_A|g\rangle_B(|j0\rangle|\mathrm{vac}\rangle_f-|0j\rangle|\mathrm{vac}\rangle_f)/\sqrt{2}$ $(j=L,R)$ being populated. Then, due to the cavity decay they will evolve into $|g\rangle_A|g\rangle_B|00\rangle|\mathrm{vac}\rangle_f$ after emitting photons, which will reduce the fidelity of transmission. As a concrete numerical example, let us assume that $g=1$ GHz and the reliable transfer can be achieved, as shown in Fig. 3, for the total adiabatic time of the order of $T=60$ ns. Considering the cavity leakage, for the initial state of atoms A and B: $|\psi\rangle=(\alpha|0\rangle_A+\beta|1\rangle_A+\gamma|g\rangle_A)\otimes|g\rangle_B$, the finial state can be approximately given by [20]

$$\rho=|\psi'\rangle\langle\psi'|+(1-|\gamma|^2)(1-P_{\mathrm{suc}})|g\rangle_A|g\rangle_B\langle g|_A\langle g|_B, \tag{13}$$

where

$$P_{\mathrm{suc}}=1-k\int_0^T\frac{2[\Omega_A(t)\Omega_B(t)]^2}{[\Omega_A(t)g]^2+2[\Omega_A(t)\Omega_B(t)]^2+[\Omega_B(t)g]^2}dt, \tag{14}$$

and

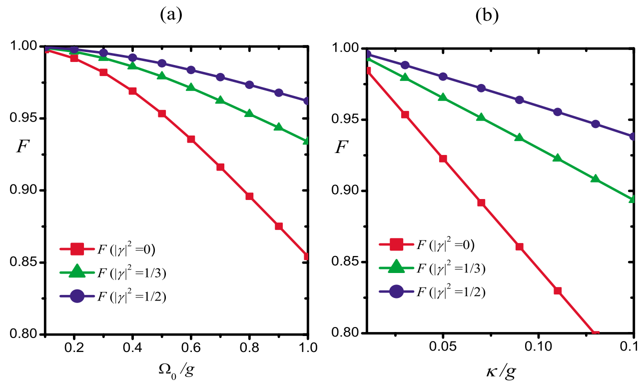

FIG. 5. (Color online) Fidelity of transmission F as a function of Ω_0/g and κ/g, respectively, for $|\gamma|^2=0$ (squares), 1/3 (triangles), and 1/2 (circles). (a) The cavity decay rate is set to be $\kappa=0.02g$. (b) The parameter of the laser-pulse shape is $\Omega_0/g=0.4$. Other parameters in Figs. 5(a) and 5(b) are the same as Fig. 3.

$$|\psi'\rangle = \alpha\sqrt{P_{\text{suc}}}|g\rangle_A|0\rangle_B + \beta\sqrt{P_{\text{suc}}}|g\rangle_A|1\rangle_B + \gamma|g\rangle_A|g\rangle_B. \quad (15)$$

Thus the fidelity is

$$F = \langle\psi'|\rho|\psi'\rangle = [(1-|\gamma|^2)\sqrt{P_{\text{suc}}} + |\gamma|^2]^2 + |\gamma|^2(1-|\gamma|^2)P_{\text{suc}}. \quad (16)$$

In Fig. 5(a) the fidelity F of our scheme is studied as a function of Ω_0/g for different transferring state with $\kappa=0.02g$ for the same parameters and laser-pulse shapes as in Fig. 3. Figure 5(b) exhibits the effect of cavity decay rate κ on fidelity F when we assume $\Omega_0=0.4g$.

The last but not the least, we discuss the experimental feasibility of the present scheme. For the atomic level structure, we can take ^{87}Rb as our choice. The ground states $|0\rangle$, $|1\rangle$, and $|g\rangle$ can be achieved in pairs of Zeeman hyperfine levels $|F=2,m=-2\rangle$, $|F=2,m=2\rangle$, and $|F=1,m=0\rangle$ of $5S_{1/2}$, respectively. The upper levels $|e_0\rangle$ and $|e_1\rangle$ correspond to $|F=2,m=-1\rangle$ and $|F=2,m=1\rangle$ of $5P_{1/2}$, respectively. The two pairs of Zeeman hyperfine levels of the ground states $|0\rangle,|g\rangle$ and $|1\rangle,|g\rangle$ are coupled via two far-detuned Raman transition, i.e., their respective transitions to the upper levels $|e_0\rangle$ and $|e_1\rangle$ have the same large detuning Δ, which is consistent with our assumption $\Delta_L=\Delta_R=\Delta$. In real experiments, it is very challenging to control atoms precisely to meet the Lamb-Dicke condition. But when all the applied laser pulses are collinear with their respective cavities axes and thus share the same spatial mode structure with their respective cavities, the systemic adiabatic dynamics evolution only depends on the ratios $g(\vec{r})/\Omega_{A(B)}(\vec{r},t)$ and $\Omega_A(\vec{r},t)/\Omega_B(\vec{r},t)$ [21,22]. So in the cavity the atomic position is not required to be accurately controlled, i.e., our scheme works robustly beyond Lamb-Dicke limit.

In summary, we have proposed a scheme for quantum transmission of three-dimensional atomic state between two remote nodes connected by optical fiber. Based on the adiabatic passages, the system of atom-cavity fiber is always on the dark states by using a sequence of pulsed laser field. During the transfer process, the fiber mode remains in vacuum states, thus, the decay of the optical fiber can be ignored. In addition, atomic transitions in our scheme are largely detuned with cavity modes so the atomic spontaneous emission can be effectively avoid. Furthermore, we encode the qubits of cavity and fiber in their polarization state instead of vacuum one-photon state, which is sensitive to phase fluctuations during the transmission. We also make an estimation on the fidelity by considering different parameters.

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[1] C. W. Gardiner and P. Zoller, *Quantum Noise* (Spring-Verlag, Berlin, 2004).
[2] C. H. Bennett, Phys. Today **48**(10), 24 (1995).
[3] C. H. Bennett, G. Brassard, C. Crepeau, R. Jozsa, A. Peres, and W. K. Wootters, Phys. Rev. Lett. **70**, 1895 (1993).
[4] H. J. Kimble, Nature (London) **453**, 1023 (2008).
[5] S. Mancini and S. Bose, Phys. Rev. A **70**, 022307 (2004).
[6] E. Knill, R. Laflamme, and G. J. Milburn, Nature (London) **409**, 46 (2001).
[7] S. J. van Enk, H. J. Kimble, and H. Mabuchi, Quantum Inf. Process. **3**, 75 (2004).
[8] S. J. van Enk, J. I. Cirac, and P. Zoller, Phys. Rev. Lett. **78**, 4293 (1997); L. M. Duan, M. D. Lukin, J. I. Cirac, and P. Zoller, Nature (London) **414**, 413 (2001).
[9] A. Serafini, S. Mancini, and S. Bose, Phys. Rev. Lett. **96**, 010503 (2006).
[10] T. Pellizzari, Phys. Rev. Lett. **79**, 5242 (1997).
[11] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Phys. Rev. Lett. **78**, 3221 (1997).
[12] V. A. Kostelecky and M. Mewes, Phys. Rev. Lett. **97**, 140401 (2006).
[13] L. M. Liang and C. Z. Li, Phys. Rev. A **72**, 024303 (2005).
[14] H. Wei, Z. J. Deng, X. L. Zhang, and M. Feng, Phys. Rev. A **76**, 054304 (2007).
[15] Y. F. Xiao, Z. F. Han, J. Gao, G. C. Guo, and X. Q. Su, Physica A **354**, 227 (2005).
[16] P. Peng and F. L. Li, Phys. Rev. A **75**, 062320 (2007).
[17] L. B. Chen, M. Y. Ye, G. W. Lin, Q. H. Du, and X. M. Lin, Phys. Rev. A **76**, 062304 (2007).
[18] U. Gaubatz, P. Rudecki, S. Schiemann, and K. Bergnann, J. Chem. Phys. **92**, 5363 (1990); H. Goto and K. Ichimura, Phys. Rev. A **70**, 012305 (2004).
[19] L. M. Duan, A. Kuzmich, and H. J. Kimble, Phys. Rev. A **67**, 032305 (2003).
[20] Note that here the local transformation $|g\rangle_A(-\alpha|0\rangle_B-\beta|1\rangle_B+\gamma|g\rangle_B)|g\rangle_B \rightarrow |g\rangle_A(\alpha|0\rangle_B+\beta|1\rangle_B+\gamma|g\rangle_B)|g\rangle_B$ has been performed.
[21] S. Y. Ye, Z. R. Zhong, and S. B. Zheng, Phys. Rev. A **77**, 014303 (2008).
[22] L. M. Duan and H. J. Kimble, Phys. Rev. Lett. **92**, 127902 (2004).