Scheme for atomic-state teleportation between two bad cavities

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A scheme is presented for the long-distance teleportation of an unknown atomic state between two separated cavities. Our scheme works in the regime where the atom-cavity coupling strength is smaller than the cavity decay rate. Thus the requirement on the quality factor of the cavities is greatly relaxed. Furthermore, the fidelity of our scheme is not affected by the detection inefficiency and atomic decay. These advantages are important in view of experiments.

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Recently, much attention has been paid to quantum information, a topical subject across a wide range of disciplines. In the classical world, we can transport the information from one system to another remote system by measuring the first system and then copying it according to the obtained information. The no-cloning theorem prohibits an arbitrary input quantum state from being cloned perfectly [1]. It seems impossible to transport quantum information from one system to another remote system since any measurement on the first system would destroy its state. Fortunately, quantum teleportation provides a way to do so [2]. Teleportation not only is important to test fundamental features of quantum mechanics, but also may be useful in quantum communication [3] and quantum computation [4].

Recently, experimental realizations of quantum teleportation have been reported by using optical systems [5]. In these experiments, the stationary qubits, whose states are to be teleported, are the light fields, not ideal for storing quantum information. Quantum teleportation has also been demonstrated using nuclear magnetic resonance [6]. Very recently, teleportation of atomic qubits in a trapped-ion system has been reported [7]; however, the interion distance is only several micrometers, which is too short to be of practical use in quantum communication. On the other hand, the cavity QED system is another qualified candidate for quantum-information processing. Schemes have been proposed for the generation of entanglement between two distant atoms by single-photon interference at photodetectors [8–11]. Schemes have also been presented for teleportation of an unknown atomic state [12]. In these schemes, the flying qubits to transfer quantum information from one place to another place are the atoms and thus are not ideal for long-distance teleportation.

Bose *et al.* have suggested a proposal to teleport the state of an atom trapped in a cavity to another atom trapped in a distant cavity by detecting the photon decays from the cavities [13]. In the protocol, the stationary qubits are the atomic states, while the flying qubits are the photon states, ideal for both storage of quantum information and long-distance teleportation. In the scheme, the atomic state (a superposition of the two ground states of a Λ-type three-level atoms) to be teleported is first mapped to the optical cavity mode, and the second atom and the second cavity are prepared in the maximally entangled state via the Raman coupling between the two atomic ground states induced by the cavity mode and the classical field. Thus, the scheme requires that the Raman coupling strength between the two atomic ground states should be much larger than the cavity decay rate. This puts a very stringent requirement on the cavity quality and thus is experimentally challenging. Furthermore, the fidelity of the scheme is affected by the detection inefficiency. Even under ideal conditions, the second atom is not exactly in the initial state of the first atom and the fidelity is state dependent.

In this paper, we propose an alternative scheme for long-distance atomic-state teleportation between two separate cavities. Our scheme does not require complete mapping between the atomic state and the photonic state, nor does it require the establishment of maximal entanglement between the second atom and the cavity. Our scheme works in the opposite regime, where the atom-cavity coupling strength is smaller than the cavity decay rate, greatly relaxing the requirement on the cavity quality. Another advantage of our scheme is that the fidelity is not affected by both the detection inefficiency and atomic decay. These features make the scheme very promising for long-distance quantum communication.

The atoms have three degenerate excited states $|e_l\rangle$, $|e_r\rangle$, and $|e_0\rangle$ and three degenerate ground states $|g_l\rangle$, $|g_r\rangle$, and $|g_0\rangle$ as shown in Fig. 1. The transitions $|g_l\rangle \rightarrow |e_0\rangle$ and $|g_0\rangle \rightarrow |e_l\rangle$ are coupled with the left-polarized, $|g_r\rangle \rightarrow |e_0\rangle$ and $|g_0\rangle \rightarrow |e_r\rangle$ are coupled with the right-polarized cavity modes, respectively. The setup is shown in Fig. 2. Two distant atoms are trapped in two two-mode optical cavities. The information encoded on the atom (atom 1) to be teleported is partially transferred to cavity 1 via the atom-cavity interaction. Meanwhile, atom 2 is entangled with cavity 2. Photons leaking from both cavities are mixed on a beam splitter, which destroys which-path information, and then transmitted through quarter-wave

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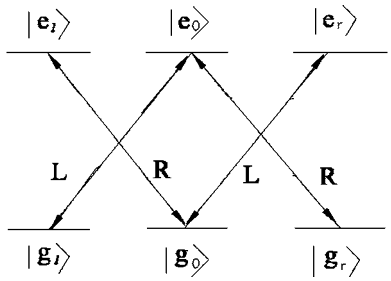

FIG. 1. The level configuration of the atoms. The transitions $|g_l\rangle \to |e_0\rangle$ and $|g_0\rangle \to |e_r\rangle$ are coupled to the left-circularly polarized cavity mode, while the transitions $|g_r\rangle \to |e_0\rangle$ and $|g_0\rangle \to |e_l\rangle$ are coupled to the right-circularly polarized cavity mode.

plates and polarization beam splitters. The photons are finally detected by photodetectors, which corresponds to a Bell-state measurement, as we will show. We here require that the cavities are one sided so that the only photon leakage occurs through the sides of the cavities facing the beam splitter.

In the interaction picture, the Hamiltonian in the kth $(k=1,2)$ cavity is

$$H_k = \sum_{j=l,r} [g_k a_{j,k}(|e_0\rangle\langle g_{j,k}| + |e_{j,k}\rangle\langle g_0|) + g_k a_{j,k}^\dagger(|g_{j,k}\rangle\langle e_0| + |g_0\rangle\langle e_{j,k}|)], \tag{1}$$

where $a_{l,k}$ and $a_{r,k}$ are the creation operators for the left-polarized and right-polarized cavity modes. We here assume that the transitions $|g_j\rangle \to |e_0\rangle$ and $|g_0\rangle \to |e_j\rangle$ have the same coupling strength g_k with the respective cavity modes in the respective cavities.

Assume the atom (atom 1) to be teleported is initially in the state

$$|\phi\rangle_1 = c_l|g_{l,1}\rangle + c_r|g_{r,1}\rangle, \tag{2}$$

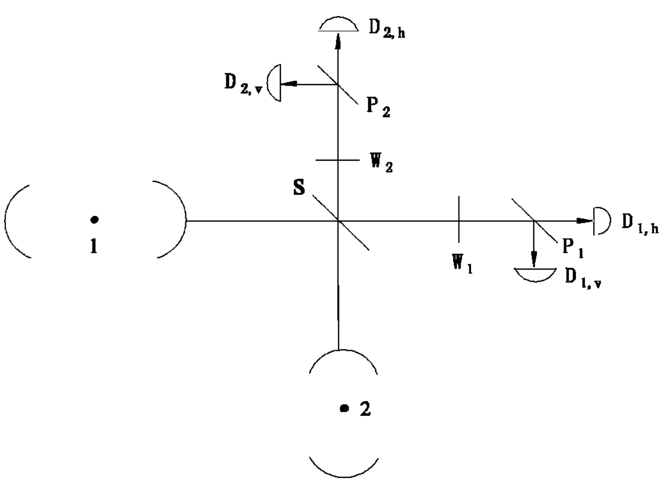

FIG. 2. The experimental setup to teleport the state of atom 1 trapped in a cavity to atom 2 trapped in another cavity. S is a beam splitter, W_1 and W_2 are quarter-wave plates, P_1 and P_2 are polarized beam splitters, and $D_{1,h}$, $D_{1,v}$, $D_{2,h}$, and $D_{2,v}$ are photodetectors.

where c_l and c_r are unknown coefficients. Perform the transformation $|g_{l,1}\rangle \to |e_{l,1}\rangle$ and $|g_{r,1}\rangle \to |e_{r,1}\rangle$, leading to

$$c_l|e_{l,1}\rangle + c_r|e_{r,1}\rangle. \tag{3}$$

The atom (atom 2) to receive the teleported state is initially prepared in the state $|e_0\rangle$. Both two cavities are initially in the vacuum states.

We now consider the cavity decay. Under the condition that no detection of a photon has been registered the system is governed by the non-Hermitian Hamiltonian

$$H_{c,k} = H_k - i\Gamma \sum_{j=l,r} a_{j,k}^\dagger a_{j,k}, \tag{4}$$

where Γ is the cavity decay rate. It should be noted that Γ may be much larger than the free-space rate, making spontaneous emission a minor problem. Assume that $g_1 \ll \kappa$, where κ is the cavity loss rate. Then the time evolution of the system combined by the first atom and the first cavity is

$$|\psi_1(t)\rangle = e^{-\Gamma t/2}\Bigg\{c_l\Bigg[\left(\frac{e^{\alpha t/2}+e^{-\alpha t/2}}{2} - \frac{\Gamma}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2}\right)|e_{l,1}\rangle|0_{l,1}\rangle \times|0_{r,1}\rangle - \frac{2g_1}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2i}|g_{0,1}\rangle|1_{l,1}\rangle|0_{r,1}\rangle\Bigg] + c_r\Bigg[\left(\frac{e^{\alpha t/2}+e^{-\alpha t/2}}{2} - \frac{\Gamma}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2}\right)|e_{r,1}\rangle|0_{l,1}\rangle|0_{r,1}\rangle - \frac{2g_1}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2i}|g_{0,1}\rangle|0_{l,1}\rangle|1_{r,1}\rangle\Bigg]\Bigg\}, \tag{5}$$

where

$$\alpha = \sqrt{-4g_1^2+\Gamma^2}. \tag{6}$$

During the evolution the atom may undergo the transition $|e_{l,1}\rangle \to |g_{0,1}\rangle$ or $|e_{r,1}\rangle \to |g_{0,1}\rangle$ accompanied by the creation of a left-polarized or right-polarized photon, respectively.

On the other hand, the time evolution of the system combined by the second atom and the second cavity is

$$|\psi_2(t)\rangle = e^{-\Gamma t/2}\Bigg[\left(\frac{e^{\beta t/2}+e^{-\beta t/2}}{2} - \frac{\Gamma}{\beta}\frac{e^{-\beta t/2}-e^{\beta t/2}}{2}\right) \times|e_{0,2}\rangle|0_{l,2}\rangle|0_{r,2}\rangle - \frac{2\sqrt{2}g_2}{\beta}\frac{e^{-\beta t/2}-e^{\beta t/2}}{2i} \times(|g_{l,2}\rangle|1_{l,2}\rangle|0_{r,2}\rangle + |g_{r,2}\rangle|0_{l,2}\rangle|1_{r,2}\rangle)\Bigg], \tag{7}$$

where

$$\beta = \sqrt{-8g_2^2+\Gamma^2}. \tag{8}$$

The whole system is in the state

$$|\psi(t)\rangle = |\psi_1(t)\rangle|\psi_2(t)\rangle. \tag{9}$$

The quarter-wave plates transform left-polarized and right-polarized photons into horizontally and vertically polarized photons, respectively. The beam splitter, quarter-wave plates, and polarization beam splitters perform the transformations

$$|1_{l,1}\rangle \to \frac{1}{\sqrt{2}}(|1_{h,1}\rangle - |1_{h,2}\rangle),$$

$$|1_{l,2}\rangle \to \frac{1}{\sqrt{2}}(|1_{h,1}\rangle + |1_{h,2}\rangle),$$

$$|1_{r,1}\rangle \to \frac{1}{\sqrt{2}}(|1_{v,1}\rangle - |1_{v,2}\rangle),$$

$$|1_{r,2}\rangle \to \frac{1}{\sqrt{2}}(|1_{v,1}\rangle + |1_{v,2}\rangle), \tag{10}$$

where $|1_{h,k}\rangle$ and $|1_{v,k}\rangle$ represent horizontally and vertically polarized photons transmitting through the kth path, respectively. Suppose that at time t_1 $D_{1,h}$ clicks. Then the system collapses to

$$|\psi_c(t_1)\rangle = \frac{1}{\sqrt{2}}(a_{l,1} + a_{l,2})|\psi(t_1)\rangle. \tag{11}$$

Using Eqs. (5), (7), (9), and (11), we have

$$\begin{aligned}|\psi_c(t_1)\rangle = {} & \frac{-1}{\sqrt{2}}e^{-\Gamma t_1}c_l\frac{2g_1}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2i}|g_{0,1}\rangle|0_{l,1}\rangle|0_{r,1}\rangle \\ & \times\left[\left(\frac{e^{\beta t_1/2}+e^{-\beta t_1/2}}{2}-\frac{\Gamma}{\beta}\frac{e^{-\beta t_1/2}-e^{\beta t_1/2}}{2}\right)|e_{0,2}\rangle|0_{l,2}\rangle\right. \\ & \times|0_{r,2}\rangle-\frac{2\sqrt{2}g_2}{\beta}\frac{e^{-\beta t_1/2}-e^{\beta t_1/2}}{2i}(|g_{l,2}\rangle|1_{l,2}\rangle|0_{r,2}\rangle \\ & \left.+|g_{r,2}\rangle|0_{l,2}\rangle|1_{r,2}\rangle)\right]-\frac{1}{\sqrt{2}}e^{-\Gamma t_1}\frac{2\sqrt{2}g_2}{\beta}\frac{e^{-\beta t_1/2}-e^{\beta t_1/2}}{2i} \\ & \times\left\{c_l\left[\left(\frac{e^{\alpha t_1/2}+e^{-\alpha t_1/2}}{2}-\frac{\Gamma}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2}\right)|e_{l,1}\rangle\right.\right. \\ & \left.\times|0_{l,1}\rangle|0_{r,1}\rangle-\frac{2g_1}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2i}|g_{0,1}\rangle|1_{l,1}\rangle|0_{r,1}\rangle\right] \\ & +c_r\left[\left(\frac{e^{\alpha t_1/2}+e^{-\alpha t_1/2}}{2}-\frac{\Gamma}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2}\right)|e_{r,1}\rangle|0_{l,1}\rangle\right. \\ & \left.\left.\times|0_{r,1}\rangle-\frac{2g_1}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2i}|g_{0,1}\rangle|0_{l,1}\rangle|1_{r,1}\rangle\right]\right\}|g_{l,2}\rangle \\ & \times|0_{l,2}\rangle|0_{r,2}\rangle.\end{aligned} \tag{12}$$

Since one cannot distinguish from which cavity the photon leaks due to the interference at the beam splitter the whole system is now in an entangled state.

Suppose that no photon is detected during the time interval $[t_1,t_2]$. Then the time evolution of the system during this time interval is governed by the conditional Hamiltonian of Eq. (4), leading to

$$\begin{aligned}|\psi'(t)\rangle = {} & \frac{-1}{\sqrt{2}}e^{-\Gamma(t_1+t)/2}c_l\frac{2g_1}{\alpha}\frac{e^{-\alpha t_1/2}-e^{\alpha t_1/2}}{2i}|g_{0,1}\rangle|0_{l,1}\rangle|0_{r,1}\rangle \\ & \times\left[\left(\frac{e^{\beta t/2}+e^{-\beta t/2}}{2}-\frac{\Gamma}{\beta}\frac{e^{-\beta t/2}-e^{\beta t/2}}{2}\right)|e_{0,2}\rangle|0_{l,2}\rangle|0_{r,2}\rangle\right. \\ & -\frac{2\sqrt{2}g_2}{\beta}\frac{e^{-\beta t/2}-e^{\beta t/2}}{2i}(|g_{l,2}\rangle|1_{l,2}\rangle|0_{r,2}\rangle+|g_{r,2}\rangle|0_{l,2}\rangle \\ & \left.\times|1_{r,2}\rangle)\right]-\frac{1}{\sqrt{2}}e^{-\Gamma(t_1+t)/2}\frac{2\sqrt{2}g_2}{\beta}\frac{e^{-\beta t_1/2}-e^{\beta t_1/2}}{2i} \\ & \times\left\{c_l\left[\left(\frac{e^{\alpha t/2}+e^{-\alpha t/2}}{2}-\frac{\Gamma}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2}\right)|e_{l,1}\rangle|0_{l,1}\rangle\right.\right. \\ & \left.\times|0_{r,1}\rangle-\frac{2g_1}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2i}|g_{0,1}\rangle|1_{l,1}\rangle|0_{r,1}\rangle\right] \\ & +c_r\left[\left(\frac{e^{\alpha t/2}+e^{-\alpha t/2}}{2}-\frac{\Gamma}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2}\right)|e_{r,1}\rangle|0_{l,1}\rangle|0_{r,1}\rangle\right. \\ & \left.\left.-\frac{2g_1}{\alpha}\frac{e^{-\alpha t/2}-e^{\alpha t/2}}{2i}|g_{0,1}\rangle|0_{l,1}\rangle|1_{r,1}\rangle\right]\right\}|g_{l,2}\rangle|0_{l,2}\rangle|0_{r,2}\rangle.\end{aligned} \tag{13}$$

If $D_{1,v}$ clicks at the time t_2 the system is collapsed to

$$|\psi_c'(t_2)\rangle = \frac{1}{\sqrt{2}}(a_{r,1} + a_{r,2})|\psi'(t_2)\rangle. \tag{14}$$

Using Eqs. (13) and (14), we obtain the state of the second atom

$$\begin{aligned}|\psi_{c,2}'(t_2)\rangle = {} & [c_l(e^{-\alpha t_1/2}-e^{\alpha t_1/2})(e^{-\beta t_2/2}-e^{\beta t_2/2})|g_{r,2}\rangle \\ & +c_r(e^{-\beta t_1/2}-e^{\beta t_1/2})(e^{-\alpha t_2/2}-e^{\alpha t_2/2})|g_{l,2}\rangle].\end{aligned} \tag{15}$$

Since two photons have been detected the two cavities are now in the vacuum states. On the other hand, the first atom is collapsed to the state $|g_{1,0}\rangle$. Suppose that $g_1=\sqrt{2}g_2$ and thus $\alpha=\beta$. Then we have

$$|\psi_{c,2}'(t_2)\rangle = c_l|g_{r,2}\rangle + c_r|g_{l,2}\rangle. \tag{16}$$

After a single-qubit transformation $|g_{r,2}\rangle \leftrightarrow |g_{l,2}\rangle$ the second atom is prepared in the initial state of the first atom.

If $D_{1,v}$ detects a vertically polarized photon at the time t_1 and $D_{1,h}$ detects a horizontally polarized photon at the time t_2 the second atom also collapses to the state of Eq. (16). If $D_{2,h}$ and $D_{2,v}$ click the second atom is again in the state of Eq. (16). On the other hand, if $D_{1,h}$ and $D_{2,v}$ click or $D_{1,v}$ and $D_{2,h}$ click the system collapses to the state

$$|\psi_{c,2}'(t_2)\rangle = c_l|g_{r,2}\rangle - c_r|g_{l,2}\rangle. \tag{17}$$

After the transformation $|g_{r,2}\rangle \to |g_{l,2}\rangle$ and $|g_{l,2}\rangle \to -|g_{r,2}\rangle$ the second atom is prepared in the initial state of the first atom.

We note that the quarter-wave platers, polarization beam splitters, and photodetectors constitute a device for the measurement of the two Bell states $(|1_{l,1}\rangle|1_{r,2}\rangle \pm |1_{r,1}\rangle|1_{l,2}\rangle)/\sqrt{2}$.

On the other hand, the two remaining Bell states $(|1_{l,1}\rangle|1_{l,2}\rangle \pm |1_{r,1}\rangle|1_{r,2}\rangle)/\sqrt{2}$ cannot be distinguished. Thus, the teleportation is probabilistic. Suppose that we wish to wait for a time t_d. The probability of success is

$$P = \frac{1}{2}\left\{1 - e^{-\Gamma t_d}\left[\frac{1}{4}e^{\beta t_d}\left(1 + \frac{\Gamma}{\beta}\right)^2 + \frac{1}{4}e^{-\beta t_d}\left(1 - \frac{\Gamma}{\beta}\right)^2 + \frac{1}{2}\left(1 - \frac{\Gamma^2}{\beta^2}\right) + \frac{2(g_2)^2}{\beta^2}(e^{\beta t_d} + e^{-\beta t_d} - 2)\right]\right\}^2. \quad (18)$$

If t_d is long enough so that $e^{-(\Gamma-\beta)t_d} \ll 1$ the probability is about $1/2$.

Suppose that the detector is imperfect. Then there is a probability that a left-polarized photon and a right-polarized photon have leaked from the cavities but only one click or no click is registered. These events are discarded. Thus, the probability of success is given by $P' = \eta^2 P$. However, the fidelity of the teleported state is not affected. We now consider the effect of the atomic spontaneous emission. The photons from the spontaneous emissions run with random directions and thus can not be detected by the photodetectors $D_{1,h}$, $D_{1,v}$, $D_{2,h}$, and $D_{2,v}$. Therefore, the atomic decay also reduces the success probability, but does not affect the fidelity.

We now briefly address the experimental feasibility of the proposed scheme. The required atomic level configuration can be achieved in ^{87}Rb. The hyperfine levels $|F=1,m=1\rangle$, $|F=1,m=0\rangle$, and $|F=1,m=-1\rangle$ of $5\,^2S_{1/2}$ can act as the states $|g_l\rangle$, $|g_0\rangle$, and $|g_r\rangle$, respectively, while the hyperfine levels $|F'=1,m'=1\rangle$, $|F'=1,m'=0\rangle$, and $|F'=1,m'=-1\rangle$ of $5\,^2P_{3/2}$ can act as $|e_l\rangle$, $|e_0\rangle$, and $|e_r\rangle$, respectively. The coupling coefficient for the transition $|5\,^2S_{1/2},F=1,m\rangle \rightarrow |5\,^2P_{3/2},F'=1,m'\rangle$ is given by $g_{m,m'} = g_0 C_{m,m'}$, where g_0 is the single-photon Rabi frequency and $C_{m,m'}$ the Clebsch-Gordan coefficient. The relevant Clebsch-Gordan coefficients are given by $C_{0,1}=C_{0,-1}=C_{1,0}=C_{-1,0}=1/\sqrt{2}$. Therefore, the transitions are equally coupled to the respective cavity modes. The transitions $|g_{l,1}\rangle \rightarrow |e_{l,1}\rangle$ and $|g_{r,1}\rangle \rightarrow |e_{r,1}\rangle$ can be achieved by applying a resonant π-polarized classical pulse [14]. As long as the classical pulse is strong enough the atom-cavity coupling is negligible during the application of the pulse. In order to perform the transformation $|g_{r,2}\rangle \rightarrow |g_{l,2}\rangle$ and $|g_{l,2}\rangle \rightarrow -|g_{r,2}\rangle$ we excite the second atom with a pair of off-resonant classical fields driving the transitions $|g_{r,2}\rangle \rightarrow |e_{0,2}\rangle$ and $|g_{l,2}\rangle \rightarrow |e_{0,2}\rangle$, respectively. The two classical fields are detuned from the respective transitions by the same amount δ. In the case that the detuning δ is much larger than the respective coupling strength the upper level $|e_{0,2}\rangle$ can be adiabatically eliminated and the two classical fields just induce the Raman transition between the states $|g_{r,2}\rangle$ and $|g_{l,2}\rangle$ [14]. We can obtain required transformations by controlling the phases of the classical fields and interaction time appropriately. The rotation $|g_{r,2}\rangle \rightarrow -|g_{r,2}\rangle$ can be achieved by applying a nonresonant right-polarized classical field to drive the transition $|g_{r,2}\rangle \rightarrow |e_0\rangle$. Suppose that the detuning Δ is much larger than the Rabi frequency Ω. In this case the state $|g_{r,2}\rangle$ undergoes a phase shift $\Omega^2 t/\Delta$, which is controllable via the Rabi frequency, detuning, and interaction time. The proposed scheme works in the Lamb-Dicke regime, i.e., the spatial extension of the atomic wave function, should be much smaller than the wavelength of the cavity mode. Recent advances in the combination of ion-trapping and cavity QED techniques allows the achievement of the Lamb-Dicke limit [15,16]. The coupling between each atom and the respective cavity mode depends upon the atomic position, $g=\Omega e^{-r^2/w^2}$, where Ω is the coupling strength at the cavity center, w is the waist of the cavity mode, and r is the distance between the atom and the cavity center [17]. We can satisfy the condition $g_2=\sqrt{2}g_1$ if we locate the first atom at the center of the first cavity and locate the other atom at the position $r=wln^{1/2}\sqrt{2}$.

In summary, we have proposed a scheme for long-distance teleportation of the state of an atom trapped in an optical cavity to a second atom trapped in another distant optical cavity via the detection of one photon leaking out from the cavities. The distinct advantage of our scheme is that it works in the regime where the cavity decay rate is larger than the atom-cavity coupling strength. Thus the requirement on the quality factor of the cavity is greatly loosened. Furthermore, the fidelity of our scheme is not affected by the detection inefficiency and atomic decay. Our scheme opens promising prospects for long-distance quantum communication as well as the test of fundamental features of quantum mechanics.

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