Teleportation of an arbitrary multipartite state via photonic Faraday rotation

Juan-Juan Chen¹,², Jun-Hong An¹, Mang Feng³ and Ge Liu²

¹ Center for Interdisciplinary Studies, Lanzhou University, Lanzhou 730000, People’s Republic of China
² Department of Modern Physics, Lanzhou University, Lanzhou 730000, People’s Republic of China
³ State Key Laboratory of Magnetic Resonance and Atomic and Molecular Physics, Wuhan Institute of Physics and Mathematics, Chinese Academy of Sciences, Wuhan 430071, People’s Republic of China

E-mail: anjhong@lzu.edu.cn

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Abstract
We propose a practical scheme for deterministically teleporting an arbitrary multipartite state, either product or entangled, using the Faraday rotation of the photonic polarization. Our scheme, based on the input–output process of single-photon pulses regarding cavities, works in low-\( Q \) cavities and only involves virtual excitation of atoms, which is insensitive to both cavity decay and atomic spontaneous emission. Besides, the Bell-state measurement is accomplished by the Faraday rotation plus product-state measurements, which could greatly reduce the experimental difficulty in realizing the Bell-state measurement by the CNOT operation.

1. Introduction
Teleportation is the faithful transfer of quantum states between spatially separated parties, based on the prior establishment of entanglement and classical communication [1]. As a practical application, teleportation has been implemented experimentally in several quantum systems [2, 3], and could be very useful for achieving quantum repeaters [4], quantum networks [5] and also quantum computing [6]. Recently, much attention has been paid to teleportation of multipartite states [7–9]. It has been shown that an arbitrary two-qubit state can be teleported using the multipartite entangled state as a quantum channel [9]. On the other hand, the cavity QED system gives a very nice platform to accomplish quantum teleportation. Many schemes have been proposed in this system. The earlier schemes use the atom as a flying qubit to transfer quantum information [10, 11], which are actually unsuitable for long-distance communication. Furthermore, the schemes mentioned above work well only in high-\( Q \) cavities, which is hard to accomplish with current technology. Besides, most of those schemes, e.g. [8, 11], are intrinsically probabilistic.

Can we achieve teleportation with low-\( Q \) cavities? More recently, schemes using photons to transfer information and atoms to store information have been proposed [12]. These schemes are based on the detection of photons leaking out of cavities. Taking the cavity damping into account, these schemes work in the low-\( Q \) cavity regime. However, the cavity damping actually plays a detrimental role in the success probability of teleportation. Thus, these schemes are intrinsically probabilistic. It shows the strong dependence of the implementation on the dissipative rate. This implies a negligible possibility of achieving a teleportation with low-\( Q \) cavities.

In this work, however, we will show this possibility with a practical scheme to teleport multipartite states using sophisticated low-\( Q \) cavities. Inspired by recent experiments of the photonic input–output process [3, 13], the key idea is to make use of the Faraday rotation produced by the single-photon-pulse input and output process regarding low-\( Q \) cavities. We had noticed a previous idea for teleportation using the Faraday rotation by single photons from microcavities at one location to those at another location in the quantum dot system [14]. While the photons’ input and output from the microcavities correspond to some decay effects, there has been little discussion about the impact from dissipation on teleportation. From our point of view, an efficient implementation of teleportation in [14] requires a high rate of the photons’ input and output with respect to the cavities. This implies the necessity of employing bad cavities with large decay rates, which goes beyond the model there.

The key point of our present scheme is to make use of a new reflection coefficient equation published recently [15], which enables us to get a rotation of the photonic polarization, named Faraday rotation [16], conditional on the atomic state.
confined in the cavity even in the low-$Q$ regime. We argue that our scheme would be advantageous over the previous schemes [12, 14] in accomplishment of teleportation in a deterministic fashion with currently available cavity QED technology, such as the microtoroidal resonator [13] or the single-sided optical cavity [17]. Furthermore, due to the disentangling action of the Faraday rotation to the photonic and atomic state at the destination side, only product-state measurement is needed in our scheme, which greatly relaxes the experimental difficulty.

This paper is organized as follows. In section 2, we first briefly review the input–output process of a single-photon pulse regarding a low-$Q$ cavity and show the Faraday rotation of the cavity and the new level structure in the cavity. On the other hand, if the atom is prepared in $|\psi\rangle$ (or $|\rho\rangle$) initially, then the photon will trigger the transition $|0\rangle \leftrightarrow |e\rangle$ (or $|1\rangle \leftrightarrow |e\rangle$). The quantum Langevin equations are

$$\begin{align*}
\dot{a}(t) &= -\frac{i}{2} [\omega_0 - \omega_p] + \frac{\kappa}{2} a(t) - g \sigma_- a(t) - \sqrt{\kappa} a_{in}(t), \\
\dot{\sigma}_-(t) &= -\frac{i}{2} [\omega_0 - \omega_p] + \frac{\kappa}{2} \sigma_- (t) - g \sigma_+(t)a(t) + \sqrt{\kappa} \sigma_+(t) b_{in}(t),
\end{align*}$$

where $\kappa$ and $\gamma$ are the cavity damping rate and atomic spontaneous emission rate, respectively. $b_{in}(t)$ is the vacuum input field felt by the atom. The contribution of the vacuum field $b_{in}(t)$ is negligible because it is much less than that of the photon pulse $a_{in}(t)$. The output field $a_{out}(t)$ relates to the input field by the intracavity field as $a_{out}(t) = a_{in}(t) + \sqrt{\kappa} a(t)$ [19]. Assume that $\kappa$ is large enough to ensure that the atom initially prepared in the ground state is only virtually excited by the photon, so we can take $|\sigma_{in}(t)\rangle = -1$. Also, under this large $\kappa$ limit, we can adiabatically eliminate the intracavity mode from the set of quantum Langevin equations and arrive consequently at the input–output relation as

$$r(\omega_p) = \frac{i[\omega_0 - \omega_p] - g^2}{[i[\omega_0 - \omega_p] + \frac{\kappa}{2}][i[\omega_0 - \omega_p] + \frac{\kappa}{2} + g^2]}$$

where $r(\omega_p) = \frac{a_{out}(t)}{a_{in}(t)}$ is the reflection coefficient of the photon to the atom–cavity system. On the other hand, if the atom is prepared in $|\rho\rangle$ (or $|\psi\rangle$) initially under the condition that an L- (or R-) circularly polarized photon pulse is input in the cavity, then no transition would be triggered. In other words, the photon only feels an empty cavity, for which the input–output relation corresponds to equation (2) with $g = 0$ [19], i.e. $r_0(\omega_p) = \frac{i[\omega_0 - \omega_p] - \frac{\kappa}{2}}{[i[\omega_0 - \omega_p] + \frac{\kappa}{2} + g^2]}$. (3)

The complex reflection coefficients (2) and (3) show that the reflected photon experiences an absorption as well as a phase shift denoted by $e^{i\phi}$ and $e^{i\phi_0}$, respectively. However, under practical experimental conditions, i.e. the strong $\kappa$ and weak $g$ and $\gamma$ [13], it has been verified that the absolute values of $r(\omega_p)$ and $r_0(\omega_p)$ are only slightly deviated from unity [15, 20]. This implies that the photon experiences a very weak absorption, and thereby we may approximately consider that the output photon only experiences a pure phase shift without any absorption.

With this basic input–output relation, the Faraday rotation can be derived readily when a linearly polarized photon is input into the cavity. Considering that the input pulse is linearly polarized in $|\psi_{in}\rangle = \frac{1}{\sqrt{2}}(|L\rangle + |R\rangle)$ and the atom initially in $|0\rangle$, the output photon would have a rotation in the polarization. In this case, the $|L\rangle$ component of the photon virtually triggers the transition of the atom from $|0\rangle$ to $|e\rangle$, and thus experiences a phase shift $e^{i\phi}$ obeying equation (2). In contrast, the $|R\rangle$ component of the photon only feels an empty cavity, which

![Figure 1](image-url)
yields a phase shift $e^{i\phi}$ obeying equation (3). So the output pulse is
$$|\Psi_{\text{out}}\rangle_+ = \frac{1}{\sqrt{2}} (e^{i\phi} |L\rangle + e^{i\phi} |R\rangle).$$

This also implies that the polarization direction of the reflected photon rotates an angle $\Theta_+^R = \frac{\phi - \phi_0}{2\pi}$ with respect to the input one, called the Faraday rotation. Similarly, if the atom is initially prepared in $|1\rangle$, then only the $R$ circularly polarized photon can sense the atom, while the $L$ circularly polarized photon only feels the empty cavity. So we have
$$|\Psi_{\text{out}}\rangle_+ = \frac{1}{\sqrt{2}} (e^{i\phi} |L\rangle + e^{i\phi} |R\rangle),$$
corresponding to a Faraday rotation with an angle $\Theta_+^R = \frac{\phi - \phi_0}{2\pi}$.

3. The realization of multipartite-state teleportation via Faraday rotation

In what follows, we show the teleportation of a multipartite state by the above Faraday rotations. As plotted in figure 1(c), to construct the entanglement between Alice and Bob, Bob has to first input $N$ single-photon pulses into his cavities to produce $N$ pairs of entangled states between the photons and atoms using the Faraday rotations of the photons. Then the emitted photons fly to Alice through the fibre. Once Alice collects the photons in her cavities, the entanglement between Alice and Bob has been established. We emphasize that the photon in our scheme acts dual roles: on the one hand, it acts as a bus to distribute entanglement; on the other hand, it also acts as a component of the entangled pair to implement teleportation.

3.1. The case of the bipartite state

To demonstrate our scheme specifically, we will take $N = 2$ below as an example. Suppose the quantum state to be teleported in Alice’s hands is
$$|\psi\rangle = \alpha|01\rangle + \beta|10\rangle + \gamma|00\rangle + \delta|11\rangle,$$
where $\alpha, \beta, \gamma$ and $\delta$ are unknown parameters with $\alpha^2 + \beta^2 + \gamma^2 + \delta^2 = 1$. In the following, we give explicitly the steps for accomplishing the quantum teleportation.

The first step: establishment of the quantum channels.

Bob inputs two linearly polarized single-photon pulses in $|\Psi_i\rangle = \frac{1}{\sqrt{2}}(|L_i\rangle + |R_i\rangle)$ generated from the single-photon source into two cavities at his side. The states of the atoms confined in Bob’s cavities are initially $|\Psi_i\rangle = \frac{1}{\sqrt{2}} (|0\rangle_i + |1\rangle_i) (i = 1, 2)$. According to the input–output relation, the atoms and photons are entangled due to the Faraday rotation:

$$|\Psi_i\rangle |\psi\rangle \rightarrow \frac{1}{\sqrt{2}} (|0\rangle_i |\Psi_{\text{out}}\rangle_{-i} + |1\rangle_i |\Psi_{\text{out}}\rangle_{+i}).$$

Then after Alice collects the two output photons via two fibres, two quantum channels are thus established. One can see that $|\Psi_{\text{out}}\rangle_+$ and $|\Psi_{\text{out}}\rangle_-$ are orthogonal when $\phi - \phi_0 = \pi/2$. It means that under this condition, the quantum channel characterized by the right-hand side of equation (7) is maximally entangled. In practice, this condition can be easily satisfied experimentally. Using the parameters in [13], $\omega_0 = \omega_a$, $\omega_p = \omega_c - \frac{\pi}{2}$ and $g = \frac{\pi}{2}$, we can verify from equations (2) and (3) that $\phi = \pi$ and $\phi_0 = \pi/2$. The Faraday rotation under the same condition has been used to realize the universal quantum gate, the entanglement generation between remote atoms and its conversion to the flying photons in [15, 21].

The second step: realization of the Bell-state measurement.

In the standard protocol of quantum teleportation, a necessary step is the Bell-state measurement on the atom and one of the entangled pairs of the quantum channel at Alice’s side [1]. The measurement would collapse the state of the total system in one of its four superposition components. Experimentally, this Bell-state measurement is ordinarily realized by the CNOT operation plus the product-state measurement. In the following, we show that the CNOT operation can actually be replaced by the Faraday rotation at Alice’s side, which much relaxes the experimental difficulty to do the CNOT operation. After collecting the photons, Alice feeds the photons to her cavities. The Faraday rotation makes the two atoms at Alice’s hands entangled with the two quantum channels. So the state of the entire system could be written as

$$\psi \rightarrow \frac{1}{\sqrt{2}} (|00\rangle |\Psi_{\text{out}}\rangle_{-1} + |11\rangle |\Psi_{\text{out}}\rangle_{+1})$$

$$\rightarrow \frac{1}{4} (|00\rangle \langle 00| + |10\rangle \langle 10| + |01\rangle \langle 01| + |11\rangle \langle 11|)$$

$$\times (|00\rangle_{12} - i|01\rangle_{12} - i|10\rangle_{12} - |11\rangle_{12})$$

$$+ |0\rangle_{12} |0\rangle_{12} - |1\rangle_{12} |1\rangle_{12} + i|1\rangle_{12} |0\rangle_{12} - i|0\rangle_{12} |1\rangle_{12} + |0\rangle_{12} |0\rangle_{12} - |1\rangle_{12} |1\rangle_{12} - i|0\rangle_{12} |1\rangle_{12} + i|1\rangle_{12} |0\rangle_{12})$$

$$\rightarrow \frac{1}{4} (|00\rangle - i|01\rangle - i|10\rangle + |11\rangle)$$

$$\times (|00\rangle_{12} - i|01\rangle_{12} - i|10\rangle_{12} - |11\rangle_{12}),$$

where the kets with and without subscripts correspond to the atomic states at Bob’s and Alice’s hands, respectively. We have used $\phi = \pi$ and $\phi_0 = \pi/2$ with the practical parameters in [13].

Next, Alice makes Hadamard gates on the photons and atoms at her side. The photonic Hadamard gating is realized by a quarter-wave plate (QWP), which makes $|R\rangle \rightarrow \frac{|L\rangle - |R\rangle}{\sqrt{2}}$ and $|L\rangle \rightarrow \frac{|L\rangle + |R\rangle}{\sqrt{2}}$. Then the right-hand side of equation (8) is converted into $\sum_{i,j,m,n} |ij\rangle |mn\rangle |f_{ij,mn}\rangle_{12}$, where $|ij\rangle$, $|mn\rangle$ and $|f_{ij,mn}\rangle_{12}$ denote the photon, and Alice’s and Bob’s atomic states, respectively, as shown in table 1.

Finally, Alice performs measurement on the states of the photons and the atoms at her side, followed by the collapse of the atomic state at Bob’s side to one of the corresponding components in above superposition.

The third step: Bob’s recovery operations.

Based on the message $(i, j, m, n)$ from Alice via the classical channel about her measurement, Bob could deterministically recover
For simplicity, we take a tripartite state as an example. With more single-photon pulses, we may generalize our results. In the last step, Bob has to choose local operations $M_{ij,mn}$ appropriately, as shown in table 2, to recover the unknown teleported state on his atom.

### 4. Discussions and summary

Our scheme, involving only virtual excitation of the atoms, is robust to spontaneous emission. Besides, as cavity decay has been considered in the reflection coefficient (2), our scheme could work with bad cavities in the case of large cavity decay and/or weak couplings. The dominant source of error in our scheme is the photon loss due to the cavity mirror absorption and scattering, the fibre absorption and the inefficiency of the detector. Nevertheless, since the accomplishment of our scheme relies on the successful detection of the photons, the photon loss does not affect the fidelity, but only the efficiency.

The Hadamard gate operations on the atoms could be done using the Raman configuration by two polarized lasers detuned from the $\ket{0} \leftrightarrow \ket{e}$ and $\ket{1} \leftrightarrow \ket{e}$ transitions. In addition, the measurement of the atomic states could be carried out by a resonant laser. For example, the successful detection of a leaking photon from the cavity after the radiation on the atom by an L-polarized laser means the population in the atomic state $\ket{0}$. Otherwise, we may try to detect the atomic state $\ket{1}$ using a R-polarized laser.

We argue that the implementation of our scheme heavily depends on the imperfection rate due to relevant techniques. Under the currently available technique, our scheme could be accomplished within a finite time if the qubit number $N$ is not very large. Specifically, supposing that the failure rate associated with the decay of the virtual atomic excitation is about 2%, the current dark count rate of the single-photon detector yields the inefficiency of $10^{-4}$, and other imperfection rate due to photon loss is about 6%; we thus have the success rate to be $(1-2\%)^2 \times 10^{-4} \times (1-6\%)^N$. Fortunately, thanks to the highly efficient single-photon generator producing 10 000 photons every second [22], we may accomplish teleportation in finite time. For example, a successful teleportation of a two-qubit state takes about 3 h.

This time-consuming implementation is mainly resulted from the low efficiency of currently available single-photon detectors, which is also a problem in any of the previously published schemes using photon interference. To shorten the implementation time, on the one hand, we should increase the resource. For example, if Bob possesses $N$ sets of a single-photon source, then the implementation time would be reduced by $N$ times. On the other hand, if the efficiency of the single-photon detector could be enhanced, then our implementation time would be much reduced by several degrees of magnitude. A simple estimate of the implementation time of our scheme is plotted in figure 2 with respect to different imperfect factors.

### Table 1. The superposition components of the final state and Bob’s corresponding operations when $N = 2$.

| $|ij, mn\rangle$ | $|f_{ij,mn}\rangle_{12}$ | $M_{ij,mn}$ |
|-----------------|----------------|----------------|
| $|LL\rangle_{00}$ | $-\alpha|0\rangle - \beta|01\rangle - \zeta|11\rangle - \delta|00\rangle$ | $-\sigma^{(1)} \otimes \sigma^{(2)}$ |
| $|LL\rangle_{11}$ | $\alpha|0\rangle + \beta|01\rangle + \zeta|11\rangle + \delta|00\rangle$ | $\sigma^{(1)} \otimes \sigma^{(2)}$ |
| $|LL\rangle_{01}$ | $-\alpha|0\rangle + \beta|01\rangle + \zeta|11\rangle - \delta|00\rangle$ | $\sigma^{(1)} \otimes \sigma^{(2)}$ |
| $|LL\rangle_{10}$ | $\alpha|0\rangle - \beta|01\rangle + \zeta|11\rangle - \delta|00\rangle$ | $\sigma^{(1)} \otimes \sigma^{(2)}$ |
| $|LR\rangle_{00}$ | $\iota |0\rangle - \iota |1\rangle + \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|LR\rangle_{11}$ | $\iota |0\rangle + \iota |1\rangle + \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|LR\rangle_{01}$ | $-\iota |0\rangle + \iota |1\rangle + \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|LR\rangle_{10}$ | $\iota |0\rangle + \iota |1\rangle - \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|RL\rangle_{00}$ | $\iota |0\rangle + \iota |1\rangle - \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|RL\rangle_{11}$ | $\iota |0\rangle + \iota |1\rangle + \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|RL\rangle_{01}$ | $-\iota |0\rangle + \iota |1\rangle - \iota |0\rangle + \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |
| $|RL\rangle_{10}$ | $\iota |0\rangle + \iota |1\rangle - \iota |0\rangle - \iota |1\rangle$ | $I^{(1)} \otimes I^{(2)}$ |

which could be further transformed into the final form $\sum_{ijk,lmn} |ijk, lmn\rangle_{12} f_{ijk,lmn}$ by local Hadamard operations applied to the photonic and to the atomic states at Alice’s side. The explicit form of the superposition components in the final state is listed in table 2. Finally, Alice performs measurement on the state of the photons and atoms and then tells Bob her results. In the last step, Bob has to choose local operations $M_{ij,mn}$ appropriately, as shown in table 2, to recover the unknown teleported state on his atom.
as quantum channels is more robust to decoherence than others.

Figure 2. The implementation time of teleportation as a function of the number of the parties involved, with different single-photon detector inefficiency \(\eta\).

In comparison with previous proposals for teleportation considering virtually excited cavity modes [10] and cavity decay [12], our scheme is advantageous to work very well not only in the case of the bad cavities, but also in perfect and deterministic fashion. Using photons as flying qubits to transfer quantum information is more suitable for long-distance communication compared to the schemes in [10, 11]. Moreover, our scheme using bipartite entanglement as quantum channels is more robust to decoherence than others based on multipartite entanglement [9].

In conclusions, we have proposed a practical scheme for teleportation of an arbitrary \(N\)-partite pure state using Faraday rotation. Besides its use in atomic systems, our idea could also be applied to the quantum-dot system after minor modification: replacing the atomic excitations by the excitonic ones [14].

As it needs no CNOT operation, only involves product-state measurements, and works perfectly and deterministically in low-\(Q\) cavities, our scheme would be useful for building quantum network and for scalable quantum computation using currently achieved microtoroidal resonators [13] or single-sided cavities [17].

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Table 2. The superposition components of the final state and Bob’s corresponding operations when \(N = 3\).

| \(|ijkl\rangle|mnm\rangle\) | \(M_{ijklmn}\) | \(f_{ijklmn}\) | \(\beta_{ijklmn}\) |
|-----------------|-----------------|-----------------|-----------------|
| \(|LLL\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |
| \(|RRR\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |
| \(|LLR\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |
| \(|RLR\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |
| \(|RLR\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |
| \(|RLR\rangle\) | \(000, 011, 101, 110\) | \(\alpha(000) - \beta(001)\) | \(\alpha(000) - \beta(001)\) |

\(0 \leq T \leq 10\) [unit of hours].

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