Controllable and fast quantum-information transfer between distant nodes in two-dimensional networks

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We construct shortcuts to adiabatic passage to achieve controllable and fast quantum-information transfer (QIT) between arbitrary two distant nodes in a two-dimensional (2D) quantum network. Through suitable designing of time-dependent Rabi frequencies, we show that perfect QIT between arbitrary two distant nodes can be rapidly achieved. Numerical simulations demonstrate that the proposal is robust to the decoherence caused by atomic spontaneous emission and cavity photon leakage. Additionally, the proposed scheme is also insensitive to the variations of the experimental parameters. Thus, the proposed scheme provides a new perspective on robust quantum information processing in 2D quantum networks.

In recent years, the QIT has been accomplished by several approaches, i.e., the resonant $\pi$ pulses, composite pulses, stimulated Raman adiabatic passage (STIRAP), and their variants1–3. Although the resonant $\pi$ pulse technique can fast transfer quantum information1, its highly sensitivity to the deviations of pulse areas restricts its extensive application in quantum information processing. The adiabatic passage techniques2,3 are robust versus variations of the experimental parameters while they usually need a long operation time. Thus the decoherence, which is one of the main obstacles in quantum information and quantum computation, would strongly affect the dynamics of the system, furthermore, may lead to the schemes become useless. The “shortcuts to adiabaticity technique”, which combines the advantages of resonant $\pi$ pulses and adiabatic techniques, has been considered as a promising venue to achieve fast and high-fidelity QIT, and has attracted much attention in recent years4–15. In view of shortcuts to adiabaticity, Chen and Muga6,7 have successfully performed fast population transfer in three-level atom systems via applying the opposite variation tendency in the time-dependent laser pulse. After that, the shortcuts to adiabaticity technique has been extended from one-atom system to two-or multi-atom system12–15.

The scalability is still another obstacles in accomplishing the quantum information and quantum computation under current cavity quantum electrodynamics technology. The emergence of coupled cavity system16 which can overcome the scalability and meet the requirement of kinds of quantum tasks, i.e., simulation of quantum many-body phenomena17–23, performing remote quantum information transfer4–26, entanglement generation27–35 and quantum gate operations between two distant nodes36,37. All such works typically focus on the cases of either two-site or one-dimensional (1D) coupled cavity arrays. Extending such researches to more complex coupled cavity arrays (i.e., two-dimensional (2D) or three-dimensional (3D)) is more significance for quantum computation. There have been several studies considering the 2D coupled cavity arrays, which have respectively considered the realization of the fractional quantum Hall system38 and 2D one-way quantum computation39. Recently, we have proposed protocols to realize the coherent coupling of multiple atoms40 and to realize two-qubits unconventional geometric phase gates in a 2D coupled cavity array41.

The quantum Zeno effect is an interesting phenomenon in quantum mechanics and has been demonstrated in many experiments42–45. It has been shown that a system can actually evolve away from its initial state, but still remain in the Zeno subspace defined by the measurements via frequently projecting onto a multidimensional subspace, which is known as quantum Zeno dynamics (QZD)46–48. In general, if a system is governed by Hamiltonian $H_{Z} = H_{\text{obs}} + KH_{\text{meas}}$, where $H_{\text{obs}}$ is the Hamiltonian of the subsystem to be investigated, $H_{\text{meas}}$ is an...
additional interaction Hamiltonian performing the “measurement”, and $K$ is a coupling constant. In a strong coupling limit $K \to \infty$, the whole system will remain in the same Zeno subspace, and is governed by the evolution operator defined as

$$\exp(−\sum_{n} i\omega_n t \mathcal{P}_n + \mathcal{P}_n t \mathcal{H}_{\text{obs}} P_n),$$

with $\mathcal{P}_n$ being the eigenvalue projection of $\mathcal{H}_{\text{obs}}$ with eigenvalues $E_n$ ($\mathcal{H}_{\text{obs}} = \sum_n E_n P_n$).

Motivated by the space division of QZD, in this paper, we construct shortcuts to adiabatic passage to achieve controllable and fast QIT between arbitrary two nodes in a 2D quantum network. Through suitably designing the time-dependent Rabi frequencies, we can controllably and fast transfer quantum-information between arbitrary two distant nodes in one-step. The distinguished advantages of the proposal are: (i) information can be controllably transferred between arbitrary two nodes; (ii) the time to accomplish the task is shorter than that in conventional adiabatic passage technique; (iii) it is robust against the parameters fluctuations and the decoherence caused by atomic spontaneous emission and cavity photon leakage. Thus it provides a new perspective on robust quantum information processing in 2D quantum networks in the future.

The theoretical model and the construction of a shortcut passage

we consider a 2D ($N \times N$) coupled cavity array, as shown in Fig. 1(a). Each cavity (denoted by $jk$) respectively couples to their neighboring ones through the $x$ and $y$ directions with intercavity photon hopping. Each cavity contains a $\Lambda$-type atom. The atoms have two ground states (labeled as $|g\rangle_{jk}$, $|f\rangle_{jk}$) and one excited state (labeled as $|e\rangle_{jk}$), as shown in Fig. 1(b). The $|g\rangle_{jk} \leftrightarrow |e\rangle_{jk}$ transition of atom couples to the cavity mode with detuning $\Delta_{jk}$, the corresponding coupling rate is $g_{jk}$. The $|f\rangle_{jk} \leftrightarrow |e\rangle_{jk}$ transition of atom is resonantly driven by a classical laser field, and the corresponding Rabi frequencies are $\Omega_{jk}$.

**Figure 1.** (a) Schematic diagram of a two-dimensional (2D) array of coupled cavities. Each node contains a $\Lambda$-type three-level atom and can respectively couple to their neighboring ones through the $x$ and $y$ directions with intercavity photon hopping. (b) The atom level scheme. The transition of the $jk$th atom $|g\rangle_{jk} \leftrightarrow |e\rangle_{jk}$ is coupled to the cavity mode with detuning $\Delta_{jk}$, the corresponding coupling rate is $g_{jk}$. The transition $|f\rangle_{jk} \leftrightarrow |e\rangle_{jk}$ of the $jk$th atom is resonantly driven by a classical laser field, and the corresponding Rabi frequencies are $\Omega_{jk}$.
The above Hamiltonian \( H_{\text{eff}} \) shows that, the atoms can resonantly interact with the nonlocal bosonic mode \( c_{\text{m,n}} \) means that the atoms resonantly interact with all the cavities simultaneously. Assume the system is initially in the state \(|\Phi_0\rangle = |f_{11}\rangle|g_{\text{NN}}\rangle|0\rangle \) (i.e., atoms in the node 11 and node NN are in the states \(|f\rangle\) and \(|g\rangle\), respectively, and the bosonic mode \( c_{\text{m,n}} \) is in the vacuum state), the whole system evolves in the subspace spanned by \(|\Phi_0\rangle\) and \(|\Phi_1\rangle\), the interaction of the atoms with the nonresonant normal modes can be neglected, the Hamiltonian reads

\[
H_{\text{eff}} = H_{1\text{eff}} + H_{2\text{eff}},
\]

where

\[
H_{1\text{eff}} = \sum_{j,k=11,\text{NN}} \Omega_{jk}|e_{jk}\rangle\langle f| + \text{ H.c.},
\]

and

\[
H_{2\text{eff}} = \sum_{j,k=11,\text{NN}} \sum_{m,n=p,q,jk=11} \frac{g_{pq}}{N} e^{-i\frac{2\pi}{N}j} e^{i\frac{2\pi}{N}k} c_{m,n}|e_{jk}\rangle\langle g| + \text{ H.c.}.
\]
$$|\Psi_1\rangle = \frac{1}{\sqrt{2}}(-|\Phi_2\rangle + |\Phi_4\rangle), \quad |\Psi_2\rangle = \frac{1}{2}(|\Phi_2\rangle - \sqrt{2}|\Phi_3\rangle + |\Phi_4\rangle).$$
$$|\Psi_3\rangle = \frac{1}{2}(|\Phi_2\rangle + \sqrt{2}|\Phi_3\rangle + |\Phi_4\rangle).$$

Next, we rewrite the Hamiltonian in Eq. (9) with the eigenvectors of \( H_{\text{eff}} \)

$$H'_{\text{eff}} = H'_{\text{eff}} + H'_{\text{eff}}^\prime,$$

where

$$H'_{\text{eff}} = \sum_{i=1}^{3} E_i |\Psi_i\rangle \langle \Psi_i|,$$

and

$$H'_{\text{eff}}^\prime = \frac{\Omega_{11}(t)}{\sqrt{2}} |\Psi_1\rangle \langle \Phi_1| + \frac{\Omega_{13}(t)}{2} (e^{i\tau \gamma \beta} |\Psi_3\rangle + e^{-i\tau \gamma \beta} |\Psi_3\rangle) \langle \Phi_3| + \frac{\Omega_{NN}(t)}{\sqrt{2}} |\Psi_5\rangle \langle \Phi_5|$$

$$+ \frac{\Omega_{NN}(t)}{2} (e^{i\tau \gamma \beta} |\Psi_3\rangle + e^{-i\tau \gamma \beta} |\Psi_3\rangle) \langle \Phi_5| + \text{H.c.}$$

It is obvious that there are four nonzero energy eigenvalues \( \pm \Omega_{11}(t) \) and \( \pm \Omega_{NN}(t) \) for the Hamiltonian \( H'_{\text{eff}} \). Defining \( H''_{\text{eff}} \) as a free Hamiltonian, and performing the unitary transformation \( U = e^{-i\tau t} \) under condition \( H''_{\text{eff}} \gg H'_{\text{eff}} \), we obtain

$$H''_{\text{eff}} = \frac{\Omega_{11}(t)}{\sqrt{2}} |\Psi_1\rangle \langle \Phi_1| + \frac{\Omega_{13}(t)}{2} (e^{i\tau \gamma \beta} |\Psi_3\rangle + e^{-i\tau \gamma \beta} |\Psi_3\rangle) \langle \Phi_3| + \frac{\Omega_{NN}(t)}{\sqrt{2}} |\Psi_5\rangle \langle \Phi_5|$$

$$+ \frac{\Omega_{NN}(t)}{2} (e^{i\tau \gamma \beta} |\Psi_3\rangle + e^{-i\tau \gamma \beta} |\Psi_3\rangle) \langle \Phi_5| + \text{H.c.}$$

Therefore, setting \( \sqrt{2} \gg \Omega_{11}(t), \Omega_{NN}(t) \), the condition \( H''_{\text{eff}} \gg H'_{\text{eff}} \) and the Zeno condition \( K \to \infty \) are satisfied. Under the rotating-wave approximation, we have a new Hamiltonian

$$H''_{\text{eff}} = \frac{\Omega_{11}(t)}{\sqrt{2}} |\Psi_1\rangle \langle \Phi_1| + \frac{\Omega_{NN}(t)}{\sqrt{2}} |\Psi_5\rangle \langle \Phi_5| + \text{H.c.}$$

Thus the Hilbert subspace splits into three invariant Zeno subspaces \( H''_{\text{eff}} = \{|\Psi_1\rangle, |\Phi_1\rangle, |\Phi_5\rangle\}, H''_{\text{eff}} = \{|\Psi_5\rangle\}, H''_{\text{eff}} = \{|\Psi_3\rangle\} \). The system can be divided into three subsystems, \( S_1 = \{|\Psi_1\rangle, |\Phi_3\rangle, |\Phi_5\rangle\}, S_2 = \{|\Psi_2\rangle, |\Phi_2\rangle, |\Phi_4\rangle\}, S_4 = \{|\Psi_3\rangle\} \). Note that the interaction between the states in subsystems \( S_1 \) and \( S_2 \) is far weaker than that in subsystem \( S_4 \), thus this weak interaction can be neglected. Then the system can be considered as a three-level atom system with two ground states \( |\Phi_1\rangle, |\Phi_5\rangle \) and an excited state \( |\Psi_3\rangle \). The Hamiltonian for STIRAP reads

$$H_{\text{eff}} = \frac{1}{\sqrt{2}} \begin{bmatrix} 0 & \Omega_{11}(t) & 0 \\ \Omega_{11}(t) & 0 & \Omega_{NN}(t) \\ 0 & \Omega_{NN}(t) & 0 \end{bmatrix}$$

The instantaneous eigenvalues are \( \chi_0 = 0, \chi_\pm = \pm \sqrt{\Omega_{11}^2 + \Omega_{NN}^2 / \sqrt{2}}, \) with the corresponding eigenstates are

$$|\Phi_0(t)\rangle = \begin{bmatrix} \cos \theta \\ 0 \\ \sin \theta \end{bmatrix}, \quad |\Phi_\pm(t)\rangle = \begin{bmatrix} -\sin \theta \\ \pm 1 \\ \cos \theta \end{bmatrix}.$$
\[ \gamma = \frac{1}{\sqrt{2}}[\Omega_{11}(t)\cos \beta - \Omega_{NN}(t)\sin \beta], \quad \beta = \frac{1}{\sqrt{2}}\tan \gamma[\Omega_{NN}(t)\cos \beta + \Omega_{11}(t)\sin \beta], \]

where the dot represents a time derivative. By inversely deriving from Eq. (17), the explicit expressions of \( \Omega_{11}(t) \) and \( \Omega_{NN}(t) \) are as follows:

\[ \Omega_{11}(t) = \sqrt{2}[\beta \cot \gamma \sin \beta - \gamma \cos \beta], \quad \Omega_{NN}(t) = \sqrt{2}[\beta \cot \gamma \cos \beta - \gamma \sin \beta]. \]

The eigenstates of the invariant \( I_{2}(t) \) are

\[ |\Phi^{g}_{0}(t)\rangle = \begin{pmatrix} \cos \gamma \cos \beta \\ -i \sin \gamma \\ -\cos \gamma \sin \beta \end{pmatrix}, \quad |\Phi^{f}_{1}(t)\rangle = \begin{pmatrix} \sin \gamma \cos \beta \pm i \sin \beta \\ i \cos \gamma \\ -\sin \gamma \sin \beta \pm i \cos \beta \end{pmatrix}, \]

corresponding to the eigenvalues \( \lambda_{0} = 0 \) and \( \lambda_{\pm} = \pm 1 \), respectively. Based on the Lewis-Riesenfeld theory, the solution of the Schrödinger equation with respect to the instantaneous eigenstates of \( I_{2}(t) \) is a superposition of orthonormal dynamical modes, \( |\Psi(t)\rangle = \sum_{n}C_{n}e^{i\alpha_{n}}|\Phi^{g}_{n}(t)\rangle \), where \( C_{n} \) is a time-independent amplitude and \( \alpha_{n} \) is the Lewis-Riesenfeld phase and obeys the form,

\[ \alpha_{n}(t_{f}) = \frac{1}{\hbar} \int_{0}^{t_{f}} \langle \Phi^{g}_{n}(t) | \frac{\partial}{\partial t} - H_{2}(t) | \Phi^{g}_{n}(t) \rangle dt'. \]

In the proposal, \( \alpha_{g} = 0 \), and

\[ \alpha(\pm) = \mp \frac{1}{\hbar} \int_{0}^{t_{f}} \beta \sin \gamma - \frac{1}{2}(\Omega_{11} \sin \beta + \Omega_{NN} \cos \beta) \cos \gamma dt'. \]

In order to get the target state \( |\Phi_{t}\rangle \) along the invariant eigenstate \( |\Phi^{g}_{0}(t)\rangle \), we suitably choose the feasible parameters \( \gamma(t) \) and \( \beta(t) \),

\[ \gamma(t) = \zeta, \quad \beta(t) = \pi t/2t_{f}, \]

where \( \zeta \) is a small value, which satisfies \( (\sin \zeta)^{-1} = 4M(M=1,2,3,...) \) for a high fidelity of the target state. And we obtain

\[ \Omega_{11}(t) = \pi t/(\sqrt{2}t_{f})\cot \zeta \sin(\pi t/2t_{f}), \quad \Omega_{NN}(t) = \pi t/(\sqrt{2}t_{f})\cot \zeta \cos(\pi t/2t_{f}). \]

Once the Rabi frequencies are specially designed, the fast QIT from initial state to the target state in subsystem \( S_{t} \) will be implemented.

**Results**

To confirm the validity of all our above derivation, we first numerically simulate the dynamics governed by the derived effective Hamiltonian in Eq. (14), and compare it to the dynamics governed by the total Hamiltonian in Eq. (1). Note that the numerical computation we performed using the python package Qutip. The validity of the model is numerically simulated by taking the evolution of the population \( P = |\langle \psi(t) | \psi(t) \rangle|^{2} \) of the proposed state \( |\psi\rangle \). We consider the case with \( N = 2 \), and set the parameters in the following way: \( \nu = 2.0 \text{ g}, \quad g_{f} = 50, \quad \xi = \arcsin(0.25) \) (the Zeno condition \( \sqrt{2}g > \Omega_{k}(t) \) can be satisfied very well). For the total Hamiltonian, a new subspace is spanned by \( \{ |\psi_{1}\rangle = |f_{11},g_{22},0\rangle, |01,01,02\rangle, |22,01,01\rangle, |22,00,00\rangle \} \), \( |\psi_{2}\rangle = |f_{11},g_{22},0\rangle, |01,02,02\rangle, |01,02,02\rangle, |22,01,01\rangle \), \( |\psi_{3}\rangle = |f_{11},g_{22},1\rangle, |01,02,02\rangle, |01,02,02\rangle, |22,01,01\rangle \), \( |\psi_{4}\rangle = |f_{11},g_{22},1\rangle, |11,01,01\rangle, |01,01,01\rangle, |22,01,01\rangle \). Thus, if the system is initially in one of these basis, the system will evolve in this subspace. In Fig. 2, the red-solid (green-solid) and blue-dashed (black-dashed) lines describe the time evolution of the population of state \( |\psi_{1}\rangle \) in \( f_{11},g_{22},0\rangle, |01,01,02\rangle, |22,01,01\rangle \) and state \( |\psi_{2}\rangle \) in \( f_{11},g_{22},0\rangle, |01,02,02\rangle, |01,02,02\rangle, |22,01,01\rangle \), respectively. It is obvious that the approximations adopted during the deriving of the effective Hamiltonian are valid, since the two curves described by the total Hamiltonian and effective Hamiltonian are nearly coincided, and their deviation is small enough as soon as the parameters are fixed.

Next, we show how the operation time is shorten when considering the shortcuts to adiabatic passage. We first numerically simulate the time dependence of the Rabi frequencies for the atoms in Fig. 3(a) when \( g_{f} = 10 \), the other parameters are set the same as those in Fig. 2. As seen from Fig. 3(a), the maximum value of \( \Omega_{k}(t) \) is 0.83, which satisfies the conditions mentioned above (the Zeno condition \( \sqrt{2}g > \Omega_{k}(t) \) is \( k = 11, NN \)). In Fig. 3(b), we plot the time evolution of the populations of states \( |\Phi_{1}\rangle = |f_{11},g_{22},0\rangle \) (blue-dash line), \( |\Phi_{2}\rangle = |g_{11},f_{22},0\rangle \) (red-dash line), and \( |\Phi_{3}\rangle = \frac{1}{\sqrt{2}}(|f_{11},g_{22},0\rangle + |g_{11},f_{22},0\rangle) \) (magenta-dash line) under effective Hamiltonian in Eq. (14). Figure 3(b) shows that a perfect and fast quantum-information transfer from the initial state \( |\Phi_{1}\rangle \) to the target state \( |\Phi_{2}\rangle \) can be achieved after reselecting the optimal value of \( \zeta \). Notice that the population of excited state \( |\Phi_{3}\rangle \) is less than 0.25 during the interaction. Thus, we can draw a conclusion that the effective model can be considered as a three-level single-atom model, as the optimal value of \( \zeta \) for the whole system faultlessly satisfy the condition. In Fig. 3(c), we plot the time evolution of the population of states \( |\psi_{1}\rangle = |f_{11},g_{22},0\rangle, |01,01,02\rangle, |22,01,01\rangle \) (blue-solid line), \( |\psi_{2}\rangle = |g_{11},f_{22},0\rangle, |01,02,02\rangle, |01,02,02\rangle, |22,01,01\rangle \) (red-solid line), and...
\[ \psi = | \psi_{11} \rangle + | \psi_{22} \rangle \] under the total Hamiltonian in Eq. (1). Also, a perfect and fast QIT from the initial state \( | \psi_1 \rangle \) to the target state \( | \psi_8 \rangle \) can be achieved at time \( t_f \). Compared to the effective Hamiltonian model, the population of excited state governed by the total Hamiltonian is larger than that governed by the effective Hamiltonian. The reason for this can be explained as follows: during the operation, the intermediate states (i.e., \( | \phi_1 \rangle, | \phi_2 \rangle, | \phi_3 \rangle, | \phi_4 \rangle, | \phi_5 \rangle, | \phi_6 \rangle, | \phi_7 \rangle \), etc.) can be slightly populated, the whole system cannot be faultlessly considered as a three-level single-atom model, and the optimal value of \( \xi \) for the whole system will not faultlessly satisfy the condition \( \xi = \arcsin(0.25) \).

In order to get more insight into dynamic of the system governed by the total Hamiltonian, we plot the population of...
states $|\psi_2\rangle$ to $|\psi_7\rangle$ versus time in Fig. 3(d). From Fig. 3(d), we can see that all the populations of these states are smaller than 0.25, especially, the probabilities to find a photon in nodes 12 and 21 are less than 0.012. We can draw a conclusion that the system can be approximately considered as a three-level atom system, although the specific procedures has small differences between the two dynamics. Thus, the information can be fast and perfect transferred between arbitrary two distant nodes under current condition.

As shown in Fig. (3), the proposal can be nearly treated as an adiabatic process which is insensitive to the fluctuations of parameters, such as the amplitude of the laser pulses $\Omega_{jk}$, the coupling constant $g$ and the parameter $\xi$. Thus, we can choose a sets of parameters to obtain high fidelity and fast QIT. In Fig. 4, we plot the fidelity of the target state $|\psi_8\rangle$ versus the value of $\xi$ and the interaction time $g_{tf}$ governed by the total Hamiltonian $H_I$ when $v = 2.0g$. The fidelity for the target state is defined as $F = \langle \psi_8 | \rho(t) | \psi_8 \rangle$, where $\rho(t)$ is the density operator of the system at the time $t_f$ by solving the equation $\dot{\rho} = i[H_I, \rho]$. As seen from Fig. 4, when $g_{tf} = 10$, the optimal value of $\xi$ for the highest fidelity (fidelity $\geq 1$) of the state $|\psi_8\rangle$ is from 0.235 to 0.265. The reason for this can be expressed as: the proposal is an adiabatic passage, thus it is robust versus variations in the experimental parameters. However, when the parameters are no longer approximately satisfied by the condition $(\sin \xi)^{-1} = 4M (M = 1, 2, 3, \ldots)$, the fidelity will show an extreme fluctuation. Figure 4 also shows that it is hardly to get high fidelity when $g_{tf} \leq 10$. Thus, in the proposed scheme, the fastest time to get the target state is $t_f = 10/g$. Therefore, it is much faster than the general adiabatic process.

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Figure 4. The fidelity of the target state $|\psi_8\rangle$ versus the value of $\xi$ and the interaction time $g_{tf}$ governed by the total Hamiltonian $H_I$ when $v = 2.0g$. The fastest time to get the high fidelity of the target state is $t_f = 10/g$ when $\xi \simeq 0.25$ (from 0.235 to 0.265).

Figure 5. (a) The fidelity of the target state $|\psi_8\rangle$ versus the ratios $\gamma_1/g$ and $\kappa/g$. (b) The fidelity of the target state $|\psi_8\rangle$ versus the ratios $\gamma_2/g$ and $\kappa/g$. The others parameters are set as $\xi = 0.25$ and $g_{tf} = 10$. The fidelity is still about 79.8% (83.3%) when $\kappa = \gamma_1 = 0.1g$ in Fig. 5(a) ($\kappa = \gamma_2 = 0.1g$ in Fig. 5(b)).
Discussion

It is necessary to discuss the influence of decoherence caused by atomic spontaneous emission and cavity photon leakage of the system. In the current model, the master equation of the whole system can be expressed by the Lindblad form:

$$\dot{\rho} = -i[H, \rho] + \sum_{l=1}^{22} \frac{N_{kl}}{2} (2a_{k}^\dagger a_{l} \rho - a_{l}^\dagger a_{k} \rho - \rho a_{l}^\dagger a_{k} + \rho a_{k}^\dagger a_{l}) + \sum_{jk=1,12}^{2} \gamma_{jk}^{ef} (2\sigma_{jk}^\dagger \rho \sigma_{jk} - \sigma_{jk}^\dagger \rho - \rho \sigma_{jk}) + \sum_{jk=1,12}^{2} \gamma_{jk}^{ls} (2\sigma_{jk}^\dagger \rho \sigma_{jk} - \sigma_{jk}^\dagger \rho - \rho \sigma_{jk}),$$

where $N_{kl}$ denotes the decay rate of cavity, $\gamma_{jk}^{ef}$ and $\gamma_{jk}^{ls}$ represent the atomic decay from level $|e\rangle$ to $|g\rangle$ and $|f\rangle$ to $|l\rangle$, respectively. For simplicity, we assume $\kappa_{0} = \kappa$ ($\kappa = 11, 12, 21, 22$), $\gamma_{jk}^{ef} = \gamma$ and $\gamma_{jk}^{ls} = \gamma_{j}$ ($j = 11, 22$). The fidelity of the target state versus the ratios $\gamma$ and $\gamma_{j}$ ($\gamma_{j} = 0.1$ and $\Gamma_{j} = 10$) is shown in Fig. 5(a,b) when $\xi = 0.25$ and $\Gamma_{j} = 10$. As seen from Fig. 5, the fidelity decreases slowly with the increasing of cavity decay and atomic spontaneous emission. Figure 5 shows that the fidelity is still about 79.8% (83.3%) when $\gamma = 0.1$ g ($\gamma_{j} = 0.1$ g). Therefore, we can draw a conclusion that the proposal is robust against the spontaneous emission and cavity photon leakage.

Finally, let us give a brief analysis of the experimental feasibility for this scheme. The proposal can be realized in solid-state qubit trapped in a 2D array of superconducting cavity system. In this system, the superconducting cavity can be strongly coupled to the solid-state qubits such as Cooper pair boxes (CPB) and the corresponding microwave photons have small loss rates. As reported in ref. 23, the coupling strength in the interaction between CPBs and the circuit cavities is $g^2 \pi \times 50$ MHz, the corresponding photon lifetime is $T_{\text{c}} = 200 \times 10^{-6}$s, the dephasing time of the spin $T_{\text{\phi}} = 1 \times 10^{-6}$s. Thus, the required time for transferring the quantum-information, in principle, is $T = 3.2 \times 10^{-6}$s, which is much shorter than $T_{\text{c}}$ and $T_{\text{\phi}}$. The proposed idea can also be used for large-scale arrays cavities in photonic crystals, in which the achievable parameters are predicted to be $(g, \kappa, \gamma) = 2 \pi \times (2.5 \times 10^4, 0.4, 1.6)$ MHz. As shown above, the required time for achieving the task is smaller than photon coherence time and the atom dephasing time. In recent experiments, a set of cavity quantum electrodynamics parameters $(g, \kappa, \gamma) = 2 \pi \times (7.6, 2.8, 3.0)$ MHz is available in an optical cavity. Thus, based on the recent cavity QED technique or the technique to be improved soon, the proposal might be realizable in the future.

In conclusion, we have proposed a promising scheme to construct shortcuts to adiabatic passage to achieve controllable and fast quantum-information transfer between arbitrary two nodes in 2D quantum networks. The proposal has several advantages. The first one is that information can be controllably transferred between arbitrary two nodes, which is the basic of quantum computation. Secondly, the operation time is shorter than that in conventional adiabatic passage technique. Third, the proposed scheme is robust against the parameter fluctuations and the decoherence caused by atomic spontaneous emission and cavity photon leakage. These are very benefit to the suppression of decoherence effect. The scheme provides a new perspective on robust quantum information processing in 2D quantum networks. In principle, the proposal can be realized in solid-state qubit trapped in a 2D array of superconducting cavity system or in large-scale arrays cavities in photonic crystals. Moreover, the proposed scheme can be extended to 3D coupled cavity system.

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