Remote Entanglement via Adiabatic Passage Using a Tunably Dissipative Quantum Communication System


1Pritzker School of Molecular Engineering, University of Chicago, Chicago, Illinois 60637, USA
2Department of Physics, University of Chicago, Chicago, Illinois 60637, USA
3Argonne National Laboratory, Argonne, Illinois 60439, USA
4Department of Physics, University of California, Santa Barbara, California 93106, USA

(Received 2 April 2020; accepted 18 May 2020; published 17 June 2020)

Remote entanglement of superconducting qubits has recently been demonstrated using both microwave photon- and phonon-mediated communication [1–6]. Many of these demonstrations are limited by loss in the communication channel, due to loss in the various microwave components or intrinsic to the channel itself [1,4,6]; similar limitations apply to, e.g., optically based quantum communication systems. Adiabatic protocols analogous to stimulated Raman adiabatic passage [7,8] can mitigate such loss by adiabatically evolving an eigenstate of the system, using states that are “dark” with respect to the communication channel. These enable the high-fidelity coherent transfer of quantum states between sender and receiver nodes, even in the presence of large channel loss. Despite their use in a number of localized systems, such protocols have not been used for the generation of remote entangled states [7,8].

In this Letter, we present a unique experimental system comprising a pair of superconducting transmon-style qubits linked by an on-chip, 0.73 m-long superconducting microwave transmission line. By changing the coupling of the transmission line to a resistive load, we can vary the energy lifetime $T_{1r}$ of the transmission line over 2 orders of magnitude. We demonstrate an adiabatic protocol for quantum communication between the qubit nodes, compare its performance to a qubit-transmission mode-qubit relay method [5,9,10], and explore the performance of both protocols as a function of transmission loss.

First, we describe the experimental device, then the two-state transfer methods. We test the performance of each protocol in the low-loss limit, then as a function of transmission loss. The adiabatic process achieves significantly improved performance compared to the relay method, especially at intermediate levels of loss in the channel.

The two quantum state transfer methods, and the device we use to test them, are shown in Fig. 1. The device comprises two frequency-tunable superconducting Xmon qubits [11,12], $Q_1$ and $Q_2$, each coupled to one end of the on-chip transmission line via an electrically controlled tunable coupler [13], $G_1$ and $G_2$, respectively [Fig. 1(b)]. We use the qubit ground $|g\rangle$ and excited $|e\rangle$ states, whose transition frequency is tunable from $\sim$3 to 6 GHz. Qubit control is via low-frequency flux tuning for Z control and quadrature-resolved microwave pulses for XY control. We read out the qubit states using standard dispersive measurements [14–16],
via a capacitively coupled readout resonator and a traveling-wave parametric amplifier. We projectively measure the excited state probability \( P_e \) of each qubit with a fidelity of 88.8 \( \pm \) 0.8%.

The tunable couplers \( G_1 \) and \( G_2 \) allow us to externally control the coupling \( g_{l,2} \) of each qubit to the individual resonant modes in the transmission line. A variable control consisting of two additional tunable couplers, \( D_1 \) and \( D_2 \), is integrated into the transmission line, 1.6 mm from the coupler \( G_1 \) and its associated qubit \( Q_1 \). This circuit element provides electrically controlled coupling between its input port and two output ports [36]. The coupler \( D_2 \) is placed inline with the transmission line and is always set to provide maximum coupling (and minimal reflection) to the remaining length of the transmission line. The other coupler \( D_1 \) connects to port 1 on the sample mount, which is terminated by a lumped 50 \( \Omega \) microwave load outside the sample box. Varying the coupling to this load allows us to set the loss in the transmission line, quantified by the energy lifetime \( T_{1\ell} \) of each resonant mode.

The transmission line of length \( \ell = 0.73 \) m supports multiple resonant modes, separated in frequency by the free spectral range (FSR) \( \omega_{\text{FSR}}/2\pi = 1/2T_e = 84 \) MHz, where \( T_e = 5.9 \) ns is the photon one-way transit time in the channel. For a sufficiently small qubit-resonator coupling, \( g_{l,2} \ll \omega_{\text{FSR}} \), each qubit can be selectively coupled to a single resonant mode in the transmission line. This is shown in Fig. 2(a), where the transition frequency \( \omega_{\text{ref}}/2\pi \) of qubit \( Q_1 \) is tuned over 400 MHz, yielding four separate vacuum Rabi swap resonances spaced by the free spectral range \( \omega_{\text{FSR}}/2\pi \). The loss coupler \( D_1 \) was set to minimum coupling, so the transmission line is limited only by its intrinsic loss. All experiments, here, were done with the mode at 5.351 GHz, just to the right of center in Fig. 2(a).

In Fig. 2(b), we demonstrate tunable control over the channel loss, using qubit \( Q_1 \) to measure the lifetime of the resonant mode at 5.531 GHz as we vary the coupler \( D_1 \) and, thus, the transmission line loss. The pulse sequence for this measurement is shown in the inset in Fig. 2(b). The mode energy decay time \( T_{1\ell} \) for each loss setting (controlled by the \( D_1 \) flux) is shown in Fig. 2(b). With no coupling through \( D_1 \), we measure the intrinsic resonant mode lifetime \( T_{1\ell} \approx 3410 \pm 40 \) ns (orange), comparable to similar transmission lines without variable loss [5].
With maximum coupling to the load, we measure a lifetime $T_{1,\rho} \approx 28.7 \pm 0.2$ ns (blue), corresponding to a loaded quality factor $Q_L = 960$, about 120 times smaller than the intrinsic quality factor of $1.1 \times 10^5$. We also measure the resonant mode’s Ramsey dephasing time $T_{2,\rho}$ at various $D_1$ flux bias points and find $T_{2,\rho} \approx 2T_{1,\rho}$, indicating the coupler $D_1$ introduces negligible additional phase decoherence. One nonideality with this system is that qubit $Q_1$, due to its close proximity to the loss coupler $D_1$, also has its lifetime reduced when the couplers $G_1$ and $D_1$ are both set to nonzero coupling, allowing energy loss from $Q_1$ to the external load; this limits the performance of $Q_1$ and is discussed further in the Supplemental Material [17–36]. This additional loss pathway could be reduced by placing the loss coupler $D_1$ in the center of the transmission line, as the transmission line would then protect both qubits from the external load.

We used two different communication protocols, adiabatic transfer and a qubit-resonant mode-qubit relay method. Both methods were used for qubit state transfer via the transmission line as well as Bell state generation, both as a function of loss in the communication channel. The relay method uses a single extended mode in the transmission line as well as Bell state generation, via the transmission line as well as Bell state generation, measured simultaneously at time $t$. Left inset: Control pulse sequence. The couplers are set so that coupling $g_2$ starts at its maximum with $g_1$ set to zero. Dissipation in the resonant channel mode is controlled using $D_1$, here, set to zero coupling. Right inset: Quantum process tomography, yielding a process fidelity $F_p = 0.96 \pm 0.01$. (b) Adiabatic remote entanglement. Right inset shows control pulse sequence: With $Q_1$ initially prepared in $|e\rangle$, $G_1$ and $G_2$ are controlled using the adiabatic protocol to share half of $Q_1$’s excitation with $Q_2$, resulting in a Bell singlet state $|\psi^\mp\rangle = (|eg\rangle - |ge\rangle)/\sqrt{2}$. Blue (orange) circles represent excited state populations of $Q_1$ ($Q_2$) measured simultaneously at time $t$. Left inset: Reconstructed density matrix of the final Bell state, yielding a state fidelity $F_s = 0.964 \pm 0.007$ and concurrence $C = 0.95 \pm 0.01$. In all panels, dashed lines are from master equation simulations accounting for channel dissipation and qubit imperfections (see [17]).

Here, we implement a simple adiabatic scheme [37,38], where we vary the couplings in time according to $g_1(t) = \tilde{g}\sin(\pi t/2t_f)$ and $g_2(t) = \tilde{g}\cos(\pi t/2t_f)$. We choose the parameters $\tilde{g}/2\pi = 15$ MHz and $t_f = 132$ ns, minimizing the impact of finite qubit coherence while maintaining sufficient adiabaticity (see [17]). We note that the adiabatic protocol supports better than 90% transfer efficiency even when $\tilde{g} = 0.4\omega_{\text{res}}$; see [17].

In Fig. 3(a), we demonstrate deterministic adiabatic state transfer from $Q_1$ to $Q_2$. With $Q_1$ in $|e\rangle$ and $Q_1$ and $Q_2$ set on resonance with a single mode in the channel, we adjust the couplers $G_1$ and $G_2$ adiabatically to complete the state transfer. We show the excited state population of each qubit as a function of time $t$, measured with the resonant mode loss at its intrinsic minimum. We observe the expected
gradual population transfer from $Q_1$ to $Q_2$, with the population of $Q_2$ reaching its maximum at $t = t_f$, with a transfer efficiency $\eta = P_{sQ_2}(t = t_f)/P_{eQ_2}(t = 0) = 0.99 \pm 0.01$. We further characterize the state transfer by carrying out quantum process tomography [39], yielding the process matrix $\chi$ shown in the inset in Fig. 3(a), with a process fidelity $F_p = 0.96 \pm 0.01$, limited by qubit decoherence. The process matrix calculated from a master equation simulation displays a small trace distance to the measured $\chi$ matrix of $D = \sqrt{\text{Tr}[(\chi - \chi_{\text{sim}})^2]} = 0.02 \pm 0.01$, indicating excellent agreement with experiment.

The adiabatic protocol can also be used to generate remote entanglement between $Q_1$ and $Q_2$. With $Q_1$ prepared in $|\psi\rangle$, we share half its excitation with $Q_2$ using the adiabatic protocol, by stopping the transfer at its midpoint $t = t_f/2$. This generates a Bell singlet state $|\psi^-\rangle = (|eg\rangle - |ge\rangle)/\sqrt{2}$. The qubit excited state population is shown as a function of time $t$ in Fig. 3(b). We further characterize the Bell state by quantum state tomography [40,41], and the reconstructed density matrix $\rho$ is shown in the inset in Fig. 3(b). We find a Bell state fidelity $F_s = \langle|\psi^-\rangle|\rho|\psi^-\rangle\rangle = 0.964 \pm 0.007$, referenced to the ideal Bell singlet state $\psi^-$, and a concurrence $C = 0.95 \pm 0.01$ (see [17]). The density matrix $\rho_{\text{sim}}$ calculated from a master equation simulation shows a small trace distance to the measured $\rho$, $\sqrt{\text{Tr}[(\rho - \rho_{\text{sim}})^2]} = 0.01$, indicating excellent agreement with experiment.

We explore the impact of loss on both the relay method and the adiabatic protocol, with results shown as a function of the resonant channel mode energy lifetime $T_1$, in Fig 4. For the highest level of dissipation, with $T_1 = 28.7$ ns, we measure an adiabatic transfer efficiency $\eta = 0.67 \pm 0.01$, even though the transfer time $t_f$ is four times the resonant mode lifetime. The efficiency is primarily limited by loss in qubit $Q_1$ due to its spurious coupling loss through $D_1$ to the 50 $\Omega$ load (see [17]), in good agreement with master equation simulations. Results from a simulation without the spurious coupling are plotted as black dashed lines in Fig 4(a), limited by a small channel occupation due to the finite adiabaticity of the sequence. We compare these results to the relay method, where we use a weak coupling $|g_{1,2}|/2\pi = 5.0$ MHz to ensure that the qubits only couple to a single transmission line mode; this results in a total transfer time $2T_{\text{swap}} = 100$ ns. We find the adiabatic protocol consistently performs better than the relay method, with a $2.6 \times$ higher transfer efficiency $\eta$ ($2.3 \times$ reduction in transfer loss) and $1.5 \times$ higher process fidelity $F_p$ ($2.3 \times$ reduction in process infidelity) compared to the relay method in the most dissipative case; the adiabatic protocol is primarily limited by spurious coupling loss in $Q_1$, while the relay method is limited by loss in the channel (see [17]).

In Fig. 4(b), we display the entanglement fidelity using the adiabatic protocol with different levels of channel loss, and compare to the relay method. The adiabatic protocol outperforms the relay method in all levels of dissipation. At the highest loss level, where $T_1 = 28.7$ ns, the adiabatic protocol achieves $1.2 \times$ higher Bell state fidelity $F_s$, $1.5 \times$ reduction in Bell state infidelity) and $1.3 \times$ higher concurrence $C$ ($1.7 \times$ reduction in concurrence infidelity) compared to the relay method; the spurious-coupling-free simulation result for the adiabatic protocol is shown by the black dashed lines, limited by a small channel occupation due to the finite adiabaticity of the sequence.

In conclusion, we describe a unique experimental system in which we can explore the performance of quantum communication protocols in the presence of controllable communication loss. We demonstrate an adiabatic protocol that realizes high-fidelity transfer of quantum states and entangled Bell states, limited mostly by spurious coupling of one qubit to the controlled transmission line loss. The platform we have developed is well-suited for exploring the impact of channel loss on other error-protecting quantum
communication protocols, such as heralding [42–44] and entanglement distillation [45–47]. The ability to introduce controlled loss dynamically into the system opens the door for studying dissipative dynamics in nonequilibrium systems, enabling approaches such as reservoir engineering [48,49]. The adiabatic protocol demonstrated here is applicable to other quantum communication systems, for example, phonon-based systems where the communication channel is significantly more lossy [6,50,51]. Future demonstrations could employ more advanced adiabatic protocols such as shortcuts to adiabaticity [52,53] and composite adiabatic passage [54,55] to further improve fidelity.

The authors declare no competing financial interests.

The authors thank A. A. Clerk, P. J. Duda, and B. B. Zhou for helpful discussions. We thank W. D. Oliver and G. Calusine at MIT Lincoln Lab for providing the traveling-wave parametric amplifier (TWPA) used in this work. Devices and experiments were supported by the Air Force Office of Scientific Research and the Army Research Laboratory. K. J. S. was supported by NSF GRFP (NSF Grant No. DGE-1144085), E. D. was supported by LDRD funds from Argonne National Laboratory; A. N. C. was supported in part by the DOE, Office of Basic Energy Sciences. This work was partially supported by the DOE, Office of Basic Energy Sciences. This work was partially supported by the UChicago MRSEC (NSF Grant No. DMR-1420709) and made use of the Pritzker Nanofabrication Facility, which receives support from SHyNE, a node of the National Nanotechnology Coordinated Infrastructure (NSF Grant No. NNCI-1542205). Correspondence and requests for materials should be addressed to A. N. Cleland (anc@uchicago.edu).

†Present address: Université Claude Bernard, CNRS, Laboratoire de Physique, F-69342 Lyon, France.
‡ Present address: Google, Santa Barbara, California 93117, USA.
§Present address: Université de Lyon, ENS de Lyon, Université Claude Bernard, CNRS, Laboratoire de Physique, F-69342 Lyon, France.
*Corresponding author.
anc@uchicago.edu


