Quantum interface between single- and dual-rail optical qubits

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Introduction. With the development of quantum technology, it is becoming clear that different physical systems are optimal for various aspects of quantum information processing. For example, superconducting circuits and trapped ions are well-suited for the implementation of quantum computation gates; spin ensembles for quantum memories; natural and artificial atoms for precise sensing. A comprehensive quantum network should enable reliable exchange of quantum information among all these systems. The primary agent of such exchange is light, as it is the only physical system able to carry quantum information over large distances. Technologies of quantum coupling between light and stationary carriers of quantum information, such as superconducting cavity modes [1], spin ensembles in solids and atomic gases [2], optomechanical devices [3] and others are being developed.

However, these technologies must address an important challenge before their broadband deployment becomes possible. In many promising quantum settings — for example, excitons in quantum dots, single atoms and superconducting qubits — the natural basis for encoding the qubit consists of two energy eigenstates. When coupled to light, such a qubit is naturally converted into the so-called single-rail qubit: the encoding in which the vacuum |0⟩ and single-photon |1⟩ states correspond to the logical one and zero, respectively. Single-rail qubits are however notoriously inconvenient when it comes to quantum information processing and communication by means of light. This is because single-qubit operations are difficult in this encoding [4]. Additionally, the information carried by a single-rail qubit can be easily distorted by optical loss, which results in |1⟩ becoming |0⟩.

A much more practical way of encoding the optical qubit is dual-rail, in which the difference between the logical states consists in the photon occupying one of the two orthogonal optical modes. In this way, the photon is present in any valid state of the qubit, thereby providing an easy way to label loss events. In many applications, the two qubit modes are the orthogonal polarisation directions, horizontal |H⟩ and vertical |V⟩. This allows easy realisation of single-qubit operations by means of polarisation rotators.

In principle, the dual-rail light qubit can be treated as a pair of single-rail qubits carried by each polarisation mode, as we can write |H⟩ = |1⟩H|0⟩V and |V⟩ = |0⟩H|1⟩V. Each mode can be directly coupled to a stationary system, as demonstrated in e.g. Ref. [5] and references therein. However, this approach would double the number of coupling units required, thus drastically increasing the quantum network’s complexity. Moreover, it would still require a method for single-to-dual rail interconversion within the stationary system, which would need to be specific for each such system. It therefore appears more practical to develop a method for such interconversion only for light.

To date, there existed methods for preparing entangled states that connect single- and dual-rail qubits with “continuous-variable” qubits carried by coherent states of opposite phases [6–8]. In principle, one could use these resources to convert between single- and dual-rail encodings through an intermediate step of continuous-variable encoding. However, this approach is quite cumbersome and prone to error. It would be much more desirable to develop a direct method for such interconversion.

This is the goal of our paper. We propose and imple-
ment a technique to prepare an entangled resource of the form

$$|\Omega\rangle = a|H\rangle|1\rangle + b|V\rangle|0\rangle.$$  

We show that this resource can be used for the interconversion between the two bases via quantum teleportation [9] from a qubit carried by a photon’s polarisation onto the single-rail encoding. Specifically, we prepared all 6 primary basis states of a dual-rail discrete variable qubit $a|H\rangle + b|V\rangle$ and teleported them onto their single-rail counterparts $a|0\rangle + b|1\rangle$. In this aspect, our experiment achieves the goal pursued in theoretical proposals [10, 11], albeit with a different method which is more general, more experimentally accessible and less vulnerable to inefficiencies.

**Theory.** To produce the entangled resource state, we begin from non-degenerate spontaneous parametric down conversion (SPDC) in crystal 1 as shown in Fig. 1. This produces the state $|\psi\rangle_{\text{cr1,CB}} = |0_H\rangle_C|0_V\rangle_B + \gamma_1|1_H\rangle_C|1_V\rangle_B + O(\gamma_1^2)$ in the vertical and horizontal polarisations of modes B and C, respectively, with $\gamma_1$ being the SPDC amplitude. We then use a polarising beam splitter (PBS) to inject a weak coherent state $|\alpha\rangle = |0\rangle + \alpha|1\rangle$ into the vertical polarisation of mode C. The collective state describing these two modes can be written as the tensor product

$$|\Omega\rangle_{CB} \equiv |\alpha\rangle_{V,C} \otimes |\psi\rangle_{\text{cr1,CB}} \approx |0\rangle_C|0\rangle_B$$

$$+ \gamma_1|1_H0_V\rangle_C|1_V\rangle_B + \alpha|0_H1_V\rangle_C|0\rangle_B$$

$$= |0\rangle_C|0\rangle_B + \gamma_1|H\rangle_C|1\rangle_B + \alpha|V\rangle_C|0\rangle_B,$$

up to the first order, where we chose the dual- and single-rail notations in modes C and B, respectively, in the last line.

Although this state is separable, entanglement of the form (1) is present in its last two terms corresponding to a photon present in mode C. Intuitively, this reflects the ambiguous situation in which this photon could originate from either crystal 1 or the weak coherent state $|\alpha\rangle$. If such photon came from crystal 1 ($|\alpha\rangle$), then Bob receives a single photon $|1\rangle$ (vacuum $|0\rangle$) in the vertical polarisation of mode B.

Theoretically, this entanglement could be recovered by means of a non-demolition (for example parity [12]), polarisation-insensitive detector in mode C that would project that mode onto the single-photon state. Regardless, even without such a detector, the state $|\Omega\rangle_{CB}$ can be used post-selectively to perform the teleportation.

For the teleportation to occur, Alice starts by

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**Figure 1.** Experimental setup. To prepare the single-dual rail entangled resource (2) in spatial modes C and B, Bob combines a photon from a pair generated in crystal 1 with a coherent state $|\alpha\rangle$ on a PBS. Alice’s source state $|\chi\rangle_A$ is obtained by applying a polarisation projection in mode D to a Bell pair $|\Psi\rangle_{AD}$ produced by crystals 2 and 3 (represented in the figure as a single crystal). HWP: half-wave plates (blue); QWP: quarter-wave plates (dark blue); PBS: polarising beam splitters (grey cubes); SPCM: single photon counting module (black boxes); BHD: balanced homodyne detector. The red (blue) lines correspond to horizontal (vertical) polarisation.
producing an entangled state $|\Psi\rangle_{AD} = |0\rangle_A |0\rangle_D + \gamma_{2,3} (|H\rangle_A |V\rangle_D + |V\rangle_A |H\rangle_D)$ between spatial modes A and D. This is achieved by means of SPDC crystals 2 and 3 (see Fig. 1 and Methods). Then, a polarization projection of $|\Psi\rangle_{AD}$ performed in mode D by one of the two single photon counting modules (SPCMs) grants Alice the heralded preparation of a dual-rail qubit of the form $|\chi\rangle_A = a |H\rangle_A + b |V\rangle_A$ in mode A which serves as the input state for quantum teleportation. Alice sends this state into a Bell state projector $\langle \Psi^+ |_{AC} = \frac{1}{\sqrt{2}} (|H\rangle_A |V\rangle_C + |V\rangle_A |H\rangle_C)$, combining it with the part of the resource state $|\Omega\rangle_{CB}$ in mode C. Upon a successful application of the Bell projector — characterised by a coincidence click from the two SPCMs in Fig. 1 — Bob finally obtains the following single-rail teleported state in his mode $B$

$$|\varphi\rangle_B = \langle \Psi^+ |_{AC} |\chi\rangle_A |\Omega\rangle_{CB} = \frac{1}{\sqrt{2}} (a |0\rangle_B + b \gamma_1 |1\rangle_B).$$

For faithful teleportation, we match the coherent state’s amplitude to that of SPDC from crystal 1 so that $\alpha = \gamma_1$. The same scheme, but without a measurement in mode D, can produce a freely-propagating single-dual entangled resource (1). This happens thanks to entanglement swapping when $|\Omega\rangle_{CB} |\Psi\rangle_{AD}$ is projected onto the Bell state $|\Psi^+\rangle_{AC}$ in modes A and C. This, however, would require a source of Bell states in modes A and D operating in a heralded [13, 14] or deterministic [15] fashion. In the absence of such a source, we can still show the viability of this procedure by reconstructing the resulting state in modes B and D post-selectively [8] and assessing its a posteriori entanglement.

**Experiment.** The technical details of our laser setup and parametric down-conversion are described in Methods. We used a picosecond Ti:Sapphire laser at 780 nm, with the pulse repetition rate of $R_L = 76$ MHz, and parametric down-conversion in periodically poled potassium titanyl phosphate crystals (PPKTP), with preliminary frequency doubling. The down-conversion amplitudes were set, for the reason explained below, to unequal values of $\gamma_1 \approx 0.20$ and $\gamma_{2,3} \approx 0.054$. The Bell state projector was implemented following the scheme of Refs. [16, 17], in which the two incoming spatial modes are overlapped on a PBS and subsequently subjected to polarisation analysis in the diagonal basis (Fig. 1). Teleportation events are heralded by a triple coincidence of the two SPCMs in Fig. 1 — Bob obtains the following single-rail teleported state in his mode $B$

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To reconstruct the teleported states in Bob’s channel, we employed a balanced homodyne detector (BHD) [18, 19]. A total of 2000 quadratures — along with their associated phases — were acquired for each teleported state. Quantum state tomography was then performed using the maximum-likelihood algorithm [20] with an efficiency correction of $\eta = 50\%$ (see Methods). The te
portation has been implemented for six input states: $|V\rangle$, $|H\rangle$, $\frac{1}{\sqrt{2}}(|H\rangle \pm |V\rangle)$ and $\frac{1}{\sqrt{2}}(|H\rangle \pm i|V\rangle)$. The results of the experiment are shown in Fig. 2. The output states show an average fidelity of $(92.8 \pm 2.2)\%$ with respect to the expected pure states in Eq. (3). Note that, for reasons of aesthetics, we only show the density matrices up to 2 photons in Fig. 2, while in fact, the states used to compute the fidelity have been reconstructed up to 4 photons. For the diagonal and circular inputs, the orthogonality between the resulting qubits, $|\pm\rangle_B = \frac{1}{\sqrt{2}}(|0\rangle_B \pm |1\rangle_B)$ and $|\pm i\rangle_B = \frac{1}{\sqrt{2}}(|0\rangle_B \pm i|1\rangle_B)$ manifests itself in the Wigner functions oriented in opposite directions.

The main source of error explaining the deviation of the experimentally observed teleported states from those predicted by Eq. (3) comes from the false positive Bell state projections [21, 22]: a coincidence click can be caused by both photons arriving from the same input mode, A or C, rather than one from each mode. The probability of such events for the two photons coming from mode A scales as $p_{bad,A} \sim 2\gamma_2^1|\gamma_2,3|^2$, and for the photons from mode C as $p_{bad,C} \sim 2\gamma_2^2|\gamma_1|^2|\gamma_2,3|^2$. These values are different because one of the photons arriving from mode A is always heralded by a click in mode D. To minimise the ratio of these events with respect to the “good” ones while keeping the latter at a reasonable rate, we set $|\gamma_2,3| \sim |\gamma_1|^2$. This results in the false positives contributing $\sim 16\%$ to all triple coincidences. Most of these events lead to the admixture of the vacuum state to Bob’s output. Because the theoretically expected output state is expected to contain, on average, one-half of the vacuum state, the effect of this admixture on the fidelity is $\sim 8\%$, consistent with our observations.

If the efficiency correction in the maximum likelihood algorithm is removed ($\eta = 1$), the reconstructed states still yield an average fidelity of $(77.7 \pm 1.5)\%$ with respect to the expected pure states in Eq. (3). Overall, these obtained fidelities are state-of-the-art compared to those achieved for various traditional teleportation experiments [23], despite a comparatively high complexity of our experiment.

The teleportation outputs for the six inputs allow us to post-selectively reconstruct the bipartite state that is produced in modes B and D due to entanglement swapping after the Bell measurement in modes A and C [8]. Comparing the result displayed in Fig. 3 with the expected maximally entangled state of Eq. (1), we find a fidelity of $(89.7 \pm 2.1)\% [(66.8 \pm 1.3)\%] with [without] applying efficiency correction to the maximum likelihood algorithm, thus certifying the a posteriori entanglement of this state since the classical limit of $1/2$ [24] is beaten.

**Conclusion.** In summary, we have proposed and experimentally demonstrated a scheme to interconvert between the dual- and single-rail encodings of an optical qubit with competitive fidelities. This scheme enables efficient exchange of quantum information between stationary carriers of different nature by means of light. The quality of the results could be improved by having number-discriminating photon detectors in order to efficiently discriminate and remove false positive Bell events. Additionally, heralded or deterministic sources of entangled photons would make the single-dual rail entangled resource readily available through entanglement swapping.

There are three primary ways to encode the qubit in the optical field: single-rail, dual-rail and continuous-variable. While previous research [6–8] established techniques to connect the two discrete-variable encodings with the continuous-variable one, the present work completes the triad to enable interconversion among all three encodings.

**Methods.**

**Experimental Setup.** The master laser is a pulsed Ti:Sapphire (Coherent Mira 900D) with a wavelength of 780 nm, mean output power of 1.3 W, repetition rate of $R_L = 76$ MHz and pulse width of 1.6 ps. We perform frequency doubling in a lithium triborate crystal with an efficiency of $\sim 30\%$. After further mode filtering, 50 mW and 5 mW of the frequency doubled wave pump crystal 1 and both crystals 2 and 3, respectively. These are periodically poled potassium titanyl phosphate crystals (PPKTP) operating in a type II spectrally and spatially de-
generate, but polarisation non-degenerate configuration, generating two-mode squeezed vacuum states of the form 
\(|0_H⟩|0_V⟩ + \gamma |1_H⟩|1_V⟩ + O(\gamma^2).\)

To generate Alice’s entangled state \(|\Psi⟩_{AD}\), we used the Mach-Zehnder interferometer architecture introduced in [23], where the outputs of two SPDC crystals interfere on a PBS. Due to different optical path lengths in the Mach-Zehnder interferometer, photons from crystals 2 and 3 experience a phase difference \(\Delta\phi_{2,3}\) which is introduced in \(|\Psi⟩_{AD}\). We use a piezoelectric transducer along with an interferometric feedback loop to ensure \(\Delta\phi_{2,3} = 0\) at all times. The quality of the entangled state \(|\Psi⟩_{AD}\) has been tested by applying a polarisation projection in mode D and measuring the probability of a coincidence click between a pair of SPCM in modes A and D as a function of the angle of a half-wave-plate (HWP) set in mode A. The resulting probabilities depend sinusoidally on that angle, with measured visibilities of 99%, 97% and 95% for projections of \(|\Psi⟩_{AD}\) onto the canonical, diagonal and circular polarisation bases in mode D, respectively.

**Bell state projector.** The central element in the projection apparatus is a PBS, which will direct the photons of different polarisations into the same spatial mode, thereby preventing coincidence clicks from the input states \(|\Phi^+⟩_{AC}\). The photons in the two PBS output modes are filtered by polarisers set at \(±\pi\) prior to detection, so that the state \(|\Phi^+⟩_{AC}\) is not transmitted. Therefore, a coincidence click projects the PBS input onto \(|Φ−⟩_{AC}\) [16, 17, 26]. However, because both input modes of the Bell projector are rotated, using HWPs, by \(\pi\) before entering the PBS, the state to which the Bell projector responds is \(|\Phi^+⟩_{AC}\). Prior to the teleportation experiment, we tested the projector’s alignment by sending a single-photon Fock state and the weak coherent state with identical polarisation into modes A and C, respectively, and observed a Hong-Ou-Mandel dip with 98% visibility.

Our implementation of the projector is advantageous with respect to the version that makes use of a sole 50:50 beam splitter [27]. First, it reduces the rate of false positive Bell detection events due to two photons present in the same input mode, being oblivious to those of them in which these two photons have orthogonal polarisations. Second, observing polarisation projections of \(|Ω⟩_{CB}\) in mode C allows us to monitor in real-time the phase in channel B necessary for quantum state reconstruction of Bob’s teleported state (see below). Third, it obviates the need to align the regular beam splitter to an exact 50% reflectivity.

**Efficiency Estimation.** We measure the quantum efficiency \(\eta\) from the quantum state tomography of a single Fock state [28]. The deviation of \(\eta\) from unity arises from three main factors: losses, mode matching between the signal and local oscillator (LO) and the quantum efficiency of the homodyne detector’s photodiodes at the laser’s wavelength. Their estimated values are 80%, 81% and 86%, respectively. Additionally, during the data acquisition runs, the high pump power impinging onto crystal 1 systematically caused the SPDC mode’s quality to degrade over time. Averaging over the duration of an acquisition batch, we have \(\eta = 50\%\).

**Phase Reconstruction.** Precise determination of the teleported state’s phase with respect to the local oscillator is essential for its reconstruction via homodyne tomography. Taking into account the difference between the phase \(\phi_\alpha\) of the state \(|\alpha⟩\) and the phase \(\phi_{\gamma_1}\) of the photons output by crystal 1, Bob’s teleported state \(3\) becomes

\[
|\varphi⟩_B = \frac{1}{\sqrt{2}} (a|0⟩_B + b\gamma_1 e^{-i(\phi_{\gamma_1} - \phi_\alpha)}|1⟩_B) .
\]

To evaluate this phase, we use one arm of the Bell projector to perform a polarisation projection of the resource state \(|Ω⟩_{CB}\) in mode C onto the diagonal basis. When a single click is recorded, we get the state

\[
|δ⟩_B = \left( \frac{|H⟩_C + |V⟩_C}{\sqrt{2}} \right) |Ω⟩_{CB} = \frac{1}{\sqrt{2}} \left( a|0⟩_B + \gamma_1 e^{-i(\phi_{\gamma_1} - \phi_\alpha)}|1⟩_B \right),
\]

whose phase equals that of the teleported state (4).

The rate of these single clicks is five orders of magnitude higher than the successful triple coincidence clicks required for teleportation. Therefore, by continuously monitoring the mean value of the quadrature observable \(X\) in Bob’s channel conditioned on these frequent single clicks, we can fit in real time the phase of \(|δ⟩_B\).

**Photon Count Rates.** Photons are measured by SPCMs from Excelitas. The typical count rates measured in both arms of the Bell projector were \(R_\alpha = R_{\gamma_1} = 22\, kHz\) for \(|\alpha⟩\) and crystal 1. For crystals 2 and 3, these rates were \(R_{\gamma_2,3} = 1.7\, kHz\). These numbers reflect the loss factor of 4 introduced by the pairs of HWPs and polarisers in the Bell projector. The coincidence rate of the photons originating from crystals 2 and 3, measured between one of the polarisation analysis SPCMs in mode D and one of the Bell projector SPCMs, was, on average, \(R_{\gamma_2,3}^{\gamma_2,3} = 51 Hz\). From these rates, we estimated a total photon detection efficiency of \(\eta_d = \frac{R_{\gamma_2,3}^{\gamma_2,3}}{R_{\gamma_2,3}} \approx 3.0\%\). This value is rather small due to the presence of narrowband (0.2 nm) filters in front of each SPCM, as well as the unavoidable optical losses present in the experiment. From here, we estimate \(\gamma_1 = \left( \frac{R_\alpha}{R_{\gamma_1}} \frac{1}{\eta_d} \right)^{1/2} \approx 0.20\) and \(\gamma_2,3 = \left( \frac{R_{\gamma_2,3}}{R_{\gamma_2,3}^{\gamma_2,3}} \frac{1}{\eta_d} \right)^{1/2} \approx 0.054\). Finally, we find for the “good” triple coincidence rate \(R_{T,good} \approx \frac{1}{2} \frac{1}{\eta_d} |\gamma_1|^2 |\gamma_2,3|^2 \approx 0.12 Hz\). The corresponding experimentally observed rate was \(R_T = (0.16 \pm 0.03) Hz\) as mentioned in the main text.
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