A controllable single photon beam-splitter as a node of a quantum network

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Abstract

A model for a controlled single-photon beam-splitter is proposed and analyzed. It consists of two crossed optical-cavities with overlapping waists, dynamically coupled to a single flying atom. The system is shown to route a single photon with near-unity efficiency in an effective ‘weak-coupling’ regime. Furthermore, two such nodes, forming a segment of a quantum network, are shown to perform several controlled quantum operations. All one-qubit operations involve a transfer of a photon from one cavity to another in a single node, while two-qubit operations involve transfer from one node to a next one, coupled via an optical fiber. Novel timing protocols for classical optical fields are found to simplify possible experimental realizations along with achievable effective parameter regime. Though our analysis here is based on a cavity-QED scenario, basic features of the model can be extended to various other physical systems including gated quantum dots, circuit-QED or opto-mechanical elements.

(Some figures may appear in colour only in the online journal)

1. Introduction

Exploring the potential of quantum systems as novel and powerful resources for information processing and communication has been a major focus of research over past three decades \cite{1–5}. Quantum information operations are usually envisioned on a network of nodes \cite{6–9} consisting of material media. The nodes in turn exchange information via a network of photonic bus transmitting single photons. Accordingly, a variety of physical systems have been explored as possible candidates towards realizing these effective nodes. These include systems like ion-chains \cite{10, 11}, photonic qubits \cite{12–15}, Josephson junction qubits \cite{16, 17}, spin and charge qubits in quantum dots \cite{18, 19} and trapped, cooled atoms in electromagnetic cavities \cite{20–22}. While single photons connect these nodes robustly, the nodes perform physical operations including generation, storage \cite{24} and routing of photons \cite{25, 26, 28, 29, 31}, entangling remote nodes \cite{9, 23} or performing basic gate operations \cite{1, 2}.

A particular focus has been to design routers for single-photons. An essential component of a quantum network, such routers should be programmable, directing a single photon in multiple directions, with minimal loses and no added photons. Much progress has been made on a one-dimensional version of such a router, including proposals for opto-mechanically routing photons in forward or backward directions \cite{25}, experimental demonstration of micro-fabricated optical cavities with single atoms efficiently coupling input-output modes of a one-dimensional tapered fiber \cite{26} or redirecting single and bi-photons in separate output modes \cite{27}. Simultaneously, exploration of schemes for directing photons in orthogonal directions to realize a higher dimensional networks is also under way. Models of coupled resonator waveguides, routing photons in orthogonal directions \cite{28} and experimental demonstration of superconducting transmon qubits redirecting microwave photons in separate outputs using quantum interference \cite{29} are two such examples.

Here we propose a model for programmable routing of a single photon in orthogonal directions, effectively acting as a controlled single-photon beam splitter. The system can act as a versatile node in a quantum network. When operated as a...
beam splitter, the splitting amplitudes can be fully controlled with a single, experimentally accessible parameter. We show that two such beam-splitters can be combined to form a node of a quantum network, with the ability to perform elementary gate operations with one and two qubits.

The model consists of a basic unit of two crossed optical cavities with overlapping cavity waists along with an atom transiting through the waist. We show that the system can transfer a single photon from one cavity mode to the mode of the other cavity. The transfer is mediated via an adiabatic rotation of a dark state, which is triggered by either passage of a single atom through the cavity waist along with timed, optical pulses or through a sequence of pulses along with a static atom, localized at the waist.

The basic unit can be repeated with optical fibers as interconnects to construct a network. An analysis for two such nodes is presented. For a range of parameters, it is found that the presence of a photon in a node allows arbitrary unitary operations on the corresponding local qubit without affecting the other node. Furthermore, there are operations through which a photon can be adiabatically transferred from one node to the other, enabling two-qubit operations such as generation of entangled atomic states. Using inter-node transfer of photon in conjunction with single node processes, one can therefore construct protocols for arbitrary one and two qubit operations by choosing appropriate classical field timing, pulse shape and peak power. We present examples of such protocols for few representative cases. These protocols are also in contrast with our earlier work on quantum network based on nodes made of single cavity coupled to a three-level static atom, localized at the waist.

We show that the system can transfer a single photon in conjunction with single node processes, the other node. Furthermore, there are operations through which a photon in a node allows arbitrary unitary operations on the corresponding local qubit without affecting the presence of a photon in a node allows arbitrary unitary operations on the corresponding local qubit without affecting the other node. Furthermore, there are operations through which a photon can be adiabatically transferred from one node to the other, enabling two-qubit operations such as generation of entangled atomic states. Using inter-node transfer of photon in conjunction with single node processes, one can therefore construct protocols for arbitrary one and two qubit operations by choosing appropriate classical field timing, pulse shape and peak power. We present examples of such protocols for few representative cases. These protocols are also in contrast with our earlier work on quantum network based on nodes made of single cavity coupled to a three-level static atom, localized at the waist.

The corresponding Hamiltonian for independent cavities and the atom (considering a single resonant mode for each cavity) is of the form:

\[ \hat{H}_0 = \sum \hbar \omega_i \left( \hat{c}_i^\dagger \hat{c}_i + \frac{1}{2} \right) + \sum \hbar \epsilon_i \langle x \rangle \langle x \rangle \]

where \( \hat{c}_i^\dagger \), \( i = 1, r \) are creation operators for cavity modes with frequencies \( \omega_i \), and \( \langle x \rangle \) denotes the atomic levels with energy \( \hbar \epsilon_i \) (which also includes the shifts due to far off-resonant dipole trapping fields for transporting the atom).

At the overlapping waist the atom interacts with the fields of each of the cavities, absorbing or emitting single cavity photons. Each of the cavity modes are resonant to an energy corresponding to an atomic transition between the ground state \( |f_i \rangle \) and excited state \( |e_i \rangle \), \( i = l, r \). Furthermore, two classical fields couple the excited states \( |e_i \rangle \) and \( |e_j \rangle \) to a single ground state \( |f_m \rangle \) (see figure 1(b)). The atom-field interaction can then be effectively described with only five of these atomic levels [33, 34]. Within the dipole approximation, the effective interaction Hamiltonian then takes the form:

\[ \hat{H}_I = -\sum \hbar |g| \hat{G}_i \langle e_i \rangle \langle f_i \rangle + \Omega_i e^{-i\omega_i t} |e_i \rangle \langle f_m \rangle + h.c. \]

The first term on the right describes interaction of the atom with modes of the cavities with coupling constant \( g_i \)’s while the last term describes the classical laser fields of Rabi frequencies \( \Omega_i \) interacting with the atomic levels \( (|f_m \rangle, |e_i \rangle) \).

With the resulting geometry, the quantization axis should be uniquely chosen for describing states of the atom along with polarizations of the classical and cavity fields. In the present geometry, since the two cavities are orthogonal, there might be ambiguity in interpreting the polarizations and realizing the interaction Hamiltonian. Next, we argue that the above interaction Hamiltonian is indeed realizable with a judicious choice of the quantization axis and classical field directions.

Quantization axis: a possible level scheme for realizing the above Hamiltonian is shown in figure 1(c). It is based on the hyperfine structure of \( D_2 \) transition (5S_{1/2} \leftrightarrow 5P_{3/2}) \) of \(^{87}\text{Rb}\). The ground states of the model correspond to the lower manifold of \( F = (1, 2) \) levels: \( |f_0 \rangle = |F = 1, m_F = -1 \rangle, \)

1.1. Physical system

The system consists of three physical elements: (a) two single mode optical cavities (labeled \( l \) and \( r \) hereafter), geometrically arranged such that the cavity axes are orthogonal to each other with overlapping waists (figure 1). (b) A single atom, prepared in a specific state of its ground-state manifold at the overlapping modes of the cavities. We consider two distinct methods for atom-cavity coupling: the coupling is either dynamically induced as the atom is physically transported (either with an optical tweezer, an optical conveyor belt or flying across the waist at a constant velocity) through the overlapping waist and a stationary atom, trapped at the cavity waist. While in the first case, the atom-cavity coupling constant, \( g \), depends explicitly on time, for the latter, it remains constant. (c) Two classical fields or lasers with specific polarizations and frequencies, near resonant with the atom.

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the photon modes are right circularly polarized and linearly polarized along the quantization axis. The excited levels are cavity states and none of these three states contain any photon in the cavity. The atom is in a ground state \( |\psi(0)\rangle \), while the cavity coupling \( \Omega_l \) is tuned between the levels \( |F = 1, m_F = \pm 1\rangle \) and \( |F' = 2, m_F = 0\rangle \). The cavity coupling \( g_l \) is near-resonant with the transition \( |F = 1, m_F = 0\rangle \leftrightarrow |F' = 0, m_F' = 0\rangle \), while the coupling \( g_r \) is near-resonant with the transition \( |F = 2, m_F = \pm 1\rangle \leftrightarrow |F' = 2, m_F' = 0\rangle \). The pumps \( \Omega_l \) and \( \Omega_r \) are Raman-resonant with the transitions \( |F = 1, m_F = 0\rangle \leftrightarrow |F' = 0, m_F' = 0\rangle \) and \( |F = 1, m_F = 0\rangle \leftrightarrow |F' = 2, m_F' = 0\rangle \) respectively. All these transitions can be realized using the proposed geometry shown in figure 1(a). Here the quantization axis, \( \hat{z} \), is chosen to be along cavity \( l \) and cavity \( r \) is along the \( \hat{x} \)-axis. Both pumps \( \Omega_l \) and \( \Omega_r \), propagating along \( y \)-axis, are \( \hat{z} \) polarized along the quantization axis (\( \hat{z} \)).

Input optical mode of cavity \( l \) then correspond to a \( \sigma_+ \) polarization, with the atom initially prepared in state \( |f_l\rangle \). The resulting linearly polarized photon along \( y \) direction, emitted from the \( r \)-cavity, leaves the atom in a state that appears as an equal superposition of the states \( |F = 1, m_F = \pm 1\rangle \), with respect to the quantization axis. No other states in the manifold are coupled. In particular, one can note that scattering cross-sections corresponding to the single photon transitions from states \( |F = 1, m_F = \pm 1\rangle \) to \( |F' = 2, m_F = \pm 2\rangle \) are negligible for large cavity detunings, compared to the two photon-resonant transitions initiated by the classical fields. Furthermore, all the relevant transitions have reasonably high Clebsch–Gordan coefficients and are therefore efficiently coupled.

2. Methods

The model is effectively described with a ‘dressed’ state

\[
|\psi(t)\rangle = \sum_{n_1,n_2} e^{-i\left((n_1+n_2)+n_1\omega_r\right)t} A_{n_1,n_2}(n_1,n_2,t)|x,n_1,n_2\rangle.
\]
Here the state $|x, n_l, n_r\rangle$ is in the joint space of the atom along with $n_l$ and $n_r$ photons in the two cavities with an amplitude $A_e(n_l, n_r, t)$. The conservation of the excitation number $M$ (number of atomic excitations in $e_l, e_r, f_m +$ total photon number) allows one to work in a specific sector with fixed number of excitations, decoupling each other completely (ignoring decay and losses).

Equations for the probability amplitudes are then
\begin{equation}
-i\dot{A}_e(n_l, n_r, t) = g_s^*\sqrt{n_l + 1}A_e(n_l, n_r, t) e^{-i(\Delta_e + \Delta_g)t},
\end{equation}
\begin{equation}
-i\dot{A}_e(n_l, n_r + 1, t) = g_s^*\sqrt{n_r + 1}A_e(n_l, n_r, t) e^{-i(\Delta_e + \Delta_g)t},
\end{equation}
\begin{equation}
-i\dot{A}_e(n_l, n_r, t) = \Omega_l^r A_e(n_l, n_r, t) e^{-i\Delta_l},
\end{equation}
\begin{equation}
-i\dot{A}_e(n_l, n_r, t) = \Omega_r^g A_e(n_l, n_r, t) e^{-i\Delta_r},
\end{equation}
\begin{equation}
-i\dot{A}_e(n_l, n_r + 1, t) = g_2\sqrt{n_r + 1}A_e(n_l, n_r, t + 1, t) e^{-i(\Delta_e + \Delta_g)t},
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\end{equation}
where $\epsilon_{e_l} - \epsilon_{e_r} - \omega = \Delta_e + \Delta_g$, $\epsilon_{e_l} - \omega = \Delta_e$, $\epsilon_{e_r} - \omega = \Delta_g$. With $\Delta_l$ and $\Delta_r$ as the detuning parameters. Using standard transformation, $A_e = A_e e^{-i\Delta_l t}$, $A_e = A_e e^{-i\Delta_r t}$ for $i = l, r$ and $\dot{A}_e = \dot{A}_e$, equations assume the form
\begin{equation}
-i\frac{\partial \hat{A}(t)}{\partial t} = \hat{M}\hat{A},
\end{equation}
where
\begin{equation}
\hat{M} = \begin{pmatrix}
\Delta_e & 0 & 0 & g_s^* \\
0 & \Delta_g & 0 & 0 \\
g_e & 0 & \Omega_l^r & \Omega_r^g \\
g_e & 0 & \Omega_l^r & 0 - \Delta_l
\end{pmatrix}.
\end{equation}

Free of any of the excited states $|e_i\rangle$. An atom in such a state therefore does not emit any photons or remains dark.

This dark state $|D_{(0,0)}\rangle$ captures some of the essential physics of our model. In particular for the single-photon sector (corresponding to $n_l = n_r = 0$), the state $|D_{(0,0)}\rangle$ can be adiabatically rotated from an initial state $|f_r, 1, 0\rangle$ to a final state $|f_r, 0, 1\rangle$ by ramping up (down) and down (up) the fields $\Omega_l (\Omega_r)$, as long as one is slower than any other relevant timescales in the system. Most importantly, one can note that for this model, the initial and the final states correspond to a photon in each of the cavities or equivalently, the adiabatic rotation can route a single photon in an orthogonal direction.

Similar dark states and corresponding adiabatic rotation with counter-intuitive pulse sequences has been well studied and applied in varying physical scenarios. However, one can note that for the present model, precise timing requirement of such dark state rotation with counter-intuitive pulse sequences can get too demanding.

On the contrary, as we show next, there is a particularly interesting and yet unexplored off-resonant regime corresponding to $\Delta_l \gg \Omega_l \gg \epsilon_g$. In this regime, the photon can not only be routed from one orthogonal mode to the other with unity efficiency but can also be put into arbitrary superposition of two cavity modes by changing a single parameter in the model and with a simple, intuitive sequence of pulses.

Since the physics in all $M$-sectors is similar, for simplicity, we present an analysis of $M = 1$ i.e. sector corresponding to single excitation.

**Large-detuning limit: adiabatic rotation** in the large detuning limit ($\Delta_l = \Delta_r = \Delta = \Omega_l / g_e$) excited states are effectively decoupled from the dynamics (equations (4) and (5)), and can be adiabatically eliminated. The system is then well described with an effective three-level Hamiltonian:
\begin{equation}
H_{\text{eff}} = \begin{pmatrix}
\Omega_l^2 & |\Omega_r|^2 & \Omega_l^r g_{li} & \Omega_r^g g_{ri} \\
\Omega_l^r g_{li}^* & |\Omega_r|^2 & 0 & 0 \\
\Omega_r^g g_{ri}^* & 0 & |\Omega_r|^2 & 0 \\
0 & 0 & 0 & |\Omega_r|^2
\end{pmatrix},
\end{equation}
in the basis $|f_r, 0, 0\rangle = |f_m\rangle$, $|f_r, 1, 0\rangle = |f_l\rangle$, $|f_r, 0, 1\rangle = |F_l\rangle$. While all the states pick up light-shifts due to their coupling to the excited states (diagonal terms), they also get coupled to each other through an effective two-photon Rabi frequency, $s_r = \frac{\Omega_l^r g_{li}}{\Delta_l}$ couples the state with one photon in the cavity, $|F_l\rangle$ to the state $|f_m\rangle$ with the atom transitioning to the state by absorbing the cavity photon.

The Hamiltonian $H_{\text{eff}}$ again has a dark state consisting of only the dressed states $|F_l\rangle = |f_r, 1, 0\rangle$, $|F_r\rangle = |f_r, 0, 1\rangle$, given by
\begin{equation}
|D_{(0,0)}\rangle = s_r |F_l\rangle - s_l |F_r\rangle.
\end{equation}

It can be heuristically argued that the evolution of the intermediate ground state, to first order, is driven by the two excited states (we verify this argument in the next section with numerical simulations). Accordingly, for $\Omega_l \ll \Delta$ one can use $A_e \approx 0$ to eliminate the intermediate state. This gives rise to an effective Hamiltonian (keeping the two
detunings $\Delta_i$ and $\Delta_e$ explicitly) of the form:

$$H_{BS} = \begin{bmatrix} -\frac{|g|^2}{{\Delta_i}} + \frac{|s|^2}{{\Delta_e}} & \frac{s^*}{\delta_m} \\ \frac{s}{\delta_m} & -\frac{|g|^2}{{\Delta_e}} + \frac{|s|^2}{{\Delta_i}} \end{bmatrix},$$  

(10)

where

$$\delta_m = \sum_j \frac{|\Omega_j|^2}{\Delta_j}. \quad \text{(11)}$$

It can be noted that this Hamiltonian is akin to a beam-splitter like interaction, mixing the two orthogonal optical modes.

3. Results

3.1. Numerical simulations

To find the effective parameter regime for the validity of the heuristic argument, the effective five-level Hamiltonian has been numerically simulated. The results are summarized in figures 2 and 3. For these and all subsequent numerical results, the frequency (time) scale is chosen to be the linewidth (life-time) of the excited state $\gamma_e$ (corresponding to $\gamma_e = 5 \text{ MHz}$). Accordingly, all frequencies can be converted to physical units. The classical fields are taken in the form of Gaussian pulses of peaks $\Omega_i$’s and widths $\sigma_i$’s, overlapping each other in time.

Transit of the atom through the overlapping waist region with a constant velocity $v_a$ leads to time-dependent cavity coupling coefficient $g_a$ taken as Gaussian pulses in simulations with widths $\sigma_a \propto 1/v_a$ (figures 2(a), (c) and (e)).

Figures 2(a) and (b) correspond to splitting of a single photon equally in two cavity modes while figures 2(c) and (d) correspond to routing a photon from cavity $i$ to $r$. Both processes occur with unity efficiency. It is interesting to note that the classical fields are turned on before the transit of the atom, such that the two field shapes along with time-dependent cavity couplings overlap in time. This greatly simplifies the procedure without a need for fine tuning atomic velocity or pulse shapes.

Furthermore, one notes that the transfer occurs in a regime of $g_a(=3) \ll \Omega_i \ll \Delta$. However, for larger $g$, i.e. in the usual far-detuned limit of $g \geq \Omega, \Delta$ there is significant population accumulation in the intermediate states leading to partial transfer (figures 2(d) and (e)). Interestingly, this seems to indicate that a ‘bad-cavity regime’ might be better suited for this model, as opposed to usual stringent ‘strong-coupling’ requirement for single atom-cavity interactions. An analysis including decays (see below) supports this claim. Hereafter, we will refer to this regime as the beam-splitter regime denoted by $g \ll \Omega \ll \Delta$.

3.2. Beam-splitter regime

The connection between a beam-splitter mixing two single mode optical fields with the Hamiltonian of equation (10) becomes evident if one compares the unitary for a beam-splitter with the effective unitary:

$$\hat{U}_{eff} = \exp\left(-\frac{i}{\hbar} \int_0^{t_f} \hat{H}(t)dt\right) = \begin{bmatrix} A & B \\ C & D \end{bmatrix}. \quad \text{(12)}$$

where $t_f$ denotes the duration of the pulse sequence. The elements $A, B, C$ and $D$ are time independent, though dependent strongly on the time dependence of $g(t)$ and $\Omega(t)$. Numerical results show that by keeping the width and height of the classical pulses same and only changing the velocity of the transiting atom, one can continuously tune the elements of the unitary. Accordingly, we fix parameters ($g = 3$, $\Omega = 20$ and $\Delta_i = \Delta_e = \Delta = 50$) along with widths of the classical pulses ($\sigma_i = 20$) and tune a single parameter, the transit velocity of the atom, thereby controlling all the elements of the unitary $\hat{U}_{eff}$. The corresponding numerical results for the matrix elements $A$ and $B$ as functions of velocity are presented in figure 3(a).

The point $P_1$ with $B = 1.0$, corresponds to the photon-router regime with a single photon redirected from one cavity to the other. For cavities with a common and equal mode waist of $10 \mu m$ this corresponds to an atomic velocity of $20 \text{ m s}^{-1}$ for the Rubidium atom. However, it is interesting to note that with a cavity waist of $100 \mu m$ one can get the same routing behaviour with a flying hot atom with a velocity as high as $200 \text{ m s}^{-1}$. Realization of such a photon-router with hot atomic ovens without laser-cooling and trapping of single atoms can greatly simplify an important component of a future quantum network.

The point $P_2$ in figure 3(a) represents a superposition of a single photon being in either of the cavities, described as

$$|1_i, 0_r \rangle \rightarrow |1_i, 0_r \rangle \pm |0_i, 1_r \rangle.$$  

The usefulness of these photonic states as a resource to perform quantum computation tasks comes due to their association with corresponding atomic states ($|f, 1, 0 \rangle$, $|f, 0, 1 \rangle$) as a marker. Accordingly, a detection of a photon out of cavity $r$ together with a measurement of the atomic state uniquely characterizes the state $|0_i, 1_r \rangle$.

We also note that the same set of operations can also be achieved when the atom is trapped in the waist, so that $g$’s do not vary. For example, an initial state $|E_i \rangle$ will adiabatically transform to $|E_i \rangle$ under the simultaneous variations of the form: $\Omega_i$ going from $\Omega_0 \rightarrow 0$ and $\Omega_i$ going from $0 \rightarrow \Omega_0$ with appropriate parameter variations (not shown here).

3.3. Effect of losses

A single excitation can be lost either via a photon leaking out of the cavity or through spontaneous decay of the excited states in free-space modes. With the cavity decay rates as $\kappa_i$ and $\kappa_e$ and the decay rates of the excited levels $|e_i \rangle$ and $|e_e \rangle$ by $\gamma_i$ and $\gamma_e$ respectively, one can incorporate dissipation by modifying the diagonal elements of matrix $M$ (equation (7)) in an effective non-Hermitian Hamiltonian.
After adiabatically eliminating the excited states, cavity states $|F_{l}\rangle$ and $|F_{r}\rangle$ pick up an additional decay channel via the cavity couplings $g_{i}$'s. In the large detuning limit, since the excited states have amplitudes $g_{i}\Delta_{i}$, the effective decay rates for the cavity states become $\Gamma_{i}\approx g_{i}\Delta_{i}$, strongly suppressing spontaneous emission. Cavity states have one more decay channel through their coupling to the intermediate state $|f_{0},0,0\rangle$, which has a linewidth. Though the intermediate state is a stable ground state, it acquires a width due to classical fields coupling it to excited states, which decay. In the large-detuning limit, this coupling is $\approx\Omega_{i}/\Delta_{i}$, which gives a width $\sum_{i}\Omega_{i}^{2}\gamma_{i}/\Delta_{i}^{2}$ to the intermediate state. Recalling that the cavity states are coupled to the intermediate state by parameters $s_{i}$'s (definition following equation (9)) one obtains the total effective linewidths $\Gamma_{i}^{\text{eff}}$, ($i=l, r$) as:

$$\Gamma_{i}^{\text{eff}} = \kappa_{i} + \frac{|g_{i}|^{2}}{\Delta_{i}^{2}}\gamma_{i} + \frac{|\Omega_{i}|^{2}|g_{i}|^{2}}{\Delta_{i}^{2}\sum_{i}\Omega_{i}^{2}\gamma_{i}}.$$

(13)

The first term on the right, $\kappa_{i}$, is the desired cavity decay channel. Furthermore, for the beam-splitter regime, $\kappa_{i}\ll\frac{g_{i}^{2}}{\Delta}$ (see the off-diagonal terms of equation (10)) or equivalently, $\kappa_{l}$, $\kappa_{r}\ll g_{i}$, $g_{r}$. Physically, this condition requires the photon to be routed from one cavity to another before it decays out of the cavity. For the parameters used in figures 2 and 3, this requires $\kappa_{i} \approx \kappa_{r} \approx 0.3$ which translates to cavities with coupling coefficient of $g_{i} \approx 5$ MHz and $\kappa_{i} \approx 500$ kHz.

**Figure 2.** Numerical simulation for the controlled beam-splitter: (a), (c) and (e) depict Gaussian pulses of $\Omega_{i}$'s and $g_{i}$'s in units of $\gamma_{i} = 5$ MHz for three cases. The time-variations of $g_{i}$'s are due to the atom transiting with different velocities. (b), (d) and (f) are the corresponding variations of photon probabilities in two cavities $l$ and $r$. (b) Corresponds to a half-transfer of population (with pulse sequences of (a)) from the dressed state $|F_{l}\rangle$, containing one photon in cavity $l$ to $|F_{r}\rangle$ with a photon in cavity $r$. The resulting photonic state is a coherent superposition of the two cavity modes. (d) Depicts routing of a photon from cavity $l$ to cavity $r$ with near unity quantum efficiency. Insets of (b) and (d) show populations (two to three orders of magnitude smaller) of the intermediate ground state and excited states $|e_{i}\rangle$ and $|f_{m}\rangle$, respectively. In the regime of $g_{i}\gg\Omega_{i}$ corresponding to pulse shapes of (e) and (f) shows poor routing efficiency, with significant occupation of the intermediate excited states, leading to spontaneous emission losses.
The next two terms are bad, scattering the useful photon out in free-space (via the excited state). However, these scale as $g^2$, implying a scattering or failure rate of one in $10^4$ operations and can therefore be ignored.

Overall, the condition on the cavities, for the validity of the beam-splitter regime including losses can now be stated as:

$$\kappa_i \ll \gamma_i \approx g_i \ll \Omega_i \ll \Delta.$$  

Accordingly, a reasonable set of parameters for Rubidium atoms and two identical cavities are: $\kappa = 500$ kHz, $g = 5$ MHz with classical field strengths $\Omega = 30$ MHz of pulse width $1 \mu$s and a transiting atom of velocity $20$ m s$^{-1}$ (for a cavity waist of $10 \mu$m).

It can be noted that fluctuations of parameters $\Omega$ and $g$ will lead to further dissipation. In particular, a thermal distribution of atomic velocity will result in fluctuations in magnitude and temporal profile of cavity coupling $g$. While we find numerically that the scheme is relatively robust to small variations in $g$, a more detailed, stochastic modeling of atomic trajectory through the cavity waist is required to account for such effects [38].

3.4. A node of a quantum network

We now show how this twin-cavity system can function as an elemental node in a quantum network. Firstly, one can note that the photon states for one node (with the two cavities) are also associated with atomic ground states $|\nu\rangle$ and $|\nu\rangle$. These two atomic states can therefore be ascribed to a qubit. The beam-splitter Hamiltonian, can then be used to produce any arbitrary superposition or qubit state.

A quantum network can be envisaged to be a lattice of such nodes connected by optical fibers in an appropriate geometry. We consider two such nodes connected with a single-mode fiber. Specifically, arrangement considered here is one in which the $l$-cavity of one node is coupled to the $l$-cavity of the other as shown in figure 4. Next, we show how the two-node system described here can, in principle, perform, on demand, general one-qubit operations and some two-qubit gates in a controlled manner.

The Hamiltonian of this two-node system can be expressed as $H_2 = H_{\text{node}} + H_f$, where $H_{\text{node}}$ is the Hamiltonian of the two separate nodes and $H_f$ couples the nodes by transfer of photon through the fiber. With excitation number conserved, we once again limit the analysis to the single excitation sector. In the large detuning-limit, the excited levels for each atom, $|\nu\rangle$ and $|\nu\rangle$, can be eliminated and we

![Figure 3. Controlled Beam-splitter](#)
take the Stark shifts $\delta_i = |g_i|^2/\Delta_i$, $i = l, r$ to be equal. The latter is not necessary but is algebraically simple for presentation. One can then write

$$H_{\text{node}} = \sum_{k=1}^{2} \left[ \delta_k \left( |F_k\rangle_k \langle F_i| + |F_i\rangle_k \langle F_k| + \delta_{\text{km}} |F_m\rangle_k \langle F_m| \right) + \sum_{k=1}^{2} \left[ s_{kl} |F_k\rangle_k \langle F_m| + s_{kl} |F_k\rangle_k \langle F_m| + \text{H.c.} \right] \right]$$

(14)

where $k$ labels the nodes and the subscript on states refer to the node number. Parameters for each node are written using obvious generalization of the earlier notation: $s_{kl} = g_{kl}\Omega_{ki}^2/\Delta_{\text{km}}$, $\delta_i = |g_i|^2/\Delta_i$ for $k = 1, 2$ and $i = l, r$. Detuning $\delta_{\text{km}}$ is same as in equation (11) for node $k$.

The fiber coupling, $H_f$ describes the transfer of the photon from the cavities to a mode in the fiber

$$H_f = w |b^\dagger (C_{li} + C_{ri}) + b(C_{li}^\dagger + C_{ri}^\dagger)|$$

(15)

where $b^\dagger$ creates a photon in the fiber mode, $w$ is the cavity-fiber transfer amplitude for the photon.

In this sector, the single photonic or atomic excitation can be in either of the nodes. So the possible states corresponding to single or zero excitations on one node is a larger set. We list these states with a shorter notation for convenience in table 1.

In table 1, the first two states have zero excitation and the last five have one excitation. The Hilbert space for the two nodes consists of zero-excitation state of one node with one-excitation node of the other. In addition there are states in which the photon is in the fiber state. The two-node Hamiltonian connects 20 such states which we label from 1 to 20 and write the two-node wave-function as

$$|\Psi_2\rangle = \sum_{i=1}^{20} A_i(\bar{\ell}) |\bar{\ell}\rangle.$$

(16)

The list of all such two-node states along with their notations is tabulated in table 2. The notation used is $|a, b\rangle$, where $a$ and $b$ correspond to the states of node one and two respectively, and the second ket gives the photon number in the fiber. For writing the one-node states we use the shorter notation of table 1.

The Schrödinger equation in terms of the coefficients can now be written straight forwardly using the Hamiltonian $H_2$ given in equations (14) and (15)

$$i\hbar \frac{dA_i}{dt} = \sum_j M_{ij} A_j.$$

(17)

With the complete equations given in appendix. It is observed that the large $(20 \times 20)$ matrix $M_{ij}$ is rather sparse. A closer inspection reveals that the states are divided into groups of states which are just connected within each group. For example, states $|2, 4\rangle$, $|6, 14\rangle$, $|17\rangle$ form such a group. Certain eigenstates of these groups have characteristics of dark states, which allow adiabatic rotations among states of interest [30, 33]. The case $\delta_1 = \delta_2 = 0$ algebraically has the

---

**Table 1.** One-node states making up the two-node wavefunction.

| State vector | Notation |
|--------------|----------|
| $|f, 0, 0\rangle$ | $|l\rangle$ |
| $|f, 0, 0\rangle$ | $|r\rangle$ |
| $|f, 1, 0\rangle$ | $|F_1\rangle$ |
| $|f, 0, 1\rangle$ | $|F_2\rangle$ |
| $|f, 0, 0\rangle$ | $|F_3\rangle$ |
| $|f, 1, 0\rangle$ | $|F_4\rangle$ |
| $|f, 1, 0\rangle$ | $|F_5\rangle$ |

**Table 2.** Labels of two-node states in single excitation sector.

| Notation | State vector |
|----------|--------------|
| $|1\rangle$ | $|f, l\rangle|0\rangle$ |
| $|2\rangle$ | $|f, r\rangle|0\rangle$ |
| $|3\rangle$ | $|f, l\rangle|0\rangle$ |
| $|4\rangle$ | $|f, r\rangle|0\rangle$ |
| $|5\rangle$ | $|F_0, l\rangle|0\rangle$ |
| $|6\rangle$ | $|F_0, r\rangle|0\rangle$ |
| $|7\rangle$ | $|l, F_1\rangle|0\rangle$ |
| $|8\rangle$ | $|r, F_1\rangle|0\rangle$ |
| $|9\rangle$ | $|l, F_2\rangle|0\rangle$ |
| $|10\rangle$ | $|r, F_2\rangle|0\rangle$ |
| $|11\rangle$ | $|l, F_3\rangle|0\rangle$ |
| $|12\rangle$ | $|r, F_3\rangle|0\rangle$ |
| $|13\rangle$ | $|l, l\rangle|1\rangle$ |
| $|14\rangle$ | $|l, r\rangle|1\rangle$ |
| $|15\rangle$ | $|r, l\rangle|1\rangle$ |
| $|16\rangle$ | $|r, r\rangle|1\rangle$ |
| $|17\rangle$ | $|l, v\rangle|0\rangle$ |
| $|18\rangle$ | $|v, l\rangle|0\rangle$ |
| $|19\rangle$ | $|r, v\rangle|0\rangle$ |
| $|20\rangle$ | $|v, r\rangle|0\rangle$ |

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**Figure 4.** Two atom-cavity system connected by a fiber. The level scheme shown is the truncated one, in which the excited levels of atoms are eliminated. It represents the Hamiltonian in equation (14).
The simplest structure and one can identify three degenerate eigenstates of zero energy for qubit operations. These are:
\[
\begin{align*}
|D_3\rangle &= \left[ |F_l, r\rangle - \frac{s_{1l}}{s_{lr}} |F_r, r\rangle - |l, v\rangle \right] |0\rangle_f, \\
|D_4\rangle &= \left[ |r, F_l\rangle - \frac{s_{2l}}{s_{2r}} |r, F_r\rangle - |v, l\rangle \right] |0\rangle_f, \\
|D_5\rangle &= \left[ \frac{1}{s_{1r}} |F_l, l\rangle - s_{1l} |F_r, l\rangle \right. \\
& \quad \left. - \frac{1}{s_{2r}} |l, F_l\rangle - s_{2l} |l, F_r\rangle \right] |0\rangle_f.
\end{align*}
\]

These states have several interesting features. For example, for $|D_4\rangle$, adiabatic rotations in only one node (say vary $s_{1l}$ and $s_{1r}$) result in an intra-node transfer of a photon from one cavity to the other, while the state of the other node remains completely unaffected. For general detuning, these states are more complex and such adiabatic-rotation analysis, serves only as a broad guide. One can accordingly, resort to numerical simulations.

We have done numerical simulations for the complete set of 20 equations equation (17). Various protocols for parameter variations were tried and we present some in which useful operations can be done with high fidelity. It can be noted that there are several free parameters in the system. For two nodes we have four laser pulses of Gaussian shape whose amplitudes ($\Omega_0$) and widths ($\sigma_t$) can be varied. Similarly there are four atom-cavity couplings, which are again taken to be Gaussian with respective amplitudes $g_0$ and widths $\sigma_i$. The time interval between the pulses can also be varied. The frequency (time) scale is the same as in the last section, i.e. spontaneous decay rate of the excited atomic levels. As far as possible, parameters are kept the same in all the simulations. These are: $\Delta_{1l} = \Delta_{2l} = \Delta_{1r} = \Delta_{2r} = \Delta = 50$. While typical laser pulse magnitudes are $\Omega_0 = 0.5$, atom-cavity coupling $g_0 = 3$ and pulse widths: $\sigma_i = 20$, $\sigma_0 = 10$. Any finer changes around these values are mentioned in figures.

3.4.1. One-qubit operations. Single qubit operations for an isolated node have already been demonstrated in the earlier sections. Here, we show that even for a two-node system with weak fiber-coupling, single qubit operations can be performed on one node without disturbing the other qubit. Figure 5(a) shows a set of laser pulses ($\Omega_{1l}$ and $\Omega_{1r}$) applied on node-1, along with a weak variation of coupling parameters $g_{1l}$. Figure 5(b) shows that only two amplitudes change, resulting in a transfer of excitation from the initial state $|F_r, r\rangle$ to $|F_l, r\rangle$, while all other state amplitudes at both nodes remaining close to zero. Furthermore, small variations in parameters for pulses on node-1 (figure 5(c)) results in another important one-qubit operation (figure 5(d)). Here one generates an equal superposition $|F_r, r\rangle + |F_l, r\rangle$ from an initial state $|F_r, r\rangle$. In the lower most panel we show a flip operation on node-2. Figures 5(e) and (f) show the pulse and the corresponding variations of amplitudes in which the initial state $|r, F_l\rangle$ is transformed to $|r, F_l\rangle$. Note with a low fiber coupling $w = 0.001$, all these operations have high fidelity $\approx 1$ (in absence of dissipation). Though with increased coupling the fidelity drops, nevertheless a large parameter space remains to be explored for optimal operations.

3.4.2. Two-qubit operation. Next we demonstrate a few two-qubit operations. As seen above, all qubit operations require presence of a photon in a node. Since we are in the single-photon sector, it is essential to be able to transfer a photon from one node to the other. The top panel of figure 6 shows such an operation. Figure 6(a) shows two sets of pulses applied on node-1 and node-2 successively. Figure 6(b) shows the variation of amplitudes for the resulting process. Here the initial state $|F_r, l\rangle$ is transformed to $|l, F_l\rangle$, but it should be noted that the amplitude for the state $|l, F_r\rangle + |F_r, l\rangle$ becomes substantial after the first set of pulses applied on node-1. This is to be expected as the photon transfers to the fiber from l-cavity of node-1 and then to the l-cavity of node-2. The photon transfer in conjunction with single node operations can yield many useful gates required in quantum processing. In the middle panel, figure 6(c) shows a sequence of two pulses applied successively on nodes 1 and 2. The corresponding variations of amplitudes in 6(d) show that the initial state $|F_r, l\rangle$ is transformed to $\sqrt{2} |l, F_r\rangle + |r, F_l\rangle$ via the intermediate state $|l, F_l\rangle + |F_l, l\rangle$. This is an example of an entangled state of two qubits. Figures 6(e) and (f) show another operation in which the first pulse on node-1 transforms the state $|F_r, l\rangle$ to a superposition $|F_r, l\rangle + |l, F_l\rangle$. The second pulse on node-2 then transforms this state to $\sqrt{2} (|F_r, l\rangle + |l, F_r\rangle) + (|l, F_l\rangle + |F_l, l\rangle)$. In terms of atomic qubits this state is $|r, l\rangle + |l, r\rangle + |l, l\rangle$. Note that in these operations, common parameters are: $\Omega_{0l} = 20$, $\Omega_{0r} = 27$, $\sigma_{1l} = 20$, $\sigma_{1r} = 10$, $w = 0.6$ and time interval between the pulses is 150. Other variations are mentioned in figures.

Figure 7 shows two protocols with the initial state $|l, F_r\rangle + |F_r, l\rangle$. In the upper panel figures 7(a) and (b) a set of two pulses applied on node-1 and node-2 successively leads to its transformation to $|F_r, l\rangle$. While in the lower panel both the qubits are flipped resulting in the state $|l, F_r\rangle + |F_r, l\rangle$ by a variation in the parameters of the two pulses. The common parameters used here are same as in figure 6. The operation shown here are some representative examples which demonstrate the flexibility and versatility of the scheme.

4. Conclusion
To conclude, we have proposed a model for a node of a quantum network, consisting of two crossed-cavities with overlapping waists. The cavities are coupled to an effective five-level atom. The node performs a number of quantum operations in a controlled manner. An application of the node is demonstrated through routing photons with unity efficiency and as a programmable beam-splitter. We demonstrate these
two operations when the coupling is provided by flying atoms through the waist. The key advantage of the scheme is that it does not place stringent requirements on pulse sequences or atom-cavity coupling, and is workable with achievable cavity and atomic parameters.

For applications relevant for quantum network, we have considered a system of two nodes coupled through an optical fiber and shown performance of few elemental one and two qubit operations. Qubits for these operations are defined with degenerate atomic states while the operations are executed via transfer of photons within two cavities of a node or via internode transfer. Explicit parameter regimes for laser powers and atom cavity couplings are presented for a unitary qubit operations on one node (without disturbing the state of the other qubit), generation of two node superposition, swap operation and for generation of two qubit entangled states. Parameters for the model are well within the reach of what is achievable with current technology. There are efforts underway to realize high finesse crossed-cavities, both in optical and microwave domains [29, 32, 37]. It can be also be noted that the proposed model can be generalized to any cavity-QED based physical systems. In particular, coupled photonic crystal cavities along with gate-induced tuning of a single coupled quantum dot [35] or tunable superconducting qubits coupled to multiple microwave resonators [36] can also lead to viable possibilities for experimental implementation of the proposed system.

Figure 5. One-qubit operations for the two-node system. (a) Shows a pulse protocol on node-1 which flips the qubit on the node. This is shown in (b) where the state |F, r⟩ is transformed to |F, r⟩. (c) Shows a pulse protocol on node-1 which as seen in (d) generates a superposition |F, r⟩ + |F, r⟩ on the node from an initial state |F, r⟩. Note parameter changes here: $\Omega_{\text{RL}} = 20$, $\Omega_{\text{LR}} = 27$ and $\sigma_{g1} = 5.25$. (e) shows a pulse-protocol on node-2 which results in flipping qubit on node-2 as shown in (f). Here the initial state |r, F⟩ is transformed to |r, F⟩.
Figure 6. Three operations involving transfer of a photon from one node to the other. In the top panel (a) shows a sequence of two pulses applied on node-1 and node-2 that results (b) in transfer of the state $|F_l, l\rangle$ to $|l, F_r\rangle$. In the middle panel (c) shows a small variation in widths of the two pulses that results (d) in transfer of the state $|F_l, l\rangle$ to $\sqrt{2} (|l, F_r\rangle + |r, F_l\rangle)$. This is an entangled state of two qubits. The bottom panel (e) and (f) show a procedure in which the first pulse on node-1 transforms the state $|F_l, l\rangle$ to a superposition $|l, F_r\rangle + |F_l, l\rangle$. The second pulse on node-2 then transforms this state to $\sqrt{2} (|F_l, l\rangle + |l, F_r\rangle + |F_l, F_r\rangle$. Have combined two or three together

\[ i\dot{A}_{1,2} = \delta_1 A_{1,2} + s_{1,2} A_{5,6} + w A_{13,14} \]
\[ i\dot{A}_{3,4} = \delta_1 A_{3,4} + s_{1,2} A_{5,6} \]
\[ i\dot{A}_{5,6} = s_{1,2}^* A_{1,2} + s_{1,2}^* A_{3,4} + \delta_{lm} A_{5,6} \]
\[ i\dot{A}_{7,8} = \delta_2 A_{7,8} + s_{2,1} A_{11,12} + w A_{13,15} \]
\[ i\dot{A}_{9,10} = \delta_2 A_{9,10} + s_{2,1} A_{11,12} \]
\[ i\dot{A}_{11,12} = s_{2,1}^* A_{7,8} + s_{2,1}^* A_{9,10} + \delta_{lm} A_{11,12} \]
\[ i\dot{A}_{13} = w (A_{1} + A_{2}) \]
\[ i\dot{A}_{14,15,16} = w (A_{2,8,19} + A_{17,18,20}) \]
\[ i\dot{A}_{17,18} = w A_{4,15} \]
\[ i\dot{A}_{19,20} = w A_{6,6} \]

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Appendix

Here we write the full form of 20 equations of motion equation (17) for the coupled two-node system. They are obtained directly using the Hamiltonian given in equation (14). The notation for the states is provided in tables 1 and 2. Since many equations are very similar, we have combined two or three together
One can immediately see that $A_{16}$, $A_{19}$ and $A_{20}$ are just coupled amongst themselves and involve states of no physical interest here. The results in the text are obtained by solving the remaining set of 17 equations numerically with time variations of the coupling coefficients as indicated in the text.

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