Letter

Scalable, chip-based optically-controlled gates for quantum information processing

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Abstract

Here we present a simple and robust method to build on-the-fly configurable quantum gates based on a photonic exchange between quantum nodes. The idea is based on a high reflectivity of Bragg grating structures near resonant wavelengths. The control is exerted by applying an external strongly off-resonant or even a static electromagnetic field and taking advantage of the Kerr effect. When the nonlinear phase shift is strong enough, the Bragg mirror disappears, thereby allowing a transmission of a wave packet from one node to another. An example of a protocol for quantum logic gates that relies on this framework is offered.

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(Some figures may appear in colour only in the online journal)
A common node cavity enables photon exchange between the nodes, giving rise to multi-qubit gates. Because we adopted the classical nonlinear Kerr switching for quantum gates, the range of material systems exhibiting the necessary nonlinearity is very broad. The enabling calculation and experimental results, then, are directly applicable to any of those materials, after scaling by their Kerr coefficient.

The advantage of this scheme is its high degree of versatility in real-time. That is, single-qubit or many-qubit gates are supported, they can be configured on demand and switched on and off as needed. Because a photon exchange between a pair of nodes is sufficient to support most common quantum gates, a discussion of multi-node photon exchange is outside the scope of this manuscript.

Here we demonstrate our concept for quantum nodes that are single two-level quantum systems, for example, quantum dots or surface-trapped atoms. In this arrangement the system ‘atom+field’ is used as a qubit, acting as storage, as single qubit gate and initialization/readout; whereas a pair of nodes offers two-qubit gates. We point out that protocols discussed here require strong coupling of nodes and cavities because we used the most basic, 2-level systems as nodes. The use of N-level systems would eliminate the need for strong coupling (see, for example, [15] and references therein). An adaptation of our method to N-level systems, however, is beyond the scope of this letter.

3. Isolated node-cavity system and one-qubit gates

Each node is surrounded by Bragg mirrors giving rise to a set of isolated cavities with an effective length \( l_0 \). A node is coupled to one of the cavity eigenmodes: \( F_0(z,l_0) = \sqrt{2l_0} \sin(\pi z l_0) \), where \( s_0 \) is the mode number. The isolated node-cavity system with one quantum of excitation undergoes vacuum Rabi oscillations between the exited state of the atom and a single excitation of the cavity mode. The coupling between the node and the cavity is given by the Rabi frequency \( \Omega_0 \), which depends on a dipole moment \( d_{eg} \), a transition wavelength \( \lambda_1 \) and a mode volume \( V \). The expression for Rabi frequency in SGC is:

\[
\Omega_0 = d_{eg} \sqrt{\frac{8\pi^2 c}{hV\lambda_1}},
\]

where \( h \) is the Planck constant and \( c \) is the speed of light in vacuum. As demonstrated by prior work on quantum dots coupled to Bragg resonators [16–18], the mode volume can be small, i.e. on the order of \( (\lambda l_0)^3 \), and the dipole moment as large as \( \approx 10^{-28} C \cdot m \approx 10 e_a \) can be achieved, yielding the vacuum Rabi frequency of \( \Omega_0 \approx 10^4 \text{rad} \cdot s^{-1} \). To achieve such a frequency [19], the transparency of Bragg mirrors in the closed channel should be less than \( 10^{-6} \), which is within experimental reach [20].

For a single node with vacuum Rabi oscillations, a natural choice of basis for qubit states are bipartite states of the system ‘atom+field’: \(|0\rangle = |\mathcal{e}\rangle|0\rangle \) and \(|1\rangle = |\mathcal{g}\rangle|1\rangle \). These states can be read out in a deterministic way by creating a transparency...
channel between a cavity and a detector with an efficiency comparable to that of a single photon detector (For systems with more than two ‘atomic’ levels, the qubit state can be measured differently: via a transition to an auxiliary exited state of the ‘atom’.) An initialization of the nodes can be done by opening a transparency channel between a single photon source and a node, i.e. similarly to read-out. A single photon source can be made with a laser and one dedicated node, through various mechanisms described in the literature. One can use a resonant classical laser $\pi/2$-pulse [21] or adiabatic rapid passage [22] or a photon blockade [23]. In the first case, the resonant external laser field applied to the system provides a deterministic initialization of the node with $\pi$-pulses, $\pi/2$-pulses, etc. With the proper choice of the intensity of the initialization classical field, the initialization time can be made several orders shorter than one period of vacuum Rabi oscillations of the node. Since initialization is an operation on a single isolated node, multiple nodes can be initialized in parallel.

Several one-qubit gates can be implemented on isolated nodes. For instance, the ‘inversion’ operator can be implemented by freezing Rabi oscillations in selected nodes by sending a set of classical $2\pi$ pulses [21] and allowing these nodes to remain in their initial quantum states while letting the other nodes evolve.

4. Common node-cavities and photon exchange

Two-qubit gates are mediated by a high-fidelity photonic exchange between the nodes. Here we describe the use of cross-phase modulation that turns the Bragg mirror that separates the two isolated nodes into a transparent waveguide. Once the control field is applied between a pair of nodes, the mirror disappears and the photon exchange occurs. Because a typical photon exchange duration is much faster than any interactions between the nodes and cavities, only a photonic part of a bipartite qubit will be affected. While the corresponding photon wavefunction is in an eigenmode of an isolated cavity, it is not in an eigenmode of a much larger common cavity. Then, the propagation of the photon is given by a solution of the wave equation, which has the form:

$$\frac{\partial^2 E}{\partial z^2} - \frac{n^2}{c^2} \frac{\partial^2 E}{\partial t^2} = 0,$$

where $n$ is the effective refractive index of the waveguide. Equation (2) is supplemented with an initial condition $E(z,t=0)$. This condition describes a photon wave-packet immediately before the control field was applied. We have:

$$E(z,t=0) = F_{s0}(z,l_0).$$

This initial field is expanded over the set of the eigenfunctions of the common cavity with an effective length $l > l_0$. Within our approximations, a propagation of the initial field has an analytical solution:

$$E(z,t) = \sum C_i F_i(z,l)e^{i\omega t},$$

where the amplitudes $C_i = \langle F_i(z,l) | F_{s0}(z,l_0) \rangle$ are the projections of the initial state on each of the eigenmodes of the common cavity and $\omega = ksc/n = \pi sc/(nl)$, $s = 1,2,...$.

The energy distribution over different eigenmodes $F_i(z,l)$ of the common cavity is given by:
into equation (4),
\[ E_{\text{ tp c n E}} \] propagates during the geometrically evanescent waveguide dispersion does not change the standing wave field shape, i.e. when the field \( \lambda_2 \) is off. Figure 3 illustrates how the wave packet with \( \lambda_1 \) propagates during the photon exchange between the two nodes, where the stop band in-between the nodes is removed via a classical control optical field \( \lambda_2 \) with dispersion \( D = 10 \) ps (nm · km)\(^{-1}\). It is evident from figure 3 that at the end of the transfer, the wavepacket has traveled to the target node with nearly unchanged amplitude and phase profiles. Therefore, upon propagation, a photon can readily interact with an 'atom' in the target node once the control field is switched off.

Our subsequent simulation illustrates the effect of dispersion on fidelity \( F \) of node-to-node photon state transfer. Because we use normalized eigenfunctions as initial conditions for (2),
\[ F = \frac{\text{Re}\left( \int_0^L E(\zeta, t = 2\Omega t / c)N^2(t) E(t, t = 0) d\zeta \right) \text{Im}\left( \int_0^L E(\zeta, t = 2\Omega t / c)N^2(t) E(t, t = 0) d\zeta \right)}{\int_0^L \|E(\zeta, t = 2\Omega t / c)N^2(t) E(t, t = 0)\|^2 d\zeta} \] we see that the higher the mode of an isolated cavity \( s_0 \) is used; the higher dispersion coefficients can be tolerated, figure 4. In particular, for a node-to-node transfer described above with \( s_0 = 1 \) the calculated fidelity \( F > 0.99 \) for \( D < 2 \) ps (nm · km)\(^{-1}\), while for \( s_0 = 10 \) dispersion coefficients up to \( D = 20 \) ps (nm · km)\(^{-1}\) yield \( F \approx 0.99 \), a ten-fold reduction in dispersion spreading. Fidelity monotonically decreases as the group velocity dispersion increases. Ripples are seen in figure 4 due to the interference effects.

5. Two-qubit gate protocols

Using isolated node manipulations and the controlled photon exchange between nodes we have developed simple and scalable protocols for basic quantum gates: SWAP, CNOT, and

| Table 1. Protocols for SWAP and CNOT operations. |
|-----------------------------------------------|
| **SWAP** |
| \( \tilde{\rho} \) |
| \( |0\rangle_1 |1\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |0\rangle_1 |1\rangle_2 |0\rangle_1 |1\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |0\rangle_1 |1\rangle_2 \) |
| \( |0\rangle_1 |1\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| **CNOT** |
| \( \tilde{\rho} \) |
| \( |0\rangle_1 |1\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |0\rangle_1 |1\rangle_2 |0\rangle_1 |1\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |0\rangle_1 |1\rangle_2 \) |
| \( |0\rangle_1 |1\rangle_2 |1\rangle_1 |0\rangle_2 \) |
| \( |1\rangle_1 |0\rangle_2 |1\rangle_1 |0\rangle_2 \) |

\[ W_2 = |C|^2 = \frac{2(-1)^{s_0} \sin(\alpha \phi)}{(\frac{1}{t})^{3/2} \sqrt{\frac{\pi}{s_0} (\pi^2 s_0^2 - \pi^2 s_0^2 s_0^2)}}. \]
phase rotation. We define the two basic operations supported by the framework of an actively-switched photonic network and required to build quantum gates. We assume that the two cubits are stored in the two adjacent isolated cavities, labeled 1 and 2. First, a photon exchange between cavities through a common cavity is denoted by an exchange operator:

$$\hat{\sigma}|a\rangle_{i}|n\rangle_{1}|a2\rangle_{2}|n2\rangle_{2} = |a\rangle_{i}|n\rangle_{1}|a2\rangle_{2}|n2\rangle_{2}$$

(6)

where $|a\rangle_{i}$ is an atomic state and $|n\rangle_{i}$ is a photon number state in the $i\text{th}$ cavity. Note that this procedure only exchanges the photonic part of the two bipartite cubits, leaving the node excitation part intact. Second, an inversion operator is naturally provided by Rabi oscillations in the nodes. When $2\pi$-pulses are applied to the selected group of nodes to freeze their evolution, the other nodes undergo an inversion:

$$\hat{\pi}|g\rangle|n\rangle = |e\rangle|n - 1\rangle$$

(7)

$$\hat{\pi}|e\rangle|n\rangle = |g\rangle|n + 1\rangle$$

(8)

where $\hat{\pi}$—is a delay equal to one half of Rabi oscillations period during the evolution of a node. In the table, $\hat{\pi}$ refers to an inversion on both the nodes, $\hat{\pi}_i$, where $i = 1,2$ denotes an inversion of just the first (second) node, while the other one is not inverted.

The above operators are sufficient to implement SWAP and CNOT gates. The associated protocols are shown in table 1.

6. Conclusions

We have introduced a method of implementing scalable light-controlled gates for quantum information processing, based on N-level systems (nodes) exchanging photons via optically-controlled nonlinear Bragg waveguides. This method for all-optical switching relies on cross-phase modulation that removes the effective Bragg resonance and creates transparency so that the control field can be detuned very far from any of the node’s resonances. Because the control field $\lambda_2$ is a classical field, multiple pairs of nodes may be controlled simultaneously. Thus, quantum information processing can be made massively parallel. This design can be implemented on a chip with the currently existing technology. A particular experimental implementation of this method is described in the supplementary material section. Note, that our proposal only requires a fast (femtosecond scale) classical switch. The search for fast, high-contrast optically controlled switches is a very active field of research, [24–28]. Therefore one should expect alternative experimental realizations of the proposed protocol.

While, in our manuscript, we assumed that the control field is applied perpendicular to the waveguide, other configurations are also possible. To implement even more exotic, multi-body based algorithms [29], nodes can be arranged in one-, two-, or three-dimensional structures. This opens the way to implement topologically protected and/or massively parallel quantum interactions, commonly studied with cold atoms in optical lattices, on a solid state chip.

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Appendix

We discussed the general idea and introduced protocols for a quantum circuit based on actively switched optical channels. Because a broad range of materials and geometries can be employed to implement these circuits [30–33] we presented general design features for such channels. Here we offer a detailed study of one possible implementation of our protocol, based on a lithium niobate waveguide. The purpose of this study is twofold. First, it provides guidance for an enabling experiment with Bragg switches. Second, it introduces certain technological enhancements that are aimed at scalability.

Lithium niobate is commonly used in optical devices and is relatively well understood. We consider it a good candidate for an experimental realization because it has one of the largest Kerr nonlinearity constants $n_2 = 83.310^{-16}$ cm$^2$ W$^{-1}$ [34].

All necessary fabrication techniques required for realization of our proposal, particularly the making of the waveguides [20], the permanent modulation of the refraction index [35, 36], the making of the dynamic gratings with a static electrical field [14] and placing the conductive nanostructures [37] were demonstrated. The making of the dynamic gratings with a light field [13] as well as field enhancement with periodic conductive nanostructures [38] were demonstrated in other materials, but are compatible with lithium niobate.

The overall geometry of the structure and the mode profile is shown in figure 1 of the manuscript. The waveguide can be implemented in lithium niobate with an average diameter of

![Figure 5. Field intensity enhancement in LiNbO3 waveguide due to plasmonic resonance of gold nanoantenna (dark blue trapezoidal regions) with air cladding.](image-url)
≈ 500 nm on a lower refractive index (RI) substrate. Periodic variation of the RI of the waveguide can be implemented statically in several different ways or dynamically induced by a control field. An RI contrast is chosen based on the nonlinear constant, available control field power and the required bandgap width.

To implement the proposed switching mechanism in a lithium niobate waveguide, a low RI contrast should be used, such that a control field of \( \approx 10^{10} \cdots 10^{11} \text{W cm}^{-2} \) can erase a static Bragg grating (BG). This intensity corresponds to a refractive index change due to a Kerr cross-phase modulation of \( \approx 10^{-3} \).

To make a high Q cavity with such a low RI contrast, the length of the grating structure should be quite long: \( \approx 3000 \) periods. Notice that the bandgap width will also decrease (\( \delta\lambda \approx 0.5 \text{ nm} \)), thus a higher longitudinal cavity mode should be coupled to the quantum node to provide adequate high reflectivity (\( \delta_0 \approx 100 \)). To erase a BG structure, the control field should have a spatial pattern that corresponds to the periodic BG structure. An array of plasmonic antennas along the waveguide could be used to enhance the local optical field and provide modulation [38, 39], thus reducing the power required and simplifying the preparation of the control beams. The effect of placing the antennas is shown in figure 5, where a numerical calculation was performed with an FDTD method for a control field with a wavelength \( \lambda = 1 \mu\text{m} \). As a result, a simple flat-top (and unmodulated) control field can be applied as a control, further aiding scalability. In addition, such an array significantly decreases the power requirements for the control field. Eliminating the BG altogether is advantageous because it reduces undesired dispersion effects. We have modeled the propagation of a wave packet through the waveguide and it is evident that dispersion in this system does not significantly distort the wavepacket, yielding a fidelity above 0.99 for the 10th mode (see figure 4 of the manuscript). As seen in that figure, the fidelity improves further at even higher mode numbers. Placing an array of conducting particles in the vicinity of the waveguide does not significantly increase the propagation loss because the \( \lambda_1 \) field is far off-resonant for gold nanoantennas.

Assuming that the control field is sufficiently detuned from the transition frequency used in the 2-level nodes (\( \Delta \lambda > 100 \text{ nm} \)), a qubit dipole moment of \( \approx e\alpha_0 \) (where \( \alpha_0 \) is Bohr radius) and a resonant energy of \( \approx 1 \text{ eV} \) [40], the resulting Rabi oscillations of the qubit due to the control field yield an amplitude that is less than \( 10^{-3} \), i.e. negligible on the time scale in question. Minor drawbacks of using shallow index contrast vs. high index contrast are the relatively long distances between nodes, longer times for gate operations, and longer Rabi periods for the coupling between nodes and isolated cavities (up to 10 times in comparison to the lowest longitudinal mode attainable with high index contrast).

On the other hand, high RI contrast, for instance \( \delta n = 0.04 \), would result in a wide band-gap (\( \delta\lambda \approx 10 \text{ nm} \), providing finesses \( f \approx 10^8 \) for 10th field mode) with only 200 periods. When a large index contrast is employed, an optically controlled bandgap shift, rather than a BG erasure, should be used to lower switching power requirements. Remarkably, a required control field pulse peak power of \( \approx 720 \text{ W} \) and a duration of \( \approx 100 \text{ ps} \) is sufficient to achieve bandgap shifts of \( \approx 0.6 \text{ nm} \) in barium fluoride [26, 41]. Due to a larger Kerr coefficient, this power is reduced more than 30-fold in lithium niobate. To reduce the role of strong dispersion as the light propagates in the medium with an optically detuned bandgap, a higher-order longitudinal isolated cavity mode should be employed (see figure 2 of the manuscript), and/or dispersion compensation techniques should be applied. Dispersion engineering can be achieved e.g. though advanced 2D patterning of lithium niobate [42].

In conclusion, we discussed general design rules using one possible experimental implementation of the platform as an example. In addition an array of plasmonic antennas along the waveguide enhances the local optical field and provides modulation, thus aids scalability.

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