Quantum Computing with Incoherent Resources and Quantum Jumps

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(Received 4 November 2011; revised manuscript received 31 December 2011; published 25 April 2012)

Spontaneous emission and the inelastic scattering of photons are two natural processes usually associated with decoherence and the reduction in the capacity to process quantum information. Here we show that, when suitably detected, these photons are sufficient to build all the fundamental blocks needed to perform quantum computation in the emitting qubits while protecting them from deleterious dissipative effects. We exemplify this by showing how to efficiently prepare graph states for the implementation of measurement-based quantum computation.

DOI: 10.1103/PhysRevLett.108.170501 PACS numbers: 03.67.Lx, 03.65.Yz, 03.67.Pp, 42.50.Lc

In the traditional circuit model of quantum computation [1], a quantum algorithm is implemented by the sequential action of entangling gates and local unitaries followed by a final measurement stage that reveals the result of the computation. An alternative approach is the measurement-based quantum computation [2,3], where an initial highly entangled multiqubit graph state is prepared and a succession of adaptive measurements on the individual qubits defines the implementation of a specific algorithm. In both scenarios, quantum computation can be carried out with a basic toolbox of three elements: measurements on the computational basis, single qubit rotations, and specific two-qubit entangling gates.

In both computational models, decoherence [4] poses a major obstacle to practical implementations, as the qubits are encoded in systems that are unavoidably coupled to the environment. Emitted and scattered photons are two notable examples of natural processes associated with decoherence in quantum systems. However, when detected, the same photons lead to the observation of quantum jumps [5–7], which can be used to read information out of the systems, to manipulate them to some extent, and even to fully and locally protect previously created entanglement between distant parts [8].

Here we show that quantum jumps actually have full and efficient quantum computing power, i.e., that not only does the ensemble of quantum trajectories produced by incoherent dissipation and absorption include any quantum computational protocol, all the way from the primary entangling gates to the final measurements, but also that it is possible, at least in principle, to unravel, in polynomial time, the dynamical evolution of the system into one of these trajectories. We demonstrate this by designing a quantum optical gedanken experiment in which the three basic elements of the quantum computing toolbox can be implemented by suitably monitoring the photons that are spontaneously emitted and inelastically scattered by qubits into their surrounding environment. We also show that the same scheme protects the system as a whole from dissipation while the computation is carried on. Furthermore, our method efficiently produces graph states for the implementation of measurement-based quantum computation [2,3]. In Ref. [9], we give a further example showing how to teleport an unknown quantum state, a protocol that already encompasses all the required elements for any quantum computation.

The first tool to be described is how to measure a qubit in the computational basis which is promptly achieved by monitoring its spontaneous emission (Fig. 1.1). If a nondegenerate qubit, prepared in a superposition $\alpha|g\rangle + \beta|e\rangle$, is allowed to decay at a rate γ_-, then it becomes entangled with a propagating photonic mode as $\alpha|g\rangle|0\rangle + \beta\sqrt{e^{-\gamma_- t}}|e\rangle|0\rangle + \beta\sqrt{1-e^{-\gamma_- t}}|g\rangle|1\rangle$. For times much larger than $T = 1/\gamma_-$, the excited (ground) component of the original state of the qubit becomes correlated with the detection (no detection) of the spontaneously emitted photon. A detection identifies the output 1 (related to state $|e\rangle$) and defines a quantum jump in the system that, mathematically, corresponds to applying the lowering operator σ_- to the state of the qubit. On the other hand, no detection outputs 0 and is associated to the so-called no-jump trajectory [10,11]. Monitoring the spontaneous decay is then equivalent to measuring the state of the qubit in the computational basis, however, always dragging the qubit into its ground state.

The nonunitary character of the spontaneous decay means that the remaining tools for quantum computation cannot be obtained solely by monitoring spontaneously emitted ("s.e.") photons. The curious solution comes from the addition of an extra incoherent process to each qubit: the inelastic scattering ("i.s.") of photons provided by an external laser field. This situation has

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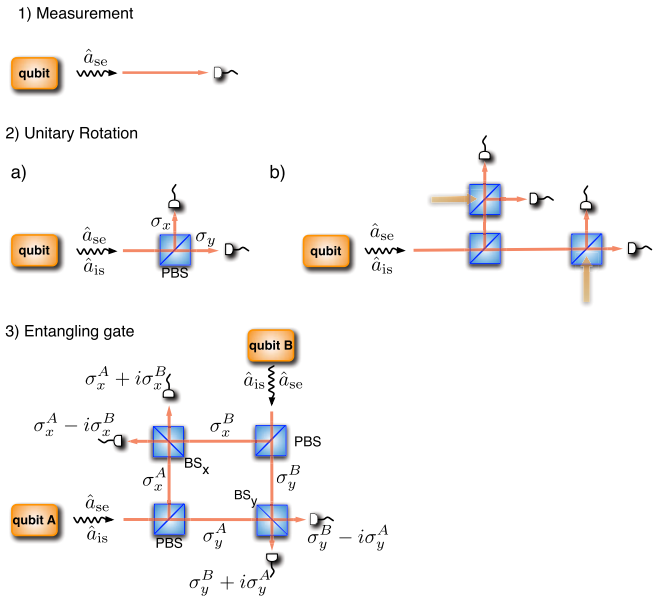


FIG. 1 (color online). Quantum-jump version of the building blocks for quantum computation. (1) Spontaneous emission (s.e.) is sufficient for measuring in the computational basis. [2(a)] The PBS erases which-process information, and a photon-detection produces a jump proportional to a unitary flip in the emitting qubit. [2(b)] More general rotations can be obtained through homodyne detection schemes [8]. (3) Entangling gates achieved through a second quantum erasing process: the standard balanced BSs erase which-qubit information, and a detection event in one of the available photocounters implements the corresponding maximally entangling operation indicated in the figure.

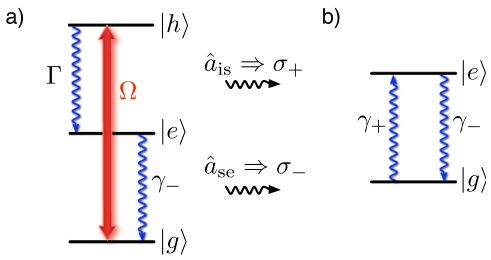


FIG. 2 (color online). (a) Each qubit is encoded in the two lower levels $|g\rangle$ and $|e\rangle$ (corresponding to the logical $|0\rangle$ and $|1\rangle$) of a discrete emitter. The third level $|h\rangle$ is introduced to produce the necessary inelastic scattering that incoherently pumps the qubit back into its excited state $|e\rangle$ and complements spontaneous emission in order to generate the building blocks for quantum computation described in Fig. 1. (b) The two decoherence processes in the effective two-level system obtained after adiabatic eliminating level $|h\rangle$.

been investigated in detail in Ref. [8] in the context of entanglement protection. For completeness, here we briefly reproduce some results of Ref. [8] that are essential to understand our computational scheme. First note that the scattering process, shown in Fig. 2(a), optically pumps an emitting qubit back into its excited state. In the limit where level $|h\rangle$ can be adiabatically eliminated [8], the dynamics is equivalent to the open system with two channels shown in Fig. 2(b) and described by the master equation

$$\dot{\rho} = \gamma_- \mathcal{D}[\sigma_-]\rho + \gamma_+ \mathcal{D}[\sigma_+]\rho, \tag{1}$$

where $\mathcal{D}[L]\rho = L\rho L^\dagger - (L^\dagger L\rho + \rho L^\dagger L)/2$ is the decoherence operator and $\gamma_+ = 4\Omega^2/\Gamma$. It is clear that when both reservoirs are monitored, i.e., when the emitted photons are detected, this equation allows for an unravelling in terms of two distinct jump processes [10,11]: the usual decay corresponding to the detection of s.e. photons and the respective application of σ_- in the system, and the complementary excitation process connected with the detection of photons in the i.s. channel (σ_+). However, this is only one of the infinite ways to unravel Eq. (1) in terms of stochastic processes [11–15]. One could, for example, use the relations $\sigma_- = (\sigma_x - i\sigma_y)/2$, $\sigma_+ = (\sigma_x + i\sigma_y)/2$ to rewrite Eq. (1) as

$$\dot{\rho} = \frac{\gamma}{2}\mathcal{D}[\sigma_x]\rho + \frac{\gamma}{2}\mathcal{D}[\sigma_y]\rho, \tag{2}$$

where we assumed equal rates for all incoherent processes by setting $\gamma_+ = \gamma_- = \gamma$. In this form, the equation suggests a stochastic evolution in terms of the jumps σ_x and σ_y. This unravelling has two crucial features: First, despite describing a detection process, the jumps correspond to unitary evolutions given by either σ_x or σ_y; second, the no-jump dynamics, given by the non-Hermitian effective Hamiltonian [10] $H_{\text{eff}} = -i\hbar\gamma(\sigma_x^\dagger\sigma_x + \sigma_y^\dagger\sigma_y)/4$, is proportional to the identity; i.e., the no-jump probability is state independent. This means that, in between clicks, the state of the emitting qubit remains the same.

Physically, we can interpret this situation with the monitoring scheme shown in Fig. 1.2(a): If both s.e. and i.s. processes are tuned to output photons that are indistinguishable in frequency and linewidth but of orthogonal circular polarizations, then, by placing polarized beam splitters (PBS) before the photocounters, the which-process information is erased [16,17]. In this case, the detection of a photon that comes out of the PBS will implement quantum jumps that are linear combinations of σ_- and σ_+ corresponding to unitary flips of the type $\cos\theta\sigma_x + \sin\theta\sigma_y$ on the qubit, with θ determined by the alignment of the PBS. Note that more general rotations are also possible if these new channels are monitored through homodyne detection [8] [Fig. 1.2(b)] or combined with an external classical photon source.

We are finally in a position to describe the remaining building block: the two-qubit entangling gate shown in Fig. 1.3. Once again, each qubit spontaneously emits and inelastically scatters photons into its own environments. If these photons are ignored, the dynamics of the qubits is properly described by a simple extension of Eq. (2) given by

$$\dot{\rho}_{A,B} = \sum_{i=A,B} \frac{\gamma}{2} \mathcal{D}[\sigma_{x_i}]\rho_{A,B} + \frac{\gamma}{2} \mathcal{D}[\sigma_{y_i}]\rho_{A,B}. \quad (3)$$

If the local reservoirs are monitored, then we get back the situation of Fig. 1.2(a) for each qubit, and we locally protect the global two-qubit state against dissipation [8].

However, in order to create entanglement between them, a second quantum erasing process is required, one that destroys the which-qubit information. This is achieved by combining the output ports of the PBS of two different qubits in a standard beam splitter (BS), which corresponds to one of the infinitely many possible unravellings of Eq. (3). In this case, as a consequence of the erasing process, clicks in the detectors of the newly defined channels will correspond to entangling jumps given by linear combinations of local unitary flips, such as $X^{\pm}_{AB} = (\sigma_{x_A} \pm i\sigma_{x_B})/\sqrt{2}$, where the sign is randomly determined by the channel where the photon is detected. For example, if the initial state of the qubits is $|g_A\rangle \otimes |g_B\rangle$, then a single detection event in these combined detectors produces a maximally entangled Bell state (Fig. 1.3). Note that, once again, in between clicks, the no-jump operator is proportional to the identity and the global state of both qubits does not change.

After the first combined detection is obtained, the standard BSs have to be removed in order to avoid undoing the entangling operation. This removal returns the reservoir monitoring setup to the configuration in which clicks just correspond to local flips, thus preserving the existing entanglement shared by the pair [8]. This removal is now achievable in very fast time scales, such as the method used in Ref. [18] to implement Wheeler's delayed choice experiment. Finally, note that, as usual in quantum-jump processes, in each realization a different trajectory and hence a different, albeit locally unitarily equivalent, entangled state will be produced. Each locally unitarily equivalent entangled state will correspond to a particular sequence of clicks, and, by keeping track of them all, one will always identify the final pure state of the global system. Naturally, the convex sum of all possible trajectories after enough photons are detected must be proportional to the identity, which is the asymptotic solution of Eq. (3).

A major advantage of this two-qubit gate over usual ways to entangle qubits by monitoring spontaneously emitted photons [19–21] is its scalability. Despite being based on detections, our gates are unitary, allowing the creation of arrays and networks of entangled qubits. This property is essential, for example, in the implementation of a teleportation protocol (see [9]) or the creation of multiqubit graph states, as we proceed to describe.

The generation of a multiqubit graph state, as shown in Fig. 3(a), follows from the parallel application of the entangling operation for every edge in the graph. Each vertex of the graph is an emitting and scattering qubit, and their reservoir channels are initially combined

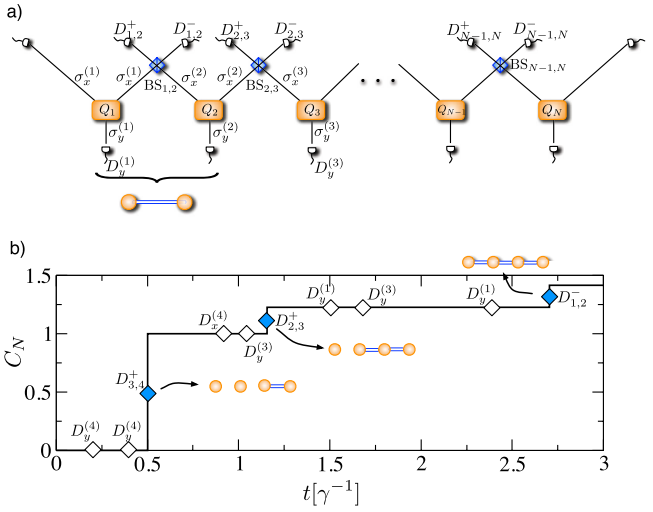


FIG. 3 (color online). The generation of a multiqubit state equivalent to a linear cluster state can be obtained by concatenating a variant of the entangling gate shown in Fig. 1.3. In (a) only the σ_x channels are combined at the BS, while the σ_y channels are detected independently (corresponding to the absence of BS_y in Fig. 1.3). (b) Single quantum trajectory simulation for the evolution of entanglement (given by the N concurrence [34]) for 4 qubits initialized in the ground state. Detection events are marked with diamonds labeled by the corresponding detector. Jumps from combined detectors (blue-filled diamonds) generate entanglement: At $t \approx 0.5/\gamma$ a Bell state between qubits 3 and 4 is formed, and the subsequent jumps at $t \approx 1.16/\gamma$ and $2.7/\gamma$ complete the generation of a four-qubit cluster. Note that the intermediate local jumps change the state of the system but not the entanglement.

according to the graph topology. However, since noncommuting operations are not desirable, only one of the channels, for example, σ_x, of each qubit should be combined. The other channel is still monitored (corresponding to Fig. 1.3 without BS_y), and its clicks do not affect the overall result (they just correspond to locally redefining the logical "0" and "1"). Each time an entangling photon is detected, the corresponding standard BS is removed, and subsequent photons in the same detectors locally flip the respective qubits. Once again, in between clicks the global state does not evolve ($H_{\text{eff}} \propto I$). An example of a quantum trajectory simulation for the generation of a four-qubit linear graph state is shown in Fig. 3(b).

We now proceed to prove that the state generated by this scheme is formally equivalent to a graph state. Let us start with all N qubits in the eigenstate $|0\rangle$ of σ_z [22]. The generate state is

$$|\tilde{G}_N\rangle = \prod_{(j,k)\in G} X^{s_{j,k}}_{jk} |0\rangle^{\otimes N}, \quad (4)$$

where $s_{j,k} = \pm$ depends on which detector clicks in the combined $BS_{j,k}$ channels and the j's and k's are chosen according to the connections present in the graph. Note that

the $X_{jk}^{\pm}$ gates are formally equivalent (up to local unitaries) to controlled-Z gates used to produce usual graph states [23] as $|G_N\rangle = \prod_{(j,k)\in G} cZ_{j,k}|+\rangle^{\otimes N}$. Indeed they can be rewritten as $X_{jk}^{\pm} = e^{\pm i\pi/4} X_j^{\mp} X_k^{\pm} cX_{j,k}$, where $X_j^{\pm} = (\mathcal{I} \pm i\sigma_{x_j})/\sqrt{2}$ is a local unitary and $cX_{j,k}$ is the X basis version of the $cZ_{j,k}$ gate. By using that $H\sigma_{x_j}H = \sigma_{z_j}$ and the fact that all gates involving σ_z commute, the state in Eq. (4) can then be rewritten as

$$|\tilde{G}_N\rangle = \prod_j U_j |G_N\rangle, \tag{5}$$

with $\prod_j U_j = H^{\otimes N} \prod_{(j,k)\in G} e^{s_{j,k} i\pi/4} Z_j^{-s_{j,k}} Z_k^{+s_{j,k}}$, $Z_j^{\pm s_{j,k}} = (\mathcal{I} \pm s_{j,k} i\sigma_{z_j})/\sqrt{2}$, and H being the usual Hadamard transformation. The state generated by our method is the usual graph state, up to local unitaries conditioned on the clicks. It is worth mentioning that all the measurements necessary to perform a universal one-way quantum computation are in the computational basis or comprised in the x-y plane of the Bloch sphere and can indeed be performed with the schemes presented in Figs. 1.1 and 1.2(a), respectively. Given that, it can be formally shown that all the different clicks in the detectors (corresponding to different local unitaries) can be treated as the usual adaptations in the measurement basis of a given one-way algorithm [2,3]. However, due to the products of Hadamards appearing in U_j, it is also necessary to introduce such transformations in our framework. This is easily achieved by the same quantum erasing principle that allows for the implementation of single qubit rotations in the x-y plane. However, in this case, an extra classical photon source is required. For example, if the propagation channel of an attenuated laser is combined in a standard BS with the σ_y output channel of a given emitting qubit k, then the detection of a photon after the BS corresponds to applying the jump operator $(\mathcal{I} \pm i\sigma_{y_k})/\sqrt{2} = \exp(\pm i\frac{\pi}{4}\sigma_{y_k})$ to qubit k, which also changes qubit k from Z basis to X basis and vice versa.

The procedure to prepare graph states described above is efficient: If τ is the average time to entangle one pair (average time for one of two qubits to emit or scatter one photon), then the parallel character of our proposal implies an overhead time of the order of $\tau \log N_e$ for the creation of an entire graph state comprised of N_e edges. In order to prove this, note that the generation process is completed after the last combined detector has clicked, in an instance of the famous coupon collector problem [24]. The mean value for the time of occurrence of this last edge is of the order of $\tau \log N_e$, where τ is the typical time for one independent detection. This reasoning is confirmed by the simulations of the average preparation time for a linear graph state with N vertices depicted in Fig. 4. It is important to stress that the unitarity of the jump operators implies that the probability of occurrence of any jump is independent of the state of the system. This means that one does not need to keep track of the state during the quantum trajectory calculations, thus allowing for simulations of a large number of qubits.

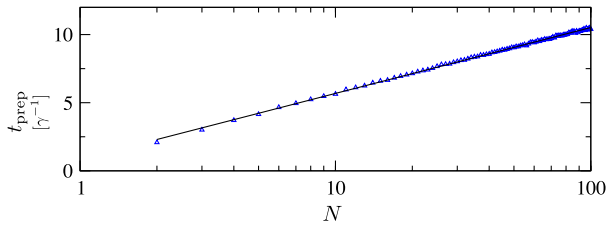


FIG. 4 (color online). Average time for the preparation of a graph state for up to 100 qubits. Simulations used 5000 trajectories and confirm the logarithmic scale with the number of vertices explained in the text.

The results here presented are fundamentally different from other proposals that use dissipation as resources for quantum computation. There, either the interactions between qubits and reservoirs have to be carefully and dynamically engineered throughout the computation [25,26], or the protocol is probabilistic [27,28]. Furthermore, the deleterious effects of natural spontaneous emission are always present. In contrast, here we consider local reservoirs and incoherent pumping that always couple to each qubit in the same way, and the computation is obtained by suitably observing the outcoming photons. Therefore, there is no need for direct action either on the qubits or on the couplings, and the computation is naturally protected from spontaneous emission. Note also that in our case the process as a whole is stochastic, but the computation is not probabilistic: Each realization produces one of many possible versions of the desired graph state, each version differing from the others by local Pauli operations that can be absorbed in the adaptation of measurement basis required by a given computational algorithm.

We have investigated the computational power of quantum trajectories and proved that one can implement any algorithm by suitably observing the photons that are spontaneously emitted and inelastically scattered by a multiqubit system. The results show that a natural reservoir that destroys quantum coherence when unobserved can be unravelled into quantum trajectories with full quantum computational power. From a fundamental point of view, one can interpret this scheme as a careful and efficient selection of the set of quantum trajectories that performs a desired computation out of the infinitude of possible unravellings allowed for that given environment interaction. Our scheme suits both circuit- and measurement-based quantum computation, and we believe that recent technological developments in 1D systems [29–31] will allow its fundamental blocks to be tested in the near future.

The concepts here presented may also significantly contribute to hybrid strategies to produce a quantum computer, for example, by combining the here introduced entangling gates with more traditional techniques to perform local operations. In that respect, the rather efficient way to

produce graph states by detecting the emitted and scattered photons is particularly encouraging.

It should be noted that, as in any scheme, the ability to perform quantum computation tasks will be degraded in the presence of imperfections. In our case, we need 100% detection efficiency to perform perfect gates or generate pure cluster states. For the two-qubit entangling gate of Fig. 1.3, we can follow the argument in Ref. [8] and calculate a 90% fidelity for a detection efficiency of 90%. Up to ten two-qubit gates could be performed with the same overall efficiency if one pushes towards the state-of-the-art single photon detectors that have recently reached efficiencies around 99% [32] at frequencies suitable for quantum optical experiments [33]. Naturally, with an increasing number of qubits we need more detectors, and the effect of inefficiencies will accumulate adversely. Two interesting directions of investigation appear in this context: the calculation of the effects of imperfections in the multiqubit scenario and the integration of standard or newly developed quantum error correcting codes to this computational approach.

Finally, the intrinsic random nature of the clicks indicates that the same scheme could be used in other interesting research directions such as the generation of random quantum circuits, t designs, and quantum walks, all useful methods to study many-body quantum systems.

This work was conceived and majorly developed during the ''Quantum Information'' session of the ''Centro de Ciencias de Benasque Pedro Pascual'' in Benasque, Spain. M. F. S. and M. T. C. thank CNPq and Fapemig for financial support. R. C. was funded by the Q-Essence project. This work is part of the INCT-IQ from CNPq, Brazil. It was also supported by the Australian Research Council Centre of Excellence for Quantum Computation and Communication Technology (Project No. CE110001027).

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