Quantum telecommunication with atomic ensembles

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Quantum mechanics provides a mechanism for absolutely secure communication between remote parties. For distances greater than 100 km, direct quantum communication via optical fiber is not viable, owing to fiber losses, and intermediate storage of the quantum information along the transmission channel is necessary. This leads to the concept of the quantum repeater, proposed by Briegel *et al.* [Phys. Rev. Lett. **81**, 169 (1998)]. Duan *et al.* [Nature **414**, 413 (2001)] have proposed to use atomic ensembles as the basic memory elements for the quantum repeater. We provide an overview of our program on the use of atomic ensembles as an interface for quantum information transfer and the prospects for long-distance quantum networks. © 2007 Optical Society of America

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1. INTRODUCTION

Quantum communication is concerned with the distribution of quantum states. An important application of quantum communication is the secure transmission of classical information between remote sites, via so-called quantum cryptographic key generation. The latter often involves the ability to entangle two distant qubits (two-level quantum systems).[1] Parametric downconversion is an established method for producing entangled photon pairs and enabling, for example, quantum teleportation of qubits.[2] Unfortunately, it is difficult to employ the approach of parametric downconversion over distances longer than a few absorption lengths (although a novel strategy of employing a satellite has been proposed,[3] and important experimental progress has been made).[4–6] As a result of the inevitable signal losses in optical fiber and the probabilistic nature of photon pair generation, the communication rate decreases exponentially with distance.

The concept of a quantum repeater was proposed to overcome this limitation and enable quantum communication over longer distances.[7,8] The idea is to insert quantum memory elements into the quantum channel every attenuation length or so and create qubits in each element. Entanglement between neighboring pairs of qubits can be generated efficiently, as light is not appreciably absorbed between them. Consider a serial network consisting of quantum memory elements connected by an optical fiber. After the entanglement of sequential pairs of atomic qubits has been established, appropriate joint measurements on neighboring internal qubit pairs are performed resulting in the entanglement of the outermost qubits at the remote ends of the network (entanglement swapping). The communication rate in this case scales polynomially with the distance.[7,9]

In general, a future quantum information network should therefore consist of spatially separated *nodes* to store and process quantum information and channels to connect the nodes. Atoms are excellent candidates for the storage and manipulation of qubits, because it is possible to isolate them from the environment and manipulate their internal states with laser light or external dc fields. Photons are ideal carriers of quantum information, because they can propagate over long distances in free space or in optical fibers. Quantum state transfer between photonic- and matter-based quantum systems forms the basis of such a scalable quantum network, as it enables the remote distribution of locally generated entanglement.

The realization of coherent quantum state transfer from a matter qubit onto a photonic qubit was achieved in 2004 using a cold Rb atomic ensemble.[10] It was quickly followed by several other significant advances: efficient generation of narrowband photon pairs deep in the regime of electromagnetically induced transparency,[11] Bell inequality violation between a collective atomic qubit and a photon,[12] storage and retrieval of single photons,[13] collapses and revivals of quantum memory,[14,15] electromagnetically induced transparency with single-photon pulses,[16] and light–matter qubit conversion and entanglement of remote atomic qubits.[17] A scheme to achieve long-distance quantum communication at the absorption minimum of optical fibers, employing atomic cascade transitions, has been proposed and its critical elements experimentally verified.[18] A deterministic single-photon source based on quantum measurement, quantum memory, and quantum feedback has been proposed and demonstrated.[19] Hong–Ou–Mandel interference of photon pairs from an ensemble has been observed.[20] In addition, important progress toward quantum networks was achieved in single-atom and single-ion experiments. In the microwave domain, single Rydberg atoms and single photons have been entangled.[21] An entangled state of an ion and a UV photon[22] and a neutral atom and a near-infrared photon[23] have been reported. In a broad sense, all these developments pave the way for the realization of a distributed network of atomic qubits, linear optical elements, and single-photon detectors.

The rest of this paper is organized as follows. We begin with an overview of the basic concepts involved in atomic ensemble-based quantum networks. We follow with a discussion of our approach to long distance quantum communication based on the production of entangled pairs of 1.53 μm and 780 nm photons using two-photon atomic cascade transitions. Next, we discuss two different methods of encoding a qubit in an atomic ensemble. We first describe our demonstration of an atomic qubit encoded in two independent atomic ensembles and then in a single ensemble. The latter provides the method for the entanglement of two qubits encoded in remote atomic ensembles. An important ingredient for a quantum network is the ability to perform certain actions based on the outcome of a measurement, i.e., quantum feedback. We discuss our protocol for a deterministic single-photon source based on an ensemble of atomic emitters, single-photon photoelectric measurements, and quantum feedback. We describe the implementation of this scheme using a cold Rb vapor. We conclude with an outlook for the future of quantum telecommunication with atomic ensembles.

2. BASICS OF THE DUAN–LUKIN–CIRAC–ZOLLER SCHEME

Collective enhancement of atom–photon interactions in optically thick atomic ensembles offers a relatively simple route toward quantum networks, as Duan, Lukin, Cirac, and Zoller (DLCZ) have demonstrated theoretically.[9] The DLCZ protocol is a probabilistic scheme based upon the entanglement of atomic ensembles via the detection of single-photon events in which the sources are intrinsically indistinguishable. The protocol generates long-distance qubit entanglement via the quantum repeater architecture.[7]

The basic mechanism proposed in Ref. 9 involves entangling a single photon (signal) with a single collective excitation of an atomic ensemble via the Raman scattering of a weak write pulse. Figure 1 illustrates the DLCZ protocol of photon pair generation for a simple three-level system.

The Raman scattered signal photon in this case is correlated with a collective atomic state $S^{\dagger}_{gg'}|0_a\rangle$. Here,

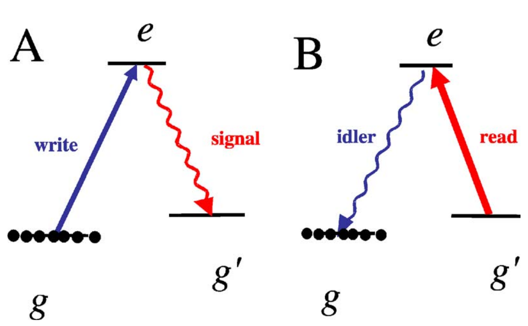

Fig. 1. (Color online) DLCZ protocol for the photon pair generation illustrating the write and read processes.

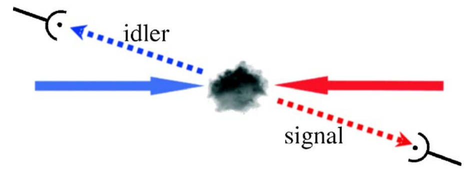

Fig. 2. (Color online) Schematic of the phase-matched off-axis geometry used to generate correlated photon pairs in the DLCZ scheme.

$$S_{gg'} = \sum_{i=1}^{N} |g\rangle_i \langle g'| e^{i\Delta\vec{k}\cdot\vec{r}_i}, \qquad (1)$$

where $\Delta\vec{k}=\vec{k}_1-\vec{k}_s$ is the difference between the first (write) and signal-photon wave vectors, $|0_a\rangle=\otimes_i|g\rangle_i$ is the atomic vacuum, and $\vec{r}_i$ is the position of the ith atom.

The quantum state of the collective atomic mode can be written as

$$|\phi\rangle = |0_a\rangle|0_p\rangle + \sqrt{p_c}S^{\dagger}_{gg'}a^{\dagger}|0_a\rangle|0_p\rangle + O(p_c), \qquad (2)$$

where p_c is the probability of scattering a signal photon into the correlated field mode. The detection of the signal photon results in the preparation of the atomic spin wave $S^{\dagger}_{gg'}|0_a\rangle$. By applying a second (read) pulse resonant on the $e\rightarrow g'$ transition, the spin excitation can be converted into another (idler) photon. The idler photon is strongly correlated with the signal photon when it is detected in a direction determined by the phase-matching condition: $\vec{k}_s+\vec{k}_i=\vec{k}_1+\vec{k}_2$, Fig. 2. Observations of nonclassical radiation produced in this manner were reported in first-generation atomic ensemble experiments.[24,25]

3. CASCADE SCHEME

The DLCZ approach to the quantum repeater relies on the atomic Raman transitions described in Section 2. Atomic and ionic Raman transitions in the UV to the near-infrared range have been successfully employed for entanglement generation,[10,12,17,22,23] making compact entanglement distribution over a range of a few kilometers conceivable. However, to extend the range of these schemes, it is vital to work at telecommunication wavelengths (1.3–1.5 μm). Unfortunately, it is difficult to identify suitable atomic or ionic Raman transitions at telecommunication wavelengths.

One possible approach to the long-distance quantum repeater could begin with the generation of narrowband, entangled photon pairs at the telecommunication wavelength using intracavity parametric downconversion.[26] Clearly, these photon pairs satisfy the need to distribute the entanglement over long distances, but it is essential to convert at least one of the entangled photons into an atomic qubit as part of the basic quantum repeater protocol described earlier. Coherent quantum state transfer between a telecommunication wavelength photon and an atomic memory would require frequency upconversion and subsequent mapping of the photon state onto an atomic qubit. Frequency upconversion of single telecommunication wavelength photons has been reported (e.g., Refs. 27 and 28 and references therein) but not under conditions of wavelength and bandwidth suitable for storage in an atomic memory.

A telecommunication wavelength quantum repeater can, however, be based on a two-photon cascade transition in alkali atomic ensembles.[18] The cascade transitions may be chosen so that the first photon (signal) emitted on the upper arm is in the telecommunication range, while the second photon (idler), emitted to the atomic ground state, is naturally suited for atomic storage.

The essential elements are illustrated in Fig. 3 and are as follows: (A) two-photon laser pulse excitation of an atomic s or d level; (B) cascade emission of two photons through intermediate p level(s). (C) Time-dependent variation of the amplitude of an auxiliary control laser field converts the idler qubit into a collective atomic qubit as discussed further below.[17] The net result is an entangled state of the signal field qubit and a long-lived collective atomic qubit.

Under conditions of collective enhancement, the atom returns to the same Zeeman state in which it was prepared,[12] and the probability the idler is emitted with wave vector $\vec{k}_i$ is weighted by the function $\varrho(\vec{k}_i+\vec{k}_s-\vec{k}_1-\vec{k}_2)$, which is the Fourier transform of the atomic density multiplied by the geometric mean of the pump beam's spatial profile and is defined as

$$\varrho(\vec{q}) \equiv \frac{1}{N}\sum_{\mu} \frac{\sqrt{I_1(\vec{r}_\mu)I_2(\vec{r}_\mu)}}{\bar{I}} e^{-i\vec{q}\cdot\vec{r}_\mu} = \left(\int d^3r \frac{\sqrt{I_1(\vec{r})I_2(\vec{r})}}{\bar{I}} \frac{n(\vec{r})}{N} e^{-i\vec{q}\cdot\vec{r}} \right) + O(1/\sqrt{N}), \quad (3)$$

where $I_1(\vec{r})$ and $I_2(\vec{r})$ are the intensities of pumps 1 and 2 at position $\vec{r}$, respectively, $n(\vec{r})$ is the atomic number density, $\vec{r}_\mu$ is the position of atom μ, and $\bar{I}$ is a characteristic intensity given by

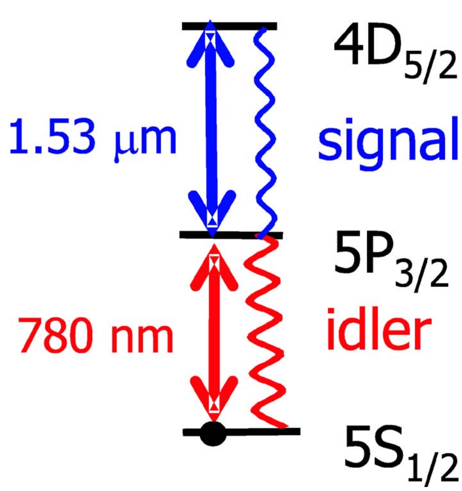

Fig. 3. (Color online) One of the atomic Rb level schemes for cascade emission involving two-photon excitation by pumps 1 and 2. The signal wavelength of 1.53 μm lies in the telecommunication wavelength range.

$$\bar{I} \equiv \int d^3r \sqrt{I_1(\vec{r})I_2(\vec{r})} \frac{n(\vec{r})}{N}. \quad (4)$$

Since the term in parentheses in Eq. (3) takes the maximum value of unity at $\vec{q}=0$, the idler photon is emitted into a distribution of plane-wave modes governed by the phase-matching condition $\vec{k}_i=\vec{k}_1+\vec{k}_2-\vec{k}_s$ and correlated with the return of the atom to its initial state. The width of this distribution depends on the overlap of the pump beams with the spatial distribution of the ensemble.

The fact that the atom begins and ends the absorption–emission cycle in the same state is essential for strong signal–idler polarization correlations. The corresponding reduced density operator for the field is given by[29]

$$\hat{\rho}(t \to \infty) \approx (1 + \sqrt{\epsilon}\hat{\Psi}_2^\dagger)\hat{\rho}_{\text{vac}}(1 + \sqrt{\epsilon}\hat{\Psi}_2), \quad (5)$$

where $\hat{\rho}_{\text{vac}}$ is the vacuum state of the field, $\hat{\Psi}_2^\dagger$ is a two-photon creation operator for the signal and idler fields, and $\epsilon \ll 1$. For linearly polarized pumps with parallel (vertical) polarizations, we find, for an unpolarized sample of ^{85}Rb,

$$\hat{\Psi}_2^\dagger = \frac{1}{\sqrt{5}}\hat{a}_H^\dagger\hat{b}_H^\dagger + \frac{2}{\sqrt{5}}\hat{a}_V^\dagger\hat{b}_V^\dagger, \quad (6)$$

where $\hat{a}_{H(V)}^\dagger$ and $\hat{b}_{H(V)}^\dagger$ are creation operators for a horizontally (vertically) polarized signal and idler photon, respectively.

Experimentally, we have demonstrated phase-matched cascade emission in an ensemble of cold Rb atoms using two different cascades: (a) at the signal wavelength $\lambda_s=776$ nm, via the $5s_{1/2}\rightarrow 5d_{5/2}$ two-photon excitation, and (b) at $\lambda_s=1.53$ μm, via the $5s_{1/2}\rightarrow 4d_{5/2}$ two-photon excitation. We have observed polarization entanglement of the emitted photon pairs and superradiant temporal profiles of the idler field in both cases.[18] With the asymmetry of the two-photon operator (6), Bell's parameter S is ideally expected to be ≈ 2.55, as opposed to $S=2.83$ for a completely symmetric state. The measured value $S=2.132\pm 0.036$ is reduced due to imperfections in the polarization optics and background counts.

For applications to quantum networks, it is necessary to perform certain interference measurements with the fields. Both signal and idler fields are labeled by polarization and wave-vector quantum numbers. High-quality interference fringes require the signal and idler wave packets to be factorizable in the latter quantum numbers.[30] This may be achieved with excitation pulses that are far detuned from two-photon resonance and with pulse lengths much shorter than the superradiant emission time t_s/d_{th} of level $|e\rangle$.

The idler field qubit is naturally suited for conversion into an atomic qubit encoded into the collective hyperfine coherence of levels $|a\rangle=|5s_{1/2},F=3\rangle$ and $|b\rangle=|5s_{1/2},F=2\rangle$. To perform such conversion, either the same or another similar ensemble–pair of ensembles could be employed, as discussed in the next section.

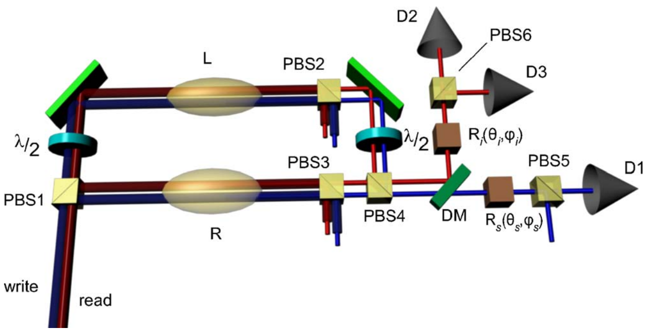

Fig. 4. (Color online) Schematic of the experimental setup used to encode a matter qubit in two atomic ensembles.

4. MATTER QUBITS

A. Matter Qubit: Two-Ensemble Encoding

The first demonstration of matter–light qubit conversion employed an atomic qubit encoded into two distinct atomic ensembles[10] (similar results were reported a year later in Ref. 31). As shown in Fig. 4, Raman scattering of a write pulse results in a single collective atomic excitation in one or the other of the two ensembles, correlated with orthogonal polarizations of the detected signal photon. These independent collective excitations form the basis states of a matter qubit, entangled with polarization degrees of freedom of the signal field. Subsequently, a read laser pulse is used to convert the matter qubit into an idler field qubit.

By observing polarization correlations of the signal and idler fields, we verified that, indeed, the signal photon and the atomic qubit were entangled, and that the matter–light qubit conversion fidelity exceeded classical bounds. An attractive feature of the two-ensemble qubit encoding is the associated capability to independently address the two basis qubit states for the purposes of quantum networking. This could lead to improved quantum communication rates compared with the single-ensemble qubit encoding discussed in Subsection 4.B.[9] On the other hand, because it is necessary to combine the light scattered from the spatially separated ensembles interferometrically, there is an inherent difficulty of phase stabilization, which may become acute for long-distance applications.

B. Matter Qubit: Single-Ensemble Encoding

The diagram in Fig. 5 illustrates schematically how one can use the ground-state degeneracy of alkali atoms to encode a qubit into orthogonal ground-state coherences of an atomic ensemble. The essence of the idea is to follow the DLCZ protocol in which a write laser field with σ^+ polarization scatters a σ^+ signal photon, thereby causing an atom to make a transition from the initially prepared $g, m=0$ hyperfine substate to the $g', m=2$ state of the other hyperfine ground level via level e, as shown in Figs. 1 and 5. The coherent superposition of the $m=0$ and $m=2$ states represents a single atomic excitation, which is, however, delocalized over the ensemble. The overall collective atomic excitation is a quasi-bosonic spin wave.[32] Alternatively, the write field results in the emission of a σ^- polarized signal photon (Fig. 5, lower panel), thereby causing an atomic transition between the $m=0$ components of the hyperfine ground levels. These two scattering processes are coherent, and the two orthogonal atomic excitations produced represent the basis states of a matter qubit.

The above arguments assume that the atoms are initially prepared in a single Zeeman state $m=0$ of hyperfine level g, whereas in practice, it is more convenient to use unpolarized atomic samples, in which all m states are equally populated (Fig. 1). In this case, perturbation theory shows that the ensemble-field density operator ρ may be written as $(1+\epsilon)^{-1}(1+\sqrt{\epsilon}\hat{\Psi}^{\dagger})\rho_{\mathrm{vac}}(1+\sqrt{\epsilon}\hat{\Psi})$, where $\hat{\Psi}^{\dagger}=\cos\eta\hat{a}_r^{\dagger}\hat{s}_{-1}^{\dagger}+\sin\eta\hat{a}_l^{\dagger}\hat{s}_{+1}^{\dagger}$, where $\varepsilon \ll 1$. Here, ρ_{vac} is the combined field and atomic vacuum density operator, the latter corresponding to an unpolarized ensemble of level g atoms. Here, $\hat{s}_{\pm1}^{\dagger}$ are collective atomic spin excitation operators, which for weak excitation, satisfy bosonic commutation relations, and $\hat{a}_{r,l}^{\dagger}$ are the creation operators for the right and left circular polarization of the detected signal field. The operator $\hat{\Psi}^{\dagger}$ clearly represents the creation of entanglement between atomic and photonic qubits. The mixing angle η is determined by the atomic couplings averaged over the unpolarized ensemble,

$$\cos^2\eta=\sum_m X_m^2(-1)\Bigg/\sum_m (X_m^2(-1)+X_m^2(+1)), \qquad (7)$$

with m summed over $\{-F_g,\ldots,F_g\}$, and $X_m(\alpha)=C_{m,1,m+1}^{F_g,1,F_e}C_{m+1,\alpha,m+\alpha+1}^{F_e,1,F_{g'}}$ is the product of the relevant Clebsch–Gordan coefficients for the transition. For our ^{85}Rb experiment $F_g=F_e=3$, $F_{g'}=2$, we find $\eta=0.81\times\pi/4$.[12]

Detection of a scattered signal photon results in the creation, in the ideal case, of exactly one atomic spin-wave excitation, and as Eq. (1) shows, the phase of the excitation is governed by the difference of the write and signal beam wave vectors. After a user-defined programmable delay time Δt (bounded by the lifetime of the ground-state atomic coherences), we convert the atomic excitation into a single photon by illuminating the atomic ensemble with a pulse of read light (Fig. 2). For an optically thick atomic sample, the idler photon will be emitted with high probability into the mode determined by the phase-matching condition $\vec{k}_i=\vec{k}_1+\vec{k}_2-\vec{k}_s$, with the atomic qubit state mapped onto a photonic one. Atomic

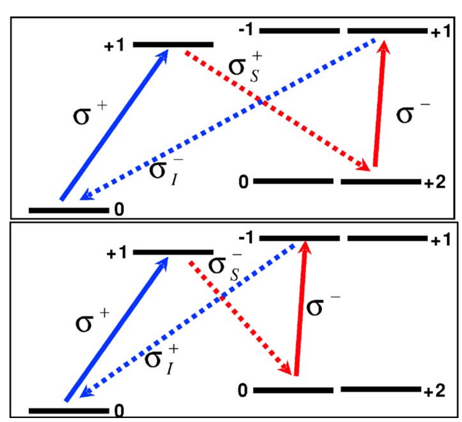

Fig. 5. (Color online) Illustration of atom–photon entanglement generation, where $\sigma_{S,I}^{\pm}$ represents the circular polarization states of the signal and idler fields, respectively.

excitations generated by $\hat{s}^{\dagger}_{\pm 1}$ map to orthogonal idler photon states up to a phase, and thus the number of correlated signal–idler counts registered by photoelectric detection provides information on atom–photon entanglement.

To infer probabilistic atom–photon entanglement, we have measured the degree of Bell inequality violation $|S| \leq 2$ (Refs. 33 and 34) and found the values $S = 2.29 \pm 0.05 \nleq 2$ (Ref. 12), and $S = 2.39 \pm 0.06 \nleq 2$.[17] These measured values of S are smaller than the ideal value of 2.77 due to nonzero counts in the minima of the interference curves. The phase of the latter curves is consistent with the theoretically predicted value of the mixing angle η.

C. Entanglement of Two Remote Matter Qubits

The entanglement of the atomic qubit and the signal field qubit discussed in subsection 4.B leads to the possibility of remote atomic qubit entanglement. The crucial new ingredient required is the conversion of the photonic qubit into a second atomic qubit, and thereby atom–photon entanglement is converted into remote atomic qubit entanglement.

In Fig. 6, we show schematically the experimental setup used to obtain remote atomic qubit entanglement using cold atomic clouds of ^{85}Rb confined at Sites A and B. These sites, separated by 5.5 m, are situated in separate laboratories and are linked by an optical fiber. An entangled state of a collective atomic qubit and a signal field is generated at Site A by Raman scattering of the write laser field. The orthogonal helicity states of the generated signal field are transmitted via an optical fiber from Site A to Site B, where they are converted to orthogonal collective atomic excitations, stored for a duration T_s, and subsequently converted into an idler field by the control field. The purpose of the control field is to subject the signal field propagation at Site B to the process of electromagnetically induced transparency (EIT). By adiabatically reducing the control field amplitude to zero, while the signal pulse lies within the cloud, the signal field is converted into a collective atomic excitation. To convert the signal field qubit into a collective atomic qubit, it is essential that the optically thick atomic sample supports EIT for both field helicities, requiring a judicious choice of the atomic level schemes. In our experiment, we optically pump the sample at Site B into the $F=2$, $m=0$ hyperfine ground state, while the sample at Site A was unpolarized as described in Subsection 4.B.

The atomic qubit at Site A is similarly converted into an idler field by a read laser pulse, counterpropagating with respect to the write pulse. By measuring polarization correlations of the idler fields generated at Sites A and B, it is possible to check for their entanglement. Since the fields are generated locally, a violation of the Bell inequality implies probabilistic entanglement of the remote atomic qubits. We find $S = 2.16 \pm 0.03 \nleq 2$, in clear violation of the Bell inequality. No corrections for background or dark counts were made to any of the experimental counting rates, and these are chiefly responsible for the reduction in the observed value of S from the ideal value of 2.60 predicted by our theoretical model.[29]

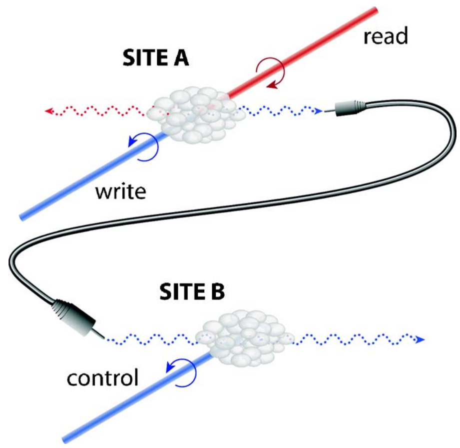

Fig. 6. (Color online) Illustration of the setup for the entanglement of two remote atomic qubits.

5. DETERMINISTIC SINGLE PHOTONS FOR QUANTUM NETWORKS

The production of single photons at specified times is an important resource for quantum information processing. Various single emitter sources of this type have been reported, including quantum dots,[35–37] color centers,[38,39] neutral atoms,[40–42] ions,[43] and molecules.[44] Related schemes using parametric downconversion have also been discussed.[45,46] Here, we describe our work on a deterministic single-photon source, which represents an important additional capability of atomic ensembles beyond reversible quantum state transfer between atoms and light described above.

A single photon can be generated at a predetermined time if one knows that the atomic ensemble contains an excitation. The presence of the latter is heralded by the measurement of a scattered photon in the write process. However, this is a probabilistic process; therefore there is a small probability $p_1 \ll 1$ of success in any given trial. Hence, a sequence of independent write trials must be performed before the excitation is heralded by the detection event. Subsequent to this event, it is necessary only to wait and to convert the excitation into the light pulse at a predetermined time. The repeated write trials and heralding measurements constitute a conditional feedback process, whose overall duration is limited by the atomic excitation's coherence time. The write trials are independent in the sense that the state of the atomic ensemble is reset after every failed write trial to avoid events from spurious atomic excitations.

The key signature of a deterministic single-photon source is sub-Poissonian statistics of the measured (unconditional) second-order coherence function $g_D^{(2)}$. The behavior of $g_D^{(2)}$ is a sensitive function of the number of write trials N in the feedback protocol and of the propagation and detection efficiencies of both the signal and idler fields. For this reason, long atomic memory time τ_c is essential. Thus, our protocol has two essential elements: (a) a high-quality probabilistic source of heralded photons, and (b) long atomic coherence times. In our experiment, we maximized τ_c by switching off the quadrupole coils of the magneto-optical trap and compensated for the ambient magnetic field with three pairs of Helmholtz coils.[12]

The measured value of $\tau_c \approx 31.5\ \mu s$ is limited by dephasing of different Zeeman components in the residual magnetic field.[13,15]

The long coherence time enables us to implement the conditional quantum evolution protocol. To generate a single photon at a predetermined time t_p, we initiate the first of a series of trials at a time $t_p - \Delta t$, where Δt is of the order of the atomic coherence time τ_c. Each trial begins with a write pulse. If D1 registers a signal photoelectric event, the protocol is halted. The atomic memory is now armed with an excitation and is left undisturbed until the time t_p when a read pulse converts it into the idler field. If D1 does not register an event, the atomic memory is reset to its initial state with a cleaning pulse, and an independent trial is repeated. The duration of a single trial $t_0 = 300$ ns. If D1 does not register a heralding photoelectric event after N trials, the protocol is halted 1.5 μs prior to t_p, and any background counts in the idler channel are detected and included in the measurement record. We emphasize that after each failed trial, one in which D1 does not register an event, the atomic state is reset to the vacuum by a cleaning pulse. We are therefore assured that after a series of failed trials, the state of the ensemble is in the vacuum, and any residual excitations corresponding to an undetected signal are wiped out.

A long atomic coherence time enables a large number of trials N, which is necessary to eliminate the classical fluctuations of the heralded single-photon source, and results in sub-Poissonian photon statistics. This can be explained most simply in the limit of infinite atomic memory in which the probability P_k that a detection event at Dk $=2,3$ (Fig. 7) is made after N trials is given by the product of the probability $1-p_{\mathrm{vac}}$ that an atomic excitation has been created (p_{vac} is the probability that the ensemble is in the vacuum state) and $p_{k|1}$ is the conditional probability that a photoelectric detection is registered on Dk given that a heralding event was recorded. Likewise, in the probability of a coincidence detection at D2 and D3, P_{23} is given in terms of the conditional joint probability $p_{23|1}$ for a coincidence given that a heralding event has been recorded, i.e.,

$$P_k = (1-p_{\mathrm{vac}})p_{k|1}, \tag{8a}$$

$$P_{23} = (1-p_{\mathrm{vac}})p_{23|1}, \tag{8b}$$

where $p_{\mathrm{vac}}=(1-p_1)^N$ for N trials, and $p_1 \ll 1$ is, as defined above, the probability of a signal photoelectric detection

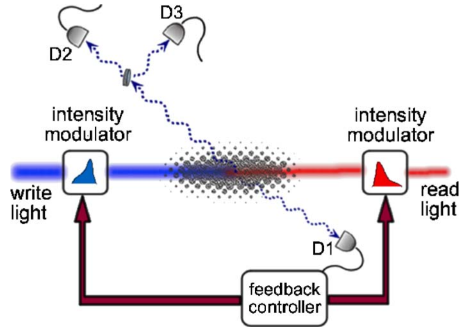

Fig. 7. (Color online) Schematic of the experimental setup for the generation of deterministic single photons.

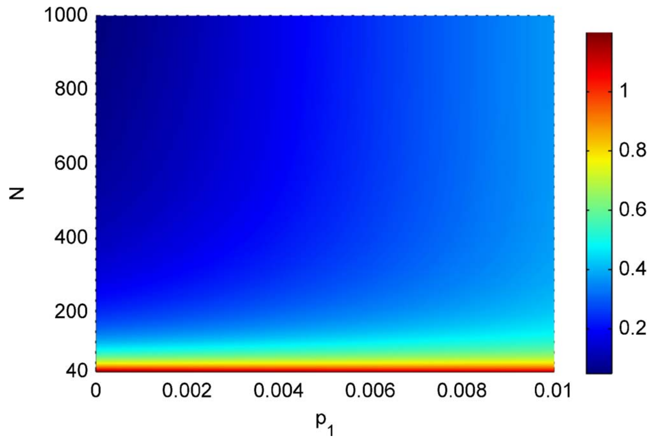

Fig. 8. (Color online) $g_D^{(2)}$ as a function of N and p_1 in the limit of infinite atomic memory.

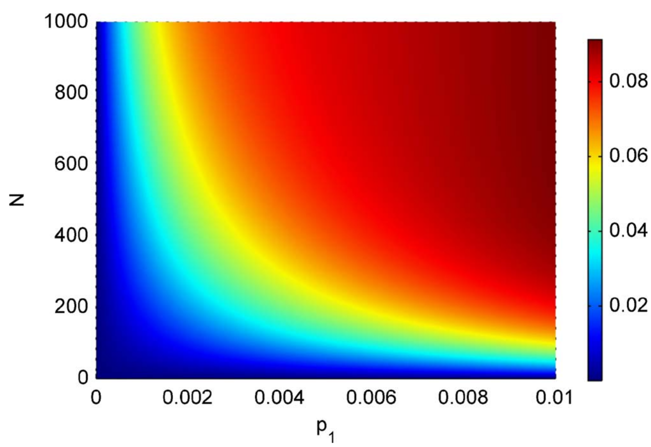

Fig. 9. (Color online) η_D as a function of N and p_1 in the limit of infinite atomic memory.

in a single write pulse trial. The calculation of the conditional probabilities and their dependence on p_1 and the propagation and detection efficiency is discussed in detail in Ref. 19. The correlation function $g_D^{(2)}$ is given by

$$g_D^{(2)} \equiv \frac{P_{23}}{P_2 P_3} = \frac{\alpha}{1-p_{\mathrm{vac}}}, \tag{9}$$

where $\alpha = p_{23|1}/p_{2|1}p_{3|1}$ is the anticorrelation parameter of Grangier *et al.*[47] When the probability p_{vac} is large, as is the case for a small maximum number of trials, N, and since the anticorrelation parameter is by definition positive $\alpha > 0$, the presence of multiple atomic excitations combined with the large vacuum component of the detected field gives super-Poissonian statistics. If N is sufficiently large, $p_{\mathrm{vac}} \rightarrow 0$, however, and we recover a value of $g_D^{(2)}$ identical to the anticorrelation parameter α, which is less than 1 for a heralded single-photon source. As a consequence, the statistics become sub-Poissonian in this limit, and the quality of the protocol is limited by the quality of the single photons produced by the heralded single-photon source (Fig. 8).

We should emphasize that $g_D^{(2)}$ does not tend to α in the limit of small N, so the feedback protocol is essential. In the limit of small N, one might expect instead that g_D^2

$\rightarrow 2$, consistent with a thermal distribution of atomic excitations. However, this expectation is correct only in the limit of perfect signal propagation and detection efficiency $\eta_s \rightarrow 1$. This can be readily demonstrated using the atomic density operator given in Ref. 19. However, for $\eta_s \ll 1$, one finds $g_D^{(2)} \gg 2$, as confirmed in our experiments.

The measured efficiency of a deterministic single-photon source η_D is defined as the probability to detect a single photon per trial of the protocol. We note that single-emitter deterministic single-photon sources typically have η_D less than 1% but in the best case may be a few times larger. In our implementation using cold Rb vapor, $\eta_D \approx 1-2\%$. In Fig. 9, we show theoretical plots of the total detection efficiency $\eta_D \equiv P_2 + P_3$ and $g_D^{(2)}$ as a function of N and p_1, for the case of an infinite atomic coherence time. Experimental data and theoretical results for finite coherence time are reported in Ref. 19.

6. CONCLUSIONS AND OUTLOOK

We have described some recent developments in the field of quantum communication with atomic ensembles carried out in our laboratory. These include the entanglement of 1.53 μm and 780 nm photons, the quantum state transfer between a matter qubit and a single-photon qubit, the entanglement of matter and light qubits, the entanglement of two remote atomic qubits, and deterministic single photons via conditional quantum evolution. This work provides basic elements for the construction of long-distance quantum networks, but before these can be realized, there are still outstanding challenges that need to be overcome.

Perhaps the most important challenge is to further increase the coherence time of one atomic memories, which is currently of the order of 30 μs.[19] Toward this end, one may implement the two-ensemble encoding of atomic qubits[9,10] using an atomic clock transition, which is insensitive to the magnetic field of first order. To eliminate the effects of atomic motion, one could confine the atoms in 3D optical lattices. These developments could result in coherence times of tens, or even hundreds, of milliseconds.

The ensemble approach has an intrinsic advantage for the rate of entanglement distribution over single-atom approaches as it should be possible to encode a large number of atomic qubits into a single cold-trapped atomic sample, with each one individually addressable. For the case of a single atomic qubit per node, the throughput is limited by the classical communication time between different quantum nodes ($\sim Ln/c$) where L and n are the length and the refractive index of the optical fiber, and c is the speed of light in vacuum. Assuming a cold sample of a 400 μm cross section and a laser beam waist of 20 μm, storage of over 100 qubits in a single atomic sample is conceivable. This could be achieved by fast (submicrosecond) 2D scanning using acousto-optic modulators, which allow for the coupling of the qubits to the same single-mode optical fiber. This should enable the multiplexed operation of quantum network protocols, vastly increasing the throughput of quantum channels in the low-memory time limit.[48] The number of atomic qubits per node can be increased to between 4 and 6 using conventional optical hardware. Larger ($\sim$100) qubit numbers should be possible using more elaborate, e.g., holographic, beam-shaping techniques similar to those that are currently being developed for single-trapped Rydberg atom approaches.[49]

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