‘Which-way’ collective atomic spin excitation among atomic ensembles by photon indistinguishability

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\textit{New Journal of Physics} \textbf{14} (2012) 063034 (11pp)

Received 15 November 2011
Published 25 June 2012
Online at http://www.njp.org/
doi:10.1088/1367-2630/14/6/063034

\textbf{Abstract.} In spontaneous Raman scattering in an atomic ensemble, a collective atomic spin wave is created in correlation with the Stokes field. When the Stokes photons from two or more such atomic ensembles are made indistinguishable, a ‘which-way’ collective atomic spin excitation is generated among the independent atomic ensembles. We demonstrate this phenomenon experimentally by reading out the atomic spin excitations and observing interference between the read-out beams. When a single-photon projective measurement is made on the indistinguishable Stokes photons, this simple scheme can be used to entangle independent atomic ensembles. Compared to other currently used methods, this scheme can be easily scaled up and has greater efficiency.
1. Introduction

Over the last decade, there have been many proposals and experimental realizations of quantum nodes involving atomic ensembles for information storage and processing in a practical quantum information network [1]. One basic element in these protocols is an entangled system involving two distinct atomic ensembles. The very first scheme for entangling two atomic ensembles is the so-called DLCZ scheme [2, 3], which is based on the complementary principle of quantum mechanics [4] and the quantum eraser concept for remote control [5, 6]: the detection of a photon erases the which-way information and puts the remote quantum systems in the entangled state.

Thereafter, many other schemes have been realized based on variations of the DLCZ scheme [7, 8]. More recently, such a strategy was used to entangle four separate atomic ensembles [9]. Furthermore, the scheme of light storage based on electromagnetically induced transparency [10, 11] was put in use to entangle remote atomic ensembles by converting photonic entanglement into atomic entanglement [12, 13].

One of the original proposed applications of the DLCZ scheme [2] is that when scaled up, it can be used as quantum repeaters to extend the quantum communication range. The DLCZ scheme, although clear and straightforward, relies on superposition by beam splitters and photon detection to project out the entangled state: there is no entanglement between the atomic ensembles before the superposition and projection. The employment of beam splitters adds a degree of complexity and reduces the efficiency, especially when scaled up [9].

In this paper, we present a new scheme for creating entanglement between two distinct atomic ensembles. We experimentally demonstrate the establishment of coherence between two atomic ensembles, a first step toward entanglement. The scheme is a remake of a mind-boggling experiment on photon interference [14], suggested by one of the authors of this work. In that experiment, the signal fields from two independent parametric down-conversion processes become coherent due to photon indistinguishability in idler fields. Here, we replace the parametric down-conversion process by a Raman scattering process and the signal photons by collective atomic spin excitations. Which-way information in atomic excitation is automatically erased and coherence is established when the corresponding Stokes photons are made indistinguishable. A similar strategy was used to entangle two macroscopic atomic ensembles in continuous variables [15, 16] as well as discrete variables [17], and the same physical principle was exhibited for atomic matter waves in [18]. Compared to the quantum eraser scheme of Duan et al [2], our scheme utilizes the same atom–photon correlation in the Raman process but completely eliminates beam splitters for the superposition of the
Figure 1. (a) Conceptual sketch for entangling two atomic ensembles (A₁, A₂). W₁, W₂: the initial Raman pump fields; S₁̂, S₂̂: the Stokes fields whose indistinguishability creates atomic entanglement between A₁ and A₂. (b) Energy levels for the atoms in A₁ and A₂; g: the ground state; m: the meta-stable state; e: the excited state.

Stokes fields. A projective measurement is not required to establish coherence among atomic ensembles. But heralding the detection of the Stokes photon is still needed to project out the entangled state among the independent atomic ensembles and to reduce the effect of losses. The scheme can be easily scaled up without suffering efficiency reduction.

2. Theoretical background

Consider the conceptual sketch in figure 1(a). Atoms in ensembles A₁ and A₂ have the energy levels shown in figure 1(b) with all the atoms prepared in the ground state g. We are interested in the Raman scattering process in atomic vapor with fields W₁, W₂ as the Raman write (pump) fields. The wiggled lines denote the generated Stokes fields. For a single ensemble of a large collection of N atoms, when the detuning Δ is large compared to the line width of level e, the coupling between the strong Raman pump field W and the Stokes field S is through the collective atomic spin excitation \( \hat{S}_{mg} = \frac{1}{\sqrt{N}} \sum_i |g\rangle \langle m_i| \). Here |g\rangle is the atomic ground state with all atoms in the g-state, i.e. |g\rangle ≡ |g...g\rangle and |m_i\rangle is the state with only the \( i \)th atom in the m-state and the remaining in the g-state, i.e. |m_i\rangle ≡ |g...g\rangle |m_i\rangle. The state \( (1/\sqrt{N}) \sum_i |m_i\rangle \) is a collective atomic state with only one atom in the m-state. The interaction Hamiltonian is given by the equation [2]

\[
\hat{H}_{SR} = i\hbar \eta \alpha_W \hat{a}_{S_w} \hat{S}_{mg}^\dagger - i\hbar \eta \alpha_W^* \hat{a}_{S_w} \hat{S}_{mg},
\]

where the strong write field is treated as a classical field with amplitude \( \alpha_W \). \( \eta = g_{eg} g_{em}^* \frac{\sqrt{N}}{\Delta} \) with \( \Delta \) being the detuning of light fields from the upper excited state e shown in figure 1(b) and \( g_{eg}, g_{em} \) the coupling coefficients between light fields and respective atomic states. Under the condition of a large atom number N and weak excitation \( (|\langle \hat{S}_{mg}^\dagger \rangle| \ll 1) \), the operators \( \hat{S}_{mg}^\dagger, \hat{S}_{mg} \) can be treated as the creation and annihilation operators for the collective atomic excitations. So, the state of the system after the interaction in perturbative expansion is given by

\[
|\psi\rangle \approx |0\rangle_A |0\rangle_{S_w} + \sqrt{p_c} |1\rangle_A |1\rangle_{S_w} + o(p_c),
\]

where \( |0\rangle_A = |g\rangle \) is the atomic ground state and \( |1\rangle_A \equiv |\hat{S}_{mg}^\dagger |0\rangle_A = (1/\sqrt{N}) \sum_i |m_i\rangle \) is the one-atom collective excited state. \( p_c \ll 1 \) is the probability of Raman scattering with \( p_c \propto |\alpha_W|^2 \). \( |1\rangle_{S_w} \) is a single-photon state for the Stokes field. Higher-order terms are ignored because \( p_c \ll 1 \). With two such processes in two independent atomic ensembles, we have an entangled state...
state between the Stokes photon and the collective atomic excitation in the sub-space of one photon or one atomic excitation:

\[
|\Psi\rangle = |\psi\rangle_1 \otimes |\psi\rangle_2 = \cdots + \sqrt{p_c}(|1\rangle_{A_1}|0\rangle_{A_2}|1\rangle_{SW_1}|0\rangle_{SW_2} + |0\rangle_{A_1}|1\rangle_{A_2}|0\rangle_{SW_1}|1\rangle_{SW_2}) + \cdots
\]

(3)

This state is the starting point for the DLCZ scheme [2] of entangling two atomic excitations. Normally, there is no coherence or entanglement in equation (3) between the two atomic ensembles before the projective measurement in the DLCZ scheme. The projective measurement of one Stokes photon collapses the state of the system into the sub-space of one photon or one atomic excitation described in equation (3). On the other hand, if we make the two states \( |1\rangle_{SW_1}, |1\rangle_{SW_2} \) indistinguishable, i.e. \( |1\rangle_{SW_1}|0\rangle_{SW_2} = |0\rangle_{SW_1}|1\rangle_{SW_2} \equiv |1\rangle_{SW} \), equation (3) becomes

\[
|\Psi\rangle = |\text{vac}\rangle + \sqrt{2p_c}|E\rangle_{A}|1\rangle_{SW} + \cdots
\]

(4)

where \( |\text{vac}\rangle \equiv |0\rangle_{A_1}|0\rangle_{A_2}|0\rangle_{SW_1}|0\rangle_{SW_2} \) and

\[
|E\rangle_{A} = (|1\rangle_{A_1}|0\rangle_{A_2} + |0\rangle_{A_1}|1\rangle_{A_2})/\sqrt{2}
\]

(5)

is an entangled state of atomic excitation of two atomic ensembles \( A_1 \) and \( A_2 \).

To make \( |1\rangle_{SW_1}, |1\rangle_{SW_2} \) indistinguishable, we can align the two Stokes fields \( SW_1, SW_2 \) together, as shown in figure 1(a), so that when a photon is detected, it is impossible to tell which cell emits it. This indistinguishability will put the atoms in an entangled state of the form in equation (5) in the sub-space of one excitation. In fact, the which-way information is erased even without actual detection of the Stokes photon: from equation (4), we see that the superposition is established irrespective of what happens to the Stokes photon. This indistinguishability-induced coherence was first discovered by Zou et al [14]. Here we replace the role of some photons with atoms.

However, caution should be exercised in claiming entanglement for the state of the entire atomic system. The existence of higher-order terms, even with a much smaller probability, may hinder the creation of entanglement between the two atomic ensembles [3, 19]. As a matter of fact, without altering the probability distribution of various numbers of atomic excitations in the current scheme, e.g. reducing the probability of multiple atomic excitations, it is impossible to achieve entanglement between the two atomic ensembles [3, 19]. One effective way of altering the probability distribution is to make a projective measurement onto the single-excitation sub-space, in a similar way as the DLCZ scheme. Such a projective measurement can be achieved by heralding the detection of a Stokes photon.

3. Experiment

Next, we describe an experiment implementing some of the above ideas. The detailed experimental arrangement is shown in figure 2(a). \( C_1, C_2 \) are 75 mm long isotopically enriched Rb-87 vapor cells without buffer gas or anti-spin-relaxation coating (paraffin). Such a cell has a relaxation time of about 3 \( \mu s \). They are heated to 72 °C using a bi-filar resistive heater. The cells were respectively mounted inside a three-layer magnetic shielding to reduce the stray magnetic fields. The energy level diagram is shown in figure 2(b). The two lower-energy levels (levels \( g, m \)) are the hyperfine splitting of the ground state: \( 5^2S_{1/2} (F = 1; 2) \). The upper excited level (level \( e \)) is \( 5^2P_{1/2} (F = 2) \). The two Raman write fields \( W_1, W_2 \) are split from a diode laser.
Figure 2. (a) Detailed experimental arrangement for establishing coherence between two atomic ensembles ($C_1$, $C_2$) and the subsequent reading of the coherence. $W_1$, $W_2$: the initial Raman pump fields; $S_{W1}$, $S_{W2}$: the Stokes fields whose indistinguishability creates atomic coherence between $C_1$ and $C_2$. $R_1$, $R_2$: the readout Raman pump fields; $S_{R1}$, $S_{R2}$: the read Stokes fields that carry coherence information of the atomic excitations in $C_1$ and $C_2$. SMF: single-mode fiber; H: half-wave plate; Q: quarter-wave plate; PZT: piezoelectric transducer. All cubes are polarization beam splitters (PBS). (b) Energy levels for the atoms in $C_1$ and $C_2$. $g$: the ground state ($5^2S_{1/2}$, $F = 1$); $m$: the meta-stable state ($5^2S_{1/2}$, $F = 2$); $e$: the excited state ($5^2P_{1/2}$, $F = 2$). (c) Time sequence. OP: optical pumping pulse.

emitting at 795 nm. The power of the write fields ranges from 0.3 to 3 mW, depending on the experimental situation. The detuning $\delta$ is set at 1 GHz for maximum Raman gain. The two read beams $R_1$, $R_2$ are from another diode laser at 795 nm. The detuning $\Delta$ is set at –1.2 GHz. The time sequence of the application of the fields is shown in figure 2(c). The atoms are initially prepared in the ground state ($g$) by an OP pulse of 80 $\mu$s length. The OP beam has a power of 100 mW, is vertically polarized and crosses the write beam at an angle of 2° in the middle of the cell. Then after approximately $\tau = 0.1 \mu$s delay, the write pulses of 1 $\mu$s length are applied to create collective atomic spin excitations in the two atomic ensembles. The Stokes field ($S_{W1}$) from the first cell is sent into the second one to make it indistinguishable from the second Stokes field ($S_{W2}$). In this way, coherence will be established between the two atomic ensembles. A silicon fast diode detector ($D_2$) is used to monitor the transmitted Stokes field ($S_{W2}$). Note that the combinations of PBS and wave plates before and after cell 2 are intended to act as beam splitters to couple in and out the read, write and Stokes fields for cell 2, because all the fields are along the same line and just opposite to each other. So, there will be some parts missing in all the fields.

Confirmation of the atomic coherence is achieved by coherently converting the atomic spin excitations into photons and observing the interference. Traditionally, this is done by a strong reading field, tuned on resonance between level $m$ and some excited state to completely empty level $m$ and convert it into anti-Stokes photon. Here we use a recently discovered effect [20] that the initially prepared atomic spin wave can significantly enhance the Raman conversion from the pump into Stokes. The converted Stokes carries information on the coherence of the atomic spin wave [21]. One advantage is that, because this process is a Raman amplification process,
the converted light field is much stronger than the traditional anti-Stokes field and easier to detect. The drawback is the existence of spontaneous noise to degrade the observed coherence.

To achieve this, we send in the read pulse to obtain Stokes fields \( S_{R1}, S_{R2} \), enhanced by the atomic spin excitations built by \( W_1, W_2 \). The read fields have a power ranging from 4 to 10 mW, depending on the strength of the spin excitations. The read Stokes fields \( S_{R1}, S_{R2} \) now carry coherence information about the atomic spin excitations prepared by \( W_1, W_2 \). We combine the two amplified Raman Stokes fields \( S_{R1}, S_{R2} \) with an ensemble of PBS and wave plates and couple the combined beam into a single-mode fiber before sending it to a detector \((D_1)\) similar to \( D_2 \). The PBS/wave plate ensemble allows us to adjust the relative intensities of \( S_{R1}, S_{R2} \) for maximum visibility of the interference. The relative phase between \( S_{R1}, S_{R2} \) is adjusted by a PZT. In figure 3, we plot the detector output as a function of time as the PZT voltage is ramped at about 10 Hz repetition rate. A clear interference fringe is observed with a visibility of 50%. The red solid curves are a sinusoidal fit to the data. The middle irregular part is due to ramping voltage return. This interference fringe is a proof of coherence between two atomic ensembles, as shown in equation (5).

The reason why we observed a lower visibility than the perfect 100% is mainly because of the mode mismatch between \( S_{W1} \) and \( S_{W2} \) and losses for \( S_{W1} \) to travel from the first cell to the second. To show this, we model the loss and mode mismatch (mode mismatch is equivalent to a loss) by a beam splitter with amplitude transmissivity \( \zeta \) between \( S_{W1} \) and \( S_{W2} \). Then the single-photon state in mode \( S_{W1} \) is transferred to mode \( S_{W2} \) by

\[
|1\rangle_{S_{W1}} = \xi |1\rangle_{S_{W2}} + \sqrt{1 - |\xi|^2} |1\rangle_b,
\]

where \( b \) is the loss output mode of the beam splitter. From equation (3), we then have

\[
|\Psi'\rangle = |\text{vac}\rangle + (\xi \sqrt{p_{c1}} |1\rangle_{A_1} + \sqrt{p_{c2}} |1\rangle_{A_2})|1\rangle_{S_{W2}} + \sqrt{1 - |\xi|^2} \sqrt{p_{c1}} |1\rangle_{A_1} |1\rangle_b \\
= |\text{vac}\rangle + \xi \sqrt{2 p_{c1}} |E\rangle_{A_1} |1\rangle_{S_{W2}} + \sqrt{1 - |\xi|^2} \sqrt{p_{c1}} |1\rangle_{A_1} |1\rangle_b.
\]

Here we assume that the probability \( p_c \) is different for the two cells and, in the second equation, we set \( \xi \sqrt{p_{c1}} = \sqrt{p_{c2}} \).

The visibility of the interference between \( S_{R1} \) and \( S_{R2} \) is directly related to the degree of first-order coherence via the collective spin excitations \( \hat{S}_{mg1}, \hat{S}_{mg2} \) between \( A_1 \) and \( A_2 \) for the

\[
\text{Figure 3. The observed interference between two Stokes fields from two reading Raman processes.}
\]
state $|\Psi'\rangle$ in equation (7) [21]:

$$V = |\gamma_{A_1A_2}| = |\langle \hat{S}_{mg1}^\dagger \hat{S}_{mg2} \rangle \psi| = |\langle 1_{SW_1} 1_{SW_2} \rangle | = |\xi|. \quad (8)$$

The maximum visibility is one if $S_{W1} \equiv S_{W2}$ or we cannot distinguish between $S_{W1}$ and $S_{W2}$. However, loss in $S_{W1}$ as well as mode mismatch between $S_{W1}$ and $S_{W2}$ makes $S_{W1}, S_{W2}$ distinguishable from each other, resulting in a reduction of visibility to $|\xi|$. In the experiment, we introduce an extra loss by a polarizer and control it by rotating the polarization of the $S_{W1}$ photon. The rotator together with a PBS acts like a beam splitter with variable transmissivity. Figure 4(a) plots the observed visibility as a function of the amplitude transmissivity $\zeta$. It can be seen that the data follow very well a linear dependence, as predicted in equation (8).

The phenomenon of induced coherence without induced emission is a purely quantum mechanical behavior, based on the fundamental quantum mechanical principle of indistinguishability. In classical theory, the coherence can only be induced by stimulated emission, which is the principle behind the coherence property of a laser. When the injected signal from $S_{W1}$ is large, the stimulated emission will dominate and it enters the classical regime. This can be done by tuning up the intensity of the write laser $W_1$. Figure 4(b) shows the measured visibility as a function of the amplitude transmission between $S_{W1}$ and $S_{W2}$. It shows a significantly different shape from that in figure 4(a) for the quantum case. We cannot use the perturbative treatment as before. On the other hand, the evolution equations of the Hamiltonian in equation (1) are the same as a parametric amplifier and have the form

$$\hat{S}_{mg1}^{\text{out}} = G_1 \hat{S}_{mg01} + g_1 \hat{a}_{SW01}^\dagger, \quad \hat{a}_{SW1} = G_1 \hat{a}_{SW01} + g_1 \hat{S}_{mg01}^\dagger$$

and for the second Raman process

$$\hat{S}_{mg2}^{\text{out}} = G_2 \hat{S}_{mg02} + g_2 \hat{a}_{SW1}^\dagger. \quad (10)$$

Here $G_{1,2}, g_{1,2}$ are the gain parameters. Now we add some losses between the Stokes field $\hat{a}_{SW1}$ generated in the first process and the input to the second process ($\hat{a}_{SW1}$). As before, this loss can be modeled as a beam splitter with amplitude transmissivity $\zeta$: $\hat{a}_{SW1} = \zeta \hat{a}_{SW1} + \hat{b}_0 \sqrt{1 - |\zeta|^2}$ with $\hat{b}_0$ being the vacuum input of the beam splitter.

Figure 4. Visibility versus the amplitude transmissivity $\zeta$ from $S_{W1}$ to $S_{W2}$ in (a) the quantum and (b) classical cases.
Figure 5. (a) Schematics for the arrangement of collinearly propagating the write and Stokes fields to the second cell in order to reduce the loss and improve the mode match. PS: the phase shifter; PZT: phase scan. (b) Interference fringe shift for various phase differences between write and Stokes fields: (i) original, (ii) $\pi/2$, (iii) $\pi$ and (iv) $3\pi/2$.

With the initial atomic system in the ground state $|g\rangle$ and $\hat{a}_{Sw1}$ in vacuum, we can deduce the degree of coherence between the two atomic systems $\hat{S}_{mg1}$ and $\hat{S}_{mg2}$ as

$$V_c = |\gamma_{A_1A_2}| = \sqrt{|\zeta|^2(1+n_{Sw1})/(1+|\zeta|^2n_{Sw1})},$$

where $n_{Sw1}$ is the average photon number in $Sw_1$. Note that equation (11) returns to equation (8) when $n_{Sw1} \ll 1$ in the quantum regime. The dashed line in figure 4(b) is a plot of $V_c$ as a function of the nominal transmissivity $\zeta$ with $n_{Sw1} = 70$. The actual transmissivity applied in equation (11) needs to multiply a factor $L = 0.65$ due to losses and mode mismatch as in figure 4(a). Furthermore, when $n_{Sw1}$ is large, the corresponding spin wave $\hat{S}_{Sw}$ experiences a large fluctuation, which results in a random readouts in the intensities of $Sr1$ and $Sr2$. Because the visibility of the interference between two fields highly depends on the intensity ratio, the fluctuations in the intensities of $Sr1$ and $Sr2$ will lead to a reduced visibility, which explains the lower values of the observed visibility in figure 4(b) compared to the dashed theoretical curve from equation (11). The solid curve fit to the experimental data is a reduction by a factor of 0.52 from the dashed line.

4. The collinear scheme and projection measurement

As demonstrated in figure 4, loss or equivalently mode mismatch is a big problem in this scheme. The low visibility is because of the fact that we separated in figure 2 the write and Stokes beams in order to introduce an extra loss and control it. If we eliminate the control elements between the two ensembles, we can propagate both write and Stokes beams through the atomic ensembles collinearly, as shown in figure 5(a). This reduces the loss and increases the mode match for higher visibility. The results are shown in figure 5(b) when the system is in the high-gain regime. The best observed visibility is nearly 90%.

Another advantage of the collinear scheme in figure 5(a) is the stable phase relationship between the atomic ensembles because it is determined by the phase difference between the...
write and Stokes beams, which are now in the same path so that any minute change in the path will not effect this phase difference. So the phase is stable even for remotely separated atomic ensembles. This basically avoids the challenging technical issues of stabilizing the phases in the DLCZ scheme [3, 7, 9]. To change this phase, we can introduce a birefringent element (PS in figure 5(a)) in the path of the write and Stokes fields, which have orthogonal polarizations and are subject to different phase shifts through the birefringent element. Figure 5(b) shows how the interference fringe shifts accordingly as we insert a quarter (ii), a half (iii) or both wave plates (iv), which correspond to $\pi/2, \pi, 3\pi/2$ phase shift, respectively.

Furthermore, because the write and Stokes fields have a $\Delta \nu = 6.8$ GHz frequency difference, the phase difference can also be introduced by a large path change $\delta L$: $\Delta \psi = (k_W - k_{SW}) \Delta L = 2\pi \Delta \nu \delta L/c$. So, for a phase difference of $2\pi$, the path difference is $\delta L = c/\Delta \nu = 44$ mm or a stage displacement of $\Delta x = \delta L/2 = 22$ mm (figure 5(a)). Figure 6(a) shows the interference fringe shift with respect to the fringe from the read fields, which are leaked from the PBS. The leaked read fields are about 1% of the $S_R$ signal. But the read pulse lasts 10 $\mu$s, while the $S_R$ signal only lasts 1 $\mu$s. So the $S_R$ signal sits on top with a long tail of the leaked read field signal. We can adjust the relative strength of the $S_R$ signal and the leaked read signal with the PBS and the half-wave plate before detector $D_1$ (figure 5(a)). From the temporal shape analysis of the detected pulses, we can thus separate the signals contributed by the $S_R$ field and the leaked read field and respectively record their interference fringes as in figure 6(a). One thing to note from figure 6(a) is that the $S_R$ interference fringe is not as stable as the read fringe even though the $S_R$ signal is much larger than the leaked read signal. This is because the Stokes field is converted from the read field by Raman scattering from the atomic spin excitation. This makes the Stokes field less stable than the read field which is directly from the laser. The phase shifts are then extracted as $(t_1 - t_2) \times 360/T$ from figure 6(a). We plot the phase shift as a function of the stage displacement $\Delta x$ in figure 6(b), from which we find that $\Delta x = 21 \pm 1$ mm for a phase shift of $360^\circ$, in agreement with the predicted value.

The scheme in figure 5 is easily scaled up to multiple atomic ensembles: we can simply put them in series, as shown in figure 7. Projective detection of the Stokes photon is achieved at the end of the series. The effect of loss can be overcome by the detection of a Stokes photon.

Figure 6. Phase shift by a large path change: (a) fringe shift (solid circles) with reference to the read with interference fringe (red solid curve) for the extraction of phase shift $= (t_1 - t_2) \times 360/T$; (b) phase shift as a function of the stage displacement. The solid red line is a linear fit.
Arrangement for scaling up to an arbitrary number of ensembles.

This can be seen in equation (7)—detection of a Stokes photon (|1⟩_S^2) will project the state to the entangled state (|E⟩_A). The loss only reduces the successful probability of heralding but does not affect the fidelity of the entanglement because the detection of Stokes photon ensures that it is not lost. Furthermore, compared with the DLCZ scheme, which suffers efficiency reduction when scaled up because of the employment of beam splitters for superposition [9], our scheme eliminates the beam splitters and thus the efficiency can reach 100% in principle even when it is scaled up, provided that caution is exercised to get rid of losses in propagation. On the other hand, as we demonstrated at the end of section 3, the overall gain of the process must be low to ensure the spontaneous nature of the process. This means that the photon production rate cannot be high for each individual cell. This may be a problem when we scale up the system. The DLCZ scheme does not have this problem.

Because our photo-detectors are not single-photon detectors, we are not able to perform the projective detection of the Stokes photons and to confirm the entanglement. Work is under way to push to the single-photon regime of operation.

5. Summary and discussion

In summary, we proposed a simple scheme for entangling atomic ensembles by photon indistinguishability. We experimentally demonstrated its feasibility by observing interference fringes. Compared to the commonly used DLCZ scheme, our scheme does not need beam splitters for superposition and thus is simpler in structure and stable in phase. This will lead to an easier way of scaling up and to a better efficiency in projective measurement for achieving entanglement. It should be noted that although we have observed coherence between the two atomic ensembles, this does not confirm the entanglement between the two ensembles. Observation of coherence is the first step towards the confirmation of entanglement, which requires the projection measurement as shown in figure 7 and a different set of criteria [3, 9].

Acknowledgments

This work was supported by the National Basic Research Program of China (973 Program grant no. 2011CB921604), the National Natural Science Foundation of China (grant numbers 11004058, 11129402, J1030309, 10828408 and 10588402) and the Program of Shanghai Subject Chief Scientist (grant number 08XD14017), the fundamental research funds for the central universities.

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