Quantum memory, entanglement and sensing with room temperature atoms

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Abstract. Room temperature atomic ensembles in a spin-protected environment are useful systems both for quantum information science and metrology. Here we utilize a setup consisting of two atomic ensembles as a memory for quantum information initially encoded in the polarization state of two entangled light modes. We also use the ensembles as a radio frequency entanglement-assisted magnetometer with projection noise limited sensitivity below femtoTesla/√Hz. The performance of the quantum memory as well as the magnetometer was improved by spin-squeezed or entangled atomic states generated by quantum non demolition measurements. Finally, we present preliminary results of long lived entangled atomic states generated by dissipation. With the method presented, one should be able to generate an entangled steady state.

1. Quantum Memory
A quantum memory for light should have two properties (i) it should be able to store a wide range of input states, and (ii) it should be able to store entanglement such as one part of an Einstein-Podolski-Rosen (EPR) pair. We demonstrate a quantum memory for a range of input states (an alphabet) which are squeezed by 6 dB and displaced in phase space [1]. A classical benchmark fidelity is calculated, and we find that the experimental fidelity is higher than the benchmark which shows that the memory is truly a quantum memory. This is an advance with respect to earlier experiments where the storage of coherent states was demonstrated in the same setup [2]. Although a memory for squeezed vacua had been realised in other setups [3, 4], the classical fidelity benchmark was never beaten in such previous attempts [5]. Before discussing our experimental results we will first define canonical variables for atoms and light which is the common language used to describe information transfer between atoms and light.

In the experiment we utilize two ensembles of cesium atoms kept in separate paraffin-coated glass cells. At room-temperature each cell contains around $N_A \approx 10^{11} - 10^{12}$ atoms. The atoms
are located in a bias magnetic field \( B = 0.92 \) Gauss pointing in the \( x \)-direction. Each ensemble can be described by the total spin vector \( \vec{J} = (J_x, J_y, J_z) \). The two ensembles (1 and 2) are optically pumped into opposite directions with macroscopic spin projections \( J_{x1} = -J_{x2} = |J_x| \) which can be considered classical. The transverse spin components \( J_y \) and \( J_z \) of each ensemble are quantum variables and will rotate at the Larmor frequency \( \Omega_L \). We define two sets \( c \) and \( s \) of non-local canonical variables by

\[
\hat{x}_{Ac} = \frac{\hat{j}_{y1} - \hat{j}_{y2}}{\sqrt{2 |J_x|}}, \quad \hat{p}_{Ac} = \frac{\hat{j}_{z1} + \hat{j}_{z2}}{\sqrt{2 |J_x|}} \quad \text{and} \quad \hat{x}_{As} = -\frac{\hat{j}_{z1} - \hat{j}_{z2}}{\sqrt{2 |J_x|}}, \quad \hat{p}_{As} = \frac{\hat{j}_{y1} + \hat{j}_{y2}}{\sqrt{2 |J_x|}},
\]

(1)

where \( \text{rot} \) refers to the fact that the spins are defined in a frame rotating at the Larmor frequency. Based on the commutation relations for the spin, \( [\hat{J}_y, \hat{J}_z] = i \hat{J}_x \), we find the commutation relations for the canonical variables \( [\hat{x}_A, \hat{p}_A] = i \) and the Heisenberg uncertainty relation \( \text{Var}(\hat{x}_A) \cdot \text{Var}(\hat{p}_A) \geq 1/4 \). The two sets \( c \) and \( s \) of canonical variables commute (for instance \( [\hat{x}_{Ac}, \hat{p}_{As}] = 0 \)). Coherent spin states, such as states which are fully pumped along the \( x \)-direction, are minimal uncertainty states with \( \text{Var}(\hat{x}_A) = \text{Var}(\hat{p}_A) = 1/2 \).

The polarization of light can be described by the Stokes vector \( \hat{S} = (\hat{S}_1, \hat{S}_2, \hat{S}_3) \). For light propagating along the \( z \)-direction and strongly polarized along the \( y \)-direction, the \( \hat{S}_1 \) operator can be replaced by a classical number. In this situation \( \hat{S}_2 \) and \( \hat{S}_3 \) describe the quantum fluctuations in the \( x \)-polarization mode. For a pulse of light with duration \( T \) we define canonical variables for a cosine mode with frequency \( \Omega \) and exponentially rising modefunction with time constant \( 1/\gamma \) by

\[
\hat{x}_{Lc}^{\text{in}} \propto \int_0^T \frac{\hat{S}_2(t)}{\sqrt{\lambda}} \cos(\Omega t)e^{+\gamma t} \, dt \quad \text{and} \quad \hat{p}_{Lc}^{\text{in}} \propto \int_0^T \frac{\hat{S}_3(t)}{\sqrt{\lambda}} \cos(\Omega t)e^{+\gamma t} \, dt,
\]

and similarly for the sine mode with canonical variables \( \hat{x}_{Ls} \) and \( \hat{p}_{Ls} \). The cosine and sine modes are orthogonal and each have the commutation relation \( [\hat{x}_L, \hat{p}_L] = i \). Coherent light states (such as vacuum) have \( \text{Var}(\hat{x}_L) = \text{Var}(\hat{p}_L) = 1/2 \), and states squeezed for instance along the \( x \)-direction have \( \text{Var}(\hat{x}_L) < 1/2 \).

The interaction between off-resonant light and two highly polarized and oppositely oriented ensembles of cesium atoms can be described as a swap and squeezing of the quantum states [6]. We can write input-output equations describing the operators before and after the interaction

\[
\hat{x}_A^{\text{out}} = \sqrt{1 - \kappa^2 Z^2} \hat{x}_A^{\text{in}} + \kappa \hat{p}_L^{\text{in}}, \quad \hat{p}_A^{\text{out}} = \sqrt{1 - \kappa^2 Z^2} \hat{p}_A^{\text{in}} - \kappa \sqrt{Z^2} \hat{x}_L^{\text{in}},
\]

\[
\hat{x}_L^{\text{out}} = \sqrt{1 - \kappa^2 Z^2} \hat{x}_L^{\text{in}} + \kappa \hat{p}_A^{\text{in}}, \quad \hat{p}_L^{\text{out}} = \sqrt{1 - \kappa^2 Z^2} \hat{p}_L^{\text{in}} - \kappa \sqrt{Z^2} \hat{x}_A^{\text{in}}.
\]

(3)

These equations are valid both for the cosine mode and the sine mode and the index \( c, s \) has therefore been omitted. Input (output) light operators are defined with exponentially rising (falling) modefunctions with a time constant \( 1/\gamma_{sw} \) where \( \gamma_{sw} \) is the swap rate proportional to number of photons and number of atoms. The coupling constant \( \kappa = Z \sqrt{1 - e^{-2\gamma_{sw} T}} \), \( \sqrt{1 - \kappa^2/Z^2} = e^{-\gamma_{sw} T} \), and \( 1/Z^2 = 14 a_2/a_1 \) is a measure of departure from the quantum non demolition (QND) interaction [7]. \( a_1 \) and \( a_2 \) are the vector and tensor polarizabilities. For cesium atoms in the \( F = 4 \) ground state hyperfine manifold and a probe blue-detuning of \( \Delta = 850 \) MHz from the \( F = 4 \rightarrow F' = 5 \) component of the \( D_2 \) line we have \( Z^2 = 6.3 \). In the QND limit \( a_2(\Delta) \rightarrow 0 \), Eq. (3) reduces to the QND-interaction where \( \hat{p}_A \) and \( \hat{p}_L \) are conserved. The interaction given by Eq. (3) can be used to measure the atomic state since the output light
operator $\hat{x}_{\text{out}}^A$ contains an atomic input operator. For long times $\gamma_{\text{dw}}T \gg 1$ the atoms and light
swap their quantum states and are squeezed by a factor $\hat{Z}$ in the variance

$$\hat{x}_{\text{out}}^A = Z\hat{p}_{in}^L, \quad \hat{p}_{A}^\text{out} = -\frac{1}{Z}\hat{z}^\text{in}_L \quad \text{and} \quad \hat{x}_{\text{out}}^L = Z\hat{p}_{A}^\text{in}, \quad \hat{p}_{L}^\text{out} = -\frac{1}{Z}\hat{z}^\text{in}_A. \quad (4)$$

This swap and squeezing interaction could be used directly as a memory; however, the swap
time $1/\gamma_{\text{dw}}$ is too long compared to the atomic decoherence time. Instead our memory protocol
consists of a light atom interaction with strength $\kappa = 1$ followed by a feedback of the measured
$\hat{x}_{\text{out}}^L$ to the atomic $\hat{p}_{A}$ operator. The final atomic operators after optimal feedback is

$$\hat{x}_{\text{fin}}^A = \sqrt{1 - \frac{1}{Z^2}}\hat{z}^\text{in}_A + \hat{p}_{in}^L \quad \text{and} \quad \hat{p}_{A}^\text{fin} = -\hat{z}^\text{in}_A. \quad (5)$$

We see that $\hat{x}_{\text{in}}^L$ is stored perfectly in $\hat{p}_{in}^L$ (up to a sign). Likewise $\hat{p}_{in}^A$ is stored into $\hat{x}_{\text{in}}^A$; however,
the initial atomic $\hat{x}_{\text{in}}^A$ is not cancelled, only slightly reduced. This will lead to additional noise
in the final atomic state. This noise can be reduced if one squeezes $\hat{x}_{\text{in}}^A$ prior to the storage.
This can be done with a near QND-measurement followed by feedback to the atoms [8]. In the
experiment, we reduced the initial variance in $\hat{x}_{\text{in}}^A$ from $1.10(8) \cdot \frac{1}{2}$ to $0.86(6) \cdot \frac{1}{2}$. Anti-squeezing
in $\hat{p}_{in}^A$ was measured to be $2.1(1) \cdot \frac{1}{2}$.

The states which are stored in the memory are two entangled sidebands of light with creation
operators $\hat{a}_+^A$ and $\hat{a}_-^A$ and frequencies $\omega_+ = \omega_0 + \Omega_L$ where $\omega_0$ is the optical carrier frequency
(see Fig. 4). Entanglement of the sidebands is equivalent to squeezing in the cosine and
sine modes which are described by the two sets of canonical operators $\hat{x}_{\text{Ac}}, \hat{p}_{\text{Ac}}$ and $\hat{x}_{\text{As}}, \hat{p}_{\text{As}}$. The Gaussian input states are characterized by the mean values and variances of these
canonical operators, and in the experiment [1] we store 18 different states with mean values
$\{\langle \hat{x}_L \rangle ; \langle \hat{p}_L \rangle = 0, 3.8, 7.6; 0, 3.8, 7.6\}$ and variances squeezed by 6 dB either along the $\hat{x}_L$- or $\hat{p}_L$-
direction (Fig. 1). The mean values are created by phase and amplitude modulation with electro-
optical modulators, and the squeezed light is produced using an optical parametric oscillator.

The fidelity of a state transfer is the overlap between the initial light state and the final atomic
state. For a given distribution of input states one can calculate the average fidelity. We
calculated the average fidelity for three different experimental input distributions parametrized
by $d_{\text{max}} = \max \{\{\langle \hat{x}_L \rangle ; \langle \hat{p}_L \rangle \} = 0, 3.8 \text{ and } 7.6\}$, see Fig. 1. For these three input distributions we
obtain average experimental fidelities of 0.58(2), 0.52(2), and 0.41(4), respectively. The choice of $\kappa = 1$ (which minimizes the noise added to the stored states) together with optical losses before
the memory had the consequence that the stored atomic mean values were smaller by a factor
0.85 than the input mean values. This is the main reason for the reduction of the experimental fidelity with larger displacements.

In order for a memory to be truly quantum, it should be able to outperform any classical
memory. For a single known quantum state the classical memory works with fidelity one. For
an unknown state selected randomly from an input distribution of states similar to the ones in
Fig. 1, an upper limit (benchmark) for the fidelity was calculated by semi-definite programming
methods [9]. For $d_{\text{max}} = 3.8$, the classical benchmark fidelity is 0.44, below the measured fidelity
of 0.52(2). The memory presented here is therefore a true quantum memory for this input set.
This means that if we choose to input one part (for instance, the upper sideband mode $\hat{a}_+$) of
an EPR-pair into the memory, the memory would be entangled with the other part (the lower
sideband mode $\hat{a}_-$). In the actual experiment, both modes of a two mode EPR entangled state
were stored. One part of the EPR pair ($\hat{a}_+$) was stored in the first ensemble and one part ($\hat{a}_-$)
of the EPR pair was stored in the second ensemble (see also Fig. 4). The two modes could in
principle have been separated spatially using narrowband filter cavities. In this case, only one
part of the EPR pair would be stored and the other would be left as a propagating mode.
Based on the performance of our memory, we would expect an EPR variance of $\Sigma_{\text{EPR}} = 0.76 < 1$
between the stored mode and the travelling mode ($\Sigma_{\text{EPR}} = 1$ is the limit for entanglement).
Figure 1. **Input States.** In total we store 18 different input states represented as ellipses in phase space. We divide the states up into three alphabets. The first alphabet (dashed line) contains 2 states (not displaced, squeezing phase $\phi = 0$ or 90 deg). The second alphabet (dotted line) contains 8 states with displacements up to $d_{\text{max}} = 3.8$. The third alphabet (solid line) contains all 18 states with a maximum displacement $d_{\text{max}} = 7.6$.

2. **Radio frequency magnetometry**

An atomic ensemble can be used as a sensor for magnetic fields [10]. Below we will describe how an atomic radio frequency (RF) magnetometer works in terms of the total spin vector $\hat{J}$ and the Stokes vector $\hat{S}$. The atomic ensemble is initially optically pumped such that the mean spin points in the $x$-direction. The transverse spin components have zero mean and variances given by the Heisenberg uncertainty relation for a minimal uncertainty state

\[ \text{Var}(\hat{J}_y) = \text{Var}(\hat{J}_z) = J_x / 2. \]

The spin can be envisioned as a long vector with an uncertainty disk on top (see Fig. 2a). The atomic ensemble is again located in a bias field leading to Larmor precession of the transverse spin components. If an RF magnetic field with amplitude $B_{\text{RF}}$ oscillating at the Larmor frequency in the transverse directions is applied for a certain duration $\tau$, the spin vector will tilt and a mean transverse component is created (see Fig. 2a)

\[ \langle \hat{J}_\perp \rangle = \Gamma B_{\text{RF}} J_x T_2 \left( 1 - e^{-\tau/T_2} \right), \]  

(6)

where $\Gamma = 2\pi \Omega_L / B = 2.2 \cdot 10^{10}$ rad/Tesla for cesium, and $T_2$ is the transverse spin decay time. The projection noise limited minimal detectable magnetic field amplitude is given by equating $\langle \hat{J}_\perp \rangle$ to the projection noise $\sqrt{J_x / 2}$, and we find

\[ B_{\text{min}} = \left[ \Gamma \sqrt{2 J_x T_2 \left( 1 - e^{-\tau/T_2} \right)} \right]^{-1}. \]  

(7)

The projection noise limited sensitivity is then $B_{\text{min}} \sqrt{T_{\text{RF}}}$, has the units Tesla/$\sqrt{\text{Hz}}$, and is the minimal detectable field using one second of measurement time.

The precessing transverse spin component $\hat{J}_\perp$ can be measured using linearly polarized light. The output Stokes operator is

\[ \hat{S}_\text{out}^2(t) = \hat{S}_\text{in}^2(t) + a S_1 \left( \hat{J}_y^{\text{rot}} \sin \Omega t + \hat{J}_z^{\text{rot}} \cos \Omega t \right), \]  

(8)

and consists of shot noise and an atomic signal. The factor $a$ is a coupling constant. For a single atomic ensemble located in a magnetic field, back-action noise from the measurement will pile up in both atomic operators $\hat{J}_y^{\text{rot}}$ and $\hat{J}_z^{\text{rot}}$. We, therefore, choose to implement the magnetometer with two atomic ensembles where a near QND measurement of the collective operators is used to cancel the backaction noise. The measured signals are

\[ \hat{S}_{2c}^{\text{out}} = \sqrt{1 - \kappa^2 / Z^2} \hat{S}_{2c}^{\text{in}} + \kappa \sqrt{\Phi / F N_A} \left( \hat{J}_y^{\text{rot}} + \hat{J}_z^{\text{rot}} \right), \]  

(9)

and $\hat{S}_{2s}^{\text{out}}$ which is given by a similar expression with $\hat{J}_y$ replaced by $\hat{J}_z$. $\Phi$ is the photon flux, $N_A$ the number of atoms, and $F = 4$ for the specific hyperfine ground state of cesium. This equation
Figure 2. (a) An illustration of the RF magnetometer. The atomic ensemble is located in a bias static field $B$ and an oscillating magnetic field $B_{RF}$. (b) Setup. The atomic ensembles are probed by off resonant light which is measured by polarization homodyning. HWP: half wave plate. (c) Results of a series of measurements with and without $B_{RF}$ applied. Each dot represents a measurement point. Dashed lines are the projection noise limit and solid lines are the measured standard deviations. Inset: Power spectrum of the signal in arbitrary units. (d) Pulse sequence for the magnetometry experiment. Firstly, atoms are polarized by optical pumping, then the magnetic RF pulse is applied and, finally, the atomic spin is probed with off-resonant light. (e) Pulse sequence for entanglement assisted magnetometry. An extra probe pulse prior to the RF pulse is used to entangle the two atomic ensembles. Figures are from [11]. ©2010 The American Physical Society

is identical to Eq. (3) but is written in terms of Stokes and spin operators instead of canonical operators. The displacement given by Eq. (6) gives a mean value in the signals $\hat{S}_{2c}^{\text{out}}$ and $\hat{S}_{2s}^{\text{out}}$ proportional to the RF field amplitude $B_{RF}$. The variance of the signal will be

$$\text{Var} \left( \hat{S}_{2c}^{\text{out}} \right) = (1 - \kappa^2/Z^2) \text{Var} \left( \hat{S}_{2c}^{\text{in}} \right) + \kappa^2 \left( \frac{\Phi}{FNA} \right) \text{Var} \left( \hat{J}_{y1}^{\text{rot}} + \hat{J}_{y2}^{\text{rot}} \right).$$

(10)

For shot noise limited light $\text{Var} \left( \hat{S}_{2c}^{\text{in}} \right) = S_x/2$ and for projection noise limited atoms in the coherent spin state $\text{Var} \left( \hat{J}_y \right) = \text{Var} \left( \hat{J}_z \right) = J_x/2$.

In the experiment a 36(3) femtoTesla RF field was applied to the atoms for a duration of 15 ms. The measured signals $S_{2c}^{\text{out}}$ and $S_{2s}^{\text{out}}$ are plotted in Fig. 2c from which we can deduce a signal to noise ratio $SNR = 12.3$. We then calculate the sensitivity $B_{RF} \sqrt{\tau}/SNR = 3.6(4) \cdot 10^{-16}$.
Tesla$/\sqrt{\text{Hz}}$ which is close to the theoretical value $2.7(4) \cdot 10^{-16}$ Tesla$/\sqrt{\text{Hz}}$ for the projection noise limited sensitivity calculated from Eq. (7) with $N_A = 2 \cdot 7.2(7) \cdot 10^{11}$ atoms and $T_2 = 32$ ms. The best experimental sensitivity obtained was $B_{RF} \sqrt{T_{tot}/SNR} = 4.2(8) \cdot 10^{-16}$ Tesla$/\sqrt{\text{Hz}}$ calculated using the duty cycle time $T_{tot}$ (optical pumping pulse of 6 ms, probe pulse of 1.5 ms and RF pulse duration of 22 ms). This sensitivity is comparable to the sensitivity of the state-of-the-art atomic magnetometer [12] which operates with $10^4$ times more atoms. This is expected since the projection noise limited measurement of the magnetic field yields the best sensitivity per atom without the use of entanglement.

The signal to noise ratio can be improved by reducing the uncertainty in the variables $\hat{J}_{y1}^{\text{rot}} + \hat{J}_{y2}^{\text{rot}}$ and $\hat{J}_{z1}^{\text{rot}} + \hat{J}_{z2}^{\text{rot}}$. This can be achieved by entangling the two atomic ensembles. The criterion for entanglement between the ensembles is that the Einstein-Podolsky-Rosen variance

$$\Sigma_{EPR} = \frac{\text{Var}(\hat{J}_{y1}^{\text{rot}} + \hat{J}_{y2}^{\text{rot}}) + \text{Var}(\hat{J}_{z1}^{\text{rot}} + \hat{J}_{z2}^{\text{rot}})}{2J_x} < 1.$$  

Such entanglement can be created with a near QND measurement followed by feedback. In the magnetometry experiment the atomic projection noise was initially $\Sigma_{EPR} = 1.10(8)$ and was reduced to $-1.5\text{dB}$ (30\%) below the projection noise level (see Fig. 3a). The entanglement decays with time constant $T_{2E} = 4$ ms. Since the $B_{RF}$ is applied after the entanglement creation, entanglement only improves the magnetometer for short RF pulses or, equivalent, broadband pulses. Experimentally we obtain a 15\% improvement in sensitivity for RF pulses with bandwidth greater than 1 kHz (Fig. 3b).

### 3. Entanglement generated by dissipation

A novel method for generating entanglement of atomic systems in a steady state has been recently proposed in [13]. Dissipation is used as the entangling mechanism, such that the atoms reach an entangled steady state independent of the initial state. This kind of steady state entanglement is highly desirable since it provides a robust entanglement source which is available on demand at any instance in time. In order to describe the entanglement generation, we implement a master equation approach detailed below.

We consider two oppositely polarized atomic ensembles with level structure as shown in Fig. 4. Atoms in ensemble 1 (2) initially occupy the level $F = 4, m_F = 4$, $(F = 4, m_F = -4)$. The
ensembles interact with strong linearly $y$-polarized light represented in the $x$-quantization axis as a sum of right hand and left hand circularly polarized light. Situated in this light field, the atoms can be excited and spontaneously emit photons in all directions. The photons are described by modes $\hat{a}_k$ where $\vec{k}$ is the wavevector. Since spontaneous emission is a resonant process and the atomic ground state levels are separated by the energy $\hbar \Omega_L$, the frequency of the spontaneously emitted photons will be $\omega_+ \text{ or } \omega_- \text{ where } \omega_{\pm} = \omega_0 \pm \Omega_L$. The atoms in ensemble 1 (2) couple to $\sigma_-$-polarized light with coefficient $\chi_1 \text{ (} \chi_2 \text{)}$ and to $\sigma_+$-polarized light with coefficient $\chi_2 \text{ (} \chi_1 \text{)}$. Upon absorption of a $\sigma_-$-polarized photon, a collective spin excitation $\hat{b}_{1\pm}^\dagger$ can be created in the first ensemble together with creation of a photon in the upper sideband (see Fig. 4). This leads to a term proportional to $\int_{\omega_{\text{us}}} d\vec{k} \chi_1 \hat{a}_{\vec{k}}^\dagger \hat{b}_1 + \text{h.c.}$ in the Hamiltonian. The integration (over $\vec{k}$) runs over modes within a narrow frequency interval $\Delta \omega_{\text{us}}$, centred around $\omega_+$ (us stands for upper sideband.) In total, we find four such terms which give rise to the total Hamiltonian

$$\hat{H} \propto \int_{\Delta \omega_{\text{us}}} d\vec{k} \left( \hat{A}_k^\dagger \hat{A}_k + \hat{A}_k^\dagger \hat{A}_k \right) + \int_{\Delta \omega_{\text{ls}}} d\vec{k} \left( \hat{B}_k^\dagger \hat{B}_k + \hat{B}_k^\dagger \hat{B}_k \right), \quad (12)$$

where $\Delta_t$, is a narrow frequency interval centred around $\omega_-$ and where we have defined the variables

$$\hat{A} = \chi_2 \hat{b}_1 + \chi_1 \hat{b}_2^\dagger \quad \text{and} \quad \hat{B} = \chi_2 \hat{b}_2 + \chi_1 \hat{b}_1^\dagger, \quad (13)$$

This Hamiltonian leads to a master equation for the density matrix $\rho$ of the ensembles [13]

$$\frac{\partial \hat{\rho}}{\partial t} = \frac{d\Gamma}{2} \left( \hat{A} \hat{\rho} \hat{A}^\dagger - \hat{A}^\dagger \hat{A} \hat{\rho} + \hat{B} \hat{\rho} \hat{B}^\dagger - \hat{B}^\dagger \hat{B} \hat{\rho} + \text{h.c.} \right) + \hat{\mathcal{L}}_{\text{noise}} \hat{\rho}, \quad (14)$$

where $\Gamma$ is the single atom radiative decay rate and $\hat{\mathcal{L}}_{\text{noise}}$ represents undesired noise such as single atom spontaneous emission. Emission in the forward propagating modes is enhanced by the optical depth $d$ as compared to the modes into other directions and give rises to the terms in Eq. (14) which are proportional to $d$. With such a master equation, the collective state of the atomic ensembles will be driven into an entangled steady state with $\Sigma_{\text{EPR}} < 1$ [13]. The emission of a forward propagating photon could arise from both the first and second ensemble, the two paths are indistinguishable, and this leads to entanglement between the two atomic ensembles. The other modes act as in the usual spontaneous emission decoherence mechanism. The generated entanglement can also be explained by the simpler model given by Eq. (4) which does not include spontaneous emission in the non-forward propagating modes. For long interaction times $\gamma_{\text{us}} T \to \infty$ the collective state of the two ensembles reaches a steady state where the operators $\hat{p}_{Ac}$ and $\hat{p}_{As}$ are squeezed by a factor $Z^2$ in the variances which means that the two atomic ensembles are entangled.

In the preliminary experiments two important things were observed. Firstly, starting from a coherent state with $\Sigma_{\text{EPR}} = 1$, entanglement is created with $\Sigma_{\text{EPR}} < 1$ for a period of more than
20 ms. Then the noise grows, $\Sigma_{\text{EPR}} > 1$, and the atomic ensembles are no longer entangled. Secondly, the mean spin $J_x$ of each ensemble decreases with time. Both the increase in noise and decrease in $J_x$ for long times are expected due to atomic depolarization. The main depolarization mechanisms are spontaneous emission into the other modes and atomic collisions. Also, when $J_x$ decreases, the collective enhancement $\propto J_x$ of the emission into the forward propagating modes is reduced and the entanglement mechanism loses strength.

In order to keep the atoms in the $F=4$ manifold well polarized, we shine resonant $\sigma_+ (\sigma_-)$ polarized light onto the atoms in ensemble 1 (2). This significantly increases the polarization $p = \frac{1}{2} \sum_{m=4}^{m=-4} m \rho_{mm}$ of the atoms in the $F=4$ manifold. Preliminary results show that also the time with entanglement could be extended beyond the 20 ms. Still, atoms are lost to the $F=3$ manifold with time and, unfortunately, we do not reach entanglement in a steady state. A full account of the results can be found in [14]. Possible improvements could be to add repump light which can recycle atoms lost to the $F=3$ manifold. According to our model, the combination of repumping and higher optical depths would yield steady state entanglement. Another possibility which may lead to steady state "forever entangled" atomic systems is a combination of the dissipation mechanism with the continuous measurement.

4. Outlook
Optically thick room temperature alkali atoms in a spin protected environment have the potential to become a robust platform for quantum technologies. Recently, a lifetime of the ground spin state exceeding several seconds with an improved spin protecting coating of the cell walls has been demonstrated in [15, 16]. Combination of improved coating with a fibre-coupled microcell design and low finesse optical cavities yields a scalable and versatile system for quantum information processing and sensing.

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References
[1] Jensen K, Wasilewski W, Krauter H, Fernholz T, Nielsen B M, Owari M, Plenio M B, Serafini A, Wolf M M and Polzik E S *Nature Physics* advance online publication, 07 November 2010 (DOI 10.1038/nphys1819)
[2] Julsgaard B, Sherson J, Cirac J I, Fiurášek J and Polzik E S 2004 *Nature* **432** 482–486
[3] Appel J, Figueroa E, Korytov D, Lobino M and Lyovsky A I 2008 *Phys. Rev. Lett.* **100** 093602
[4] Honda K, Akamatsu D, Arikawa M, Yoko Y, Akiba K, Nagatsuka S, Tanimura T, Furusawa A and Kozuma M 2008 *Phys. Rev. Lett.* **100** 093601
[5] Adesso G and Chiribella G 2008 *Phys. Rev. Lett.* **100** 170503
[6] Wasilewski W, Fernholz T, Jensen K, Madsen L S, Krauter H, Muschik C and Polzik E S 2009 *Opt. Express* **17** 14444–14457
[7] Hammerer K, Sørensen A S and Polzik E S 2010 *Rev. Mod. Phys.* **82** 1041–1093
[8] Julsgaard B, Rozhekin A and Polzik E S 2001 *Nature* **413** 400–403
[9] Owari M, Plenio M B, Polzik E S, Serafini A and Wolf M M 2008 *New J. Phys.* **10** 113014
[10] Budker D and Romalis M 2007 *Nat. Phys.* **3** 227–234
[11] Wasilewski W, Jensen K, Krauter H, Renema J J, Balabas M V and Polzik E S 2010 *Phys. Rev. Lett.* **104** 133601
[12] Lee S K, Sauer K L, Seltzer S J, Alem O and Romalis M V 2006 *Appl. Phys. Lett.* **89** 214106
[13] Muschik C A, Polzik E S and Cirac J I 2010 (Preprint arXiv:1007.2209v1)
[14] Krauter H, Muschik C A, Jensen K, Wasilewski W, Petersen J M, Cirac J I and Polzik E S 2010 (Preprint arXiv:1006.4344v1)
[15] Balabas M V, Jensen K, Wasilewski W, Krauter H, Madsen L S, Müller J H, Fernholz T and Polzik E S 2010 *Opt. Express* **18** 5825–5830
[16] Balabas M V, Karaulanov T, Ledbetter M P and Budker D 2010 *Phys. Rev. Lett.* **105** 070801

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