Photon-echo-based quantum memory for optical squeezed states

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Abstract
The ability to efficiently realize storage and readout of optical squeezed states plays a key role in continuous-variables quantum information processing. Here we study the quantum memory for squeezed state of propagating light in atoms based on the hybrid photon echo re-phasing. The optical quantum state is recorded in two sublevels of the ground state of an atomic ensemble to realize long-lived quantum memory. Taking into account the noise effect due to atomic decay, our estimation indicates that high fidelities larger than the classical fidelity threshold 81.5% are obtainable even with currently available techniques. Moreover, our result shows that the decay rate of atoms restricts the maximal fidelity. Our work provides some practical guidance for the realization of efficient and faithful photon-echo-based memory for squeezed light.

Keywords: quantum memory, photon echo, optical squeezed states

1. Introduction
Quantum memory for light plays a key role in quantum-information processing and has been considered as a basic ingredient for quantum repeaters [1] and scalable all-optical quantum computers [2]. Currently, much effort is being devoted to such memory through different approaches, including Faraday rotation [3, 4], electromagnetically induced transparency (EIT) [5, 6], off-resonant Raman transitions [7, 8] and photon echoes [9–12]. Photon-echo-based quantum memories [13–15] are currently attracting particular attention, not only for their abilities to implement in solid-state materials, but also for their successful achievements in storage efficiency [16, 17] and multimode-memory capacity [18]. Many impressive results have now been achieved in this area. Gisin et al showed [19] that photon-echo techniques can be used to store and retrieve time-bin qubits [20]. Clausen et al demonstrated [21] the storage of entangled photons using the atomic frequency comb technique. Recently, highly efficient storage of optical coherent state (about 87%) based on gradient echo memory has also been demonstrated [17]. Despite these achievements, however, storage and revival of optical squeezed-state light based on photon-echo techniques, as far as we know, has not yet been studied both theoretically and experimentally.

Squeezed state light, which has a reduction in one of its quadratures below the standard quantum limit, has been shown to be especially valuable for interferometry [22], high-precision measurement [23], quantum information with continuous variables [24], quantum illumination [25] and quantum reading [26]. Therefore, realization of the efficient storage and revival of squeezed light would be very appealing. Over the last decade, a variety of approaches concerning the memory for squeezed light have been proposed. One of them proposes to transfer the squeezed state from non-classical light to the excited atomic states [27]. Another approach to map the squeezed state of light onto atomic ensemble based on stimulated Raman scattering has been developed by Kozhekin et al [8]. In particular, storage and retrieval of squeezed light with EIT was also investigated [6]. Very recently, Jensen et al demonstrated [4] the quantum memory for entangled two-mode squeezed states via Faraday rotation.

In this paper, we study the long-lived quantum memory for squeezed state of light based on hybrid photon echo re-phasing (HYPER)—a variant of the two-pulse photon echo (2PE) technique [28]. 2PE is well known to be unsuitable for...
serving as a quantum memory technique, because of the irritating collective emission of photons (amplified spontaneous emission). In order to overcome this problem, McAuslan et al developed a new photon-echo technique called HYPER, which uses electrical field gradients to eliminate the unwanted gain created by the strong re-phasing pulse ($\pi$-pulse [28] or frequency-chirped pulse [29]). Such a technique has the outstanding characteristic of avoiding spectral hole-burning, which enables us to implement large-bandwidth and high-efficiency quantum memories. Here, utilizing this approach, we carefully derived the processes for storage and retrieval of optical squeezed-state light. We find that, for the ideal case, the squeezed-quadrature variance of the retrieved optical field only relies on the optical depth of atomic ensemble. The larger the optical depth, the more the squeezing of the input mode will be preserved. Taking into account the noise effect due to atoms coupled to thermal reservoirs, the squeezed-quadrature variance now depends not only on the optical depth but also on the decay rate. We show that substantial squeezing of the input mode can also be preserved even with currently available techniques [28]. The rest of this paper is organized as follows. In section 2, we present the basic equations governing the evolution of the quantized atomic and photonic fields based on photon-echo techniques. Next, storage and retrieval of optical squeezed-state is considered. In section 3, we discuss the possibility of implementing the technique in solid-state materials. Section 4 gives our conclusions.

2. The storage process

2.1. The basic evolutions

Our storage material consists of a large number of atoms with an excited state $|e\rangle$ and two lower states $|g\rangle$, $|s\rangle$, see figure 1(a). Let us first study the absorption process. Consider a quantized optical pulse being injected into the storage medium to couple $|g\rangle$ and $|e\rangle$. In the one-dimensional light propagation model, the negative frequency part of the optical field can be decomposed into forward and backward modes

$$
\hat{a}(z, t) = \epsilon \hat{a}_f(z, t)e^{-i(\omega t - kz)} + \epsilon \hat{a}_b(z, t)e^{-i(\omega t + kz)},
$$

(1)

where $\epsilon = \sqrt{\hbar \omega_0/(\varepsilon_0 V)}$, $\omega_0$ is the central frequency of optical pulse, $\varepsilon_0$ is the vacuum permeability, and $V$ is the quantized volume. $\hat{a}_k(z, t)$ (where $k = f, b$) represents the optical operators with the commutation relation $[\hat{a}_k(z, t), \hat{a}^\dagger_k(z, t)] = \delta(t - t')$. For atomic operators, the mean field per atom is defined as [30]

$$
\delta_{ge}(z, t, \Delta) = \frac{1}{N(\Delta, z)} \sum_{n=1}^{N(\Delta, z)} |g\rangle_n \langle e|,
$$

(2)

where $\Delta = \omega_{eg} - \omega_0$ is the detuning from resonance, with the atomic transition frequency $\omega_{eg}$. In the above sum, the index $n$ runs over all atoms $N(\Delta, z) = \rho \delta z g(\Delta) \delta \Delta$ within the infinitesimal slice $\delta z$ and $\delta \Delta$. The notation $\rho$ denotes the number density of atoms, and $g(\Delta)$ is the normalized atomic spectral distribution. In analogy with the optical field, the negative-frequency part of the atomic operator can be described by the two counter-propagating contributions

$$
\delta_{ge}(z, t, \Delta) = \delta_f(z, t, \Delta)e^{-i(\omega t - k z)} + \delta_b(z, t, \Delta)e^{-i(\omega t + k z)}.
$$

(3)

In the low-excitation limit, if the atoms are initially prepared in the ground state, the operators associated to the atomic coherence $\delta_{ge}$ have the following commutation relation [31]

$$
\left[\delta_{ge}(z, t, \Delta), \delta^\dagger_{ge}(z', t, \Delta')\right] = \frac{1}{\rho g(\Delta)}\delta(z - z')\delta(\Delta - \Delta').
$$

(4)

With these definitions, the dynamic of the optical absorbing process can be described by the following equations

$$
\left\{\frac{1}{c}\partial_t + \partial_z\right\}\hat{a}_f(z, t) = ig\rho \int_{-\infty}^{\infty} d\Delta g(\Delta) \hat{\delta}_f(z, t, \Delta),
$$

(5)

$$
\partial_t \hat{\delta}_f(z, t, \Delta) = \left(\frac{\Delta}{2} - \frac{g}{2}\right)\hat{\delta}_f(z, t, \Delta) + ig\hat{a}_f(z, t) + \hat{\delta}_f(z, t, \Delta),
$$

(6)
where $g$ is the coupling constant of the light–atoms interactions, and we have considered the decoherence of atomic system with $\Gamma$ the decay rate of the excited state. $\hat{f}(z, t, \Delta)$ is the Langevin noise operator, satisfying the following correlation functions \cite{32}

$$
\langle \hat{f}(z, t, \Delta) \hat{f}^\dagger(z', t', \Delta') \rangle = \frac{\Gamma}{\rho g(\Delta)} \delta(z - z') \delta(t - t') \delta(\Delta - \Delta'),
$$

and

$$
\langle \hat{f}^\dagger(z', t', \Delta') \hat{f}(z, t, \Delta) \rangle = 0.
$$

The solution of equation (6) can easily be found to be

$$
\hat{\sigma}(z, t, \Delta) = \hat{\sigma}(z, t_0, \Delta) e^{(i/\gamma) \Delta \int_{t_0}^{t} ds \hat{f}(s, t, \Delta)} + ig \int_{t_0}^{t} ds \hat{\sigma}(z, s, \Delta) e^{(i/\gamma) \Delta \int_{t_0}^{s} ds \hat{f}(s, t, \Delta)},
$$

where $\hat{\sigma}(z, t_0, \Delta)$ is the initial atomic operator. To derive the time evolution of optical field, we insert equation (9) into equation (5) to generate

$$
\partial_t \hat{\alpha}(z, t) = -g^2 \rho \int_{t_0}^{t} ds \hat{\sigma}(z, s, \Delta) e^{-\frac{(t-s)(t-s)}{2}},
$$

$$
\times \sigma(\Delta) e^{i \Delta \int_{t_0}^{t} ds \hat{f}(s, t, \Delta)} + ig \sigma(\Delta) \sigma(\Delta) e^{i \Delta \int_{t_0}^{t} ds \hat{f}(s, t, \Delta)} + ig \sigma(\Delta) \sigma(\Delta) e^{i \Delta \int_{t_0}^{t} ds \hat{f}(s, t, \Delta)} + \frac{\Gamma}{\rho g(\Delta)} \delta(z - z') \delta(t - t') \delta(\Delta - \Delta'),
$$

(10)

Here, we have neglected the temporal derivative in equation (5), since we consider the regime $\tau \gg L/c$ ($\tau$ is the temporal length of the incident optical pulse; $L$ is the length of medium), which allows us to ignore the temporal retardation effect of light. In order to find the solution of equation (10), we assume the atomic spectral distribution is uniform, that is $g(\Delta) \sim 1/\gamma$, where $\gamma$ is the bandwidth of the ensemble, which is valid for relatively narrow bandwidth (\delta\omega_0) of the input signal light pulse $\delta\omega_0 < \gamma$ \cite{9}. Furthermore, the duration of the absorption process is essentially governed by the temporal length of the signal optical pulse, which ensures the values of $t - s$ of order $\tau$. Hence, in the limit $\tau \gg 1$ the Fourier transform $\int_{-\infty}^{\infty} d\Delta g(\Delta) e^{i \Delta (t-s)}$ in equation (10) will act like a $\delta$ function \cite{9, 33, 34}. With these assumptions, equation (10) can be directly solved to give

$$
\hat{\alpha}(z, t) = \hat{\alpha}(0, t) e^{-\Gamma t/2} + ig \rho \int_{t_0}^{t} ds \int_{-\infty}^{\infty} d\Delta g(\Delta) \hat{f}(s, t, \Delta)
$$

\times e^{i \Delta (t-s)} e^{-\frac{\gamma}{2} (t-s)^2} + \frac{\Gamma}{\rho g(\Delta)} \delta(z - z') \delta(t - t') \delta(\Delta - \Delta'),
$$

(11)

where $\hat{\alpha}(0, t)$ is the input photonic field, and we have defined the absorption coefficient $\alpha = 2 \gamma g^2 / \gamma$. The first term describes the exponential decay of the incoming optical pulse.

### 2.2. Long-lived memory for squeezed state of light

Now, we consider a long-time storage protocol based on the HYPER technique. This protocol consists of three separate time regions, as shown in figure 1(b). In region 1 ($t_0 < t < t_1$), just after absorbing the signal optical field, a linear electrical field gradient is applied to the sample to cause an additional atomic phase change depending on the position of the atoms. Soon after this field, a strong control optical field coupling to $|e\rangle$, $|s\rangle$ is turned on to transfer the population from $|e\rangle$ to $|s\rangle$. Obviously, information about the signal field is now stored in the spin wave. After a time interval ($t_1 < t < t_2$), another strong optical field propagating in the same direction is applied to drive the population back onto the excited state. In region 2 ($t_2 < t < t_5$), two $\pi$ pulses sandwiched with an electrical field are applied to the atoms to rephase the input optical mode. In region 3 ($t > t_5$), the ensemble will emit an expected echo at $t_6$.

During the storage process, we assume for simplicity that the first electrical field is short enough, which enables us to ignore the decay of the excited state during the time interval ($t_1 < t < t_2$). We also neglect spontaneous decay of the state $|s\rangle$, which is reasonable since we assume the time interval ($t_3 - t_2$) is much smaller than the lifetime of $|s\rangle$. Furthermore, the time spent in the region 2 is also temporally short, which allows us to reduce the dynamics of this region to an instantaneous operation at $t_5$ \cite{28}. Besides, if the twopulses used in region 2 propagate in opposite directions, the system will satisfy the condition of the backward retrieval resulting in the capability for achieving 100% quantum memory efficiency \cite{9, 34–38}. After applying the secondpulse in region 2, the atoms start to rephase. At $t_6 = 2t_5 - 2t_4 + t_1 - t_2 + t_1$, an echo propagating along negative $z$-axis through the medium occurs, which can be described by the following equations

$$
-\partial_t \hat{\delta}(z, t) = ig \rho \int_{-\infty}^{\infty} d\Delta g(\Delta) \hat{\delta}(z, t, \Delta),
$$

(12)

$$
\partial_t \hat{\delta}(z, t, \Delta) = \left( i \Delta - \frac{\gamma}{2} \right) \hat{\delta}(z, t, \Delta) + ig \hat{\alpha}(z, t) + \hat{f}(z, t, \Delta),
$$

(13)
where \( \hat{F}(z, t, \Delta) \) is the Langevin noise of the retrieval process. In deriving the above equations, we have used the boundary condition of the atomic field created at \( t_5 \)

\[
\hat{a}_b(z, t = 0; \Delta) = \hat{F}(z, t = 0; \Delta)e^{-i\Delta T},
\]

(14)

where we have set \( t_5 = 0 \), leading to the quantum memory time \( T = 0 - t_1 \). In accordance with (9), (12), and (14), one can easily derive the output optical field

\[
\hat{a}_b(0, t) = \hat{a}_{b}(L, t)e^{-\alpha L/2} + ae^{-iT/2}
\]

\[
\times \int_0^L dz \hat{f}(z, t - T)e^{-\alpha z/2}
\]

\[
- igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \hat{d}_f(z, t_0, \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}
\]

\[
+ igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_{t_0}^{0} dz \hat{f}(z, s; \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}
\]

\[
- \int_{t_0}^{0} dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_{t_0}^{0} dz \hat{f}(z, s; \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}.
\]

(15)

where \( \hat{a}_b(L, t) \) is the vacuum input of the backward retrieval optical mode, and we have used the fact that the optical field is out at \( z = 0 \). Next, equation (11) is substituted into equation (15), and after interchanging the order of the associated double integral \( \int_0^L dz \int_{-\infty}^{\infty} d\omega dz' \rightarrow \int_0^L dz \int_{-\infty}^{\infty} dz' \), we will obtain

\[
\hat{a}_b(t) = \hat{a}_b(0, t - T)
\]

\[
+ e^{-iT/2}(e^{-\alpha L} - 1)\hat{d}_f(0, t - T)
\]

\[
+ igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \hat{d}_f(z, t_0, \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}e^{-\alpha L}
\]

\[
+ igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_{t_0}^{0} dz \hat{f}(z, s; \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}
\]

\[
+ igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_{t_0}^{t - T} dz \hat{f}(z, s; \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2}
\]

\[
+ igp \int_0^L dz \int_{-\infty}^{\infty} d\Delta g(\Delta) \int_{t_0}^{t} dz \hat{f}(z, s; \Delta)
\]

\[
\times e^{i(\Delta - T/2)(t - t_0)}e^{-i\Delta T}e^{-\alpha z/2},
\]

(16)

where \( t - T \) is the time when the incident pulse is completely absorbed. The second term shows that information about the input optical mode is now printed onto the output mode. The last three noises terms are independent of each other, since they correspond to the Langevin noises at different times. Integrating equation (16) over the times of the duration of the input pulse, one obtains the collective optical operator [39]

\[
\hat{a}_b = \tau^{-1/2} \int_{0}^{T} dt \hat{d}_a(T)
\]

\[
= \tau^{-1/2} e^{-\alpha L/2} \int_{0}^{T} dt \hat{d}_a(L, t)
\]

\[
+ \tau^{-1/2} e^{-IT/2} (e^{-\alpha L} - 1) \int_{0}^{T} dt \hat{d}_a(0, t - T)
\]

\[
+ i \sqrt{\alpha} \int_{0}^{T} dz \left\{ \hat{D}(z) + \hat{F}_{11}(z) \right\} e^{i\Delta z/2} e^{-\alpha L}
\]

\[
+ \left[ \hat{F}_{12}(z) + \hat{F}_{2}(z) \right] e^{-\alpha z/2}
\]

(17)

where we defined the new operators for the initial atomic operator and Langevin noise operators

\[
\hat{D}(z) = \frac{\langle \rho \rangle}{2\pi} \int_{-\infty}^{\infty} dt \int_{-\infty}^{\infty} d\Delta g(\Delta)
\]

\[
\times \hat{d}(z, t_0, \Delta) e^{i(\Delta - T/2)(t - t_0)} e^{-i\Delta T},
\]

\[
\hat{F}_{11}(z) = \frac{\langle \rho \rangle}{2\pi} \int_{-\infty}^{\infty} dt \int_{-\infty}^{\infty} d\Delta g(\Delta)
\]

\[
\times \int_{t_0}^{T - t} dz \hat{f}(z, s; \Delta) e^{i(\Delta - T/2)(t - t_0)} e^{-i\Delta T},
\]

\[
\hat{F}_{12}(z) = \frac{\langle \rho \rangle}{2\pi} \int_{-\infty}^{T} dt \int_{-\infty}^{\infty} d\Delta g(\Delta)
\]

\[
\times \int_{t_0}^{t - T} dz \hat{f}(z, s; \Delta) e^{i(\Delta - T/2)(t - t_0)} e^{-i\Delta T},
\]

\[
\hat{F}_{2}(z) = \frac{\langle \rho \rangle}{2\pi} \int_{-\infty}^{T} dt \int_{-\infty}^{\infty} d\Delta g(\Delta)
\]

\[
\times \int_{t_0}^{0} dz \hat{f}(z, s; \Delta) e^{i(\Delta - T/2)(t - t_0)} e^{-i\Delta T},
\]

According to equations (4), (8) and (9), the above operators obey

\[
\left\langle \hat{D}(z) \hat{D}^\dagger(z') \right\rangle = A e^{i\Delta_0} \delta(z - z'),
\]

(18)

\[
\left\langle \hat{F}_{11}(z) \hat{F}_{11}^\dagger(z') \right\rangle = (e^{-\chi} - A e^{i\Delta_0}) \delta(z - z'),
\]

(19)

\[
\left\langle \hat{F}_{12}(z) \hat{F}_{12}^\dagger(z') \right\rangle = (A - e^{-\chi}) \delta(z - z'),
\]

(20)

\[
\left\langle \hat{F}_{2}(z) \hat{F}_{2}^\dagger(z') \right\rangle = (1 - A) \delta(z - z').
\]

(21)

Here we have defined the dimensionless factor

\[
A = \frac{1}{\Gamma T} (1 - e^{-\Gamma T})
\]

\[
\chi = IT.
\]

Equation (17) is the main result of this paper. Obviously, in contrast to the input light state the output one is severely distorted. In order to measure how well the output state compares to the input one, we introduce the fidelity \( F = \langle \rho | \tilde{\rho}_{\text{out}} | \rho \rangle \) for light, where |\( \rho \rangle \rangle \) is the input state, and \( \tilde{\rho}_{\text{out}} \) is the output state. Note that the atomic and optical states involved here are all pure Gaussian states. For such states, the fidelity is given by

\[
F = \frac{1}{2} \sqrt{\det(M_{\text{in}} + M_{\text{out}})^{-1}} \cdot \exp \left[ -\frac{1}{2} \xi^T (M_{\text{in}} + M_{\text{out}})^{-1} \xi \right]
\]

(40-42), where \( M \) represents the covariance matrix containing the elements \( M_{ij} = \left\langle \hat{y}_i \hat{y}_j \right\rangle \), and \( \xi = m_{\text{in}} - m_{\text{out}} \) with m standing for the mean values. As an
example, we consider storing and retrieving such an optical squeezed state whose mean values center around zero, such that \( \xi = 0 \). The fidelity then reduces to a simple form \( F = \frac{1}{2} \sqrt{\det \left( M_m + M_n \right)} \). In accordance with equations (17)–(21), one can directly derive the covariance matrices, and we finally get

\[
F = \frac{2}{\sqrt{\left( 1 + e^{-2\tau} \right) + e^{-\chi} \left( 1 - e^{-aL} \right)^2 \left( e^{2\tau} - 1 \right)}} \times \left[ \left( 1 + e^{2\tau} \right) + e^{-\chi} \left( 1 - e^{-aL} \right)^2 \left( e^{2\tau} - 1 \right) \right].
\]

In deriving equation (22), we have used the fact that \( \langle \hat{X}_f(t - T) \hat{X}_f(t' - T) \rangle = e^{-2\tau} \delta(t - t')/4 \), which means that the x-quadrature of the input optical mode are initially squeezed with the squeezing parameter \( r \). Note that, for the ideal case \( \chi = 0 \) and the certain squeezing input, the fidelity depends only on the optical depth. The larger the optical depth, the higher the fidelity. For the case of \( \chi \neq 0 \), in figure 2 we show the fidelity in its dependence on the optical depth for a different decay parameter \( \chi \), where the dashed line denotes classical fidelity threshold (CFT) [43]. One can see from the figure that, for an optical depth large than 2, the fidelity is well above CFT, and the maximal fidelity is mainly limited by the atomic decay. More specifically, for an optical depth large than 7, the achievable fidelities are nearly maximal but can be enhanced by decreasing the atomic decay rate. Fidelity versus optical depth for different squeezing is also depicted in figure 3, showing that the more highly squeezed the input mode, the less we can tolerate the losses of the memory process. It should be noted that, for low-squeezed light, its fidelity is much higher than CFT even within a small optical depth, which reflects that the quantum features of low-squeezed states are much more easily reconstructed via a measure-and-prepare strategy [43].
both the storage and the retrieval processes, while most of the previous proposed scheme mainly focused on the storage process [8, 27]. In conclusion, our work provides hopeful possibilities for the practical realization of efficient quantum memory for squeezed state of light based on photon echoes.

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