Entropy **2013**, *15*, 434-444; doi:10.3390/e15020434

OPEN ACCESS

entropy

ISSN 1099-4300

www.mdpi.com/journal/entropy

Article

All-Optically Controlled Quantum Memory for Light with a Cavity-Optomechanical System

Jinjin Li, Wen Bin and Ka-Di Zhu *

Department of Physics, Shanghai Jiao Tong University, 800 Dong Chuan Road, Shanghai, 200240, China

* Author to whom correspondence should be addressed; zhukadi@sjtu.edu.cn; Tel.: 86-21-3420-3864.

Received: 29 November 2012; in revised form: 26 December 2012 / Accepted: 6 January 2013 / Published: 24 January 2013

Abstract: Optomechanics may be viewed as a light-mechanics interface to realize hybrid structures for (classical or quantum) information processing, switching or storage. Using the two-laser technique, in this paper, we theoretically devise a protocol for quantum light memory via a cavity optomechanical system composed of a Fabry–Perot cavity and a mechanical resonator. Due to the long-lived mechanical resonator, this quantum memory for light based on optomechanically induced transparency (OMIT) can serve as a long-term memory that can store the full quantum light contained in an optical pulse. It is shown that, with the tunable pump laser, the quantum signal light can be reaccelerated and converted back on demand. Our presented work could open the door to all-optical routers for light memory devices and have a guide to actual experiments.

Keywords: all-optically; quantum memory; optomechanical system

1. Introduction

Storing and retrieving quantum light on demand, without corrupting the information it carries, have attracted wide attention in the field of quantum optics and quantum information processing [1–3]. So far, successful demonstrations of quantum light storage have been used in atomic vapor [4], doped solid [5], and quantum cavity system [6,7]. Equally exciting is the possibility that these results will lead to useful applications in the fields of telecommunications and optical buffer [8]. There are several possible techniques of implementing quantum light memory such as electromagnetically induced transparency (EIT) [9], coherent population oscillation (CPO) [10], Raman and Brillouin

amplification [11], among others [12]. EIT is the earliest technique to achieve the light storage in three-level system, which has been proved to be a powerful technique for eliminating the effect of a medium on a propagating beam of electromagnetic radiation [13–16]. Since EIT can retain the large and desirable nonlinear optical properties associated with the resonant response of a material system, it has been subjected to increasing investigation, especially in the realm of modern quantum optics and quantum information processing [17–21].

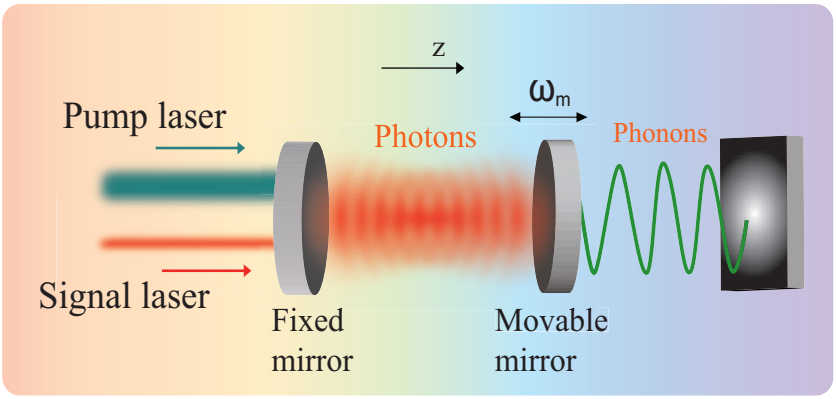
Moveover, based on EIT effect, a novel interference optical technique called optomechanically induced transparency (OMIT) has emerged recently, which leads to destructive interference for the excitation of an intracavity probe field, inducing a tunable transparency window for the probe beam [22,23]. Kippenberg et al. [23] have experimentally demonstrated that OMIT is equivalent to EIT in a cavity optomechanical system operated in the resolved sideband regime. Since then, the unique features of optomechanics system such as the on-chip integration, large bandwidth of operation, high Q factor and relatively long delay times have attracted many theoretical and experimental researchers' attention. For example, Teufel *et al.* [24] have realized the circuit cavity electromechanics in the strong-coupling regime. Most recently, Safavi-Naeini *et al.* [25] have theoretically achieved the light storage involving an optical waveguide coupled to an optomechanical crystal array, where light in the waveguide can be dynamically and coherently transferred into long-lived mechanical vibrations of the array. Some other work includes the report of storing optical information as a mechanical excitation in a silica optomechanical resonator [7]. All these works have experimentally demonstrated EIT and tunable optical delays in a nanoscale optomechanical crystal [26].

With this backdrop, in the present letter, we use the same protocol with Kippenberg *et al.* [23] and theoretically predict that in such system the quantum light memory can be achieved without any signal distortion. This is an all-optical controlled processing, in which the cavity photons interfere with the quantum signal pulse while dressing with the mechanical resonator mode (phonon mode). Recently, we theoretically demonstrated the nonlinear optical mass sensor with an optomechanical microresonator, which will provide a route towards the use of cavity optomechanical system in biomedical sensors, deposition monitors, chemical reaction monitors and other nonlinear measurement fields [27]. In the present article, theoretical analysis shows that light in the Fabry–Perot cavity can be dynamically and coherently transferred into long-lived vibrations of mechanical resonator, which provides a metastable state for photons. This dark-state polariton can be reaccelerated and converted back into a photon pulse by switching off and on the pump pulse on demand. Thanks to the long-lived mechanical resonator, the signal pulse can stay for a long time compared with other semiconductor systems. Our proposed quantum memory for light in an optomechanics system combines many of merits over other previous approaches, in that it allows the delay and release of pulse to be rapidly and all-optically controlled and achieves the long-term light memory due to the long lifetime of mechanical resonator. This scheme gives rise to the speculation of the ability of photons dressed by phonons, which can enable new possibilities for manipulating light through the manipulation of sound.

2. Theory

We consider the canonical situation in which a driven high-finesse optical cavity is coupled by momentum transfer of the cavity photons to a micro-mechanical resonator. The physical realization is shown in Figure 1, which is a Fabry–Perot-type cavity with the right mirror movable and perfectly reflecting and the left mirror fixed. The cavity is driven by a strong pump field and a weak signal field with frequencies of ω_p and ω_s, respectively. The circulating light inside the cavity deflects the movable mirror, and therefore, via radiation pressure, the driven high-frequency cavity mode is coupled to the high-Q, low-frequency mechanical mode. This mechanical mode can be treated as phonon mode. In the adiabatic limit, the input laser drives only one cavity mode ω_c, and the cavity-free spectrum range $L/2c$ (c is the speed of light in the vacuum and L is the effective cavity length [28,29]) is much larger than the frequency of movable mirror (ω_m). Hence, we can ignore the scattering of photons to other cavity modes.

Figure 1. Schematic diagram of an optomechanics system consisting of a Fabry–Perot cavity with a movable mirror in the presence of a strong pump field and a weak signal field.



In the simultaneous presence of a strong pump field and a weak signal field, the Hamiltonian of the total system can be written as [30,31]

$$H_t = H_c + H_m + H_{c-m} + H_{c-o} \tag{1}$$

$$H_c = \hbar\omega_c a^+ a \tag{2}$$

$$H_m = \frac{1}{2}\hbar\omega_m(p^2 + q^2) \tag{3}$$

$$H_{c-m} = \hbar G_0 a^+ a q \tag{4}$$

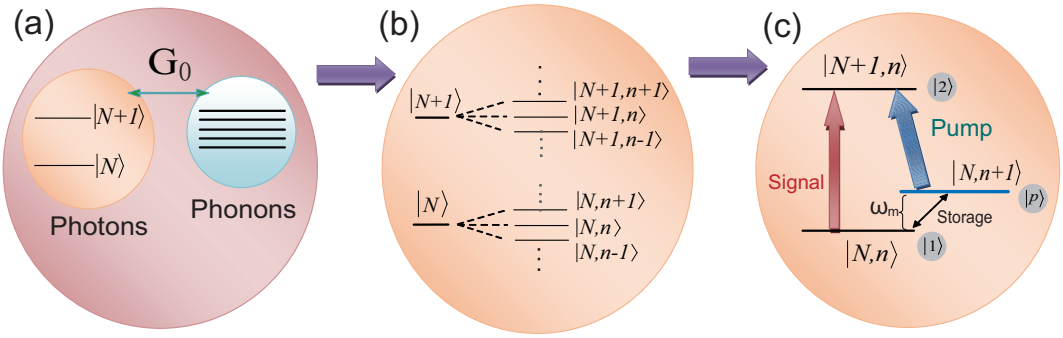
$$H_{c-o} = i\hbar E_p(a^+ e^{-i\omega_p t} - a e^{i\omega_p t}) + i\hbar E_s(a^+ e^{-i\omega_s t} - a e^{i\omega_s t}) \tag{5}$$

where H_c is the energy of the cavity mode, a and a^+ denote the annihilation and creation operators with commutation relation $[a, a^+] = 1$. H_m gives the energy of the mechanical mode, which is modeled as a harmonic oscillator and described by dimensionless position and momentum operators q and p with commutation relation $[q, p] = i$. The Hamiltonian H_{c-m} corresponds the coupling of the movable mirror to the cavity field via radiation pressure. The parameter $G_0 = (\omega_c/L\sqrt{\hbar/m\omega_m})$ is the coupling rate between the cavity and the oscillator (m is the effective mass of the mechanical mode [32]). H_{c-o} shows the interaction between cavity and classical inputs with frequencies ω_p and ω_s, E_p and E_s are slowly

varying envelope of the pump field and signal field, respectively, which related to the laser power P by $|E_p| = \sqrt{2P_p\kappa/\hbar\omega_c}$ and $|E_s| = \sqrt{2P_s\kappa/\hbar\omega_s}$ (κ is the cavity amplitude decay rate).

When the cavity photons couple to the mechanical resonator via the radiation pressure, the cavity photons will be dressed by an infinite number of possible phonon states as shown in Figure 2(b). Based on some experimental studies [23,26], we can only select the three-level as a lambda structure as shown in Figure 2(c). All the metastable states other than $|N, n+1\rangle$ are not considered due to the transition energy of $|N, n\rangle \rightarrow |N+1, n\rangle$ is much larger than that in $|N, n\rangle \rightarrow |N, n+1\rangle$ ($|N\rangle$ and $|n\rangle$ denote the number states of the photon and phonon, respectively).

Figure 2. The process of quantum memory for light (**a**) The initial model of the cavity photons and a mechanical resonator. G_0 is the coupling between the optical cavity and mechanical resonator. The vibration modes of mechanical resonator are treated as phonons modes; (**b**) The split energy levels of the cavity photon when dressing the vibration mode of mechanical resonator. $|N\rangle$ and $|n\rangle$ denote the number states of the photon and phonon, respectively; (**c**) The process of quantum memory for light, where the $|N, n+1\rangle \rightarrow |N+1, n\rangle$ and $|N, n\rangle \rightarrow |N+1, n\rangle$ transitions can be induced by the pump beam and signal beam, respectively. $|N, n+1\rangle$ is the metastable state caused by mechanical vibration. The signal pulse is temporarily stored in the state between $|N, n+1\rangle$ and $|N, n\rangle$.



The mechanism of quantum memory by adiabatic passage is well-known when the medium is made of lambda three-state atoms. The storage is then realized through the coherence of the atomic ground states. The idea of present work is to replace the lambda atoms by the photons dressed by the mechanical resonator as done in [23,26] and to take advantage of the long coherence induced by the resonator. In the following, in order to investigate the propagation properties of the quantum signal pulse in this optomechanics system, we follow the treatment outlined in [33,34] and use a quasi-one-dimensional model, consisting of one propagating beam $\hat{\varepsilon}(z,t)$ passing through a cavity of length L. $\hat{\varepsilon}(z,t)$ is a weak signal field that couples the state $|N, n\rangle$ and $|N+1, n\rangle$, and is related to the positive frequency part of the electric field by

$$\hat{E}_s^+(z,t) = \sqrt{\frac{\hbar\omega_c}{2\varepsilon_0 V}}\hat{\varepsilon}(z,t)e^{i(\omega_c/c)(z-ct)} \tag{6}$$

where ω_c is the frequency of the $|N, n\rangle \rightarrow |N+1, n\rangle$ transition (*i.e.*, the frequency of the cavity photon), V is the quantization volume of the electromagnetic field, c is the light velocity in vacuum, ε_0 is the free space permittivity, z is the direction along the cavity of length L. As compared with the three-level

atomic systems, here the $|N, n+1\rangle \rightarrow |N+1, n\rangle$ transition is driven by the strong pump laser (see Figure 2(c)). One can change the pump detuning simply by a tunable pump laser.

To perform a quantum analysis of the quantum signal light interacting with the cavity photons dressed by the resonator mode, it is useful to introduce locally averaged operators. $\aleph_z$ is the photon number in cavity along with z direction ($\aleph_z \gg 1$). Assuming the slowly-varying amplitude $\hat{\varepsilon}(z_j, t)$ does not change much, we can introduce the locally-averaged, slowly-varying operators

$$\hat{\sigma}_{\mu\nu}(z_j, t) = \frac{1}{\aleph_{z_j}} \sum_{z_j \in N_z} \hat{\sigma}^j_{\mu\nu}(t) e^{i(\omega_{\mu\nu}/c)(z_j - ct)} \tag{7}$$

where $\hat{\sigma}^j_{\mu\nu}(t) = |\mu^j(t)\rangle\langle\nu^j(t)|$ is the jth dressed photon state in cavity(μ,$\nu = 1, 2, p$, where for simplification we have used $|1\rangle$, $|p\rangle$, $|2\rangle$ to replace $|N, n\rangle$, $|N, n+1\rangle$, $|N+1, n\rangle$).

Going to the continuum limit, the effective interaction Hamiltonian for the reduced three-level system can be written in terms of the locally-averaged operators as [15]

$$\hat{H} = -\int \frac{\aleph\hbar}{L}[g\hat{\sigma}_{21}(z,t)\hat{\varepsilon}(z,t) + \Omega_p\hat{\sigma}_{2p}(z,t) + \mathrm{H.c.}]dz \tag{8}$$

where $\aleph$ is the photon number density in Fabry–Perot cavity, $g = \mu\sqrt{\omega_c/2\varepsilon_0 V\hbar}$ is the cavity photon-quantum signal field coupling constant. $\Omega_p = \langle n+1, N|\ \mu_{2p}\ \mathbf{E(t)}\ |N+1, n\rangle/2\hbar$ describes the coupling to the pump field and the transition, **E(t)** is the amplitude of the pump field, the "dipole moment" μ_{2p} corresponds to the transition between the state $|N+1, n\rangle$ and the state $|N, n+1\rangle$.

The evolution of the Heisenberg operator corresponding to the quantum signal field can be described in a slowly varying amplitude approximation by the propagation equation

$$(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z})\hat{\varepsilon}(z,t) = ig\aleph\hat{\sigma}_{12}(z,t) \tag{9}$$

The dressed photon state evolution is governed by a set of Heisenberg–Langevin equations

$$\frac{\partial}{\partial t}\hat{\sigma}_{\mu\nu} = -\gamma_{\mu\nu}\hat{\sigma}_{\mu\nu} + \frac{i}{\hbar}[\hat{H}, \hat{\sigma}_{\mu\nu}] + \hat{F}_{\mu\nu} \tag{10}$$

where $\gamma_{\mu\nu}$ are the cavity decay rates and $F_{\mu\nu}$ are δ-correlated Langevin noise operators.

Making the approximation that the quantum signal field intensity is much less than that of pump field and assuming all the photon are initially in the state $|N, n\rangle$, we can solve Equation (10) perturbatively to the first order in $g\hat{\varepsilon}/\Omega_p$ to obtain a pair of two equations

$$\frac{\partial}{\partial t}\hat{\sigma}_{12} = -\gamma_{12}\hat{\sigma}_{12} + ig\hat{\varepsilon} + i\Omega_p\hat{\sigma}_{1p} + \hat{F}_{12} \tag{11}$$

$$\frac{\partial}{\partial t}\hat{\sigma}_{1p} = -\gamma_{1p}\hat{\sigma}_{1p} + +i\Omega_p\hat{\sigma}_{12} + \hat{F}_{1p} \tag{12}$$

where γ_{12} and γ_{1p} are the decay rates of $|N+1, n\rangle \rightarrow |N, n\rangle$ and $|N, n+1\rangle \rightarrow |N, n\rangle$, respectively. In the above equations if we assume a sufficiently slow change of pump laser power and neglect the decay rate γ_{1p} of the photon, and then [33]

$$\hat{\sigma}_{12}(z,t) \approx -\frac{i}{\Omega_p}\frac{\partial}{\partial t}\hat{\sigma}_{1p}(z,t) \tag{13}$$

$$\hat{\sigma}_{1p}(z,t) \approx -g\frac{\hat{\varepsilon}(z,t)}{\Omega_p} \tag{14}$$

By combining with Equation (9), we can obtain the propagation equation of the quantum light pulse in the optomechanics system in the perturbative and the adiabatic limit as follows,

$$(\frac{\partial}{\partial t}+c\frac{\partial}{\partial z})\hat{\varepsilon}(z,t)=-\frac{g^2\bar{\aleph}}{\Omega_p}\frac{\partial}{\partial t}\frac{\hat{\varepsilon}(z,t)}{\Omega_p} \tag{15}$$

where $\bar{\aleph}$ denotes the mean photon numbers in the cavity. The group velocity of the signal field is given by

$$v_g=\frac{c}{1+(\bar{\aleph}g^2/\Omega_p^2)} \tag{16}$$

We can see from the above equation that for finite pump field and large $\bar{\aleph}$, the incident signal field can be slowed down significantly as shown in [26,35]. One can obtain an analytical solution of Equation (15) by introducing a new quantum field operator $\hat{\Psi}(z,t)$,

$$\hat{\Psi}(z,t) \;=\; \cos\theta(t)\hat{\varepsilon}(z,t)-\sin\theta(t)\sqrt{\bar{\aleph}}\hat{\sigma}_{1p}(z,t) \tag{17}$$

with

$$\cos\theta(t) \;=\; \frac{\Omega_p(t)}{\sqrt{\Omega_p^2(t)+g^2\bar{\aleph}}} \tag{18}$$

$$\sin\theta(t) \;=\; \frac{g\sqrt{\aleph}}{\sqrt{\Omega_p^2(t)+g^2\bar{\aleph}}} \tag{19}$$

$$\tan^2\theta(t) \;=\; \frac{g^2\bar{\aleph}}{\Omega_p^2(t)} \tag{20}$$

Introducing the adiabaticity parameter $\varepsilon\equiv(g\sqrt{\bar{\aleph}}T)^{-1}$ with T being a characteristic time, one can expand the equations of motion in power of ε. In the lowest order, such as the adiabatic limit, we can obtain [33]

$$\hat{\varepsilon}(z,t)=\cos\theta(t)\hat{\Psi}(z,t) \tag{21}$$

$$\sqrt{\bar{\aleph}}\hat{\sigma}_{1p}=-\sin\theta(t)\hat{\Psi}(z,t)e^{-i\Delta kz} \tag{22}$$

Furthermore, the new operator $\hat{\Psi}(z,t)$ obeys the following equation of motion

$$[\frac{\partial}{\partial t}+c\cos^2\theta(t)\frac{\partial}{\partial z}]\hat{\Psi}(z,t)=0 \tag{23}$$

which describes a shape-preserving propagation with velocity $v=v_g(t)=c\cos^2\theta(t)$. It is should be noted here that $\hat{\sigma}_{1p}(z,t)$ in Equations (17) and (22) corresponds to the creation operator of the $|N,n\rangle\rightarrow|N,n+1\rangle$, which is different from atomic spin operator in three-level atomic systems [36]. In the linear limit, this new operator satisfies the Bosonic commutation relation and we can refer this new Bosonic particle to phonon-polariton.

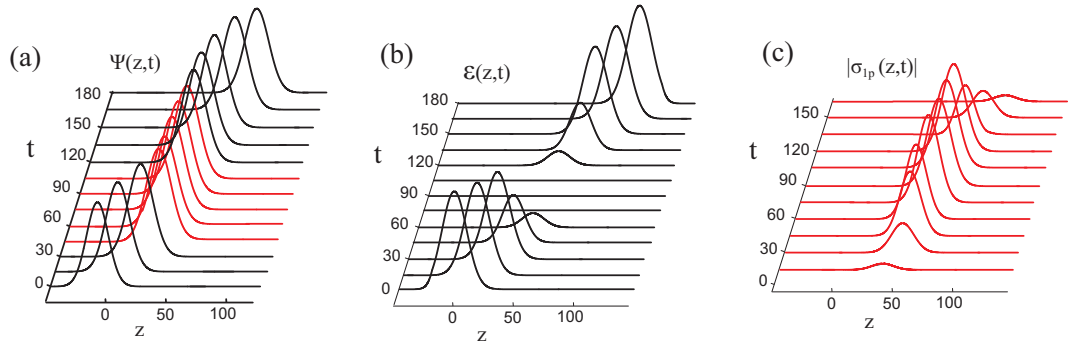
3. Results and Discussions

Equation (23) illustrates a shape- and quantum-state preserving propagation

$$\hat{\Psi}(z,t) = \hat{\Psi}(z,t)(z - c\int_0^t d\tau \cos^2\theta(\tau), 0) \tag{24}$$

It is obvious that by adiabatically rotating θ from 0 to $\pi/2$ via a tunable pump laser, one can decelerate and stop the signal pulse on demand. When $\theta \to 0$, $\Omega^2 \gg g^2\bar{\aleph}$, which corresponds to the strong external pump field, the polariton has purely photonic character $\hat{\Psi}(z,t) = \hat{\varepsilon}(z,t)$ and the signal propagation velocity equals to the vacuum speed of light. For $\theta \to \pi/2$, the polariton becomes phonon-like, $\hat{\Psi}(z,t) = -\sqrt{\bar{\aleph}}\hat{\sigma_{1p}}e^{i\Delta kz}$, and its propagation velocity approaches zero. In this process, the quantum signal field are mapped onto a phonon-polariton which are different from the atomic spins in three-level atomic systems. This phonon-polariton can then be reaccelerated to the vacuum speed of light in which the stored quantum states are transferred back to the photonic state. Figure 3 displays the coherent amplitude of a dark-state phonon-polariton, which results from an initial light pulse, as well as the corresponding field and matter components.

Figure 3. Quantum light propagation of a dark-state polariton with envelope $exp\{-(z/10)^2\}$, according to the expression $\cot\theta(t) = 100(1 - 0.5\tanh[0.1(t-15)] + 0.5tanh[0.1(t-125)])$. (**a**) The coherent amplitude of the polariton $\hat{\Psi}(z,t) = \langle\hat{\Psi}(z,t)\rangle$ as functions of time t and length z; (**b**)The process of light storage. $\varepsilon(z,t) = \langle\hat{\varepsilon}(z,t)\rangle$; (**c**)The matter component part $\mid \sigma_{1p} \mid = \mid \langle\hat{\sigma_{1p}}\rangle \mid$,where $c = 1$.



Here, we show an example of the signal laser storage and recall on demand. First, under constant pump laser E_p, the signal optical pulse E_s completely enters the optomechanics system consisting of a Fabry–Perot cavity and a mechanical resonator (Figure 1). Once the signal pulse inside, the amplitude of E_s will turn to reduce gradually. However, when the pump laser is switched off rapidly, the signal laser information is stored in the coherence between the $|N,n\rangle$ and $|N,n+1\rangle$ subbands. The maximum storage time is set by the mechanical decay rate, $\sim 1/\gamma_m$. Otherwise, turning on the pump laser results in a retrieved signal pulse. There is no additional distortion during the storage because the width of the signal pulse spectrum is very much less than the width of the OMIT window [26,35].

As we know, when storing light in atomic systems using EIT, characteristics of the field are recorded as a spin wave in the atomic ensemble. The storage time is determined by the coherence times of the hyperfine transitions. However, in optomechanical systems, the term "$storage$" implies the conversion

of signal pulse into the phonon coherence σ_{1p}, whose lifetime is determined by the lifetime of mechanical resonator. Towards the realistic optomechanical systems, Gröblacher *et al.* [37] recently experimentally showed that in optomechanical system, the mechanical quality factor $Q = 6700$ and the decay rate $\gamma_m = 2\pi \times 140Hz$, which corresponds to the lifetime of resonator is $7ms$. In this case, the light storage time can reach an ideal situation if we select a long vibration lifetime of mechanical resonator.

During the theoretical simulations, the noise processes are neglected. Actually, there are various noise sources that will limit the performance of the cavity optomechanical system for quantum light memory. Among these noise sources, the thermal noise of the mechanical motion in the linear regime is the dominant noise of error in our device, which through the optomechanical coupling can be mapped into noise power in the cavity optomechanical system. This is essentially coursed by the conversion of signal pulse into the phonon coherence to yield longer storage times. When $\Gamma_{opt} \gg \gamma_m$ the optical noise power is given by [25]

$$P_{nois} \approx \frac{\hbar\omega_c}{2}\frac{\kappa_{ex}}{\kappa}(\gamma_m \overline{n}_{th} + \Gamma_{opt}(\frac{\kappa}{4\omega_m})^2) \tag{25}$$

$\Gamma_{opt} = 4\Omega_p/\kappa$ is defined as an effective optically induced cooling rate, $\overline{n}_{th} = (e^{\hbar\omega_m/k_B T_b} - 1)^{-1}$ is the Bose occupation number at the mechanical frequency and T_b is the bath temperature. The first term in the brackets corresponds to the up-conversion of thermal mechanical noise, which can be suppressed by working at lower temperatures, while the second term corresponds to optically-induced Stokes scattering, which can be suppressed with good sideband resolution, $\kappa \ll \omega_m$. At room temperature, $\overline{n}_{th} \approx k_B T_b/\hbar\omega_m$ is large and thermal noise will dominate. However, as the thermal noise scales inversely with mechanical frequency, it makes noise remain tolerable even at room temperature to employ high-frequency mechanical oscillators. At very low temperature, these high-frequency oscillators can be thermally cooled to the ground state, which enables the optomechanical system we proposed here as a suitable choice for quantum memory for photons.

4. Conclusions

In conclusion, we have theoretically achieved the quantum light memory by tuning another pump laser in an optomechanics system consisting of a Fabry–Perot cavity and a mechanical resonator. Theoretical analysis shows that it is possible to store and retrieve the quantum signal light by switching off and on the pump laser. This light storage is based on the coherence between quantum light and mechanical resonator via cavity photons. The storage time of signal pulse can be determined by the vibration lifetime of mechanical resonator. Taking advantages of this optical device, it has prepared the stage of all optical buffers, quantum optical processing and other related biological [38], chemical applications [39]. Finally, we hope that our predictions in the present work can be testified by experiments in the near future.

Acknowledgments

The part of this work has been supported by National Natural Science Foundation of China (No.10974133 and No.11274230), the National Ministry of Education Program for Ph.D.

References

1. De Riedmatten, H. Quantum optics: Light storage at record bandwidths. *Nat. Photonics* **2010**, *4*, 206.
2. Hedges, M.P.; Longdell, J.J.; Li, Y.M.; Sellars, J. Efficient quantum memory for light. *Nature* **2010**, *465*, 1052.
3. Karpa, L.; Nikoghosyan, G.; Vewinger, F.; Fleischhauer, M.; Weitz, M. Frequency matching in light-storage spectroscopy of atomic raman transitions. *Phys. Rev. Lett.* **2009**, *103*, 093601.
4. Dao, T.L.; Kollath, C.; Carusotto, I.; Köhl, M. All-optical pump-and-probe detection of two-time correlations in a Fermi gas. *Phys. Rev. A* **2010**, *81*, 043626.
5. Beil, F.; Buschbeck, M.; Heinze, G.; Halfmann, T. Light storage in a doped solid enhanced by feedback-controlled pulse shaping. *Phys. Rev. A* **2010**, *81* , 053801.
6. Elshaari, A.W.; Aboketaf, A.; Preble, S.F. Controlled storage of light in silicon cavities. *Opt. Express* **2010**, *18*, 3014.
7. Fiore, V.; Yang, Y.; Kuzyk, M.C.; Barbour, R.; Tian, L.; Wang, H. Storing optical information as a mechanical excitation in a silica optomechanical resonator. *Phys. Rev. Lett.* **2011**, *107*, 133601.
8. Lobino, M.; Kupchak, C.; Figueroa, E.; Lvovsky, A.I. Memory for light as a quantum process. *Phys. Rev. Lett.* **2009**, *102*, 203601.
9. Tu, Y.; Zhang, G.; Zhai, Z.; Xu, J. Angular multiplexing storage of light pulses and addressable optical buffer memory in $Pr^{3+} : Y_2SiO_5$ based on electromagnetically induced transparency. *Phys. Rev. A* **2009**, *80*, 033816.
10. Xue, W.Q.; Sales, S.; Capmany, J.; Mørk, J. Wideband 360° microwave photonic phase shifter based on slow light in semiconductor optical amplifiers. *Opt. Express* **2010**, *18*, 6156.
11. Kalosha, V.P.; Li, W.H.; Wang, F.; Chen, L.; Bao, X.Y. Frequency-shifted light storage via stimulated Brillouin scattering in optical fibers. *Opt. Lett.* **2008**, *33*, 2848.
12. Elshaari, A.W.; Aboketaf, A.; Preble, S.F. Controlled storage of light in silicon cavities. *Opt. Express* **2010**, *18*, 3014.
13. Harris, S.E.; Field, J.E.; Imamoğlu, A. Nonlinear optical processes using electromagnetically induced transparency. *Phys. Rev. Lett.* **1990**, *64*, 1107.
14. Harris, S.E. Electromagnetically induced transparency. *Phys. Today* **1997**, *50*, 36.
15. Lukin, M.D. Colloquium: Trapping and manipulating photon states in atomic ensembles. *Mod. Rev. Phys.* **2003**, *75*, 457.
16. Fleischhauer, M.A. Imamoglu, A.; Marangos, J.P. Electromagnetically induced transparency: Optics in coherent media. *Rev. Mod. Phys.* **2005**, *77*, 633.
17. Hu, X.M.; Sun, H.; Wang, F. Scalable network of quadrangle entanglements via multiple phase-dependent electromagnetically induced transparency. *Phys. Rev. A* **2010**, *82*, 045807.
18. Hammerer, K.; Sørensen, A.S.; Polzik, E.S. Quantum interface between light and atomic ensembles. *Rev. Mod. Phys.* **2010**, *82*, 1041.
19. Arikawa, M.; Honda, K.; Akamatsu, D.; Nagatsuka, S.; Akiba, K.; Furusawa, A.; Kozuma, M. Quantum memory of a squeezed vacuum for arbitrary frequency sidebands. *Phys. Rev. A* **2010**, *81*, 021605.

20. Litvinov, A.N.; Kazakov, G.A.; Matisov, B.G. Electromagnetically induced transparency in nanocells. *J. Phys. B* **2009**, *42*, 165402.
21. Ding, J.L.; Hou, B.P.; Wang, S.J. Squeezing survival and transfer in single and double electromagnetically induced transparency. *J. Phys. B* **2010**, *43* , 225502.
22. Huang, S.; Agarwal, G.S. Normal-mode splitting and antibunching in Stokes and anti-Stokes processes in cavity optomechanics: Radiation-pressure-induced four-wave-mixing cavity optomechanics. *Phys. Rev. A* **2010**, *81*, 033830.
23. Weis, S.; Rivière, R.; Deléglise, S.; Gavartin, E.; Arcizet, O.; Schliesser, A.; Kippenberg, T.J. Optomechanically induced transparency. *Science* **2010**, *330*, 1520.
24. Teufel, J.D.; Li, D.; Allman, M.S.; Cicak, K.; Sirois, A.J.; Whittaker, J.D.; Simmonds, R.W. Circuit cavity electromechanics in the strong-coupling regime. *Nature* **2011**, *471*, 204.
25. Chang, D.E.; Safavi-Naeini, A.H.; Hafezi, M.; Painter, O. Slowing and stopping light using an optomechanical crystal array. *New J. Phys.* **2011**, *13*, 023003.
26. Safavi-Naeini, A.H.; Alegre, T.P.M.; Chan, J.; Eichenfield, M.; Winger, M.; Lin, Q.; Hill, J.T.; Chang, D.E.; Painter, O. Electromagnetically induced transparency and slow light with optomechanics. *Nature* **2011**, *472*, 69.
27. Li, J.J.; Zhu, K.D. Nonlinear optical mass sensor with an optomechanical microresonator. *Appl. Phys. Lett.* **2012**, *101*, 141905.
28. Kippenberg, T.J.; Rokhsari, H.; Carmon, T.; Scherer, A.; Vahala, K.J. Analysis of radiation-pressure induced mechanical oscillation of an optical microcavity. *Phys. Rev. Lett.* **2005**, *95*, 033901.
29. Schliesser, A.; Del'Haye, P.; Nooshi, N.; Vahala, K.J.; Kippenberg, T.J. Radiation pressure cooling of a micromechanical oscillator using dynamical backaction. *Phys. Rev. Lett.* **2006**, *97*, 243905.
30. Genes, C.; Vitali, D.; Tombesi, P.; Gigan, S.; Aspelmeyer, M. Ground-state cooling of a micromechanical oscillator: Comparing cold damping and cavity-assisted cooling schemes. *Phys. Rev. A* **2008**, *77*, 033804.
31. Giovannetti, V.; Vitali, D. Phase-noise measurement in a cavity with a movable mirror undergoing quantum Brownian motion. *Phys. Rev. A* **2001**, *63*, 023812.
32. Pinard, M.; Hadjar, Y.; Heidmann, A. Effective mass in quantum effects of radiation pressure. *Eur. Phys. J. D* **1999**, *7*, 107.
33. Fleischhauer, M.; Lukin, M.D. Effective mass in quantum effects of radiation pressure. *Phys. Rev. A* **2002**, *65*, 022314.
34. Peng, A.; Johnsson, M.; Bowen, W.P.; Lam, P.K.; Bachor, H.A.; Hope, J.J. Squeezing and entanglement delay using slow light. *Phys. Rev. A* **2005**, *71*, 033809.
35. He, W.; Li, J.J.; Zhu, K.D. All-optical control of light group velocity with a cavity optomechanical System. **2010**, ArXiv: 1011.4993v1.
36. Fleischhauer, M.; Lukin, M.D. Quantum memory for photons: Dark-state polaritons. *Phys. Rev. Lett.* **2000**, *84*, 5094.
37. Gröblacher, S.; Hammerer, K.; Vanner, M.R.; Aspelmeyer, M. Observation of strong coupling between a micromechanical resonator and an optical cavity field. *Nature* **2009**, *460* , 724.

38. Dos Santos, D.S.; Bassi, A., Jr.; Rodrigues, J.J.; Misoguti, L., Jr.; Oliveira, O.N., Jr.; Mendonça, C.R. Light-induced storage in layer-by-layer films of chitosan and an Azo dye Biomacromolecules. *Biomacromolecules* **2003**, *4*, 1502.
39. Yu, J.X.; Mathew, S.; Flavel, B.S.; Johnston, M.R.; Shapter, J.G. Ruthenium porphyrin functionalized single-walled carbon nanotube arrayss a step toward light harvesting antenna and multibit information storage. *J. Am. Chem. Soc.* **2008**, *130*, 8788.