Quantum storage via refractive-index control

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Abstract

Off-resonant Raman interaction of a single-photon wave packet and a classical control field in an atomic medium with controlled refractive index is investigated. It is shown that a continuous change of refractive index during the interaction leads to the mapping of a single photon state to a superposition of atomic collective excitations (spin waves) with different wave vectors and visa versa. The suitability of refractive index control for developing multichannel quantum memories is discussed and possible schemes of implementation are considered.

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I. INTRODUCTION

During the past decade the optical quantum memories have became one of the active areas of research in the field of quantum optics and quantum information (see the Reviews [1–4]). Such devices are considered as basic ingredient for scalable linear-optical quantum computers and efficient quantum repeaters. For practical quantum information applications, it is necessary to develop memories which could store quantum states of light with close to 100% efficiency and fidelity, and provide long and controllable storage times or delay-bandwidth products. In this respect significant experimental progress has been achieved in demonstration of optical quantum storage using electromagnetically induced transparency [5–8], photon echo induced by controlled reversible inhomogeneous broadening [9–13] or by atomic-frequency comb [14–20], and off-resonance Raman interaction [21–23]. Optical quantum memories are usually assumed to store and recall optical pulses, such as single-photon wave packets, exploiting inhomogeneous broadened transitions or modulated control fields. In the present work we suggest one more possibility. By considering quantum storage based on off-resonant Raman interaction, we show that manipulation of refractive index in a three-level resonant medium allows one to store and recall single-photon wave packets without using inhomogeneous broadening of the atomic transitions or manipulating the amplitude of the control field. A single-photon wave packet may be reversibly mapped to a superposition of atomic collective excitations with different wave vectors, which is analogues to that of orthogonal subradiant states created in an extended atomic ensemble [24].

As well as providing an interesting possibility for storage, the refractive index control may also be useful for optimizing multiplexing regimes of multimode quantum memories, development of which is important in the prospect of both quantum communication [25, 26] and computation [27]. Particularly, multimode memories can significantly increase the quantum communication rate for short storage times. Different ways of multiplexing has been suggested [28–31], among which time domain multiplexing is currently most demanded from the view point of fiber optical communication. Being combined with any approach to quantum storage mentioned above, the refractive index control provides an additional degree of freedom for multiplexing thereby improving capacity of a multimode quantum storage device or allowing operation in a multichannel regime. Such a multiplexing method is closely connected with the angular [28] or holographic [31] ones since it also resorts to

phase-matching conditions in an extended atomic ensemble, but it does not exploit different spatial modes of the field. In effect, the additional multiplexing capacity is based on the possibility to use frequency and wavelength of the field in a storage material as independent parameters.

The paper is organized as follows. In Sec. II, we analyze the storage and retrieval of single-photon wave packets via refractive index control during off-resonant Raman interaction. In Sec. III, suitability of refractive index control for developing multichannel quantum memories is discussed. In Sec. IV, we consider possible ways of refractive index manipulation and some implementation issues.

II. STORAGE AND RETRIEVAL OF SINGLE-PHOTON WAVE PACKETS

As a basic model we consider cavity-assisted quantum storage, which is motivated by the following reasons. First, enclosing an atomic ensemble in a cavity makes it possible to achieve high efficiency of quantum storage with optically thin materials. This may be especially useful for considered off-resonant Raman interaction since the cross-section of the two-photon transition is usually small. Second, there is no need for backward retrieval when optically thin materials are used, which relieves one of having to perform phase conjugation of the atomic states used for storage.

We consider a system of $N \gg 1$ identical three-level atoms which are placed in a single-ended ring cavity and interact with a weak quantum field (single-photon wave packet) to be stored and with a strong classical control field (Fig. 1). The atoms have a Λ-type level structure, the fields are Raman resonant to the lowest (spin) transition, and the cavity is resonant to the quantum field. We assume that the atoms are stationary like impurities embedded in a solid state material or cold atoms in the magnetic trap. The interaction volume is supposed to have a large Fresnel number, which allows us to take advantage of one-dimensional approximation. The Hamiltonian of the three-level system in the dipole and rotating wave approximations is

$$H = H_0 + V, \tag{1}$$

where

$$H_0 = \hbar\omega a^\dagger a + \sum_{j=1}^{N} \left(\hbar\omega_2 \sigma_{22}^j + \hbar\omega_3 \sigma_{33}^j \right), \tag{2}$$

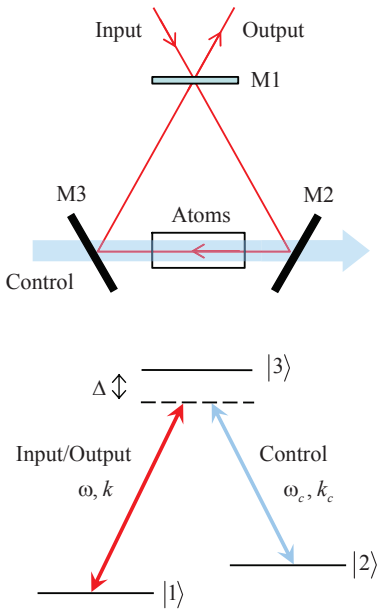

FIG. 1: (Color online) Schematic of a quantum memory device (above) and atomic level structure (below). Mirrors M2 and M3 are perfectly reflecting for the single-photon field and fully transmitting for the control field, M1 is a partially transmitting mirror. The difference of wave vectors $k - k_c$ is modulated via refractive index control during off-resonant Raman interaction.

$$V = -\hbar \sum_{j=1}^{N} \left(\Omega(t) \sigma_{32}^{j} \, e^{ik_c z_j - i\omega_c t} + ga\sigma_{31}^{j} \, e^{ikz_j} \right) + \text{H.c.} \tag{3}$$

Here $\sigma_{mn}^{j} = |m_j\rangle\langle n_j|$ are the atomic operators, $|n_j\rangle$ is the nth state ($n = 1, 2, 3$) of jth atom with the energy $\hbar\omega_n$ ($\omega_1 = 0 < \omega_2 < \omega_3$), z_j is the position of the jth atom, a is the photon annihilation operator in the cavity mode, $k_c = \omega_c n_c/c$ and $k = \omega n/c$ are the wave vectors of the classical and quantum fields, respectively, n_c and n are refractive indices at the frequencies ω_c and ω, $\Omega(t)$ is the Rabi frequency of the classical field, and g is the coupling constant between the atoms and the quantized field mode. The values of n_c and n are considered below as parameters changing in time. Since variations of them are supposed to be much less than 1, we leave such time dependence only in phase factors and ignore it in the factors Ω and g as functions of refractive indices.

In the Heisenberg picture, we define the following slowly varying atomic operators: $P_j = \sigma_{13}^{j} \, e^{i\omega t}$, $S_j = \sigma_{12}^{j} \, e^{i(\omega - \omega_c)t}$, and cavity field amplitude $\mathcal{E} = a \, e^{i\omega t}$. From the input-output

relations for the cavity field [32, 33] we have

$$\mathcal{E}_{\rm out}(t) = \sqrt{2\kappa}\mathcal{E}(t) - \mathcal{E}_{\rm in}(t), \tag{4}$$

where 2κ is the cavity decay rate and $\mathcal{E}_{\rm in}$ ($\mathcal{E}_{\rm out}$) is the input (output) field (a single-photon wave-packet). From the Heisenberg—Langevin equations, assuming that all the population is in the ground state initially and taking into account that the quantum field is weak, we find

$$\dot{P}_j = -(\gamma_P + i\Delta)P_j + i\Omega S_j\, e^{ik_c z_j} + ig\mathcal{E}\, e^{ikz_j}, \tag{5}$$

$$\dot{S}_j = -(\gamma_S + i\Delta_S)S_j + i\Omega^* P_j\, e^{-ik_c z_j}, \tag{6}$$

$$\dot{\mathcal{E}} = -\kappa\mathcal{E} + ig\sum_j P_j\, e^{-ikz_j} + \sqrt{2\kappa}\,\mathcal{E}_{\rm in}. \tag{7}$$

Here γ_P and γ_S are the rates of dephasing, which in general case include both homogeneous and inhomogeneous broadening of the resonant transitions, $\Delta = \omega_3 - \omega$ is a one-photon detuning, and $\Delta_S = \omega_2 + \omega_c - \omega$ is a two-photon detuning. We have not included the Langevin noise atomic operators since they make no contribution to normally ordered expectation values in consistence with the approximation that almost all atoms remain in the ground state (see, e.g., [34] for discussion). Finally, in the Raman limit, when the single-photon detuning is sufficiently large, adiabatically eliminating P_j in Eqs. (5) and (6), and going to the collective atomic operators

$$S_q = \frac{1}{\sqrt{N}}\sum_j S_j\, e^{-iqz_j}, \tag{8}$$

we obtain

$$\dot{S}_q = -\gamma' S_q + ig'\sqrt{N}\phi(k - k_c - q)\mathcal{E}, \tag{9}$$

$$\dot{\mathcal{E}} = -\kappa\mathcal{E} + \sqrt{2\kappa}\,\mathcal{E}_{\rm in} + ig'^*\sqrt{N}\sum_q \phi(q + k_c - k)S_q. \tag{10}$$

Here $\phi(q) = 1/N\sum_{j=1}^{N}\exp(iqz_j)$ is the diffraction function, $\gamma' = \gamma_S + \gamma_P|\Omega|^2/\Delta^2$, $g' = g\Omega^*/\Delta$, and the resulting frequency shift $\Delta' = \Delta_S + |\Omega|^2/\Delta$ has been compensated by tuning the coupling field frequency. The wave vectors q are multiples of $2\pi/L$, where L is the length of the atomic medium.

The phase mismatching factors $\phi(q)$, which are usually ignored on the assumption that a single spatial mode of the spin coherence is excited and phase-matching is perfect, are now considered. Suppose that we can manipulate the difference $q + k_c - k$ by refractive index control without changing frequencies and propagation directions of the interacting fields. We discuss possible ways of implementation below. For now it is sufficient to consider the case when one of the wave vectors, say k_c of the control field, is changed linearly in time during the interaction so that $q + k_c - k = (\omega_c/c)\dot{n}_c(t - t_q)$, where t_q is the moment when $q + k_c - k = 0$ for a given q. Then

$$\phi[\pm(q + k_c - k)] = e^{\pm i\beta(t-t_q)}\,\mathrm{sinc}[\beta(t - t_q)], \tag{11}$$

where $\mathrm{sinc}(x) = \sin(x)/x$ and $\beta = (\omega_c/c)(L/2)\dot{n}_c$. The phase factors $e^{\pm i\beta t}$ may be compensated by linear or sawtooth phase modulation of the control field. In such a situation, the phase of the control field remains constant at the point $z = L/2$ during the refractive index change. As a result, Eqs. (9) and (10) take the form

$$\dot{S}_q = -\gamma' S_q + ig'\sqrt{N}\, e^{i\beta t_q}\,\mathrm{sinc}(\beta(t - t_q))\,\mathcal{E}, \tag{12}$$

$$\dot{\mathcal{E}} = -\kappa\mathcal{E} + \sqrt{2\kappa}\,\mathcal{E}_{\mathrm{in}} + ig'^{*}\sqrt{N}\sum_q e^{-i\beta t_q}\,\mathrm{sinc}(\beta(t - t_q))\,S_q. \tag{13}$$

Now it is possible to consider storage and retrieval of a single-photon wave packet. Let the atomic system interacts with the quantum field during the time interval $[-T, 0]$ with the initial condition $S_q(-T) = 0, \forall q$. Then from Eqs. (12) and (13) we have

$$S_q(t) = ig'\sqrt{N}F_q(-T, t, \mathcal{E})\, e^{i\beta t_q - \gamma' t}, \tag{14}$$

$$\dot{\mathcal{E}}(t) = -\kappa\mathcal{E}(t) + \sqrt{2\kappa}\mathcal{E}_{\mathrm{in}}(t) - |g'|^2 N\sum_q \mathrm{sinc}[\beta(t - t_q)]F_q(-T, t, \mathcal{E})\, e^{-\gamma' t}, \tag{15}$$

where $F_q(-T, t, \mathcal{E}) = \int_{-T}^{t}\mathcal{E}(\tau)\,\mathrm{sinc}[\beta(\tau - t_q)]\, e^{\gamma'\tau}d\tau$. If the cavity field $\mathcal{E}$ varies slowly than $\delta = \pi/\beta$, and $\gamma'\delta \ll 1$, then Eq. (15) takes the form

$$\dot{\mathcal{E}} = -\kappa\mathcal{E} + \sqrt{2\kappa}\mathcal{E}_{\mathrm{in}} - \Gamma\mathcal{E}, \tag{16}$$

where $\Gamma = |g'|^2 N\delta/2$. Then the cavity field can be adiabatically eliminated provided that $\Gamma + \kappa$ is much greater than the bandwidth of the input field, which gives $\mathcal{E} = \sqrt{2\kappa}(\kappa + \Gamma)^{-1}\mathcal{E}_{\mathrm{in}}$,

and from Eq. (14) we find

$$S_q(0) = \frac{ig'\sqrt{N}\sqrt{2\kappa}}{\kappa+\Gamma} F_q(-T, 0, \mathcal{E}_{\rm in})\, e^{i\beta t_q}$$
$$= \frac{ig'\sqrt{N}\sqrt{2\kappa}}{\kappa+\Gamma} \frac{\pi}{\beta} \mathcal{E}_{\rm in}(t_q)\, e^{(i\beta+\gamma')t_q}. \quad (17)$$

Equation (17) describes the mapping of an input single-photon wave packet to a superposition of collective excitations (spin waves) with different wave vectors. The quantum storage mechanism may be explained qualitatively in the following way. During absorption of the input field via off-resonant Raman interaction a spin wave is created with a wave vector $q = k - k_c$, which is the difference between wave vectors of the input and control fields. If we change refractive index during the interaction, the input field amplitude at different moments of time is mapped to the amplitude of the spin waves with different wave vectors q. This is possible due to phase matching condition which allows us to switch the collective interaction between the input field and atoms from one spin wave to another. As a result, by the end of the storage process a coherence grating is created on the spin transition, which is a superposition of the spin waves. If the rate of refractive index change is sufficiently large so that δ is smaller than the fastest time scale of the input pulse, then the temporal shape $\mathcal{E}_{\rm in}(t)$ is imprinted on the amplitude $S_q(0)$ as a function of wave vector q. Retrieval is achieved by off-resonant interaction of the atomic system with the control field when the values of n_c that used for storage are scanned again. The retrieved pulse is actually a superposition of pulses created via Raman interaction from different spin waves satisfying phase matching condition for different wave vectors of the control field. Suppose, e.g., that the time dependence of refractive index during the time interval $[0, T]$, when $\mathcal{E}_{\rm in}(t) = 0$, is reversed. In this case, instead of Eq. (15) we have

$$\dot{\mathcal{E}}(t) = -\kappa\mathcal{E}(t)$$
$$- \frac{|g'|^2 N\sqrt{2\kappa}}{\kappa+\Gamma} \sum_q \mathrm{sinc}[\beta(t+t_q)] F_q(-T, 0, \mathcal{E}_{\rm in})\, e^{-\gamma' t}. \quad (18)$$

For slow-varying $\mathcal{E}_{\rm in}$ this equation takes the form

$$\dot{\mathcal{E}}(t) = -\kappa\mathcal{E}(t) - \frac{\Gamma\sqrt{2\kappa}}{\kappa+\Gamma} \mathcal{E}_{\rm in}(-t)\, e^{-2\gamma' t}. \quad (19)$$

Finally, after adiabatic elimination of the cavity field and using Eq. (4) we obtain

$$\mathcal{E}_{\rm out}(t) = -\frac{2\Gamma}{\kappa+\Gamma} \mathcal{E}_{\rm in}(-t)\, e^{-2\gamma' t}. \quad (20)$$

The solution Eq. (20) is exactly the same as that in the cavity-assisted storage with inhomogeneous broadening [35, 36]. The output field becomes time-reversed replica of the input field provided that the duration of wave packet is much smaller that the decay time $1/\gamma'$, and the efficiency of the storage followed by retrieval is maximum under impedance-matching condition $\kappa = \Gamma$. The latter corresponds to the situation when a directly reflected field and a transmitted field coming from the circulating one inside the cavity cancel each other, which allows one to put all input signal into the cavity, thereby achieving the maximum efficiency of the storage. The only difference in the solutions is the collective absorption/emission rate Γ, which in our case takes the form

$$\Gamma = \frac{g^2 N |\Omega|^2}{\Delta^2} \frac{\pi}{2\beta}. \tag{21}$$

It means that the time interval $\delta = \pi/\beta$, which is actually the time interval between two adjacent orthogonal spin states created upon refractive index control, is analogous to inhomogeneous life-time, i.e., reversal inhomogeneous linewidth. We see that a single-photon wave packet can be effectively stored and reproduced via refractive index control in a three-level system without inhomogeneous broadening and without modulating the Rabi frequency of the control field during the interaction. It is also important that the time dependence of the refractive index need not be reversed during the retrieval. If the values of n_c during retrieval are ordered like those during storage, Eq. (20) is replaced by

$$\mathcal{E}_{\rm out}(t) = -\frac{2\Gamma}{\kappa + \Gamma} \mathcal{E}_{\rm in}(t - T)\, e^{-\gamma' T}. \tag{22}$$

Thus a single-photon wave packet may be reconstructed without time reversal so that its temporal shape be not deformed by the dephasing process.

The total change of refractive index during storage or retrieval is

$$\Delta n = \dot{n}_c T = \left(\frac{T}{\delta}\right) \frac{\lambda}{L}, \tag{23}$$

where $\lambda = 2\pi c/\omega_c$. Numerics show that a Gaussian pulse with the duration (FWHM) as short as 2δ can be stored and recalled with the efficiency 0.99 provided that γ' is small enough (Fig. 2). Therefore, taking $T/\delta \sim 1$ and $\lambda/L \sim 10^{-5}$ we have $\Delta n \sim 10^{-5}$, which may be considered as the minimum refractive index increment needed for storage of a single pulse under typical experimental conditions. The ratio between the total accessible range of refractive index change and this minimum value determines the number of pulses can be stored in a series, i.e., the mode capacity of quantum memory.

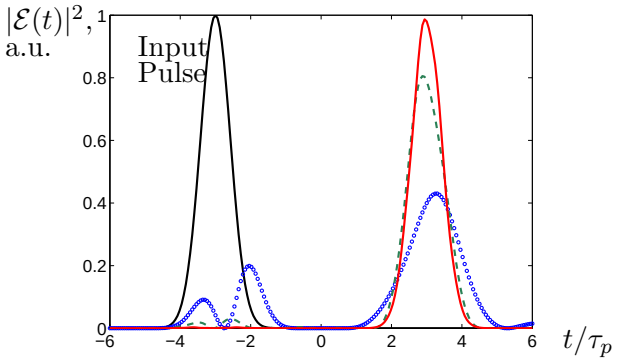

FIG. 2: (Color online) Storage and retrieval of a Gaussian pulse for different values of $\delta = \pi/\beta$. The black solid curve is the input pulse of duration τ_p. Other curves represent the output field for the cases $\delta = 2\tau_p$ (blue circles), $\delta = \tau_p$ (green dashed line), $\delta = \tau_p/2$ (red solid line). These curves were obtained by numerically solving Eqs. (12) and (13), treated as complex number equations, with the condition $\kappa = \Gamma$ and $\gamma'\tau_p \ll 1$. The sign of dk_c/dt is changed at the moment $t = 0$.

The impedance matching condition $\Gamma = \kappa$ may be written in the following form:

$$C\gamma'\delta/2 = 1, \tag{24}$$

where $C = g^2N/\kappa\gamma = |g'|^2N/\kappa\gamma'$ is the cooperativity parameter, which could also be expressed in terms of the cavity finesse $\mathcal{F}$ and resonant absorption coefficient α as $C = \alpha L\mathcal{F}/2\pi$, provided that the resonant medium fills the cavity. Thus to store and recall pulses broader than the Raman linewidth γ', we need to increase C above 1 appropriately. In general case, γ' consists of homogeneous and inhomogeneous contributions. If we do nothing with the latter, the delay-bandwidth product of quantum storage (or multimode capacity) proves to be of the order of C. It may be increased by choosing materials with a narrower line-width or shorter pulses with appropriate increasing the cooperative parameter. But if we are able to reverse the inhomogeneous broadening, the delay-bandwidth product may be increased significantly depending on a residual noncontrollable broadening, say $\gamma'_{\rm hom}$, of the Raman transition. As a result, the mode capacity of quantum storage may be C times larger than that achievable without refractive index control, which is determined by the ratio $\gamma'/\gamma'_{\rm hom}$.

III. MULTICHANNEL QUANTUM STORAGE

In the previous Section, we were interested in the situation when the bandwidth of the photon is larger than the inhomogeneous linewidth of the Raman transition. Now we consider the opposite case, when the storage and retrieval are implemented, for example, by manipulating the inhomogeneous broadening, which should be larger than the photon bandwidth, and refractive index control is used for realizing multichannel regime of quantum storage. The idea is that different wave vectors of the spin waves correspond to different channels of storage and retrieval just as in the case of angular multiplexing. For example, consider the Raman echo memory scheme with controlled reversible inhomogeneous broadening that is switched in time periodically [11]. The multichannel regime can be achieved by assigning different wave vectors k_c to the control field (in our case — by refractive index control) during different dephasing/rephasing circles. In a similar way, we can consider memory schemes based on resonant interaction. Let the storage and retrieval be implemented using atomic frequency comb [30], which dephases and then rephases after a time T, and π-pulses transferring the optical coherence to/from the spin coherence are used [15]. Then we can make different n for different π-pulses thereby creating orthogonal spin waves on different storage/retrieval cycles or we can change refractive index for the weak field to be stored so that $\delta = T$, which leads to the same result. In any case, such multichannel regimes enable one to process new quantum states while preserving those stored before and provides access to all states kept in store in any order. It is also important that the phase modulation of the control field, which is required to store and recall pulses without resort to inhomogeneous broadening, becomes needless in the case of such channel division. Regarding the impedance matching condition, it takes the form $C = 1$ since $\delta/2$ is replaced by inhomogeneous life-time $T_2^* = 1/\gamma'$.

Apart from the finite storage time due to irreversible relaxation, the maximum number of channels is also limited by the signal-to-noise ratio. The latter can be estimated in the following simple way. Consider two channels corresponding to wave vectors with the difference $\Delta k_m = 2\pi m/L$, where m is an integer. During retrieval from one of them the probability of retrieval from another is determined by phase mismatching and proportional to $P_m = [\mathrm{sinc}(\Delta k_m L/2)]^2$, which may be equal to zero only for monochromatic field. Now it is necessary to take into account the bandwidth $\delta k = \delta\omega n/c$ of the retrieved signals (here

n is the average value of refractive index during storage or retrieval). If $\delta k \ll k$, then

$$P_m \approx \frac{1}{\delta k} \int_{-\delta k/2}^{\delta k/2} \left[\frac{Lx}{2\pi m}\right]^2 dx = \frac{1}{12} \left[\frac{\delta\omega}{\omega} \frac{L}{\lambda} \frac{n}{m}\right]^2 . \tag{25}$$

We see that the noise from another channel is quadratically proportional to the bandwidth of the photons $\delta\omega$ and length of the sample L. As an example, let $L/\lambda = 10^5$, $\omega/2\pi = 2{\cdot}10^{14}$ Hz, and $n = 2$. If $\delta\omega/2\pi \sim 50$ MHz, then $P_m \lesssim 10^{-4}$ for $m \geq 1$ so that 100 channels provide total signal-to-noise ratio of the order of 100. In order to maintain this ratio for broader signals, the value of m should be increased proportionally to the bandwidth, which means increasing refractive index difference between adjacent channels.

IV. REFRACTIVE INDEX CONTROL

Let us discuss possible ways of refractive index manipulation that are suited to quantum storage devices. First, if a doped nonlinear crystal is used as a storage medium, we can take advantage of the linear electro-optic effect. For example, the maximum value of the index change in $LiNbO_3$, which is limited by the breakdown electric field, is of the order of 10^{-3}. Although attaining this value by applying a moderate voltage is possible only for a waveguide configuration, the doped nonlinear materials, particularly $LiNbO_3$:Er^{3+}, hold promise in quantum storage applications, and therefore traditional electro-optic techniques of refractive index manipulation might be of helpful. Second, the resonant enhancement of the refractive index with vanishing absorption can be realized via quantum interference effects [37–42]. Although the maximum enhancement demonstrated experimentally in a gaseous medium is 10^{-4} [43, 44], there is good reason to expect much larger values, such as 10^{-2}, for solid state materials [45, 46]. It should be noted that manipulating refractive index either on the frequency of the quantum field to be stored or on the frequency of the strong control field should not only guarantee small losses but also no amplification. One way of avoiding the gain is to use excited state absorption [47]. Finally, we would like to note that some possibilities of refractive index control are provided even by a frequency shift of an absorption structure relative to the Λ-type one. At cryogenic temperatures optical transitions of impurity crystals and especially of rare-earth ion-doped ones have very narrow homogeneous lines so that in the neighborhood of an inhomogeneous profile there might be rather strong frequency dependence of refractive index and yet very small absorption. In this

respect, spectral hole-burning techniques can be very useful for preparing inhomogeneous profiles with sharp edges as well as for creating Λ-systems on a nonabsorbing background. A more detailed analysis of such an approach, which invokes specific information about resonant materials, will be presented elsewhere.

V. CONCLUSION

It is shown that single-photon wave packets can be stored and recalled in a resonant three-level medium by means of refractive index control without recourse to inhomogeneous broadening and modulating the amplitude of control fields. Such a scheme for quantum storage can be combined with other techniques to increase the mode capacity of quantum memories or to develop multichannel storage devices. In particular, quantum storage via refractive index control may be considered as a kind of a more general three-dimensional approach, wherein an input pulse is mapped to a superposition of collective atomic states corresponding to orthogonal macroscopic polarization modes of a finite interaction volume. By continuous changing refractive index we project the input pulse shape on a subsystem of the modes which differ in length of their wave vector. In a similar manner, by changing direction of propagation of the control field we can project the input pulse on the modes with different direction of the wave vector. It is the combination of both methods that allows one to exploit all three dimensions of the wave vector space for quantum storage. Although implementation of refractive index control is still a challenging experimental problem, we hope that significant progress can be made in this field by taking advantage of solid-state materials, which are also promising for quantum storage applications.

VI. ACKNOWLEDGMENTS

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