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Generalized dark-state polaritons for photon memory in multilevel atomic media

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We generalize the concept of dark-state polariton in multilevel electromagnetically induced transparency systems. We show that the quantum states of light for the pulses can be mapped onto more than one collective atomic polarization states of the multilevel atomic system, which can act as a quantum state copier or divider. Such dark-state polaritons potentially have applications in quantum information processing.

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Optical memories using classical techniques of spin and photon echoes are quite powerful as far as the high capacity of data storage is concerned [1–6]. However, direct applications of these techniques for quantum memories are very much limited as the number of photons required is larger than number of atoms in the system. Quantum memories are essential in quantum information processing, which involves physically transporting or communicating quantum states between different nodes of a quantum network [7]. It has been well recognized [8] and now experimentally demonstrated [9] that photons can be used as ideal carriers of quantum information and ensembles of atoms can act as long lived storage and processing units. The basic requirement for a reliable quantum memory system is its capability of storing and releasing quantum states on demand at the level of individual qubits, which puts stringent conditions on coherent transfer of information among photons and atoms.

Among many methods for coherently controlling the photon-atom interactions, the techniques based on electromagnetically induced transparency (EIT) [10] have been shown to be very promising in quantum state manipulations [11] and enhanced nonlinearity applications [12]. EIT systems can be transparent for the probe light beam at certain frequencies due to quantum interference and provide a large variation in linear dispersion within the transparency window [13]. The sharp change in linear dispersion can lead to substantial group-velocity reduction of light [14] to preserve quantum states of the slowed down light pulses, thus allowing the atomic medium to act as a temporary storage or buffering device for the quantum states of light. Physically, the photon storage in an EIT system is due to the formation of so called dark-state polariton (DSP), which is a mixture of electromagnetic signal field and the collective polarization of the atomic system, controlled by the strong coupling beam. It is the mixing angle of the signal field and the collective atomic polarization, which determines the group velocity of the signal pulse propagating in the atomic medium. By adiabatic following of DSP it is possible to reduce the group velocity of light pulses and to convert the photon states into the metastable atomic-polarization state. The stored light field in the collective atomic polarization can be recovered by manipulating the mixing angle of the DSP components by controlling the coupling beam. The concept of DSP was first introduced in Raman adiabatic passage by Mazets and Matisov [15] and then formulated for the three-level EIT system by Fleischhauer and Lukin [16]. Recent experiments have demonstrated photon storage using the EIT system consisting of three-level atoms in Λ configuration due to dynamic group-velocity reduction via adiabatic following in DSP [9].

In this work we extend the concept of a simple DSP to a generalized dark-state polariton (GDSP) for a multilevel atomic system having more than one dark state. One such example is the inverted-Y configuration involving four atomic levels that can easily be realized in rubidium atoms [17]. As one will see in the following paragraphs that such a system can preserve quantum states of light in two different collective atomic polarization states of the same atomic medium and thus provide channelization or bifurcation of photon memory. The two different collective atomic polarization channels contributing to GDSP can be used to retrieve back the light pulses on demand at different times or at the same time. In a way this system acts as a copier or divider of the quantum memory.

A schematic diagram of the closed four-level atomic system in inverted-Y configuration is depicted in Fig. 1. Levels $|a\rangle$, $|b\rangle$, and $|c\rangle$ are in a three-level Λ-type configuration and levels $|b\rangle$, $|a\rangle$, and $|d\rangle$ form a three-level ladder-type configuration. The transition from $|b\rangle$ to $|a\rangle$ (with frequency ω_{ab}) interacts with a signal field of amplitude $\hat{\varepsilon}(z,t)$ [defined below in Eq. (4)] and frequency ν. A coupling field of frequency ν_{ac} (considered to be classical) drives the transition $|c\rangle$ to $|a\rangle$ (frequency ω_{ac}) with Rabi frequency $\Omega_c(z,t)$, while a pumping field of frequency ν_{ad}, which is also considered to be classical, drives the transition $|a\rangle$ to $|d\rangle$ (frequency ω_{ad}) with Rabi frequency $\Omega_d(z,t)$. If there is no coupling (pump-

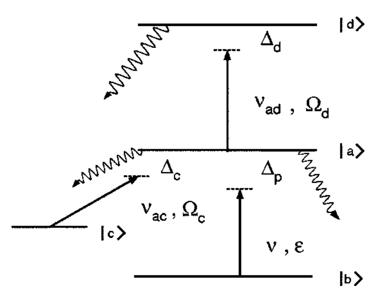

FIG. 1. Schematic diagram of a four-level atomic system in inverted-Y configuration. Here, ε, Ω_c, and Ω_d define signal, coupling, and pumping fields, respectively.

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ing) field, i.e., $\Omega_c(z,t)=0$ $[\Omega_d(z,t)=0]$, then this system reduces to a standard ladder-type (Λ-type) three-level EIT configuration driven by the signal and the pumping (coupling) fields. This four-level inverted-Y configuration acts as a double EIT system with independent controls of the two EIT windows by Ω_c and Ω_d, respectively.

We consider a collection of N four-level atoms in inverted-Y configuration interacting with three single-mode optical fields as described above. The system can be modeled by the interaction Hamiltonian in the dipole and the rotating-wave approximations under the condition that the carrier frequencies of the electromagnetic fields ν, ν_{ac}, and ν_{ad} coincide with the atomic resonances ω_{ab}, ω_{ac}, and ω_{ad}, respectively, and it is given by

$$\hat{V} = -\wp \sum_j [\hat{\sigma}^j_{ab}\hat{E}^{(+)}(z_j) + \text{H.a.}] - \hbar \sum_j [\hat{\sigma}^j_{ac}\Omega_c(z_j,t)e^{i(k^p_c z_j - \nu_{ac}t)} + \text{H.a.}] - \hbar \sum_j [\hat{\sigma}^j_{ad}\Omega_d(z_j,t)e^{i(k^p_d z_j - \nu_{ad}t)} + \text{H.a.}], \tag{1}$$

in which $\wp$ is the dipole-matrix element of the signal transition and the jth atom is located at the position z_j. The atomic ladder operators are defined as

$$\hat{\sigma}^j_{xy} = |x_j\rangle\langle y_j|. \tag{2}$$

The projections of the wave vectors ($\vec{k}_c$, $\vec{k}_d$), corresponding to the coupling and pumping fields on the propagation axis of the quantum (signal) field, are given by

$$k^p_c = \vec{k}_c \cdot \vec{e}_z = (\nu_{ac}/c)\cos(\phi_c), \quad k^p_d = \vec{k}_d \cdot \vec{e}_z = (\nu_{ad}/c)\cos(\phi_d), \tag{3}$$

where ϕ_c (ϕ_d) is the angle between wave vector $\vec{k}_c$ ($\vec{k}_d$) and $\vec{e}_z$, and $\vec{e}_z$ is the unit vector in the signal propagation direction. Next, we write down the signal field and the ladder operators in slowly varying variables

$$\hat{E}^{+}(z,t) = \sqrt{\frac{\hbar\nu}{2\epsilon_0 V_Q}}\hat{\varepsilon}(z,t)e^{i(\nu/c)(z-ct)},$$
$$\hat{\sigma}^j_{xy} = \tilde{\sigma}^j_{xy}e^{-i(\omega_{xy}/c)(z-ct)}, \tag{4}$$

where V_Q represents the quantization (interaction) volume. If the amplitudes of the slowly varying variables remain constant in the propagation direction in an interval Δz containing $N_z(\gg 1)$ atoms then we have

$$\tilde{\sigma}_{xy}(z,t) = \frac{1}{N_z}\sum_{z_j \epsilon N_z}\tilde{\sigma}^j_{xy}(t), \tag{5}$$

and the sum can be replaced by the integral in Eq. (1), i.e., $\Sigma^N_{j=1} \rightarrow (N/L)\int dz$. Therefore, the interaction Hamiltonian can be recast in the continuous form

$$\hat{V} = -\int \frac{dz}{L}[\hbar g N\tilde{\sigma}_{ab}(z,t)\hat{\varepsilon}(z,t) + \hbar\Omega_c(z,t)\tilde{\sigma}_{ac}(z,t)e^{i\Delta k_c z} + \hbar\Omega_d(z,t)\tilde{\sigma}_{ad}(z,t)e^{i\Delta k_d z} + \text{H.a.}], \tag{6}$$

in which we have used $g=\wp\sqrt{\nu/2\hbar\epsilon_0 V_Q}$ as the atom-signal field coupling constant and $\Delta k_i = k^p_i - k_i = (\omega_{ai}/c)[\cos(\phi_i)-1]$, $(i=c,d)$. The evolution of the atomic variables can be given in terms of Heisenberg-Langevin equations as

$$\dot{\tilde{\sigma}}_{dd} = -\gamma_d\tilde{\sigma}_{dd} - i(\Omega_d e^{i\Delta k_d z}\tilde{\sigma}_{ad} - \text{H.a.}) + \hat{F}_d,$$
$$\dot{\tilde{\sigma}}_{aa} = \gamma_d\tilde{\sigma}_{dd} - \gamma_a\tilde{\sigma}_{aa} - \gamma_c\tilde{\sigma}_{aa} + i(\Omega_d e^{i\Delta k_d z}\tilde{\sigma}_{ad} - \text{H.a.}) - i(\Omega^*_c e^{-i\Delta k_c z}\tilde{\sigma}_{ca} - \text{H.a.}) - ig(\hat{\varepsilon}^\dagger\tilde{\sigma}_{ba} - \text{H.a.}) + \hat{F}_a,$$
$$\dot{\tilde{\sigma}}_{bb} = \gamma_b\tilde{\sigma}_{aa} + ig(\hat{\varepsilon}^\dagger\tilde{\sigma}_{ba} - \text{H.a.}) + \hat{F}_b,$$
$$\dot{\tilde{\sigma}}_{cc} = \gamma_c\tilde{\sigma}_{aa} + i(\Omega^*_c e^{-i\Delta k_c z}\tilde{\sigma}_{ca} - \text{H.a.}) + \hat{F}_c,$$
$$\dot{\tilde{\sigma}}_{ba} = -\gamma_{ba}\tilde{\sigma}_{ba} + ig\hat{\varepsilon}(\tilde{\sigma}_{bb} - \tilde{\sigma}_{aa}) + i\Omega_c e^{i\Delta k_c z}\tilde{\sigma}_{bc} + i\Omega_d e^{i\Delta k_d z}\tilde{\sigma}_{bd} + \hat{F}_{ba},$$
$$\dot{\tilde{\sigma}}_{ca} = -\gamma_{ca}\tilde{\sigma}_{ca} + i\Omega_c e^{i\Delta k_c z}(\tilde{\sigma}_{cc} - \tilde{\sigma}_{aa}) + ig\hat{\varepsilon}\tilde{\sigma}_{cb} + i\Omega_d e^{i\Delta k_d z}\tilde{\sigma}_{cd} + \hat{F}_{ca},$$
$$\dot{\tilde{\sigma}}_{bc} = i\Omega^*_c e^{-i\Delta k_c z}\tilde{\sigma}_{ba} - ig\hat{\varepsilon}\tilde{\sigma}_{ac},$$
$$\dot{\tilde{\sigma}}_{cd} = -i\Omega_c e^{i\Delta k_c z}\tilde{\sigma}_{ad} + i\Omega^*_d e^{-i\Delta k_d z}\tilde{\sigma}_{ca},$$
$$\dot{\tilde{\sigma}}_{bd} = -ig\hat{\varepsilon}\tilde{\sigma}_{ad} + i\Omega^*_d e^{-i\Delta k_d z}\tilde{\sigma}_{ba},$$
$$\dot{\tilde{\sigma}}_{ad} = -\gamma_{ad}\tilde{\sigma}_{ad} + i\Omega^*_d e^{-i\Delta k_d z}(\tilde{\sigma}_{aa} - \tilde{\sigma}_{dd}) - ig\hat{\varepsilon}^\dagger\tilde{\sigma}_{bd} - i\Omega^*_c e^{-i\Delta k_c z}\tilde{\sigma}_{cd} + \hat{F}_{ad}, \tag{7}$$

in which $\gamma_a = \gamma_b + \gamma_c$. γ_d, γ_b, γ_c are longitudinal decay rates of levels $|d\rangle$ to $|a\rangle$, $|a\rangle$ to $|b\rangle$, $|a\rangle$ to $|c\rangle$, respectively, and γ_{xy} are transverse decay rates related to longitudinal rates. The frequency detunings are kept zero. The terms $\hat{F}_x$ and $\hat{F}_{xy}$ correspond to δ-correlated noise operators (with correlation functions $\langle F_\alpha(t)F_\beta(t')\rangle = D_{\alpha\beta}\delta(t-t')$, where $D_{\alpha\beta}$ are time-independent functions).

In order to solve the atomic equations we assume that photon number density of the signal pulse is much smaller than the number density of atoms and the Rabi frequency associated with the signal field is also much weaker than Ω_c and Ω_d. With these approximations, one can consider $\hat{\varepsilon}$ as perturbation in the atomic equations and these equations can be treated perturbatively in $\hat{\varepsilon}$. In zeroth-order approximation, $\tilde{\sigma}_{bb} \cong 1$ and all other atomic elements are set to zero. In the first-order approximation, we have

$$\Omega_c e^{i\Delta k_c z}\tilde{\sigma}^{(1)}_{bc} + \Omega_d e^{i\Delta k_d z}\tilde{\sigma}^{(1)}_{bd} = -g\hat{\varepsilon}(z,t) - i\left(\frac{\partial}{\partial t} + \gamma_{ba}\right)\tilde{\sigma}^{(1)}_{ba} + \hat{F}_{ba},$$
$$\frac{\partial}{\partial t}(\tilde{\sigma}^{(1)}_{bc}) = i\Omega^*_c e^{-i\Delta k_c z}\tilde{\sigma}^{(1)}_{ba},$$
$$\frac{\partial}{\partial t}(\tilde{\sigma}^{(1)}_{bd}) = i\Omega^*_d e^{-i\Delta k_d z}\tilde{\sigma}^{(1)}_{ba}. \tag{8}$$

The propagation equation of the signal pulse in the atomic medium takes the form

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{\varepsilon}(z,t) = igN\tilde{\sigma}^{(1)}_{ba}(z,t), \tag{9}$$

where $\tilde{\sigma}^{(1)}_{ba}(z,t)$ on the right hand side of Eq. (9) can be expressed in terms of $\tilde{\sigma}^{(1)}_{bc}$ and $\tilde{\sigma}^{(1)}_{bd}$. Under the adiabatic condition, i.e., Ω_c and Ω_d change sufficiently slow (their time variations are small) or more explicitly $L_P \gg \sqrt{\gamma_b cL/g^2 N}$ and

$T \gg \gamma_b v_g / g^2 N c$, where L_p is the length of pulse in medium and L is the total length of medium, T is the characteristic time corresponding to turn on and turn off times of the coupling/pump field, and v_g is the initial group velocity of the signal pulse [defined in Eq. (12)], which brings dependence of the adiabatic condition on the Rabi frequencies' quadrature sum [18]. Normalizing the time to $\tilde{t}=t/T$, then expanding the right hand side of the first equation of (8) in powers of $1/T$ and retaining the lowest nonvanishing order, we obtain:

$$\Omega_c e^{i\Delta k_c z}\tilde{\sigma}_{bc}^{(1)} + \Omega_d e^{i\Delta k_d z}\tilde{\sigma}_{bd}^{(1)} = -g\hat{\varepsilon}(z,t),$$
$$\tilde{\sigma}_{bc}^{(1)} = -\frac{\Omega_c^*}{|\Omega_c|^2+|\Omega_d|^2} g e^{-i\Delta k_c z}\hat{\varepsilon}(z,t),$$
$$\tilde{\sigma}_{bd}^{(1)} = -\frac{\Omega_d^*}{|\Omega_c|^2+|\Omega_d|^2} g e^{-i\Delta k_d z}\hat{\varepsilon}(z,t). \tag{10}$$

The Langevin noise operators do not contribute and the pulse propagation Eq. (9) under the adiabatic and perturbative conditions [using Eqs. (8) and (10)] is given by

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{\varepsilon}(z,t) = -\frac{g^2 N}{\Omega_c^* + \Omega_d^*}\left[\frac{\partial}{\partial t}\left(\frac{\Omega_c^*}{|\Omega_c|^2+|\Omega_d|^2}\hat{\varepsilon}(z,t)\right) + \frac{\partial}{\partial t}\left(\frac{\Omega_d^*}{|\Omega_c|^2+|\Omega_d|^2}\hat{\varepsilon}(z,t)\right)\right]. \tag{11}$$

With the assumption that $\Omega_c(z,t)$ and $\Omega_d(z,t)$ remain constant in time, we can have an expression for the modified group velocity v_g of the signal field as

$$v_g = \frac{c}{1+n_g}, \quad \left[n_g = \frac{g^2 N}{|\Omega_c|^2+|\Omega_d|^2}\right]. \tag{12}$$

The slowing down of the signal pulse does not change its temporal profile. Also, the photon storage time (τ_d) in such an EIT medium is set by the finite lifetimes of the two dark states or the dephasing times of the two transitions. The four-level inverted-Y atomic double EIT system can be characterized by the susceptibility of the medium [19]

$$\chi \sim \frac{g^2 N/kc}{(\gamma_a - i\Delta_p) - \frac{|\Omega_c|^2}{i(\Delta_p - \Delta_c)} + \frac{|\Omega_d|^2}{\gamma_d - i(\Delta_p + \Delta_d)}}, \tag{13}$$

where Δ_p, Δ_c, and Δ_d are the detunings of probe, coupling, and pumping lasers from their respective atomic transitions. Under resonance conditions $\Delta_c = 0$, $\Delta_d = 0$, we have

$$\chi \approx \frac{n_g}{kc}\frac{(|\Omega_c|^2+|\Omega_d|^2)\Delta_p}{(|\Omega_c|^2+|\Omega_d|^2) - i\gamma_a\Delta_p - \Delta_p^2}, \tag{14}$$

in which $\Delta_p = \nu - \omega_{ab}$. The transparency function under the homogeneous field condition is

$$T(\Delta_p, z) = \exp(-kz\,\mathrm{Im}[\chi]) \sim \exp\left(-\frac{\Delta_p^2}{\Delta\omega_{tr}^2}\right),$$
$$\Delta\omega_{tr} = \left[\frac{c}{\gamma_a l}\frac{|\Omega_c|^2+|\Omega_d|^2}{n_g}\right]^{1/2} = \frac{|\Omega_c|^2+|\Omega_d|^2}{\gamma_a\sqrt{\beta}}, \tag{15}$$

where β is related to absorption without EIT. From Eq. (15) above, it is clear that the transparency width decreases as n_g (which is controlled by Ω_c and Ω_d) increases. For the EIT medium, $\tau_d = n_g l/c$, then it is easy to write

$$\Delta\omega_{tr} = \sqrt{\beta}/\tau_d. \tag{16}$$

Clearly, a large delay due to large n_g will cause the transparency window to be narrower than the spectral width of the pulse, violating the adiabatic condition, and part of the signal pulse will be absorbed [16]. This limitation can be overcome by making use of the time-dependent coupling and pumping fields. For spatially homogeneous and time-dependent [$\Omega_c = \Omega_c(t)$ and $\Omega_d = \Omega_d(t)$] fields we can solve the propagation equation in a quasi particle picture using a GDSP in such a system following the procedure used in Ref. [16].

We define GDSP (assuming Ω_c and Ω_d to be real and dropping superscripts on atomic polarization components) as

$$\hat{\Psi}(z,t) = \cos\theta(t)\hat{\varepsilon}(z,t) - \sin\theta(t)\sqrt{N}\left(\frac{\Omega_c}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bc}(z,t)e^{i\Delta k_c z} + \frac{\Omega_d}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bd}(z,t)e^{i\Delta k_d z}\right). \tag{17}$$

Essentially, the quantum (signal) field $\hat{\varepsilon}(z,t)$ and the atomic spin coherence operators $\tilde{\sigma}_{bc}(z,t)$ and $\tilde{\sigma}_{db}(z,t)$ form a joint space of evolution with a mixing angle $\theta(t)$ defined as

$$\cos\theta(t) = \frac{\sqrt{\Omega_c^2+\Omega_d^2}}{\sqrt{\Omega_c^2+\Omega_d^2+g^2N}}, \quad \sin\theta(t) = \frac{g\sqrt{N}}{\sqrt{\Omega_c^2+\Omega_d^2+g^2N}}, \tag{18}$$

and

$$\tan^2\theta(t) = \frac{g^2N}{\Omega_c^2+\Omega_d^2} = n_g. \tag{19}$$

We call $\hat{\Psi}(z,t)$ GDSP and another quantum field $\hat{\Phi}(z,t)$ a generalized bright-state polariton (GBSP) can be defined as

$$\hat{\Phi}(z,t) = \sin\theta(t)\hat{\varepsilon}(z,t) + \cos\theta(t)\sqrt{N}\left(\frac{\Omega_c}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bc}(z,t)e^{i\Delta k_c z} + \frac{\Omega_d}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bd}(z,t)e^{i\Delta k_d z}\right). \tag{20}$$

The components of $\hat{\Psi}(z,t)$ and $\hat{\Phi}(z,t)$ are the quantum (signal) electromagnetic field $\hat{\varepsilon}(z,t)$ and the collective atomic spin operators $\tilde{\sigma}_{bc}(z,t)$ and $\tilde{\sigma}_{bd}(z,t)$, which can be controlled by adjusting $\theta(t)$ through coupling and pumping field strengths. Both of these generalized polaritons can be decomposed using the plane waves, i.e., $\hat{\Psi}(z,t) = \Sigma_k \hat{\Psi}_k(t)e^{ikz}$ and $\hat{\Phi}(z,t) = \Sigma_k \hat{\Phi}_k(t)e^{ikz}$. In the limit of photon number density to be much smaller than the atomic density, i.e., $\tilde{\sigma}_{bb}^j \cong 1$, $\tilde{\sigma}_{cc}^j \cong 0$, $\tilde{\sigma}_{dd}^j \cong 0$, and $\tilde{\sigma}_{dc}^j \cong 0$, etc. It is straightforward to show that the mode operators obey the following commutation relations:

$$[\hat{\Psi}_k, \hat{\Psi}_{k'}^\dagger] \cong [\hat{\Phi}_k, \hat{\Phi}_{k'}^\dagger] \cong \delta_{kk'}, \quad [\hat{\Psi}_k, \hat{\Phi}_{k'}^\dagger] \cong 0. \tag{21}$$

Therefore, the operators $\hat{\Psi}_k$ and $\hat{\Phi}_k$ behave like bosonic operators. One can transform propagation Eq. (11) in terms of new variables $\hat{\Psi}$ and $\hat{\Phi}$ easily in a low intensity approximation (where we assume $\partial\theta/\partial z=0$). In the adiabatic limit it can be shown using these equations that $\hat{\Phi}\approx 0$ and $\hat{\Psi}$ obey the simple equation of motion

$$\left[\frac{\partial}{\partial t}+c\cos^2\theta(t)\frac{\partial}{\partial z}\right]\hat{\Psi}(z,t)=0. \quad (22)$$

This equation gives a shape preserving, as well as quantum state preserving, solution propagating with a velocity $v=v_g(t)=c\cos^2\theta(t)$.

After obtaining our central results [Eqs. (17)–(22)] we can discuss the "stopping" (converting photons into collective spin polarizations of atoms) and "reacceleration" (converting collective spin polarization of atoms back into optical photons) of the photon wave packet, i.e., quantum memory by using GDSP as defined in Eq. (17), in this four-level atomic system by manipulating the mixing angle θ. When $\theta\cong 0$, i.e., $\cos\theta\cong 1$, meaning strong coupling and pumping fields ($\Omega_c^2+\Omega_d^2\gg g^2N$), then GDSP has a purely photonic character, i.e., $\hat{\Psi}\sim\hat{\varepsilon}$, which propagates with the velocity of light in vacuum. On the contrary, if $\theta\cong\pi/2$, i.e., $\cos(\theta)\cong 0$, then GDSP becomes spin-wave like and is composed of two spin components as

$$\hat{\Psi}\sim-\sqrt{N}\left(\frac{\Omega_c}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bc}e^{i\Delta k_c z}+\frac{\Omega_d}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bd}e^{i\Delta k_d z}\right).$$

This is achieved when $\Omega_c\to 0$ and $\Omega_d\to 0$, and the pulse propagation velocity goes down to zero. Mapping the traveling quantum field and the atomic spin polarization components can be stated as

$$\hat{\varepsilon}\leftrightarrow\frac{\Omega_c}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bc}(z')e^{i\Delta k_c z'}+\frac{\Omega_d}{\sqrt{\Omega_c^2+\Omega_d^2}}\tilde{\sigma}_{bd}(z')e^{i\Delta k_d z'}, \quad (23)$$

where $z'=z+\int_0^\infty d\tau c\cos^2\theta(\tau)$. Now from photon wave packets traveling with the speed of light we can convert information to stationary atomic excitation in two different channels just by adiabatically changing θ from 0 to $\pi/2$. The state of the polariton and its spatial shape do not change because of the linearity of Eq. (22) [16]. The stored light pulse in the two stationary spin wave channels can be simultaneously retrieved to signal electromagnetic field by applying Ω_c and Ω_d simultaneously and adiabatically, so that GDSP will be "reaccelerated" to the vacuum speed of light to give back the pure signal pulse in the forward direction. In such a system, the signal field is bifurcated and stored in two stationary spin waves with a weight factor governed by the two classical fields as given in Eq. (10) for the respective spin polarization components. Interestingly, if one puts the pumping (coupling) field $\Omega_d(\Omega_c)=0$, then the system reduces to a three-level Λ (ladder)-type system as evident from Eqs. (10) and (17) in which $\hat{\sigma}_{bc}$ ($\hat{\sigma}_{bd}$) provides the quantum storage channel. The main advantage of the four-level system is the ability to channelize the quantum memory into two different spin waves which can then be retrieved on demand either simultaneously or with a certain time delay (by adjusting the time delay between Ω_c and Ω_d pulses). One can also simultaneously monitor what is being stored as $\Omega_c\to 0$ adiabatically so the spin wave $\hat{\sigma}_{ac}$ will be activated for quantum memory purposes. However, by keeping $\Omega_d\neq 0$ (but small in magnitude) one can monitor the fraction of the signal pulse.

To summarize, we have formulated the generalized dark-state polariton in the four-level inverted-Y atomic system, which can be used for two-channel quantum memory. This system behaves like a copier, divider, or multiplexer for quantum states of photons, which are essential for quantum information processing and quantum networking. The above derivations can be easily modified to apply to more than four-level systems to form multiple memory channels.

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