Correlated quantum memory: Manipulating atomic entanglement via electromagnetically induced transparency

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We propose a feasible scheme of quantum state storage and manipulation via electromagnetically induced transparency in flexibly *united* multiensembles of three-level atoms. For different atomic array configurations, one can properly steer the signal and the control lights to generate different forms of atomic entanglement within the framework of linear optics. These results opens up new possibility for the future design of quantum memory devices by using, e.g., an atomic grid.

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The remarkable demonstration of ultraslow light speed in a Bose-Einstein condensate in 1999 [1] have stimulated rapid advances in both experimental and theoretical works on exploring the mechanism and its fascinating applications of electromagnetically induced transparency (EIT) [2,3]. In 2000, Fleischhauer and Lukin presented a dark-state polaritons (DSPs) theory, which shows an elegant mapping-and-readout quantum memory technique by exchanging the quantum state information between the two components of DSPs, the quantized light field and the collective atomic excitations [4,5]. The crucial condition of adiabatic passage for the dark states was then fully confirmed by Sun *et al.* by revealing the dynamical symmetry of a single EIT medium consisting of three-level atoms [6]. As an extension, the quantum memory process in an atomic ensemble composed of complex N-level ($N>3$) atoms was also studied [7,8], which shows more freedom for the quantum state control. Most recently, the storage of two entangled lights in two *independent* ensembles of three-level atoms was also proposed [9], motivated by building a quantum communication network in which the stored entanglement in each of and/or among the atomic nodes should be conveniently manipulated [10].

In this paper, we present a feasible technique of correlated quantum memory via EIT mechanism in many ensembles composed of Λ-type three-level atoms, which has build-up flexible ability to generate and manipulate entangled states of atomic ensembles. This differs from previous entanglement schemes (even the EIT-based ones) [5,10] since it inserts new freedoms of manipulations between the mapping and readout processes in itself. We resolve the general quantized DSPs model of m-atomic-ensemble system and find that, for two atomic ensembles, the input probe light can be stored in either the first or the second atomic ensemble or even both of them (in correlated manner) by steering the control fields. Particularly, by preparing the initial probe light in a coherent superposition state, the entangled atomic states of two or three ensembles can be created within the framework of linear optics. Further manipulations of the atomic entanglement are manifested under two configurations for three ensembles. This scheme may have an impact on future research in quantum information science or even designing an atomic grid as light storage medium by combining the EIT and other methods [11].

GENERAL MODEL

The system we are considering is an atomic array composed of m ensembles of Λ type three-level-atoms [Fig. 1(a)]. Atoms of the σth ($\sigma=1,2,\ldots m$) atomic ensemble interact with the input single-mode quantized field with coupling constants g_σ, and one classical control filled with time-dependent Rabi-frequencies $\Omega_l(t)$. Generalization to multimode probe pulse case is straightforward. Considering all transitions at resonance, the interaction Hamiltonian of the system is given by

$$\hat{V}=\sum_{\sigma=1}^{m} g_\sigma\sqrt{N_\sigma}\hat{a}\hat{A}_\sigma^\dagger+\sum_{\sigma=1}^{m}\Omega_\sigma(t)\hat{T}_\sigma^{+}+\text{H.c.},\qquad(1)$$

where a is the photon annihilation operator and N_σ denotes the atom number of the σth ensemble, the collective excitation operators $\hat{A}_\sigma=1/\sqrt{N_\sigma}\sum_{j=1}^{N_\sigma}e^{-i\mathbf{k}_{ba}\cdot\mathbf{r}_\mathbf{j}^{(\sigma)}}\hat{\sigma}_{ba}^{j(\sigma)}$, $\hat{T}_\sigma^{-}=(\hat{T}_\sigma^{+})^\dagger=\sum_{j=1}^{N_\sigma}e^{-i\mathbf{k}_{ca}\cdot\mathbf{r}_\mathbf{j}^{(\sigma)}}\hat{\sigma}_{ca}^{j(\sigma)}$ and $\hat{C}_\sigma=1/\sqrt{N_\sigma}\sum_{j=1}^{N_\sigma}e^{-i\mathbf{k}_{bc}\cdot\mathbf{r}_\mathbf{j}^{(\sigma)}}\hat{\sigma}_{bc}^{j(\sigma)}$; as usual, the flip operator $\hat{\sigma}_{\mu\nu}^{i(\sigma)}=|\mu\rangle_{ii}\langle\nu|$ ($\mu,\nu=a,b,c$) is defined for the ith atom between states $|\mu\rangle$ and $|\nu\rangle$ of the σth ensemble, and $\mathbf{k}_{ba}$ ($\mathbf{k}_{ca}$) is the wave vectors of the quantized (classical) optical field with $\mathbf{k}_{bc}=\mathbf{k}_{ba}-\mathbf{k}_{ca}$. In large-$N_\sigma$ limits and low excitation conditions [6,7], it follows that $[\hat{A}_i,\hat{A}_j^\dagger]=\delta_{ij}$, $[\hat{C}_i',\hat{C}_j^\dagger]=\delta_{ij}$ and $[\hat{T}_i^{+},\hat{T}_j^{-}]=\delta_{ij}\hat{T}_j^{z}$, $[\hat{T}_i^{z},\hat{T}_j^{\pm}]=\pm\delta_{ij}\hat{T}_j^{\pm}$, where $\hat{T}_\sigma^{z}=\sum_{j=1}^{N_\sigma}(e^{-i\mathbf{k}_{aa}\cdot\mathbf{r}_\mathbf{j}^{(\sigma)}}\hat{\sigma}_{aa}^{j(\sigma)}-e^{-\mathbf{k}_{cc}\cdot\mathbf{r}_\mathbf{j}^{(\sigma)}}\hat{\sigma}_{cc}^{j(\sigma)})/2$. Meanwhile, $[\hat{T}_i^{+},\hat{A}_j]=-\delta_{ij}\hat{C}_j$, $[\hat{T}_i^{-},\hat{C}_j]=-\delta_{ij}\hat{A}_j$. Thereby the dynamical symmetry of this system is governed by a semidirect product

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Lie group $[\otimes_\sigma SU(2)]\bar{\otimes} h_{2m}$; h_{2m} denotes the Heisenberg algebra generated by $(\hat{A}_i, \hat{A}_i^\dagger, \hat{C}_i, \hat{C}_i^\dagger; i=1,2,\ldots,m)$.

The quantum memory process via the EIT technique requires that the total system should be kept in the dark-state subspace, thus the key point is to obtain the dark states of the multi-atomic-ensemble system which can be realized by the following general definition of DSPs operator

$$\hat{d} = (\cos\theta)\hat{a} - \sin\theta \prod_{j=1}^{m-1} (\cos\phi_j)\hat{C}_1 - \sin\theta \sum_{k=2}^{m} \sin\phi_{k-1} \prod_{j=k}^{m-1} (\cos\phi_j)\hat{C}_k, \qquad (2)$$

where the mixing angles θ and ϕ_j are defined through

$$\tan\theta = \frac{\sqrt{g_1^2 N_1 \Omega_2^2\Omega_3^2 \ldots \Omega_m^2 + g_2^2 N_2 \Omega_1^2\Omega_3^2 \ldots \Omega_m^2 + \ldots + g_m^2 N_m \Omega_1^2\Omega_2^2 \ldots \Omega_{m-1}^2}}{\Omega_1\Omega_2 \ldots \Omega_m} \qquad (3)$$

and

$$\tan\phi_{m-1} = \frac{g_m\sqrt{N_m}\Omega_1\Omega_2 \ldots \Omega_{m-1}}{\sqrt{g_1^2 N_1 \Omega_2^2\Omega_3^2 \ldots \Omega_m^2 + g_2^2 N_2 \Omega_1^2\Omega_3^2 \ldots \Omega_m^2 + \ldots + g_{m-1}^2 N_{m-1} \Omega_1^2\Omega_2^2 \ldots \Omega_{m-2}^2\Omega_m^2}}. \qquad (4)$$

Since $[\hat{d},\hat{d}^\dagger]=1$, $[\hat{V},\hat{d}]=0$, we can obtain the general atomic dark states as $|D_n\rangle = [n!]^{-1/2}(\hat{d}^\dagger)^n|0\rangle$, where $|0\rangle = |b^{(1)}, b^{(2)}, \ldots, b^{(m)}\rangle_{atom} \otimes |0\rangle_{photon}$ ($|0\rangle_{photon}$ denotes the vacuum state of the probe field and $|b^{(\sigma)}\rangle = |b_1, b_2, \ldots, b_{N_\sigma}\rangle$ is the collective ground state of the σth atomic ensemble).

In order to get more insights on quantum state control of this system, we first consider the special case of two ensembles which has the dark states (just as the standard EIT procedure as in Ref. 4)

$$|D_n\rangle = \sum_{k=0}^{n}\sum_{j=1}^{n-k} \sqrt{\frac{n!}{k!l!j!}} (\cos\theta)^k(-\sin\theta)^{n-k} \times (\sin\phi_1)^j(\cos\phi_1)^l |c_{(1)}^l, c_{(2)}^j\rangle_{spin} \otimes |k\rangle_{photon}, \qquad (5)$$

where $l=n-k-j$ and the state $|c_{(i)}^j\rangle$ $(i=1,2)$ represents the collective excitations of the ith ensemble with j excitations. We will show that not only the quantum memory process still can be revealed in this quite general system but also for the build-up ability to generate and manipulate the atomic entanglement in a highly extensible style. In fact, if we initially prepare the total state of the quantized light and atomic ensembles in $|\Psi_0\rangle = \Sigma_n P_n(\alpha_0)|b^{(1)}, b^{(2)}\rangle_{atom} \otimes |n\rangle_{photon}$, where $P_n(\alpha_0) = \alpha_0^n/\sqrt{n!}e^{-|\alpha_0|^2/2}$ is the probability of distribution function, then the mixing angle θ is rotated from 0 to $\pi/2$ while keeping the ratio Ω_1/Ω_2 by, e.g., an acoustic-optical modulator (AOM) [12] and switching them off adiabatically ($\tan\phi_1 = \lim_{\Omega_1,\Omega_2\to 0}[g_2\sqrt{N_2}\Omega_1/g_1\sqrt{N_1}\Omega_2]$), the system evolves into

$$|\Psi_e\rangle = |\alpha_1\rangle_{spin1} \otimes |\alpha_2\rangle_{spin2} \otimes |0\rangle_{photon} \qquad (6)$$

with $\alpha_1 = \alpha_0\cos\phi$ and $\alpha_2 = \alpha_0\sin\phi$. Clearly, the input optical state can be converted into the atomic coherences via manipulating two control fields. Particularly, (i) if $\phi=0$, the injected light is fully stored in the first ensemble ($\alpha_2=0$); (ii) if $\phi=\pi/2$, the input pulse is now stored in the second ($\alpha_1=0$). This mechanism can be extended to any nonclassical or entangled state of the input light.

The important issue of entanglement generation in the macroscopic atomic ensembles hardly can be overestimated due to its practical applications in quantum information processing [11,12]. For the present scheme, if the injected quantized field is in a coherent super-position state, e.g., for the initial state $|\Psi_0\rangle^\pm = 1/\sqrt{\mathcal{N}_\pm(\alpha_0)}(|\alpha_0\rangle \pm |-\alpha_0\rangle)_{photon} \otimes |b^{(1)}, b^{(2)}\rangle_{atom}$ [here $\mathcal{N}_\pm(\alpha_0)$ is a normalized factor [2]], a two-ensemble entangled state would be created as ($|\Psi_0\rangle^\pm \to |\Psi_e\rangle^\pm$)

$$|\Psi_e\rangle^\pm = \frac{1}{\sqrt{\mathcal{N}_\pm(\alpha_0)}}|0\rangle_{photon} \otimes (|\alpha_1, \alpha_2\rangle \pm |-\alpha_1, -\alpha_2\rangle)_{spin}, \qquad (7)$$

and the entanglement of atomic coherences [13] $E^\pm(\alpha_1,\alpha_2) = -\mathrm{tr}(\rho_{\alpha_1}^\pm \ln \rho_{\alpha_1}^\pm)$, with the reduced density matrix $\rho_{\alpha_1}^\pm = \mathrm{tr}^{(\alpha_2,atom)}(|\Psi_e\rangle\langle\Psi_e|)^\pm$, can be controlled by the two control fields. In particular, if we start from an initial state $|\Psi_0\rangle^-$ and choose $\phi=\pi/4$, we then obtain an EPR-type entangle state: $(|+\rangle|-\rangle + |-\rangle|+\rangle)/\sqrt{2}$, where $|\pm\rangle = (|\alpha_0/\sqrt{2}\rangle \pm |-\alpha_0/\sqrt{2}\rangle)_{spin}/\sqrt{\mathcal{N}(\alpha_0/2)}$. This process may be viewed as a simple linear optical circuit which transforms a standard basis to an entangled one [14]. Following the method developed in Ref. 8 it is straightforward also here to confirm the condition of adiabatic evolution and therefore the robustness of the system [7,15]. In addition, as a noticeable analogy, it deserves further explorations by recalling our scheme of generating two entangled *lights* from an initial optical superposition state via a *single* four-state atomic medium [7].

These remarkable properties can be readily extended to the general m-atomic-ensemble case. For an interesting example, we consider the case of $m=3$ in which the dynamical

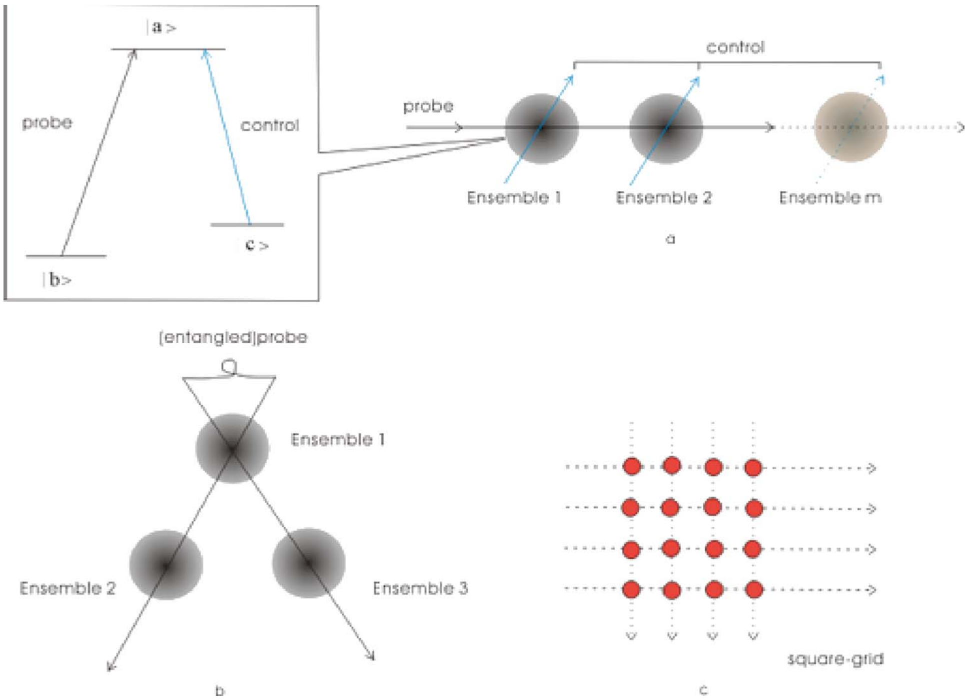

FIG. 1. (Color online) EIT process for three or many ensembles of Λ type atoms located in the (a) straight-line configuration; (b) cross-line configuration; (c) expanded square-grid configuration (the arrow-lines represent the probe beams).

symmetry is governed by the Lie group $\mathrm{so}(4)\otimes\mathrm{su}(2)\bar{\otimes}h_6$ [see also Fig. 1(a)]. Now, if the probe light is in a coherent superposition state, e.g., $|\Psi_0\rangle^{\pm}=1/\sqrt{\mathcal{N}_{0\pm}}(|\alpha_0\rangle\pm|\beta_0\rangle)_{photon}\otimes|b^{(1)},b^{(2)},b^{(3)}\rangle_{atom}$ with a normalized factor $\mathcal{N}_{0\pm}=2\pm 2e^{-|\alpha_0-\beta_0|^2/2}$, the achieved entangled state between three ensembles reads ($|\Psi_0\rangle^{\pm}\rightarrow|\Psi_e\rangle^{\pm}$)

$$|\Psi_e\rangle^{\pm}=\frac{1}{\sqrt{\mathcal{N}_{0\pm}}}|0\rangle_{photon}\otimes(|\alpha_1,\alpha_2,\alpha_3\rangle\pm|\beta_1,\beta_2,\beta_3\rangle)_{spin}, \tag{8}$$

where $\alpha_1=(\cos\phi_1\cos\phi_2)\alpha_0$, $\alpha_2=(\sin\phi_1\cos\phi_2)\alpha_0$, $\alpha_3=(\sin\phi_2)\alpha_0$, and $\beta_1=(\cos\phi_1\cos\phi_2)\beta_0$, $\beta_2=(\sin\phi_1\cos\phi_2)\beta_0$, and $\beta_3=(\sin\phi_2)\beta_0$. If we choose $\phi=\pi/4$ and $\varphi=\tan^{-1}(\sqrt{2}/2)$, we can get a "GHZ-like" entangled state: $(|\alpha,\alpha,\alpha\rangle\pm|\beta,\beta,\beta\rangle)_{spin}/\sqrt{\mathcal{N}_{0\pm}}$, where $\alpha=\alpha_0/\sqrt{3}$, $\beta=\beta_0/\sqrt{3}$. Notably, with the orthogonal basis $|\pm\rangle\propto(|\alpha\rangle\pm|\beta\rangle)_{spin}$ (for a normalized factor), this state also can be put into: $\Phi_{123}(\pm)=\xi_{\pm}|\pm\rangle|\pm\rangle|\pm\rangle+\zeta_{\pm}|W_{\pm}\rangle$, where $\xi_{\pm}$ and $\zeta_{\pm}$ are the normalized factors, and $|W_{\pm}\rangle=|\pm\rangle|\mp\rangle|\mp\rangle+|\mp\rangle|\pm\rangle|\mp\rangle+|\mp\rangle|\mp\rangle|\pm\rangle$ is the W-state [14]. It is fascinating to see that the two-ensemble state is *still* entangled after reducing the third one, and the general feature of coherent entanglement oscillations of Ramsey fringes for a highly entangled array also should be observed [12].

EXPANDED ILLUSTRATION

Finally we prove that the characteristics in the above scheme can be flexibly extended to other different atomic configurations. As a concrete example, we consider the three ensembles of Λ type atoms with such an atomic array [see Fig. 1(b)]: one probe light beam is injected to interact with the atoms of the first and second ensembles with the coupling constants g_1 and g_1', while another beam interacts with the first and third ones with the coupling g_2 and g_2', and the three classical control fields couple the transitions $|c\rangle\rightarrow|a\rangle$ with time-dependent Rabi-frequencies $\Omega_{\sigma}(t)$ $(\sigma=1,2,3)$ [12]. For simplicity, we still consider the single-mode probe lights; the generalizations to the multimode case is straightforward. The interaction Hamiltonian of the total system now is

$$\hat{V}=g_1\sqrt{N_1}\hat{a}_1\hat{A}_1^{\dagger}+g_1'\sqrt{N_2}\hat{a}_1\hat{A}_2^{\dagger}+g_2\sqrt{N_1}\hat{a}_2\hat{A}_1^{\dagger}+g_2'\sqrt{N_3}\hat{a}_2\hat{A}_3^{\dagger}+\Omega_1(t)\hat{T}_1^{+}+\Omega_2(t)\hat{T}_2^{+}+\Omega_3(t)\hat{T}_3^{+}+\mathrm{H.c.} \tag{9}$$

and the operators here are in the same definitions as the above. The dynamical symmetry of this system is governed by the Lie algebra $\mathrm{so}(4)\otimes\mathrm{su}(2)\bar{\otimes}h_6$ and now the DSPs operator is given by

$$\hat{d}=\hat{d}_1+\hat{d}_2, \tag{10}$$

where $\hat{d}_1=(\cos\theta_1)\hat{a}_1-(\sin\theta_1\cos\phi_1)\hat{C}_1-(\sin\theta_1\sin\phi_1)\hat{C}_2$, $\hat{d}_2=(\cos\theta_2)\hat{a}_2-(\sin\theta_2\cos\phi_2)\hat{C}_1-(\sin\theta_2\sin\phi_2)\hat{C}_3$ and the mixing angles $\theta_{1,2}$ and $\phi_{1,2}$ are defined just as Eq. (3) and (4), e.g., $\tan\phi_1=g_1'\sqrt{N_2}\Omega_1/(g_1\sqrt{N_1}\Omega_2)$ and $\tan\phi_2=g_2'\sqrt{N_3}\Omega_1/(g_2\sqrt{N_1}\Omega_2)$. It can be readily verified that $[\hat{d},\hat{d}^{\dagger}]=1$ and $[\hat{d},\hat{V}]=0$. Now we consider the initial state of the system with two entangled probe lights: $|\Psi_0\rangle^{\pm}=1/\sqrt{\mathcal{N}_{0\pm}}(|\alpha_0\rangle_1|\alpha_0\rangle_2\pm|\beta_0\rangle_1|\beta_0\rangle_2)_{photon}\otimes|b^{(1)},b^{(2)},b^{(3)}\rangle_{atom}$ ($\mathcal{N}_{0\pm}$ is a normalized factor), then it turns out that the atomic entangled state of the three ensembles as Eq. (8) can be obtained again by properly steering the control fields, but with *different* parameters ($\eta=\alpha,\beta$): $\eta_1=(\cos\phi_1+\cos\phi_2)\alpha_0$, $\eta_2=(\sin\phi_1)\alpha_0$ and $\eta_3=(\sin\phi_2)\alpha_0$. In particular, if $\phi_1=\phi_2=\pi/2$ or the Rabi frequency Ω_1 is kept much smaller than both Ω_2 and Ω_3 during the process that the three control fields are turned off, we can arrive at the reduced two-ensemble entanglement of the three atomic ensembles

$$\Phi_{atom}(\pm) = \frac{1}{\sqrt{\mathcal{N}_{0\pm}}}|b^{(1)}\rangle \otimes (|\alpha_0\rangle_1|\alpha_0\rangle_2 \pm |\beta_0\rangle_1|\beta_0\rangle_2)_{spin}. \tag{11}$$

The entanglement of two probe lights are fully transferred into that of the second and third ensembles. Clearly, this way to create the atomic entanglement differs itself from previous schemes [5,10,11].

In conclusion, we first proposed a nontrivial model for light storage in multi-atomic-ensemble and then easily obtained the quantum entanglement between many atomic ensembles via EIT techniques, which will be widely useful in quantum information science. The essential feature of the present scheme is that both the storage style and the quantum entanglement form can be controlled by just using the standard EIT procedure in a "single" memory node of a possible quantum network. We should remark here that, just as the quantum memory for photons in an atomic ensemble is much different from that in single atoms, the model for a multi-atomic-ensemble system is also quite different from that of single-atomic-ensemble EIT system. Obviously, the atomic-ensemble entanglement can be obtained only in the former system but not for the later system (for example, we illustrated that by preparing an initial probe light in coherent superposition state through, e.g., a beam-splitter-based catalysis technique [16], an atomic entanglement of two ensembles can be created by linear optics). Furthermore, similar differences on the Lie algebra structure exist between the multiensemble system and the single-ensemble system just like that between an atomic ensemble system and the single atom [for example, the so(4) algebra in two ensemble is different from the su(2) in single system]. The manipulations of atomic entangled state, for the case of three ensembles in particular, are explored under two different configurations and it can be flexibly extended to more complex structures like a tree or square array [see e.g., Fig. 1(c)]. This may have a strong impact on future research in quantum entanglement manipulations between many atomic ensembles or many propagating atom lasers and quantum information science by using even an atomic *grid* based on the EIT-type experiment.

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