Tightly localized stationary pulses in a multilevel atomic system

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We show that the pulse matching phenomenon can be obtained in the general multilevel system with electromagnetically induced transparency. For this we find a different way to create tightly localized stationary pulses by using counterpropagating pump fields. The present process is a spatial compression of excitation so that it allows us to shape and further intensify the localized stationary pulses, without using standing waves of pump fields or spatially modulated pump fields.

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Recently, an important progress in electromagnetically induced transparency (EIT) [1–7] is that experimental realization for the coherent control of stationary pulses was achieved by using standing waves of pump fields in the three-level system [8,9]. The creation of stationary pulses can enhance the nonlinear couplings between photons or collective excitations corresponding to stored photons, both of which are useful for deterministic logic operations. The key point in the creation of stationary pulses is expressed by the pulse matching phenomenon between the forward (FW) and backward (BW) propagating probe fields [8–10]. For a three-level system, the technique to generate tightly localized stationary pulses involves the use of standing waves of pump fields with a frequency comb or spatially modulated pump fields [9]. However, such tight localization is created by a filtering process rather than a compression of excitation. Thus the three-level technique cannot be applied directly to applications in quantum nonlinear optics.

On the other hand, coherent manipulation of probe lights has been studied in the four-level double Λ-type system [11,12] and also in the general multilevel atomic system that interacts with many probe and pump fields [13,14]. It has also been shown in Ref. [13] that one can convert different probe lights into each other by manipulating the external pump fields based on such general EIT method, indicating a sort of pulse matching phenomenon between all applied probe fields. This observation also motivates us to probe into a new technique of creating localized stationary pulses based on a multilevel atomic system. In this paper, we shall show that the tightly localized stationary pulses can be obtained through a spatial compression of excitation the general multilevel EIT system.

We consider the quasi-one-dimensional system shown in Fig. 1(a) for an ensemble of m-level atoms interacting with $m-2$ quantized probe fields which couple the transitions from the ground state $|b\rangle$ to excited state $|e_\sigma\rangle$ $(1\leq\sigma\leq m-2)$ with coupling constants g_σ, and $m-2$ classical pump fields which couple the transitions from the state $|c\rangle$ to excited ones $|e_\sigma\rangle$ with Rabi frequencies $\Omega_\sigma(z,t)$. All probe and pump fields are co-propagating in the $+z$ or $-z$ direction [Fig. 1(b)], and

$$E_\sigma(z,t)=\sqrt{\frac{\hbar\nu_\sigma}{2\epsilon_0 V}}\hat{\mathcal{E}}_\sigma(z,t)e^{i(k_{p\sigma}z-\nu_\sigma t)},$$

$$\Omega_\sigma(z,t)=\Omega_{\sigma 0}e^{i(k_{c\sigma}z-\omega_\sigma t)}, \tag{1}$$

where $\sigma=1,2,\ldots,m-2$, $\hat{\mathcal{E}}_\sigma$ and $\Omega_{\sigma 0}$ are slowly varying amplitudes, $k_{p\sigma}$ and $k_{c\sigma}$, respectively z-component wave vectors of probe and pump fields, can be positive or negative. For $k_{p\sigma}>0$ and $k_{c\sigma}>0$ ($k_{p\sigma}<0$ and $k_{c\sigma}<0$), it means the σth pair of probe and pump fields propagate in the $+z$ ($-z$) direction. We consider all transitions to be at resonance. Under the rotating-wave approximation, the interaction Hamiltonian can be written as

$$\hat{\mathcal{V}}=-\int\frac{dz}{L}\Bigg(\hbar N\sum_{\sigma=1}^{m-2}g_\sigma\tilde{\sigma}_{e_\sigma b}(z,t)\hat{\mathcal{E}}_\sigma(z,t)+\hbar N\sum_{\sigma=1}^{m-2}\Omega_{\sigma 0}(t)\tilde{\sigma}_{e_\sigma c}(z,t)+\text{H.c.}\Bigg), \tag{2}$$

where N is the total atom number, L is the length of the

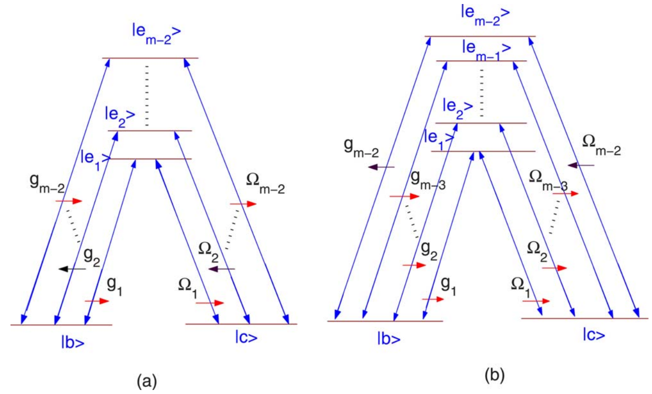

FIG. 1. (Color online) (a) General m-level atomic system coupled to $m-2$ quantized probe and classical pump fields which propagate in $+z$ or $-z$ directions. (b) No. 1 to no. $m-3$ pump and probe pulses propagate in the $+z$ direction, while no. $m-2$ pump and probe pulse propagates in the $-z$ direction.

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medium in the z direction, and the continuous atomic variables $\tilde{\sigma}_{\mu\nu}(z,t)=\frac{1}{N_z}\Sigma_{z_j\in N_z}\hat{\sigma}^j_{\mu\nu}(t)$ are defined by a collection of $N_z\gg 1$ atoms in a very small length interval Δz [3]. $\hat{\sigma}^j_{e_\sigma b}=|e^j_\sigma\rangle\langle b^j|e^{-i(k_{p\sigma}z-\omega_{e_\sigma b}t)}$ and $\hat{\sigma}^j_{e_\sigma c}=|e^j_\sigma\rangle\langle c^j|e^{-i(k_{c\sigma}z-\omega_{e_\sigma c}t)}$ are the slowly varying parts of the jth atomic flip operator. We note that an essential difference between our model and the three-level case is that for the case of multifrequency optical pulses, here the one- and two-photon detunings can be avoided for all optical transitions, and no standing waves of pump fields or spatially modulated pump fields are used.

The evolution of the slowly varying amplitudes $\hat{\mathcal{E}}_\sigma(z,t)$ can be described by the propagation equations

$$\left(\frac{\partial}{\partial t}+\frac{\nu_\sigma}{k_{p\sigma}}\frac{\partial}{\partial z}\right)\hat{\mathcal{E}}_\sigma(z,t)=ig_\sigma N\tilde{\sigma}_{be_\sigma}(z,t), \tag{3}$$

where we note $\nu_\sigma/k_{p\sigma}=\pm c$ for the $\pm z$ directional propagation field. Under the condition of low excitation, i.e., $\tilde{\sigma}_{bb}\approx 1$, the atomic evolution governed by the Heisenberg-Langevin equations can be obtained by

$$\dot{\tilde{\sigma}}_{be_\sigma}=-\gamma_{be_\sigma}\tilde{\sigma}_{be_\sigma}+ig_\sigma\hat{\mathcal{E}}_\sigma+i\Omega_{\sigma 0}\tilde{\sigma}_{bc}+F_{be_\sigma}, \tag{4}$$

$$\dot{\tilde{\sigma}}_{bc}=i\sum_{\sigma=1}^{m-2}\Omega_{\sigma 0}\tilde{\sigma}_{be_\sigma}-i\sum_{\sigma=1}^{m-2}g_\sigma\hat{\mathcal{E}}_\sigma\tilde{\sigma}_{e_\sigma c}+F_{bc}, \tag{5}$$

$$\dot{\tilde{\sigma}}_{ce_\sigma}=-\gamma_{ce_\sigma}\tilde{\sigma}_{ce_\sigma}+i\sum_{\sigma=1}^{m-2}g_\sigma\hat{\mathcal{E}}_\sigma\tilde{\sigma}_{cb}+F_{ce_\sigma}, \tag{6}$$

where $\gamma_{\mu\nu}$ are the transversal decay rates that will be assumed $\gamma_{be_\sigma}=\Gamma$ in the following derivation and $F_{\mu\nu}$ are δ-correlated Langevin noise operators. From Eq. (4) we find in the lowest order: $\tilde{\sigma}_{be_\sigma}=(ig_\sigma\hat{\mathcal{E}}_\sigma+i\Omega_{\sigma 0}\tilde{\sigma}_{bc}+F_{be_\sigma})/\Gamma$. Substituting this result into Eq. (5) yields $\tilde{\sigma}_{bc}=\Gamma^{-1}\Omega_0^2\tilde{\sigma}_{bc}-\Gamma^{-1}\Sigma_{\sigma=1}^{m-2}g_\sigma\Omega_{\sigma 0}\hat{\mathcal{E}}_\sigma-i\Sigma_{\sigma=1}^{m-2}g_\sigma\hat{\mathcal{E}}_\sigma\tilde{\sigma}_{e_\sigma c}$, where $\Omega_0=\sqrt{\Sigma_{\sigma=1}^{m-2}\Omega_{\sigma 0}^2}$. The Langevin noise terms are neglected in the present results, since under the adiabatic condition the Langevin noise terms have no effect on EIT quantum memory technique [3]. For our purpose we shall calculate $\tilde{\sigma}_{bc}$ to the first order, neglecting the small time derivatives of $\Omega_{\sigma 0}$, thus

$$\tilde{\sigma}_{bc}\approx-\frac{1}{\Omega_0^2}\sum_{\sigma=1}^{m-2}g_\sigma\Omega_{\sigma 0}\hat{\mathcal{E}}_\sigma-\frac{1}{\Omega_0^4}\sum_{jk\sigma}g_jg_kg_\sigma\Omega_{\sigma 0}\hat{\mathcal{E}}_j^\dagger\hat{\mathcal{E}}_k\hat{\mathcal{E}}_\sigma+\frac{\Gamma}{\Omega_0^4}\sum_{\sigma=1}^{m-2}g_\sigma\Omega_{\sigma 0}\partial_t\hat{\mathcal{E}}_\sigma. \tag{7}$$

The second term in the right-hand side of the above equation represents the nonlinear couplings between the probe pulses.

The dark-state polaritons (DSPs) in the general multilevel EIT system is first obtained in [13], where the single-mode probe pulses are considered. Accordingly, the dark- and bright-state polaritons (BSPs) in the present general multilevel system can be defined by

$$\hat{\Psi}(z,t)=\cos\theta\prod_{j=1}^{m-3}\cos\phi_j\hat{\mathcal{E}}_1+\cos\theta\sum_{l=2}^{m-2}\sin\phi_{l-1}\prod_{j=l}^{m-3}\cos\phi_j\hat{\mathcal{E}}_l-\sin\theta(t)\sqrt{N}\tilde{\sigma}_{bc}(z,t), \tag{8}$$

$$\hat{\Phi}(z,t)=\sin\theta\prod_{j=1}^{m-3}\cos\phi_j\hat{\mathcal{E}}_1+\sin\theta\sum_{l=2}^{m-2}\sin\phi_{l-1}\prod_{j=l}^{m-3}\cos\phi_j\hat{\mathcal{E}}_l+\cos\theta(t)\sqrt{N}\tilde{\sigma}_{bc}(z,t), \tag{9}$$

which are superpositions of the atomic coherence and the $m-2$ probe fields. The mixing angles θ and ϕ_j in the new quantum fields are defined through

$$\tan\theta=\frac{g_1g_2\cdots g_{m-2}\sqrt{N}}{\left[\sum_{j=1}^{m-2}\left(\Omega_{j0}^2\prod_{l=1,l\neq j}^{m-2}g_l^2\right)\right]^{1/2}}$$

and

$$\tan\phi_j=\frac{\prod_{l=1}^{j}g_l\Omega_{j+1,0}}{\left[\sum_{l=1}^{j}\left(\Omega_{l0}^2\prod_{s=1,s\neq l}^{j+1}g_s^2\right)\right]^{1/2}}.$$

Using the definitions above, one can transform the equations of motion for the probe fields and the atomic variables into the new field variables. With the low-excitation approximation and neglecting the nonlinear effects we find that the DSP field satisfies

$$\begin{aligned}(\partial_t+c\cos^2\theta\cos\alpha_{m-2}\partial_z)\hat{\Psi}&=-\dot{\theta}\hat{\Phi}+\sum_{j=1}^{m-2}\dot{\phi}_j\cos\theta\hat{s}_j-\frac{c}{2}\sin 2\theta\cos\alpha_{m-2}\partial_z\hat{\Phi}\\&\quad+c\cos\theta\sum_{j=1}^{m-2}\prod_{l=j}^{m-3}\cos\phi_l\sin 2\phi_{j-1}\\&\quad\times\left(\frac{1}{2c}\frac{\nu_j}{k_{pj}}+\frac{\cos\alpha_{j-1}}{2}\right)\partial_z\hat{s}_j,\end{aligned} \tag{10}$$

where we have defined

$$\cos\alpha_\sigma=c\frac{\sum_{j=1}^{\sigma}\frac{k_{pj}}{\nu_j}\Omega_{j0}^2\prod_{l=1,l\neq j}^{\sigma}g_l^2}{\sum_{j=1}^{\sigma}\Omega_{j0}^2\prod_{l=1,l\neq j}^{\sigma}g_l^2},\quad\sigma=1,2,\ldots,m-3$$

and $\hat{s}_j=\partial_{\phi_j}\hat{\Psi}/\cos\theta$. It then follows that $\hat{s}_1=\Pi_{j=2}^{m-3}\cos\phi_j\times(-\sin\phi_1\mathcal{E}_1+\cos\phi_1\mathcal{E}_2)$, $\hat{s}_2=\Pi_{j=3}^{m-3}\cos\phi_j[-\sin\phi_2(\cos\phi_1\mathcal{E}_1+\sin\phi_1\mathcal{E}_2)+\cos\phi_2\mathcal{E}_3]$, and generally

$$\hat{s}_k=\prod_{j=k+1}^{m-3}\cos\phi_j\hat{f}_k,\quad k=1,2,\ldots,m-3, \tag{11}$$

with

$$\hat{f}_k = \cos\phi_k \mathcal{E}_{k+1} - \sin\phi_k \sum_{m=2}^{k} \left(\prod_{l=m}^{k-1} \cos\phi_l \right) \sin\phi_{m-1} \mathcal{E}_m - \sin\phi_k \prod_{l=1}^{k-1} \cos\phi_l \mathcal{E}_1. \quad (12)$$

On the other hand, the equation of BSP field can be obtained as

$$\Phi = \frac{\Gamma}{\sqrt{N}} \left[\sum_{j=1}^{m-2} \left(\frac{\Omega_{j0}/\Omega_0}{g_j} \right)^2 \right]^{1/2} \frac{\cos\theta}{\Omega_0} \partial_t (\sin\theta\Psi - \cos\theta\Phi) + \cos\theta \left(g_1 \Omega_{01} \sin\phi_1 \hat{s}_1 - \sum_{l=2}^{m-3} g_l \Omega_{l0} \cos\phi_{l-1} \hat{s}_{l-1} \right). \quad (13)$$

By comparing Eqs. (10) and (13) with the corresponding DSP and BSP fields in the three-level system, one can see a key difference is the appearance of $\hat{s}_j(z,t)$ from the probe fields in our model. The adiabatic condition in the present case can be fulfilled only if $\hat{s}_j(z,t)=0$ for all j. However, the input probe pulses are generally independent of each other so that the fields $\hat{s}_j$ need not be zero. To study the dynamics of the DSP field, we should therefore investigate first the pulse matching between all the probe fields needed for adiabatic condition. Bearing these ideas in mind, we next explore the evolution of a set of normal fields by introducing

$$\hat{G}_{j,j+1} = -\sin\phi_{j,j+1} \hat{\mathcal{E}}_j(z,t) + \cos\phi_{j,j+1} \hat{\mathcal{E}}_{j+1}(z,t), \quad (14)$$

where $j=1,2,\ldots,m-3$ and $\tan\phi_{j,j+1}=g_j\Omega_{j+1,0}/g_{j+1}\Omega_{j0}$. From the Eq. (3) and together with the results of $\tilde{\sigma}_{be_\sigma}$ and $\tilde{\sigma}_{bc}$ one can verify that the field $\hat{G}_{j,j+1}$ satisfies the equation

$$\begin{aligned}(\partial_t - c\cos^2\beta \cos 2\phi_{j,j+1}\partial_z)\hat{G}_{j,j+1} &= -\frac{(g_j^2\Omega_{j+1}^2 + g_{j+1}^2\Omega_j^2)N}{\Gamma}\frac{\cos^2\beta}{\Omega_0^2}\hat{G}_{j,j+1} \\ &\quad - \frac{1}{2}g_j g_{j+1}\sqrt{N}\sin 2\beta\,\partial_t\hat{\mathcal{E}}_{j,j+1} \\ &\quad + c\cos^2\beta\sin 2\phi_{j,j+1}\partial_z\hat{\mathcal{E}}_{j,j+1} + F(\hat{\mathcal{E}}_\sigma, \sigma \neq j, j+1)\end{aligned} \quad (15)$$

with

$$\tan^2\beta = \frac{N\Omega_j^2\Omega_{j+1}^2}{g_j^2\Omega_{j+1}^2 + g_{j+1}^2\Omega_j^2}\frac{(g_j^2 - g_{j+1}^2)^2}{\Omega_0^4},$$

and $\hat{\mathcal{E}}_{j,j+1}=\cos\phi_{j,j+1}\hat{\mathcal{E}}_j(z,t)+\sin\phi_{j,j+1}\hat{\mathcal{E}}_{j+1}(z,t)$. $F(\hat{\mathcal{E}}_\sigma)$ includes no $\hat{\mathcal{E}}_j$ or $\hat{\mathcal{E}}_{j+1}$. The first term in the right-hand side of Eq. (15) reveals a very large absorption of $\hat{G}_{j,j+1}$, which results in a large decay in the field $\hat{G}_{j,j+1}$ and then the system satisfies the pulse matching condition [13,15,16] $\hat{\mathcal{E}}_{j+1} \to \tan\phi_{j,j+1}\hat{\mathcal{E}}_j$. For an explicit numerical estimation, we set some typical values [2,4]: $g_j \approx g_{j+1} \sim 10^5\ \mathrm{s}^{-1}$, $N \approx 10^8$, $\Gamma \approx 10^8\ \mathrm{s}^{-1}$, so that the lifetime of field $\hat{G}_{j,j+1}(z,t)$ is about $\Delta t < 10^{-8}$ s which is much smaller than the storage time [4]. Furthermore, by introducing the adiabatic parameter $\tau \equiv [\sum_j (1/g_j)^2]^{1/2}/\sqrt{N}T$ where T is the characteristic time scale, we can calculate the lowest order in Eq. (13) and obtain $\hat{\Phi} \approx 0$, $\hat{G}_{j,j+1} \approx 0$. On the other hand, under the condition of pulse matching, one can verify that $\hat{s}_j(z,t) \propto \hat{G}_{j,j+1}=0$. Thus Eq. (10) is reduced to the shape- and state-preserving case

$$(\partial_t + c\cos^2\theta\cos\alpha_{m-2}\partial_z)\hat{\Psi}(z,t) = 0. \quad (16)$$

The formula (16) is the main result of the present work. The group velocity of the DSP field is

$$V_g = \cos^2\theta \frac{\sum_{j=1}^{m-2} \frac{\nu_j}{k_{pj}} \Omega_{j0}^2 \prod_{l=1,l\neq j}^{m-2} g_l^2}{\sum_{j=1}^{m-2} \Omega_{j0}^2 \prod_{l=1,l\neq j}^{m-2} g_l^2}. \quad (17)$$

One should bear in mind that the wave vectors k_{pj} can be positive (in the $+z$ direction) or negative (in the $-z$ direction). So, by adjusting the Rabi frequencies for external pump fields properly under the adiabatic condition so that $\cos\alpha_{m-2}=0$, we can obtain a zero velocity for the DSP field. In particular, one may set no. 1 to no. $m-3$ pump and probe pulses in the $+z$ direction, while no. $m-2$ pump and probe pulse is in the $-z$ direction [Fig. 1(b)] and $\Omega_{m-2,0} = \sum_{j=1}^{m-3} \frac{g_{m-2}^2}{g_j^2} \Omega_{j0}^2$, in an experiment so that the group velocity $V_g=0$. In this way, we create the multifrequency stationary pulses with each component:

$$\hat{\mathcal{E}}_1 = \cos\theta \prod_{j=1}^{m-3} \cos\phi_j \Psi(z,t),$$

$$\hat{\mathcal{E}}_l = \cos\theta\sin\phi_{l-1} \prod_{j=l}^{m-3} \cos\phi_j \Psi(z,t),$$

$$l = 2, \ldots, m-2. \quad (18)$$

These components interfere to create a sharp spatial envelope. It is helpful to present a comparison of our results with those obtained for a three-level system [8,9]: (i) In the present system, all optical pulses can couple in resonance to the corresponding atomic transitions, thus all the applied probe fields with different frequencies contribute to the generation of stationary pulses. This means the present process is a spatial compression of excitation, which allows us to shape and intensify the localized stationary pulses. Our technique can therefore be expected to enhance further nonlinear couplings and be applied in the most straightforward manner, e.g., to applications in quantum nonlinear optics [17]. Numerical results in Fig. 2 show how tight localization of stationary pulses can be readily obtained when the multilevel system is applied. In contrast, for a three-level system, a frequency comb is used to create a localized pulse, filtering the off-resonant input probe pulses and retaining only the resonant one for creation [9]. Generally, the total number of

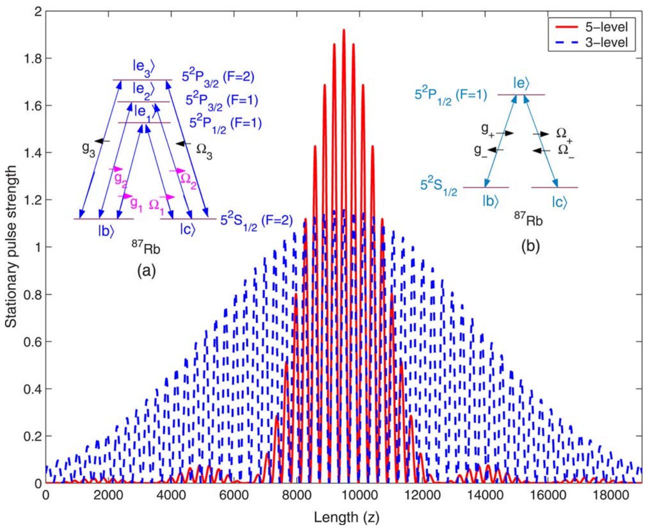

FIG. 2. (Color online) Localization of created stationary pulses for five-level (red solid line) cases, where three input probe lights are used and the parameters are set as $\omega_{e_3e_2}=\omega_{e_2e_1}\approx\nu_2/100$ (a). As a comparison, the blue dashed line shows the stationary pulses created in a three-level system (b) by employing one standing wave of pump fields. The probe lights are used with the envelope $\exp(-z^2)$.

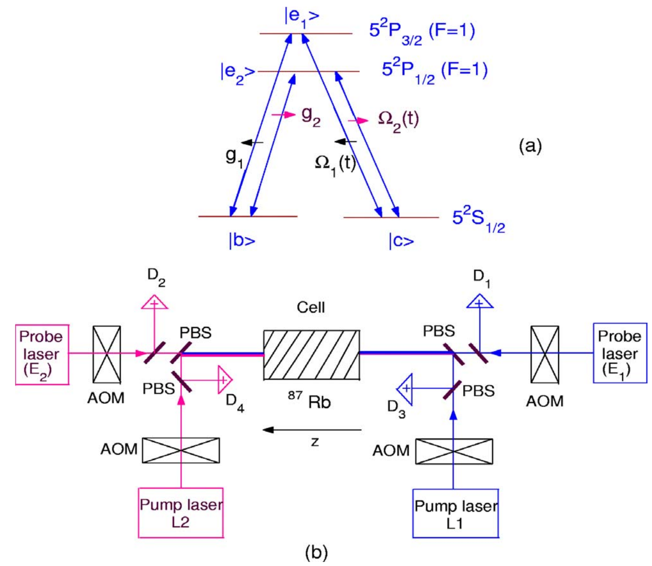

FIG. 3. (Color online) (a)(b) Schematic of experimental realization of stationary pulses with four-level double-Λ-type ^{87}Rb atoms coupled to two single-mode quantized and two classical pump fields that propagate in $+z$ and $-z$ directions, respectively.

probe photons created by a frequency comb in a three-level system is much less than in the current model. (ii) The present technique can be freely controlled. For example, based on our results, we see that the pulse matching in the present case is between all of the probe pulses with different frequencies, say, $\hat{\mathcal{E}}_\sigma=\Pi_{j=l}^{\sigma}\tan\phi_{j,j+1}\hat{\mathcal{E}}_l(l\geqslant 1,\ \sigma\leqslant m-2)$. Thus, in principle, one can use just one pump field to achieve the stationary pulse by adjusting its intensity to match those of the other pump fields. (iii) It requires no standing waves in the pump fields or spatially modulated pump fields to create localized stationary pulses.

Experimentally, the simplest multilevel system is an ensemble of four-level double-Λ-type ^{87}Rb atoms. The schematic diagram for experimental realization is shown in Fig. 3. All the atoms first are trapped in state $|b\rangle$ $(5^2S_{1/2})$ and only the $\pm z$ directional propagation pump fields (Ω_1 and Ω_2) are applied to couple the transitions from $|c\rangle$ $(5^2S_{1/2})$ to $|e_1\rangle$ $[5^2P_{1/2}(F=1)]$ and $|e_2\rangle$ $[5^2P_{3/2}(F=1)]$, respectively. We then input the probe pulses $(\hat{\mathcal{E}}_{1,2})$ and allow the system to achieve adiabatic condition. Finally, by adjusting Ω_1 or Ω_2 so that $g_1\Omega_{20}=g_2\Omega_{10}$, we can create the stationary pulses for probe fields $\hat{\mathcal{E}}_1(z,t)=\cos\theta\cos\phi\hat{\Psi},\hat{\mathcal{E}}_2(z,t)=\cos\theta\sin\phi\hat{\Psi}$, where $\hat{\Psi}$ is determined by Eq. (8) with $m=4$. It can be expected that when the level number m becomes larger, the more tightly localized stationary pulses can be created. According to the numerical results in Fig. 2, the effect becomes substantial when $m\geqslant 5$.

In conclusion, we have shown the tightly localized stationary pulses can be obtained in the general multilevel EIT system. We have examined the dynamics of DSPs in detail and found that, all the applied probe pulses with different frequencies contribute to the stationary pulses. The present process is therefore a spatial compression of excitation, which may be able to enhance further nonlinear optical couplings and will have interesting applications in quantum nonlinear optics [17]. In particular, our technique may open up a different way towards the spatial compression of many probe photons with small losses. According to the results in [13], if initially input is a nonclassical probe pulse, e.g., a quantum superposition of coherent states, we may also generate entangled stationary pulses within our model.

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