Role of degenerate Zeeman states in the storage and retrieval of light pulses

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We have studied the role of degenerate Zeeman states in the storage and retrieval of light pulses in the three-level Λ-type system of electromagnetically induced transparency. A probe pulse with the σ_+ polarization is stored in the medium and subsequently released with the σ_- polarization. However, the manipulation of the retrieved polarization causes an energy loss. Our study shows that the incompatibility between the stored ground-state coherence and the ratio of the probe and coupling Clebsch-Gordan coefficients is responsible for the energy loss. Such energy loss can be avoided as long as all population accumulates in a single Zeeman state. We provide the useful knowledge for the light-storage technique and demonstrate that the degenerate Zeeman states should be considered in the relevant applications.

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Light is an ideal carrier of information, as it possesses ultimate propagation speed, provides large communication bandwidth, and hardly interacts with the environment. We can foresee that optical photons will be employed to transport quantum information among the logic gates in quantum computation and among the devices in quantum communication. Based on the effect of electromagnetically induced transparency (EIT), the recent developments of slow light, the storage and retrieval of light pulses, and low-light-level nonlinear optics have attracted great attention [1–13]. These developments may lead to applications in the manipulation of quantum information. Utilizing the storage of light pulses, Chanellière et al. [14] and Eisaman et al. [15] have demonstrated the transfer of quantum states between atoms and single photons; Chen et al. have studied low-light-level cross-phase modulation and shown that a phase shift of the order of π with single photons is feasible [16]; Wang et al. have explored low-light-level all-optical switching and shown that a single photon switched by another is attainable [17].

A quantum computer or a quantum network may consist of different types of quantum devices made by various kinds of media. Manipulating the wavelength or frequency, polarization, and pulse width of stored photonic information can bridge different quantum devices, each of which only interacts with light of specific properties. Zibrov et al. demonstrated that the stored light pulse is retrieved into a different wavelength and a different propagation direction [18]. Wang et al. reported similar experimental results [19]. Chen et al. [20] experimentally studied the manipulation of the retrieved pulse width proposed by Patnaik et al. [21] and showed that this manipulation process does not introduce any phase shift. Some theoretical studies of the manipulation of the retrieval can be found in Refs. [22–26]. In Ref. [27], we reported the experimental demonstration that a stored light pulse is released with a different polarization, but discovered an energy loss caused by the manipulation process. In this work, we provide a detailed analysis to explain the energy loss and show that degenerate Zeeman states play an important role in the manipulation of retrieved light pulses and should be taken into consideration in the relevant applications.

We employed laser-cooled 87Rb atoms in the experiment of Ref. [27]. The coupling fields drove the |5S_{1/2},F=2⟩ → |5P_{3/2},F'=2⟩ transition resonantly and the carrier frequency of the probe pulse was tuned to the |5S_{1/2},F=1⟩ → |5P_{3/2},F'=2⟩ transition frequency. Both the writing coupling field and the input probe pulse were σ_+ polarized. The reading coupling field was either σ_+ or σ_- polarized in different measurements. Figure 1 shows the relevant 87Rb energy levels and the laser excitations in the experiment. Figure 2 shows the experimental demonstration which is redrawn from the data of Fig. 3 of Ref. [27]. In Figs. 2(a) and 2(b), the reading fields are σ_+ and σ_- polarized, respectively. The probe pulses with the σ_+ and σ_- polarizations were directed through different paths after leaving the atoms and separately detected by two photo detectors. Please see Fig. 1 of Ref. [27] for the experimental setup. Figure 2(b) demonstrates that a probe pulse with the σ_+ polarization was stored in the medium and subsequently released with the σ_- polarization. The retrieved probe signal in Fig. 2(a) and twice of that in Fig. 2(b) are plotted together in Fig. 2(c). It is clearly shown that an attenuation of about 50% is caused by the process of retrieving the stored light pulse into a different polarization.

In the consideration of the degenerate Zeeman states, there are five subsystems as shown in Fig. 1, three EIT subsystems, a two-level subsystem, and a subsystem containing only one ground state. The Maxwell-Schrödinger equation of the probe field and the optical Bloch equation of the density-matrix operator are given by

\[
\frac{1}{c} \frac{\partial}{\partial t} \hat{\Omega}_p + \frac{\partial}{\partial x} \hat{\Omega}_p = i \eta \sum_j a_{p,j} \rho_{31,j},
\]

(1)

\[
\frac{\partial}{\partial t} \rho_{31,j} = \frac{i}{2} a_{p,j} \hat{\Omega}_p (\rho_{11,j} - \rho_{31,j}) + \frac{i}{2} a_{c,j} \hat{\Omega}_c \rho_{21,j} - \frac{\Gamma}{2} \rho_{31,j},
\]

(2)

\[
\frac{\partial}{\partial t} \rho_{32,j} = \frac{i}{2} a_{c,j} \hat{\Omega}_c (\rho_{22,j} - \rho_{33,j}) + \frac{i}{2} a_{p,j} \hat{\Omega}_p \rho_{21,j} - \frac{\Gamma}{2} \rho_{32,j},
\]

(3)

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\[ \frac{\partial}{\partial t} \rho_{21,j} = -i 2 \left( a_{c,j} \hat{\Omega}_c (\rho_{31,j} - \rho_{31,j}^*) - a_{p,j} \hat{\Omega}_p (\rho_{32,j} - \rho_{32,j}^*) - \gamma \rho_{21,j} \right) \]  

\[ \frac{\partial}{\partial t} \rho_{11,j} = \frac{i}{2} a_{p,j} \hat{\Omega}_p (\rho_{31,j} - \rho_{31,j}^*) + \sum_k (a_{3,k-1,j})^2 \Gamma \rho_{33,k}, \]  

\[ \frac{\partial}{\partial t} \rho_{22,j} = \frac{i}{2} a_{c,j} \hat{\Omega}_c (\rho_{32,j} - \rho_{32,j}^*) + \sum_k (a_{3,k-2,j})^2 \Gamma \rho_{33,k}, \]  

In the above equations, \( j \) represents each of the five subsystems, \( \eta = \frac{3 \lambda^2 N \Gamma}{(4\pi)} \) where \( \lambda \) is the wavelength of the light and \( N \) is the number density of the atoms, \( a_{p,j} \) and \( a_{c,j} \) are the Rabi frequencies of the probe and coupling transitions, \( a_{p,j} \hat{\Omega}_p \) and \( a_{c,j} \hat{\Omega}_c \) are the Rabi frequencies of the probe and coupling transitions, \( \rho_{31,j} \) is the amplitude of the optical coherence of each probe transition and \( \rho_{32,j} \) is that of each coupling transition, \( \rho_{21,j} \) is the amplitude of each ground-state coherence, \( \rho_{11,j} \) and \( \rho_{22,j} \) are the populations of the ground states driven by the probe and the coupling fields, \( \rho_{33,j} \) is the population of each excited state, \( \Gamma \) is the total spontaneous decay rate of the excited states, \( (a_{3,k-1,j})^2 \Gamma \) is the rate of the spontaneous decay from the excited state in the subsystem \( k \) to the ground state \( |1\rangle \) in the subsystem \( j \), \( (a_{3,k-2,j})^2 \Gamma \) is similar but to the ground state \( |2\rangle \) and \( \gamma \) is the ground-state decoherence rate. The probe field is not treated as the perturbation in the deduction of these equations. Since we only consider the case that the coupling and probe fields are tuned to the resonance frequencies, \( \hat{\Omega}_c, \hat{\Omega}_p \), and \( \rho_{21,j} \) must be real numbers and \( \rho_{31,j} \) and \( \rho_{32,j} \) must be imaginary numbers. The two-level subsystem only has Eqs. (3), (6), and (7) with \( \hat{\Omega}_c = 0 \). The ground state without any driving field only has Eq. (6) with \( \hat{\Omega}_c = 0 \).

We numerically solve Eqs. (1)–(7) and assume that initially all population is equally distributed among the three ground states driven by the probe field. This assumption is practical for the degenerate states. In the calculation, \( \hat{\Omega}_c = 0.64 \Gamma \) for the writing and reading fields, \( \hat{\Omega}_p(t) = \hat{\Omega}_p \exp(-t^2/\tau^2) \) for the input probe pulse, \( \hat{\Omega}_p = 0.19 \Gamma, \tau = 79/\Gamma, \gamma = 0.001 \Gamma, \) and \( \eta = 14 \Gamma/L \) where \( L \) is the optical path length of the medium. These values are consistent with the experimental parameters.

In Fig. 3(a), the black solid line is the probe transmission under the condition that the writing and reading fields have same polarization; the red line is the probe transmission under the condition that the two fields have opposite polarizations. In the calculation of the latter condition, all \( a_{p,j}, a_{c,j}, a_{3,k-1,j} \) and \( a_{3,k-2,j} \) in Eqs. (1)–(7) are changed from the Clebsch-Gordan coefficients in the sys-
FIG. 3. (Color online) The main plot (a) and the inset (b) are the theoretical predictions calculated from Eqs. (1)–(7) and from Eqs. (8)–(10), respectively. Black and red (gray in printed journal) solid lines are the probe transmissions under the conditions that the writing and reading fields are both \( \sigma_+ \) polarized and that they have the opposite \( \sigma_- \) and \( \sigma_+ \) polarizations. The left parts and the gaps of the two solid lines overlap precisely. The dashed line is the combination of the writing and reading fields. The dotted line is the input probe pulse with the size reduced to half. The left and right vertical axes indicate \( \dot{\Omega}_p^2 \) (proportional to the probe intensity) and \( \dot{\Omega}_p^{-2} \) (proportional to the writing or reading field intensity). All the horizontal and vertical scales and the axis labels in (b) are the same as those in (a).

In order to verify whether the energy loss arises from the population transfer among the subsystems due to spontaneous decays, we treat the probe field as the perturbation and obtain the following equations:

\[
\frac{1}{c} \frac{\partial}{\partial t} \hat{\Omega}_p + \frac{\partial}{\partial x} \hat{\Omega}_p = i \eta \sum_{j=1}^{3} a_{p,j} \rho_{31,j},
\]

\[
\frac{\partial}{\partial t} \rho_{31,j} = \frac{i}{2} a_{p,j} \hat{\Omega}_p \hat{\Omega}_p + \frac{i}{2} a_{c,j} \hat{\Omega}_c \rho_{21,j} - \frac{\Gamma}{2} \rho_{31,j},
\]

\[
\frac{\partial}{\partial t} \rho_{21,j} = \frac{i}{2} a_{c,j} \hat{\Omega}_c \rho_{31,j} - \gamma \rho_{21,j},
\]

where \( j = -1, 0, \) and 1 represent the three EIT subsystems containing the \( |m=-1\rangle, |m=0\rangle, \) and \( |m=1\rangle \) ground states,

respectively. The coefficient of 1/3 between \( a_{p,j} \) and \( \hat{\Omega}_p \) in Eq. (9) indicates that all population is equally distributed in the three ground states driven by the probe field. Due to the nature of the first-order perturbation, the population distribution does not change in the calculation. We numerically solve the above equations with the same parameters as those of Fig. 3(a). The predictions from the perturbation calculation shown in Fig. 3(b) are very similar to those from the non-perturbation calculation shown in Fig. 3(a). There is also an energy loss of about 49\% caused by the manipulation of the polarization of the probe pulse in the perturbation calculation. Hence, the population transfer among the subsystems due to spontaneous decays contributes little to the energy loss. The theoretical predictions shown by the dotted lines in Figs. 2(a) and 2(b) are calculated from Eqs. (8)–(10), in good agreement with the experimental data. See Ref. [27] for the details of the calculation and the experiment.

We design a more ideal condition in order to understand the origin of the energy loss. In this condition, \( \gamma = 0 \); the square probe pulse has the slowly rising and falling edges to well satisfy the EIT bandwidth; the optical density is large enough such that about half of the probe pulse can be stored in the medium; the writing and reading fields are switched adiabatically to avoid any effect from the rapid changes of their intensities. With all other parameters the same as those of Fig. 3, we use Eqs. (8)–(10) to calculate the probe transmissions. The black and red solid lines in Fig. 4(a) show the probe transmissions without and with manipulation of the retrieved polarization. There is also an energy loss of about 49\% caused by the manipulation process. The following discussion will show that the ratio of \( a_{p,j} \) to \( a_{c,j} \) in each three-
state subsystem is the crucial factor for the energy loss. The amplitude of the ground-state coherence is determined by the ratio of the Rabi frequencies of the probe and the coupling fields—i.e.,

\[ \rho_{21,j} = \frac{1}{3} \left( \frac{a_{p,j}}{a_{c,j}} \right) \left( \hat{\rho}_p - \hat{\Omega}_c \right). \]  

(11)

This equation is also the consequence of the adiabatic condition for the storing and retrieving processes [4,22]. The 1/3 in the right-hand term represents the population in each of the three ground states driven by the probe field. If the three \((a_{p,j}/a_{c,j})\)'s change during the retrieval, the equation may not be simultaneously satisfied by all the subsystems without the change of \(\rho_{21,j}\). For the \(\sigma_x\) writing field and the \(\sigma_z\) reading field, \((a_{p,j}/a_{c,j})\)'s are \(\sqrt{1/3}, 1, \text{and } \sqrt{3}\) during writing and \(-\sqrt{3}, -1, \text{and } -\sqrt{1/3}\) during reading, where \(j=1, 0, \text{and } 1\), respectively. The red lines in Fig. 4(b) show that Re[\(\rho_{21,j=1}\)] decreases and Re[\(\rho_{21,j=0}\)] and Re[\(\rho_{21,j=1}\)] increases. According to Eq. (10), the change of Re[\(\rho_{21,j}\)] must be accompanied by non-negligible Im[\(\rho_{21,j}\)] as demonstrated by the red lines in Fig. 4(c). The three \(\rho_{21,j}\)'s eventually evolve to the final values which all satisfy Eq. (11). These final values determine the amount of the energy loss of the retrieved probe pulse. Without the change of Re[\(\rho_{21,j}\)] shown by the black lines in Fig. 4(b), Im[\(\rho_{21,j}\)] is negligible as shown by the black lines in Fig. 4(c); the probe energy is intact. Hence, the incompatibility between the stored ground-state coherence and the ratio of the probe and coupling Clebsch-Gordan coefficients causes the energy loss in the manipulation of the retrieved polarization.

The values of \(\rho_{21,j}\)'s due to the change of the Clebsch-Gordan coefficients can be determined analytically. In the following description, we define \(R_j = a_{p,j}/a_{c,j}\) and use the superscripts of \((w)\) and \((r)\) to denote the values before writing and after reading, respectively. Because the reading field is turned on adiabatically and \(c \eta\) is usually very large in the light-storage studies, the sum of \(\rho_{21,j}\) weighted by \(d_{p,j}^{(r)}\) on the right-hand side of Eq. (8) should be nearly zero during the retrieval as demonstrated by the light gray line in Fig. 4(c)—i.e.,

\[ \sum_j d_{p,j}^{(r)} \rho_{21,j}(t) = 0. \]  

(12)

We multiply Eq. (10) of each subsystem by \(d_{p,j}^{(r)}\) divide it by \(d_{c,j}^{(r)}\), and drop the negligible term of \(-\gamma \rho_{21,j}\). The sum of Eq. (10) of all subsystems is given by

\[ \frac{\partial}{\partial t} \left( \sum_j R_j^{(r)} \rho_{21,j}(t) \right) = \frac{i}{2} \hat{\Omega}_p \sum_j d_{p,j}^{(r)} \rho_{21,j}(t) \approx 0. \]  

(13)

Therefore, \(\sum_j R_j^{(r)} \rho_{21,j}(t)\) during the reading process is constant as demonstrated by the light gray line in Fig. 4(b) and

\[ \sum_j R_j^{(r)} \rho_{21,j} = \sum_j R_j^{(w)} \rho_{21,j}. \]  

(14)

The \(\rho_{21,j}^{(w)}\) in the right-hand term represents the stored coherence or the value of \(\rho_{21,j}(t)\) just before the reading; the \(\rho_{21,j}^{(r)}\) in the left-hand term represents the value of \(\rho_{21,j}(t)\) just after the reading. All the \(\rho_{21,j}^{(r)}\)'s after the retrieval satisfy Eq. (11) and the population is equally distributed among the ground states driven by the probe field. We then obtain

\[ \frac{\rho_{21,j}^{(r)}}{R_{j}^{(r)}} = \frac{\rho_{21,j}^{(w)}}{R_{j}^{(w)}} = \cdots \Rightarrow \frac{\rho_{21,j}^{(r)}}{R_{j}^{(r)}} = \cdots . \]  

(15)

Similarly, all the \(\rho_{21,j}^{(w)}\)'s also satisfy Eq. (11) and

\[ \frac{\rho_{21,j}^{(w)}}{R_{j}^{(w)}} = \frac{\rho_{21,j}^{(w)}}{R_{j}^{(w)}} = \cdots \Rightarrow \frac{\rho_{21,j}^{(w)}}{R_{j}^{(w)}} = \cdots . \]  

(16)

In Eq. (14), we can replace all \(\rho_{21,j}^{(r)}\)'s by a single \(\rho_{21,1}^{(r)}\) and all \(\rho_{21,j}^{(w)}\)'s by a single \(\rho_{21,1}^{(w)}\) according to Eqs. (15) and (16), and obtain

\[ \sum_j R_j^{(r)} \rho_{21,j}^{(r)} = \sum_j R_j^{(w)} \rho_{21,j}^{(w)}. \]  

(17)

Hence,

\[ \rho_{21,1}^{(r)} = \rho_{21,1}^{(w)}. \]  

(18)

In general, the value of each \(\rho_{21,j}\) is given by

\[ \rho_{21,j}^{(r)} = \sum_k R_{jk}^{(w)} \rho_{21,k}^{(w)}. \]  

(19)

The \(\rho_{21,1}^{(w)}\)'s and \(\rho_{21,1}^{(r)}\)'s in the manipulation of the retrieved polarization shown in Fig. 4(b) are in good agreement with the above results. The ratio of the retrieved and stored energy densities \(u_{\text{out}}\) and \(u_{\text{in}}\) can be written as

\[ \frac{u_{\text{out}}}{u_{\text{in}}} = \frac{[\hat{\Omega}_p^{(r)}]_2^2 u_g^{(r)}}{[\hat{\Omega}_p^{(w)}]_2^2 u_g^{(w)}}. \]  

(20)

where \(\hat{\Omega}_p^{(w)}\) is the value of \(\hat{\Omega}_p\) before the writing, \(\hat{\Omega}_p^{(r)}\) is that after the reading, and \(u_g^{(w)}\) and \(u_g^{(r)}\) are the group velocities. Since \(1/\sqrt{u_g^{(w)}} \approx \sum_j R_j^{(w)}/[\hat{\Omega}_j^{(w)}]^2\) and the ratio of \(\hat{\Omega}_p^{(r)}\) to \(\hat{\Omega}_p^{(w)}\) can be replaced by the ratio of \(\rho_{21,1}^{(r)} \hat{\Omega}_j^{(r)}/R_j^{(r)}\) to \(\rho_{21,1}^{(w)} \hat{\Omega}_j^{(w)}/R_j^{(w)}\) for any \(j\), the above equation becomes

\[ \frac{u_{\text{out}}}{u_{\text{in}}} = \frac{[\sum_j R_j^{(r)}/R_j^{(w)}]_2^2 \sum_j [R_j^{(w)}]_2^2}{[\sum_j [R_j^{(r)}]_2^2] [\sum_j [R_j^{(w)}]_2^2]}. \]  

(21)

This indicates that the retrieved energy can only be less than or equal to the stored energy when the Clebsch-Gordan coefficients change. With the \(R^{(w)}\)'s and \(R^{(r)}\)'s in our experiment, Eq. (21) gives the energy-density ratio of 48% which is consistent with the experimental data and the predictions of the numerical calculations.

The precise criterion for the energy preservation during the manipulation of the retrieval is described below. According to Eqs. (19) and (21), as long as \(R_j\)'s change by the same factor in all subsystems, all \(\rho_{21,j}\)'s remain the same during reading and the energy of the retrieved probe pulse is intact.
We provide two examples to illustrate the criterion. The first example is the manipulation of the wavelength from the 87Rb lines of 87Rb atoms and all the laser fields have the same $\sigma_+ \leftrightarrow \sigma_-$ polarization. The input probe pulse and the writing field drive the $|S_{1/2}, F=1\rangle \rightarrow |P_{3/2}, F'=2\rangle$ and $|S_{1/2}, F=2\rangle \rightarrow |P_{3/2}, F'=2\rangle$ transitions, and $R_{ij}^{(\omega)}$'s are $\sqrt{1/3}$, 1, and $\sqrt{3}$. In the manipulation of the retrieved wavelength, the output probe pulse and the reading field drive the $|S_{1/2}, F=1\rangle \rightarrow |P_{3/2}, F'=2\rangle$ and $|S_{1/2}, F=2\rangle \rightarrow |P_{3/2}, F'=2\rangle$ transitions, and $R_{ij}^{(\omega)}$'s are $\sqrt{1/3}$, $-1$, and $\sqrt{3}$. Although each $R_{ij}$ changes sign, there is no energy loss caused by the retrieval as shown in Fig. 5(a). The second example is the manipulation of the wavelength from the D2 to D1 lines of 87Rb atoms. All the laser fields also have the same $\sigma_+ \leftrightarrow \sigma_-$ polarization. The input probe pulse and the writing field drive the $|S_{1/2}, F=2\rangle \rightarrow |P_{3/2}, F'=3\rangle$ and $|S_{1/2}, F=3\rangle \rightarrow |P_{3/2}, F'=3\rangle$ transitions, and $R_{ij}^{(\omega)}$'s are $\sqrt{16/125}$, $\sqrt{8/25}$, $\sqrt{16/25}$, $\sqrt{32/25}$, and $\sqrt{16}/5$. In the manipulation of the retrieval, the output probe pulse and the reading field drive the $|S_{1/2}, F=2\rangle \rightarrow |P_{3/2}, F'=3\rangle$ and $|S_{1/2}, F=3\rangle \rightarrow |P_{3/2}, F'=3\rangle$ transitions, and $R_{ij}^{(\omega)}$'s are $-(5/4)\sqrt{16/25}$, $-(5/4)\sqrt{8/25}$, $-(5/4)\sqrt{32/25}$, and $-(5/4)\sqrt{16}/5$. Since all $R_{ij}$'s change by the same factor, the probe energies with and without the manipulation of the retrieved wavelength are the same as shown by the red and black lines of Fig. 5(b). Due to the manipulation, the Rabi frequency of the reading field is effectively reduced by a factor of 5/4. Consequently, the retrieved pulse has a larger width and smaller amplitude. This has been explained in Ref. [20], and the retrieved amplitude and width are exactly as expected. In the above two examples, all $\text{Re}[\rho_{21,ij}]$'s maintain constant and all $\text{Im}[\rho_{21,ij}]$'s are negligible during the reading process in the wavelength manipulation.

Once the origin of the energy loss is understood, it can be avoided. In Ref. [27], we optically pumped all population to the $|F=1, m=0\rangle$ ground state and observed very little energy loss in the manipulation of the retrieved polarization. This solution came from the idea that the Clebsch-Gordan coefficients of the $\sigma_+ \leftrightarrow \sigma_-$ transitions of the $|F=1, m=0\rangle$ and $|F=2, m=0\rangle$ ground states are symmetrical. Since only the ground-state coherence between $|F=1, m=0\rangle$ and $|F=2, m=0\rangle$ exists during storage of the probe pulse, the Clebsch-Gordan coefficients are unchanged in the manipulation of the retrieved polarization. Nevertheless, the symmetry of the Clebsch-Gordan coefficients is not the necessary condition for avoiding energy loss. With the detailed analysis in this work, it is now clear that one can manipulate the retrieved polarization without energy loss as long as all population accumulates in a single Zeeman state.

In conclusion, the incompatibility between the stored ground-state coherence and the ratio of the probe and coupling Clebsch-Gordan coefficients causes the energy loss in the manipulation of the retrieval. When the ratios of the probe and coupling Clebsch-Gordan coefficients change by the same factor in all subsystems, the incompatibility does not occur and the energy of the retrieved probe pulse is intact. One can also manipulate the retrieved polarization without energy loss as long as all population accumulates in a single Zeeman state. We have demonstrated that degenerate Zeeman states play an important role in the manipulation of the retrieval of stored light pulses and should be taken into consideration in practical applications utilizing the slow light and the light-storage techniques.

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