Superluminal propagation of an optical pulse in a Doppler-broadened three-state single-channel active Raman gain medium

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Using a single-channel active Raman gain medium, we show a 220 ± 20 ns advance time for an optical pulse of $\tau_{FWHM} = 15.4$ μs propagating through a 10 cm medium, a lead time that is comparable to what was reported previously using a two-mode pump field. In addition, we have verified experimentally all the features associated with this single-channel Raman gain system.

In recent work [1] on room-temperature Cs atoms with two close Raman resonances created by two pump frequencies, a 3.7 μs optical pulse was observed to travel superluminally with about 62 ns lead time, corresponding to about 1.6% of the full width at half maximum (FWHM) of the probe pulse used in the experiment. In this paper, we demonstrate experimentally that a single Raman resonance achieved by a *single-mode* pump field can also bring about the same superluminal propagation. Our experiments are based on the theoretical work of Payne *et al.* [2,3], which showed analytically [4] that the reason for the observed apparent superluminal propagation is the virtual population produced by a pump field, and it is not necessary to have a two-mode pump field.

Before describing our experiments, we first cite several pioneering works on superluminal propagation of optical pulses in anomalous-dispersion regions of optical media. These include the early work of Garrett and McCumber [5] on the propagation of a Gaussian light pulse through an anomalous dispersive medium, the study by Casperson and Yariv [6] on pulse proration in a high-gain medium, and the work by Chu and Wong [7] on linear pulse propagation in an absorbing medium. More recently, Chiao [8] investigated the superluminal propagation of wave packets in transparent media with an inverted atomic population, Bolda *et al.* [9] studied optical pulse propagation with negative group velocity due to a nearby gain line, and Steinburg *et al.* [10,11] investigated the single-photon tunneling time and superluminal propagation in a medium with a gain doublet. Stenner *et al.* [12] studied the information speed in superluminal propagation.

We use ^{85}Rb to demonstrate the superluminal propagation of an optical pulse in a single-frequency-channel active Raman gain medium (Fig. 1). An external cavity diode laser system produces nearly 400 mW power and is locked to a rubidium reference cell. The output of the laser passes through a 3 GHz acousto-optical modulator (AOM) to generate two laser beams with the frequency difference of the ground-state hyperfine levels $|1\rangle = |5S_{1/2}, F=2\rangle$ and $|3\rangle = |5S_{1/2}, F=3\rangle$ of ^{85}Rb atoms. The AOM is driven by a rf synthesizer which allows accurate adjustment so that the two-photon resonance can be precisely maintained. A frequency shifter is then added to the path of the probe field, allowing accurate and independent tuning of the small two-

photon detuning. Figure 1 shows the energy level diagram and relevant laser couplings for the present experiment. The probe and the pump fields couple the $|1\rangle = |5S_{1/2}, F=2\rangle \rightarrow |2\rangle = |5P_{1/2}\rangle$ and $|3\rangle = |5S_{1/2}, F=3\rangle \rightarrow |2\rangle = |5P_{1/2}\rangle$ transitions, respectively. Both lasers are detuned, however, on the high-energy side of the $5P_{1/2}$ hyperfine manifold by a large (≈ 3 GHz) one-photon detuning. In this experiment the strong pump field serves two purposes: (1) it is on the $|1\rangle = |5S_{1/2}, F=2\rangle \rightarrow |2\rangle = |5P_{1/2}\rangle$ resonance (see the discussion below), and (2) it also pumps the $|3\rangle = |5S_{1/2}, F=3\rangle \rightarrow |2\rangle = |5P_{1/2}\rangle$ transitions with a large one-photon detuning. The former serves as the optical pumping field that keeps nearly all population in the state $|3\rangle$, whereas the latter provides the
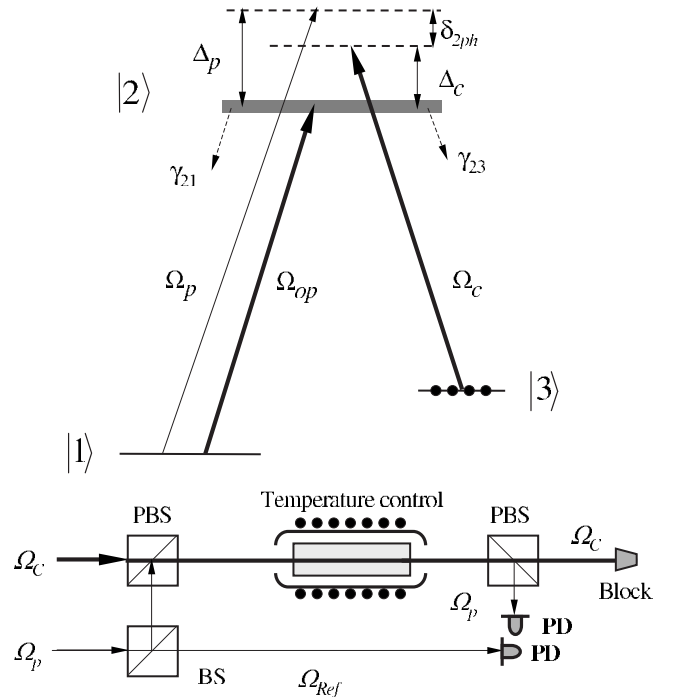

FIG. 1. Top panel: Energy level diagram and relevant laser couplings. Level assignment: $|1\rangle = |5S_{1/2}, F=2\rangle$, $|2\rangle = |5P_{1/2}\rangle$, and $|3\rangle = |5S_{1/2}, F=3\rangle$. Because of the large one-photon detuning, the hyperfine splitting, Doppler broadening, and small ac Stark shift induced in the state $|2\rangle$ have been neglected. Lower panel: Schematics of the experimental setup.

active Raman gain to a weak probe field tuned near the $|1\rangle = |5S_{1/2}, F=2\rangle \rightarrow |2\rangle = |5P_{1/2}\rangle$ transition. It is in this active Raman gain region where we study the superluminal propagation of the weak probe pulse. Contrary to the two-frequency pump field scheme used in Ref. [1], there is no second pump frequency component that can lead to the interference effect discussed in [1]. Finally, the cell for the medium is 100 mm in length and 25 mm in diameter with antireflection coating on both ends, and is placed in a temperature-controlled enclosure. Experiments were carried out in the temperature region of $T = 60-90$ °C.

Theoretically, the system described above can be understood using a simple lifetime broadened three-state model where nearly all population remains in the initial state $|3\rangle$. This is because, under our experimental conditions, the one-photon detuning is sufficiently large in comparison with the Doppler-broadened linewidths and the ac Stark shift produced in the $5P_{1/2}$ hyperfine states by the optical pumping field. Consequently, the Doppler broadenings of the hyperfine states are unimportant and the contributions by the two hyperfine states can easily be included. Thus, we neglect the optical pumping field, the Doppler broadening, and the hyperfine splitting, and assume that nearly all population is maintained in the state $|3\rangle$. Since the probe field is very weak [13,14], there is never appreciable population that can be transferred to the state $|1\rangle$.

The equations of motion for the relevant density matrix elements are given as (assuming $\rho_{33} \approx 1$)

$$\frac{\partial \rho_{12}}{\partial t} \approx -i\Omega_c^* e^{i\Delta_c t} \rho_{13} - \gamma_{12}\rho_{12}, \quad (1a)$$

$$\frac{\partial \rho_{13}}{\partial t} \approx i\Omega_p^* e^{i\Delta_p t} \rho_{23} - i\Omega_c e^{-i\Delta_c t} \rho_{12} - \gamma_{13}\rho_{13}, \quad (1b)$$

$$\frac{\partial \rho_{23}}{\partial t} \approx i\Omega_p e^{-i\Delta_p t} \rho_{13} + i\Omega_c e^{-i\Delta_c t} - \gamma_{23}\rho_{23}, \quad (1c)$$

where $\Omega_p = D_{21}E_{p0}/(2\hbar)$ and $\Omega_c = D_{23}E_{c0}/(2\hbar)$ are the half Rabi frequencies of the probe (E_{p0}) and pump (E_{c0}) fields for the corresponding transitions, and D_{ij} and γ_{ij} are the electric dipole moment and the decoherence rate of the respective transitions. In addition, Δ_p and Δ_c are one-photon detunings to the state $|2\rangle$ by the probe and pump fields.

Equation (1) must be solved self-consistently together with the wave equation for the pulsed probe field in order to correctly predict the propagation characteristics of the probe field. Payne *et al.* have shown that the positive frequency part of the probe field amplitude, in the case of an unfocused plane wave and within the slowly varying amplitude and adiabatic approximations, must satisfy the wave equation which, in the Fourier transform space, is given by

$$\frac{\partial \Lambda_p}{\partial z} - i\frac{\omega}{c}\Lambda_p = \frac{i\kappa_{12}|\Omega_c|^2 W(\omega)}{(\Delta_c - i\gamma_{23})(\Delta_c + i\gamma_{21})}\Lambda_p, \quad (2a)$$

$$W(\omega) = \frac{1}{\omega - (\Delta_c - \Delta_p - i\gamma_{31}) - |\Omega_c|^2/(\Delta_c + i\gamma_{21})}, \quad (2b)$$

where Λ_p is the Fourier transform of the probe field Rabi frequency Ω_p and ω is the Fourier-transform variable. In addition, $\kappa_{12} = 2\pi N_0 \omega_p |D_{21}|^2/c$, where N_0 and ω_p are the atom number density and frequency of the probe field, respectively. We note that the last term in the denominator in Eq. (2b) leads to a small shift of the Raman gain peak and a shift of the minimum group velocity. This term has its origin in the small ac Stark shift due to the pump field.

Payne *et al.* have carried out a non-steady-state calculation and shown the group velocity of a probe pulse to be (for the two-photon detuning $|\delta_{2ph}| = |\Delta_p - \Delta_c| \gg \gamma_{31}$, and $|\delta_{2ph}\tau| \gg 1$),

$$V_g \approx -\frac{\Delta_c^2 \delta_{2ph}^2}{\kappa_{12}|\Omega_c|^2}, \quad (3)$$

where $|\Delta_p|, |\Delta_c| \gg \gamma_{21}$ (assuming $\gamma_{21} \approx \gamma_{23}$), $|\Delta_p|, |\Delta_c| \gg |\Omega_c|$, and $|\Omega_c\tau|^2/(\Delta_c\tau)| \ll |\delta_{2ph}\tau|$ have been used. We note that the negative sign indicates superluminal propagation. Equation (3) clearly demonstrates that it is the virtual population generated by the pump field, i.e., $|\Omega_c/\Delta_c|^2$, and the two-photon detuning that are responsible for apparent superluminal propagation, not some interference effect.

Our experiment is aimed at demonstrating three predictions given in Eq. (3): (1) the existence of nondistorted apparent superluminal propagation in a single-pump-frequency-channel active Raman gain medium; (2) the characteristic inverse parabolic dependence of the group velocity on the pump field Rabi frequency; and (3) quadratic dependence of the group velocity on the two-photon detuning (neglecting the small ac Stark shift).

In Fig. 2 we show typical data for the advanced propagation of a Gaussian probe pulse with FWHM pulse width of $\tau_{FWHM} = 15.4$ μs. In order to show the advanced time clearly we have plotted only a small portion of the probe pulse profile. The lower left panel shows the advance of the front edge (solid line) in comparison with a reference Gaussian pulse (dashed line) that traverses through the air, whereas the lower right panel depicts the advance of the rear edge in comparison with the same reference pulse. We observed excellent signal-to-noise ratio [15], and the fitting of the data to a Gaussian using a standard statistical routine yielded the advanced time of $\delta t = 220 \pm 20$ ns, or about 1.4% of the FWHM pulse width of the Gaussian probe pulse. This is comparable to the advanced time of the two-pump-frequency experiment, where 1.6% advanced time was reported [1].

In Fig. 3(a) we show the group velocity of a Gaussian probe pulse as a function of the single-frequency pump field Rabi frequency. The inverse quadratic behavior (notice the negative sign) of the group velocity as a function of the pump field Rabi frequency is in accord with the prediction based on Eq. (3) when the pump field Rabi frequency is relatively small and therefore the ac Stark shift can be neglected. As the pump field Rabi frequency is strong, the ac Stark shift term in Eq. (3) must be considered, together with the ground-state population depletion. In particular, a sizable
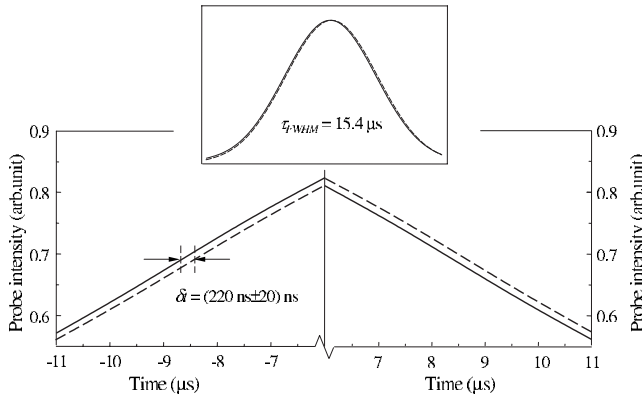
FIG. 2. Typical data (solid line) showing the advanced propagation of a Gaussian probe pulse of full width at half maximum $\tau_{FWHM}=15.4$ μs (top panel; the dashed line is for a reference pulse). Here, we have plotted portions of the front (lower left) and rear (lower right) edges of the probe pulse (solid line) in comparison with a reference Gaussian pulse that travels in the air (dashed line). The vertical axis is the normalized optical intensity of the Gaussian probe pulse. The probe pulse has a slightly higher intensity because of the Raman gain. There is, however, no detectable pulse spreading or distortion when the probe field amplitude is normalized. Parameters: $|\Omega_c|=2\pi\times25$ MHz, $\delta_{2ph}\approx400$ kHz, $T=60$ °C, and $\Delta_c=2\pi\times3$ GHz.

*ac Stark shift leads to a transition from inverse quadratic behavior to a quadratic increase of the magnitude of the group velocity, as can be seen from Eq. (3), in the large-Ω_c limit. In this limit, however, one should rederive Eqs. (1)–(3) by including the equation of motion for the ground-state population and the possible effects due to "accidental" opti-

cal pumping and the four-wave mixing (FWM) field generated by the allowed transition $|2\rangle\rightarrow|3\rangle$ due to this accidental optical pumping field. This complicated situation will be studied later.

In Fig. 3(b) we show the group velocity of the probe pulse as a function of the two-photon detuning $\delta_{2ph}=\Delta_p-\Delta_c$. Equation (3) predicts that the group velocity is a quadratic function of the two-photon detuning. In addition, when $|\delta_{2ph}|\gg\text{Max}[|\Omega_c|^2/|\Delta_c|,\gamma_{31}]$, the group velocity should be symmetric with respect to the two-photon detuning. This is indeed what has been observed experimentally. Here, we plot the red two-photon detuning portion of the group velocity dispersion curve. It is worth noting that the vertical dashed line indicates the approximate two-photon detuning where the propagation characteristics of the probe pulse propagation change from subluminal (above the horizontal dashed line) to superluminal (below the horizontal dashed line). This behavior is in complete agreement with the predictions based on Eq. (2b). Indeed, from Eq. (2b) it is seen that, when $\delta_{2ph}\approx|\Omega_c|^2/\Delta_c$, the $i\gamma_{31}$ term dominates the denominator [16], resulting in subluminal propagation. For red-detuned two-photon detuning, this point occurs at about

$$\delta_{2ph}\approx\frac{|\Omega_c|^2}{\Delta_c}\approx-2\pi\times200\text{ kHz}.$$

This is very close to what can be seen from Fig. 3(b).

It is also worth pointing out another observation that is in accord with the prediction of Eqs. (2a) and (2b) and Ref. [2]. We have observed about 1% *narrowing* of the probe pulse for $|\delta_{2ph}|\leq2\pi\times400$ kHz. Superluminal propagation without pulse narrowing, as depicted in Fig. 2, is observed for $|\delta_{2ph}|>2\pi\times400$ kHz, as expected from Eq. (3) and Ref. [2]. With a shorter probe pulse such as $\tau_{FWHM}=5$ μs, we have observed nearly 10% pulse narrowing in addition to a 295 ± 20 ns lead time. This lead time is about 5% of the FWHM width of the probe pulse used.

We now discuss the possible effect of the accidental optical pumping field and the likelihood of possible FWM generation due to the allowed transition $|2\rangle\rightarrow|3\rangle$ because of the presence of this field. In our simple three-state treatment given before, the optical pumping field is not included. Consequently, the possible generation of a FWM field is also neglected. We neglect the accidental optical pumping field and the FWM field in our treatment for the following reasons. (1) Theoretically, with all four fields included, one could not find a clean analytical solution to the problem; extracting the essential physics is nontrivial. In addition, a steady-state solution is highly questionable. This is because the equations of motion contain fast-oscillating factors that cannot be removed by simple phase transformation. (2) Although the accidental optical pumping field is on resonance with the $|1\rangle\rightarrow|2\rangle$ transition, and therefore causes perturbation to the dispersion properties of state $|2\rangle$, this perturbation is not very significant in the gain process and the propagation of the probe field simply because the detunings from state $|2\rangle$ are so large. It should be pointed out that the contribution by FWM is small in our case even with a sizable gain. Theoretically, however, the gain increases with the pump field
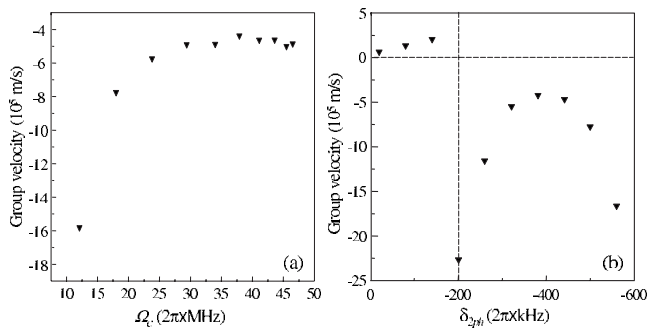

FIG. 3. (a) Group velocity of a Gaussian probe pulse as a function of the pump field Rabi frequency. As expected, the group velocity exhibits inverse quadratic dependence on the Rabi frequency of the single-channel pump field, as predicted in Eq. (3). For this plot $\delta_{2ph}\approx2\pi\times400$ kHz, $\Delta_c=2\pi\times2.2$ GHz, and temperature $T=60$ °C. (b) Group velocity of a Gaussian probe pulse as a function of the two-photon detuning δ_{2ph}. In the superluminal region where $|\delta_{2ph}|\gg\gamma_{31}$, the group velocity exhibits quadratic dependence on the two-photon detuning, as predicted from Eq. (3). When $|\delta_{2ph}-|\Omega_c|^2/\Delta_c|\ll\gamma_{31}$ the propagation characteristics changes from superluminal to subluminal. For this plot $|\Omega_c|=2\pi\times25$ MHz, $\Delta_c=2\pi\times3$ GHz, and temperature $T=60$ °C. The Raman gain coefficient of $G_{Raman}\simeq0.05$ cm^{-1} used in our experiment is similar to the gain coefficient of $G\simeq0.04$ cm^{-1} reported in Ref. [1].

intensity, and ground-state depletion will occur at sufficient gain. In this strong-pumping regime, the inclusion of the accidental optical pumping field and FWM generation becomes necessary in order to accurately predict the propagation dynamics.

In conclusion, we have verified experimentally all predictions based on Eq. (3) and Ref. [2]. Using a Raman gain coefficient similar to that reported in Ref. [1] (see [15]) we have demonstrated experimentally similar gain-assisted superluminal propagation with only a single pump frequency mode. We stress that the probe field is a transform-limited Gaussian pulse with a single longitudinal frequency mode. It does not, in every sense of slowly varying amplitude and phase approximation, have any other frequency components

or amplitude excursions. Thus, the observed superluminal propagation should not be attributed to any interference effect between different frequency components [1]. In addition, we have demonstrated four features of a single-channel active Raman gain medium for superluminal propagation predicted before: (1) the inverse quadratic dependence of the group velocity as a function of the pump field Rabi frequency; (2) the quadratic and symmetric dependence of the group velocity as a function of the two-photon detuning; (3) a group velocity dispersion region where the propagation characteristics change from superluminal to subluminal; and (4) probe pulse narrowing when $|\delta_{2ph}| < \gamma_{31}$. None of these features has been demonstrated before to our knowledge with a narrowband gain medium.

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[3] We note that Ref. [2] also predicts an ultraslow propagation of the probe field when $\delta_{2ph}=0$. This prediction agrees with experimental observations. See S. Inouye, R. F. Low, S. Gupta, T. Pfau, A. Gorlitz, T. L. Gustavson, D. E. Pritchard, and W. Ketterle, Phys. Rev. Lett. **85**, 4225 (2000); W. Ketterle and S. Inouye, C. R. Acad. Sci., Ser. IV Phys. Astrophys. **2**, 336 (2001).

[4] It is widely believed that certain interference effects lead to the observed superluminal propagation [1]. We note, however, that it is incorrect to attribute the origin of the observed superluminal propagation to interference between the two pump frequency components or between the multiple frequency components of the pulsed probe field. It is the gain nature of the anomalous dispersion that is at the root of the apparent superluminal propagation observed.

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[15] In our experiment, we choose the Raman gain coefficient to be about $G_{Raman} \approx 0.05$ cm^{-1}. This is comparable to the gain coefficient of $G \approx 0.04$ cm^{-1} reported in Ref. [1].

[16] That is, when the denominator in Eq. (2b) changes from mostly real to mostly imaginary, the propagation characteristics change from superluminal to subluminal.