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 \((220 \pm 20)\)ns advance time for an optical pulse of \(\tau_{FWHM} = 15.4\) ns propagating through a 10 cm medium, a lead time that is comparable to what was reported previously. In addition, we have verified experimentally all the features associated with this single channel Raman gain system. Our results show that the reported gain-assisted superluminal propagation should not be attributed to the interference between the two frequencies of the pump field.

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In a recent work [1] in room-temperature Cs atom and with two nearby Raman resonances created by two pump frequencies, a 3.7 \(\mu\)s optical pulse is 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 their conclusion the authors have attributed the observed superluminal propagation to the interference effect between different frequency components of the pump field (in the anomalous dispersion region provided by the two resonances) and also to the gain that occurs at the frequency of the probe pulse. Payne et al. [2] have shown analytically, however, that in this case the superluminal propagation is not based on any interference effect. Indeed, they have shown that with a pump field of single frequency, the superluminal propagation should still be present. In this work, we demonstrate experimentally that an active Raman gain medium enabled by a single frequency pump field can indeed bring about the superluminal propagation, as predicted in Ref. [3]. Thus, while the attribution of the observed superluminal propagation is assisted by Raman gain it is correct it is incorrect to attribute the origin of the superluminal propagation to an interference between the two pump frequency components. It is the gain nature of the anomalous dispersion that is at the root of the apparent superluminal propagation observed.

Before describing our experiments we first cite several pioneer works on superluminal propagations of optical pulses in anomalous dispersion regions of optical media. This include the early work of Garrett and McCumber [4] on the propagation of a Gaussian light pulse through an anomalous dispersive medium, the study by Casperson and Yariv [5] on pulse proration in a high-gain medium, the work by Chu and Wong [6] on the linear pulse propagation in an absorbing medium. More recently, Chiado [7] investigated the superluminal propagation of wavepackets in transparent media with inverted atomic population, Bolda et al. [8] studied optical pulse propagation with negative group velocity due to a nearby gain line, and Steinburg et al. [9, 10] investigated single photon tunneling time and superluminal propagation in a medium with a gain doublet. Stenner et al. [11] have studied the information speed in superluminal propagation.

We use \(^{85}\text{Rb}\) 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 frequency difference of the ground state hyper-fine levels \(|1\rangle = |5S_{1/2}, F = 2\rangle\) and \(|3\rangle = |5S_{1/2}, F = 3\rangle\) of \(^{85}\text{Rb}\) 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 \(|1\rangle = |5S_{1/2}, F = 2\rangle\) to \(|2\rangle = |5P_{1/2}\rangle\) and \(|3\rangle = |5S_{1/2}, F = 3\rangle\) to \(|2\rangle = |5P_{1/2}\rangle\) transitions, respectively. Both lasers are detuned, however, on the high energy side of the 5\(P_{1/2}\) hyper fine manifold by a large (\(\geq 3\)GHz) one-photon detuning. In this experiment the strong pump field serves two purposes: (1) it is on the \(|1\rangle = |5S_{1/2}, F = 2\rangle\) to \(|2\rangle = |5P_{1/2}\rangle\) resonance (see discussion below), and (2) it also pumps the \(|3\rangle = |5S_{1/2}, F = 3\rangle\) to \(|2\rangle = |5P_{1/2}\rangle\) 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 active Raman gain to a weak probe field tuned near the \(|1\rangle = |5S_{1/2}, F = 2\rangle\) to \(|2\rangle = |5P_{1/2}\rangle\) 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 the 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 anti-reflection coating on both ends, and is placed in a temperature controlled enclosure. Experiments were carried out in the temperature region of $T = 60\, ^\circ \text{C}$ to $T = 90\, ^\circ \text{C}$.

Theoretically, the system described above can be understood using a simple life time broadened three-state model where nearly all population remains in the initial state [3]. This is because under our experimental conditions the one-photon detuning is sufficiently large in comparison with the Doppler broadened line widths 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 un-important and the contributions by the two hyperfine states can be easily 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 [12, [3], there is never appreciable population can be transferred to the state $|1\rangle$.

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

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

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

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

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 respective transitions, $D_{ij}$ and $\gamma_{ij}$ are the electric dipole moment and the decoherence rate of the respective transitions. In addition, $\Delta_p$ and $\Delta_c$ 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 un-focused 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

$$\frac{\partial \Lambda_p}{\partial z} - \frac{\omega}{c} \Lambda_p = -i \frac{\kappa_{12} |\Omega_c|^2 W(\omega)}{(\Delta_c - i \gamma_{23}) (\Delta_c - i \gamma_{31})} \Lambda_p, \quad (2a)$$

$$W(\omega) = \frac{1}{\omega - (\Delta_c - \Delta_p - i \gamma_{31}) - \frac{|\Omega_p|^2}{(\Delta_c + i \gamma_{23})}}, \quad (2b)$$

where $\Lambda_p$ is the Fourier transform of the probe field Rabi frequency $\Omega_p$ and $\omega$ is the Fourier transform variable. In addition, $\kappa_{12} = 2\pi N_0 \omega_p |D_{21}|^2/c$ where $N_0$ and $\omega_p$ 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 the 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 as (for the two-photon detuning $|\delta_{2ph}| = |\Delta_p - \Delta_c| \gg |\gamma_{31}|$),

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

where $|\Delta_p|, |\Delta_c| \gg |\gamma_{21}$ (assuming $\gamma_{21} \approx \gamma_{23}$), and $|\Delta_p|, |\Delta_c| \gg |\Omega_c|$ have been used. We note that the negative sign indicates the superluminal propagation, thus proving the interference between the two-frequency components of a pump field is not the cause of the superluminal propagation. It is the active Raman gain that is responsible for apparent superluminal propagation.

Our experiment is aimed at demonstrating three predictions given in Eq. (3): (1) the existence of the non-distorted 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 a typical data of the advanced propagation of a Gaussian probe pulse with FWHM pulse width of $\tau_{FWHM} = 15.4\, \mu s$. 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 rare edge in comparison with the same reference pulse. We have observed excellent S/N ratio [14] and the fitting of the data to a Gaussian using a standard statistical routine has 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 advance time of the two-pump-frequency experiment where $1.6\%$ advanced time has been reported [1].

In Fig. 3a 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 the function of the pump field Rabi frequency is in accord to the prediction based on Eq. (3) when the pump field Rabi frequency is small relative small 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 ac Stark shift leads to the transition from the inverse quadratic behavior to the quadratic increase of magnitude of the group velocity as can be seen from Eq. (3) at large $\Omega_{c}$ limit. In this limit, however, one should re-derive Eqs. (1-3) by including the equation of motion for the ground state population and the possible effects due to the “accidental” optical pumping and a four-wave mixing (FWM) field generated by the allowed transition $|2\rangle \rightarrow |3\rangle$ due to the “accidental” optical pumping field. This complicated situation shall be studied later.

In Fig. 3b we show the group velocity of a 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 $|\delta_{2ph}| > \max[|\Omega_{c}|^{2}/|\Delta_{c}|, |\gamma_{31}|]$, 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 changes 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 the when $\delta_{2ph} \approx |\Omega_{c}|^{2}/|\Delta_{c}|$, the $i\gamma_{31}$ term dominates the denominator [13], resulting in a subluminal propagation. This point occurs at about, for red detuned two-photon detuning,

$$\delta_{2ph} \approx \frac{|\Omega_{c}|^{2}}{|\Delta_{c}|} \approx -2\pi \times 200 \text{ kHz}.$$ 

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

It is also worth pointing out another observation that is in accord with the prediction of Eq. (2a,2b) and ref. [3]. We have observed about 1% narrowing of the probe pulse for $|\delta_{2ph}| \leq 2\pi \times 400 \text{ kHz}$. Superluminal propagation without pulse narrowing, such as those depicted in Fig. 2, are observed for $|\delta_{2ph}| > 2\pi \times 400 \text{ kHz}$, as expected from Eq. (3) and ref. [3]. With a shorter probe pulse such as $\tau_{FWM} = 5 \mu s$, we have observed nearly 10% pulse narrowing in addition to a $(295 \pm 20)$ 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 likely hood of the possible FWM generation due to the allowed transition $|2\rangle \rightarrow |3\rangle$ because of the presence of the “accidental” optical pumping 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 because (1), theoretically with all four fields included, one could not find a clean analytical solution to the problem, therefore the extract of the essential physics is non-trivial. In addition, 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 to 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 intensity, and the ground state depletion will occur at sufficient gain. In this strong pumping regime, the inclusion of the “accidental” optical pumping field and the FWM generation become necessary in order to accurately predict the propagation dynamics.

In conclusion we have verified experimentally all predictions based on Eq. (3) and ref. [3]. We have demonstrated experimentally that the superluminal propagation of an optical pulse in an active Raman gain medium is not the result of the interference between two frequency components in the pump field. Indeed, with only a single frequency pump field where no second frequency component exists to allow the possibility of interference between the pump frequencies but with nearly the same Raman gain coefficient ($G_{Raman} \approx 0.05 \text{ cm}^{-1}$ in our experiment and $G \approx 0.04 \text{ cm}^{-1}$ reported in Ref. [4]) we have shown comparable advance times as reported before by other group. In addition, we have demonstrated four features of the single channel active Raman gain medium for superluminal propagation: (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) the group velocity dispersion region where the propagation characteristics changes from superluminal to subluminal; and (4) probe pulse narrowing when $|\delta_{2ph}| < |\gamma_{31}|$. None of these features has been demonstrated before with a narrowband gain medium.

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

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

**Figure captions**

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

**Figure 2** Typical data (solid line) showing the advanced propagation of a Gaussian probe pulse of full width at half maximum (FWHM) $\tau_{\text{FWHM}} = 15.4 \mu s$ (top panel, the dashed line is for a reference pulse). Here, we have plotted portions of the front (lower left) and rare (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: $|\Omega_c| = 2\pi \times 25 \text{ MHz}$, $\delta_{2\text{ph}} \approx 400 \text{ kHz}$, $T = 60^\circ \text{C}$, and $\Delta_c = 2\pi \times 3 \text{ GHz}$.

**Figure 3** (a) Plot of the group velocity of a Gaussian probe pulse as a function of the pump field Rabi frequency ($\Omega_c$). As expected, the group velocity exhibits the inverse quadratic dependence of the Rabi frequency of the single channel pump field, as predicted in Eq. (3). For this plot $\delta_{2\text{ph}} \approx 2\pi \times 400 \text{ kHz}$, $\Delta_c = 2\pi \times 2.2 \text{ GHz}$, and temperature $T = 60^\circ \text{C}$. (b) Plot of the group velocity of a Gaussian probe pulse as a function of the two-photon detuning $\delta_{2\text{ph}}$. In the superluminal region where $|\delta_{2\text{ph}}| \gg \gamma_{31}$, the group velocity exhibits the quadratic dependence on the two-photon detuning, as predicted from Eq. (3). When $|\delta_{2\text{ph}} - |\Omega_c|^2/\Delta_c| << \gamma_{31}$ the propagation characteristics changes from superluminal to subluminal. For this plot $|\Omega_c| = 2\pi \times 25 \text{ MHz}$, $\Delta_c = 2\pi \times 3 \text{ GHz}$, and temperature $T = 60^\circ \text{C}$. The Raman gain coefficient of $G_{\text{Raman}} \approx 0.05 \text{ cm}^{-1}$ used in our experiment is similar to the gain coefficient of $G \approx 0.04 \text{ cm}^{-1}$ reported in Ref. [1].
\[ \tau_{FWHM} = 15.4 \mu s \]

\[ \delta t = (220 \text{ ns} \pm 20) \text{ ns} \]
