Spectral Narrowing via Quantum Coherence

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We have studied the transmission through an optically thick $^{87}\text{Rb}$ vapor that is illuminated by monochromatic and noise broaden laser fields in $\Lambda$ configuration. The spectral width of the beat signal between the two fields after transmission through the atomic medium is more than 1000 times narrower than the spectral width of this signal before the medium.

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The resonant interaction of an quasi-monochromatic electromagnetic field with atomic media is very important because of its applications to spectroscopy, magnetometry, nonlinear optics and quantum information and computing. The interaction of a phase noise broaden electromagnetic field with media possessing a resonant absorption or transmission has been studied both experimentally and theoretically. In particular, several experiments involving noisy laser fields transmitted through a cell containing alkali atomic vapor were performed. For example, the study of the conversion of phase-noise to amplitude noise in dense Cs and Rb confirmed the theoretical prediction in \cite{14}. In addition, there are experimental studies of intensity fluctuations and correlations between the drive and probe fields in the electromagnetically induced transparency (EIT) regime \cite{9, 10, 11, 12}.

In this letter, we show that the spectral width of the beat signal between two lasers is modified by transmission through a medium to give a very narrow spectral feature that is more than 1000 times narrower than the spectral width of the probe laser radiation. This can find broad applications to spectroscopy to develop light sources that have a very narrow (less than natural) spectral line ranging their carry frequency from optics to gamma-rays where experimental demonstration of EIT has been recently reported \cite{14}. It is important to note here that the spectral width of the line and the coherent time of radiation is controlled by auxiliary external laser field. It is also worth to mention that the current technique can be applied to a single photon source (see for example \cite{10}) and allows one to transfer the flux of single photons with a particular or not defined coherent time to the flux of single photons with a given coherent time.

Our experimental setup is shown in Fig. 1. The laser is tuned to the transition $|a\rangle \rightarrow F=2 \rightarrow |c\rangle \rightarrow F=2$, see in Fig. 1, and the output laser beam is modulated by an electro-optic modulator (EOM) which is driven at the frequency of the ground level splitting (6.835 GHz, see Fig. 1). Two sidebands are generated, one with the frequency of the probe field and another with frequency down-shifted by 6.835 GHz with respect to the drive field. This down-shifted field is far from resonance and has a negligible effect on the experiment. The power and frequency of the sidebands is varied by changing the frequency and amplitude of the microwave field driving the EOM. If the EOM is driven by a spectrally narrow microwave source, then we observe a narrow EIT spectrum, just as for the case when two phase-locked lasers are used \cite{14}.

After the EOM, optical fields are sent through a single mode optical fiber to make a clean spatial intensity distribution with diameter 0.7 cm. The optical fields are circularly polarized by a quarter-wave plate. Then the drive and phase-noise broaden probe beams propagate through a cell containing atomic $^{87}\text{Rb}$ at temperature 67.7C and a buffer gas neon at pressure 30 Torr; the length of the cell is 2.5 cm; the diameter is 2.2 cm. The power at the entrance of the cell is 0.58 mW; after the cell power is 0.32 mW. We use heterodyne detection of the probe by mixing the transmitted light with an additional field that is frequency shifted by an acousto-optic modulator (AOM) at 60 MHz with respect to the drive field. This field does not propagate through the $^{87}\text{Rb}$ medium. This detection technique has been described in \cite{14}. A spectrum analyzer tuned in the vicinity of the beat-note frequency of the probe and additional fields is used to record the spectrum of beat-note which coincides with the spectral density of the phase-noise broaden transmitted probe field for the case of monochromatic drive field.

In the current experiment, the EOM is modulated with a broad (“white”) noise spectrum centered about the selected probe frequency. The total power in each side-band is about 10% of the power in the drive field, and the spectral width of the beat signal between the probe and drive is $\sim$ 1 MHz. We can characterize
FIG. 1: (a) Experimental setup. (b) Atomic level structure for $^{87}$Rb.

FIG. 2: Spectral density of the phase-noise broadened probe laser after passing the $^{87}$Rb cell. (a) shows the spectrum with the lasers far from resonance, so that the interaction with the atoms is negligible. (b) shows the spectrum with the lasers on resonance. The solid line in plot (a) is a Gaussian fit.

FIG. 3: Comparison of the EIT resonance obtained two ways. Data points marked with squares are transmission versus probe laser frequency for spectrally narrow drive and probe fields. Data points marked with stars are the spectrum of the transmitted probe light when the probe is phase-broadened.

the modulation of the EOM as a time dependent frequency: $f(t) = f_0 + \Delta f(t)$ where $f_0$ is the carrier frequency (6.835 GHz), and $\Delta f(t)$ is the noise driven time dependent frequency shift. Since the phase of such an oscillation is equal to $\varphi(t) = \varphi_0 + f_0 t + \int_{-\infty}^{t} \Delta f(t) dt$, we have a phase-noise broadened microwave signal that drives the EOM. Thus we have a phase-noise broadened probe field instead of a monochromatic probe field which is observed without the phase noise modulation. This is equivalent to using two lasers, one of which is spectrally narrow, and the other of which is has a phase-noise broadened spectrum.

We note that the spectral density of the probe signal does not coincide exactly with the spectrum of driving microwave signal because of the limited bandwidth of the EOM response to the modulation frequency. A typical plot of spectral density is depicted in Fig. 2. Before the cell, the spectral FWHM (full width at half maximum) of the probe field is 980 kHz, whereas after the cell we see significant narrowing of the spectrum with a FWHM of 4.6 kHz (see also Fig. 3). Before the cell, we have a Gaussian distribution of the spectral density

$$f_{in}(\omega) = e^{-\frac{(\omega - \omega_0)^2}{\gamma_n^2}}$$

where $\omega_0$ is the average probe field frequency and $\gamma_n$ is the width of the spectral density spectrum. However, after the cell, the probe field has a Lorentzian distribution

$$f_{out} = \frac{\gamma_n^2}{(\omega - \omega_0)^2 + \gamma_n^2}$$

where $\gamma_n$ is the width of the transmitted spectral density spectrum.

In Fig. 3 we compare transmission spectra taken under two experimental conditions: First, we measure the usual EIT resonance with coherent drive and probe fields by scanning the two-photon detuning between the two coherent fields. Second, we use the phase-noise broadened probe and measure the spectral density of the transmit-
be written as depicted in Fig. 1. We let levels $a$ and $b$ correspond to $^{87}$Rb atomic levels $5P_{1/2}F = 2$, $5S_{1/2}F = 1$, and $5S_{1/2}F = 2$, respectively. The driving field couples levels $a$ and $c$ and the probe field couples levels $a$ and $b$. The interaction picture Hamiltonian for the system can be written as

$$\hat{H} = \hbar \Omega_{ab} e^{i\Delta t}|a\rangle\langle c| + \hbar \Omega_{p} e^{i\Delta t}|a\rangle\langle b| + \text{h.c.}$$  \hspace{1cm} (3)$$

where $\Delta_p = \omega_{ab} - \nu_p$ and $\Delta_d = \omega_{ac} - \nu_d$ are the detunings of probe and drive fields from the atomic transitions $a-b$ and $a-c$, respectively, $\nu_p$ and $\nu_d$ are the frequencies of the probe and drive fields, $\Omega_d = |\psi_{ac}\rangle\langle \psi_{p}\rangle/\hbar$ and $\Omega_p = |\psi_{ab}\rangle\langle \psi_{p}\rangle/\hbar$ are the Rabi frequencies, and $\psi_{ab}$ and $\psi_{ac}$ are the dipole moments of the $a-b$ and $a-c$ transitions.

We are interested in finding the spectrum of the beat signal $S(t, z) = \Omega_{p}(t, z)\Omega_{d}^{*}(t, z)$ that is governed by the propagation equation along $z$ axis ($0 < z < L = 2.5 \text{ cm}$)

$$\frac{\partial S}{\partial z} = -i\eta \rho_{ab} \Omega_{d}^{*} + i\eta \rho_{ca} \Omega_{p},$$ \hspace{1cm} (4)$$

where $\rho_{ab}$ and $\rho_{ca}$ are the density matrix elements; $\eta = 3\lambda^{2}N\gamma_{r}/8\pi; N$ is the atomic density; $\gamma_{r}$ is the radiative decay from level $a$ to level $b$; $\lambda = 2\pi e/\omega_{ab}$.

The adiabatic approximation can be used to find the coherences $\rho_{ab}$ and $\rho_{ca}$, because the bandwidth of the phase-noise is less than relaxation rate of these coherences. Under these conditions, the propagation of the correlation function for the beat signal, $R(t, z) = \langle S(t, z)S(t + \tau, z) \rangle_{\tau}$, and time dependence of the correlation of beat signal and spin coherence $G(t, z) = \langle S(t, z)\rho_{cb}(t + \tau, z) \rangle$ whose space and time evolution is given by

$$\frac{\partial}{\partial z} R(t, z) = 2\eta \left( \frac{n_{ab}}{\Gamma_{ab}} - \frac{n_{ca}}{\Gamma_{ca}} \right) R(t, z) - 2\eta \left( \frac{\Omega_{d}^{2}}{\Gamma_{ab}} + \frac{\Omega_{p}^{2}}{\Gamma_{ca}} \right) G(t, z),$$ \hspace{1cm} (5)$$

$$\frac{\partial}{\partial t} G(t, z) = \left( \frac{n_{ab}}{\Gamma_{ab}} - \frac{n_{ca}}{\Gamma_{ca}} \right) R(t, z) - \left( \frac{\Omega_{d}^{2}}{\Gamma_{ab}} + \frac{\Omega_{p}^{2}}{\Gamma_{ca}} \right) G(t, z).$$ \hspace{1cm} (6)$$

Here $\Gamma_{cb} = \gamma_{ab} + i\Delta_p$ and $\Gamma_{ca} = \gamma_{ac} - i\Delta_{ac}$ where $\gamma_{ab}$ and $\gamma_{ac}$ are the relaxation rates of atomic coherences $\rho_{ab}$ and $\rho_{ac}$ respectively. Also $n_{ab} = \rho_{ab} - \rho_{bb}$ and $n_{ca} = \rho_{ca} - \rho_{cc}$ are the population differences. We have assumed a slow variation of atomic populations such that they do not change appreciably during propagation. A detailed analysis justifying this approximation will be presented elsewhere. It is clear that under these approximations, the correlation function of the beat signal (and hence the corresponding spectral density) is independent of the phase fluctuations of the the drive and probe fields.

Let us note the important role of atomic coherence in the discussed processes. Before we proceed to solving a set of coupled Eqs. (5,6), note that the atomic coherence plays an important role in these processes, once induced it gives rise a term $G(t, z)$. In the adiabatic limit, one can obtain transparency behavior governed by

$$\frac{\partial}{\partial z} R(t, z) = 2\eta N \Gamma_{cb} R(t, z),$$ \hspace{1cm} (7)$$
if condition $\vert \Omega_d \vert^2 \gg \Gamma_{ab}\Gamma_{cb}$ is met. Here we introduce 
\[ \Gamma_{cb} = \Gamma_{cb} + \Omega_{cb}^2 \overline{\Gamma_{cb}} + \Omega_{cb}^2 \overline{\Gamma_{cb}} \text{ and } \mathcal{N} = \frac{\overline{\Gamma_{cb}}}{\overline{\Gamma_{cb}}} - \frac{1}{\overline{\Gamma_{cb}}}. \]

In order to determine the spectral density of the beat signal $I_\omega$, we recall the definitions

\[ R(\tau, z) = \int I_\omega(z) e^{-i\omega \tau} d\omega, \quad G(\tau, z) = \int \rho_\omega(z) e^{-i\omega \tau} d\omega. \]

Substituting these expressions for $R(\tau, z)$ and $G(\tau, z)$ into Eqs. (5) and (6) we find

\[ \rho_\omega = \mathcal{N} I_\omega \quad \text{and} \quad \frac{\partial I_\omega}{\partial z} = 2\eta \left( \frac{\Gamma_{cb} - i\omega}{\Gamma_{cb} - i\omega} \right) I_\omega, \quad (8) \]

For the simplest case, with a weak probe and a strong drive, $\vert \Omega_d \vert \gg \vert \Omega_p \vert$ and all population remains in state $b$, so that $n_b = 1, n_a = n_c = 0$. Taking Doppler broadening into account by integration over velocity distribution leads to changing the homogeneous width $\gamma (2 \cdot 10^7 s^{-1})$ to the Doppler width $\Delta W (2 \pi 500 MHz)$. Also, there is narrowing because the medium is optically thick [10]. The spectral density of the beat signal is then given by

\[ I_\omega(z) = I_\omega(0) \exp \left[ \frac{-\eta \omega^2}{\Delta W \left( \left( \frac{\Omega^2}{\Delta W} \right)^2 + \omega^2 \right)} \right], \quad (9) \]

with the following expression for the spectral width of beating signal:

\[ \Delta \omega_{bs} = \frac{\vert \Omega \vert^2}{\Delta W \sqrt{\frac{\eta \omega}{\Delta W} - 1}}, \quad (10) \]

These expressions agree very well with our experimental results, namely, it gives a linear dependence of the spectral width on the driving power, and the same slope for the experimental parameters and a density of atoms of the order of $3 \times 10^{11} \text{ cm}^{-3}$ which corresponds to the cell temperature. Note, for the case of monochromatic drive and broadened probe field, the spectral width of the beating signal coincides with the spectral width of the EIT for a monochromatic probe field. The results obtained here can be interpreted in some sense that the probe field passing the cell and the atomic coherence in the cell are strongly correlated (see Eq.(8)) [17]. It is also worth to mention that there are interesting aspects of the problem was studied recently in where it has been demonstrated that, when probe and drive intensities are the same order of magnitude, the noise of probe and drive fields are strongly correlated [11], and phase noise of one lase is transfered to the second laser [12]. Also we would like to underline that the results obtained here go beyond classical EIT treatment [15] where the broadening of EIT resonance was predicted.

In conclusion, we have experimentially observed that the width of the beat signal between a coherent drive and broad-band probe field is greatly reduced (more than 1000 times) by propagation through a cell containing $^{87}\text{Rb}$ vapor. We analyzed the modification of a broad emission spectral line after transmission through a coherentlly prepared resonant medium. The final spectral lineshape is defined by the spectral shape of EIT resonance. The applications of the obtained results can be light sources (including a single photon sources) [10] with controllable coherent time of radiation, and notch filters with naturally narrow transmission band that are in great demand for many practical applications to background suppression in imaging including astrophysics and environmental imaging [13].

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