Width of the electromagnetically induced transparency resonance in atomic vapor

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The width of the electromagnetically induced transparency resonance is studied in rubidium vapor. Nonlinear dependence of the width on drive intensity that is caused by Doppler broadening and optical pumping is found. Density-matrix analysis supports the observed dependence. The discrepancies between experimental data and theoretical prediction are discussed.

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Atomic coherence and interference have changed the face of the many well-studied nonlinear processes, such as parametric- and second-harmonic generations as well as four-wave mixing [1]. The efficiency of these processes is very high and the input power is extremely low (it can reach a level of a few photons) [2]. The electromagnetically induced transparency (EIT) plays a key role in these processes [3,4]. Gain and lasing without inversion in a medium with EIT are achievable too [5]. A new flash of interest in EIT emerges since the proposal to use EIT for the storage of quantum information [6]. The application of EIT in the new generation of compact, small size, and totally optical atomic clock also attracts significant interest [7].

The EIT resonance is, practically, a unique product of atomic coherence that may be easily measured experimentally. The study of the EIT width (Γ) may lead to a better understanding of the influence of the different processes on atomic coherence. Among them are the relaxation and collisional rates, optical pumping, multicoherences [8], spontaneous transfer of coherence [9], radiation trapping [10], and especially, Doppler broadening.

Experimentally the EIT width has been studied in several papers. Linear [11–14] and nonlinear [15,16] behaviors of EIT width on laser intensity have been found. Theoretically the EIT has been studied in Refs. [3,12,13,17,18]. Linear dependence on light intensity is predicted in Refs. [3,12,13], whereas in Refs. [17,18]—nonlinear one.

The goals of our study are (i) to verify the validity of the Refs. [17,18] approach, in which nonlinear dependence of EIT width due to optical pumping and Doppler broadening has been obtained, and (ii) to explain why in previous papers mentioned above both linear and nonlinear EIT width dependencies have been observed.

In this paper, we present the study of EIT width in Doppler-broadened rubidium vapor in a wide range of intensities of the drive field. The EIT width reveals its nonlinear behavior on drive intensity. We find the qualitative agreement between the experimental data and theory [17,18], meanwhile quantitative difference is quite significant. The role played by the optical pumping and Doppler broadening in the shaping EIT resonance is discussed. Numerical calculations, fulfilled by density-matrix equation technique, give clear illustrations of the EIT-width dependence on Doppler broadening and the rate of optical pumping.

We perform our experiment with the ^{87}Rb atomic vapor contained in a 3-cm long cylindrical (diameter 2.5 cm) pyrex cell. The extended cavity diode lasers are tuned to the D_2 line and form Λ interaction scheme. The probe (with Rabi frequency α) and drive (with Rabi frequency Ω) laser fields excite $5S_{1/2}(F=1)\rightarrow 5P_{3/2}(F'=0,1,2)$ and $5S_{1/2}(F=2)\rightarrow 5P_{3/2}(F'=1,2,3)$ transitions, respectively, as shown in Fig. 1. To be accurate, we choose the drive frequency around $5S_{1/2}(F=2)\rightarrow 5P_{3/2}(F'=1)$ transition at which the EIT resonance exhibits the symmetrical shape at low intensities $\sim$0.1 W/cm^2. We combine the orthogonally polarized beams before cell by polarization beamsplitters. To separate the beams after the cell the identical beamsplitter is used. The beams are overlapping in the cell and the angle between them is $\varphi \sim 0.6\times 10^{-2}$ rad. The cell is covered by three layers μ-metal shield, that suppresses the laboratory magnetic field to the value of $\sim$0.1 mG.

The measured absorption of the probe field is about 30% at 40 °C cell temperature in the absence of the drive. The application of the drive field results in the change of the absorption spectrum of probe field (Fig. 2) in the following: (i) The probe absorption [curve (b)] is increased by two times compared with the absorption indicated by curve (a). This results from the optical pumping, by which most of population is pumped into $F=1$ from $F=2$ hyperfine levels of the state $5S_{1/2}$ via the upper state $5S_{3/2}$. (ii) EIT resonance, which occurs at the two-photon resonance, enhances the transparency. This transparency is not complete [only about 20%, see curve (b)] and cannot prevail the absorption from optical pumping. The reason for this nonsuppressed

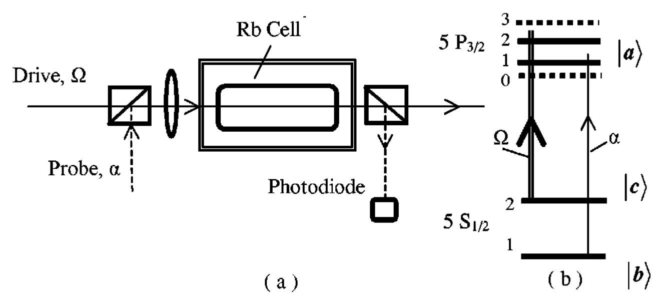

FIG. 1. Experimental setup (a) and energy-level diagram of ^{87}Rb (b). Double and single lines stand for the drive and probe laser fields, respectively. Dashed lines depict uncoupled states in the energy diagram.

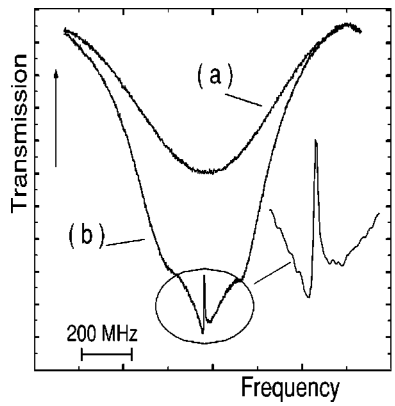

FIG. 2. Transmission spectra of the probe field without (a) and with drive (b). The inset shows the EIT feature of the spectrum.

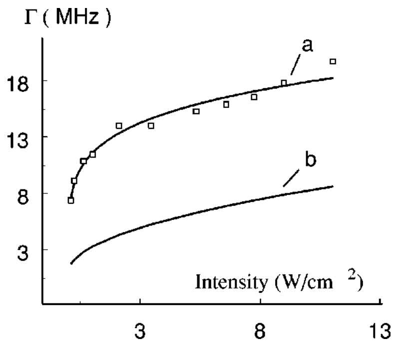

FIG. 3. Experimental (a) and analytical (b) dependence of EIT width on the drive-field intensities. The solid curve (a) is drawn as a guide to the eye. Curve (b) is plotted according to Eq. (1).

absorption is that the drive and probe fields cannot couple the ground states via the hyperfine levels $F'=0$ and 3 of $5P_{3/2}$ state due to the selection rules. These states are marked by the dashed lines in Fig. 1(b).

To measure the EIT width, we modulate the injection current of the probe laser (similar to Ref. [15]). The spectrum of the probe laser has two side bands under modulation, that are used as reference marks. The modulation frequency lies in the range of 1–50 MHz.

There are two types of errors of the measurement of the EIT width: asymmetric shape of EIT resonance as shown in Fig. 2 and poor signal-to-noise ratio. The first error dominates mostly at high intensities when Rabi frequency becomes comparable with Doppler width. The distortion of the shape in this case results from Autler-Townes doublet asymmetry originating from the detuned hyperfine components $F'=1$ or 2 of $5P_{3/2}$ state [19] and their different light shifts. At low intensities, the EIT resonance has small amplitude that gives low signal-to-noise ratio. As a result, the measurement error of the EIT width is about ~25% in our experiment.

The drive laser with power less than 10 mW is focused into a 0.3 mm diameter spot size. The transit broadening is $\sim 0.98u/\pi D = 0.3\,\mathrm{MHz}$ [20], where u is thermal speed of atoms and D the diameter of the laser beam. The power of the probe light is 10 μW.

The experimental dependence of the EIT width on the intensity of the drive field is shown in Fig. 3 [curve (a)]. Similar dependence of the EIT width on drive intensities in Doppler-broadened vapor has been predicteded theoretically in Refs. [17,18]. To compare our results with theory, we use formula from Ref. [17].

According to the Ref. [17], the width of the EIT resonance should follow the square-root dependence on the drive intensity.

$$\Gamma = \Omega \sqrt{2\gamma_{bc}/\gamma}, \tag{1}$$

when $\gamma_{bc}\gamma \ll \Omega^2 \ll \Delta_D^2(2\gamma_{bc}/\gamma)$. At higher intensities, this dependence implies linear relationship with drive intensity.

$$\Gamma = \Omega^2/\Delta_D, \tag{2}$$

where γ_{bc} is the relaxation rate. γ—the natural width of the P—state (~6 MHz), and Δ_D is a Doppler width (~540 MHz). The Rubi frequency for Rb atoms is estimated from formula $\Omega = 2\pi\gamma\sqrt{I/8}$, where I is the drive intensity (mW/cm^2) [21],

Equation (2) has simple physical content in term of "dressed-state" atom presentation [22] (see Fig. 4). The splitting between dressed-state levels $|+\rangle$ and $|-\rangle$ equals 2Ω when the probe field is tuned to the center of the Doppler resonance. The scale of EIT width in the case of immobile atoms $v_z=0$ is defined by the Rabi splitting 2Ω. The moving atoms see the coupling field at shifted frequency $\omega' = \omega_{ab} \pm |\vec{\kappa}\cdot\vec{v}|$ and the responding splitting increases $2\Omega' = \sqrt{\Omega^2 + |\vec{\kappa}\cdot\vec{v}|^2}$, where $\vec{\kappa}$ and $\vec{v}$ are the wave and velocity vectors, respectively. One may see that for atoms moving in the same direction as the propagation of the light, the "dressed" state $|-\rangle$ is attracted to center frequency ω_{ab} resulting in the narrowing of the EIT resonance [Fig. 4(b)]. The same is true for atoms moving in opposite to light propagation [Fig. 4(c)]. So the EIT width becomes $2\Gamma' \leqslant 2(\sqrt{\Omega^2 + |\vec{\kappa}\cdot\vec{v}|^2} - |\vec{\kappa}\cdot\vec{v}|)$. In the case where $|\vec{\kappa}\cdot\vec{v}| \gg \Omega$, the width of the EIT resonance is

$$\Gamma \simeq \Omega^2/\kappa v. \tag{3}$$

Replacing $|\vec{\kappa}\cdot\vec{v}|$ by Δ_D, we obtain relation between the EIT width and Doppler broadening $\Gamma \simeq \Omega^2/\Delta_D$. To demonstrate this relation, we perform numerical calculation (see the Appendix). The EIT width for atoms with different Doppler broadening is shown in Fig. 5. The inset shows that the broader Doppler width leads to the smaller EIT width.

Now let us pay attention to the role played by optical pumping in the shaping of EIT resonance. The EIT occurs

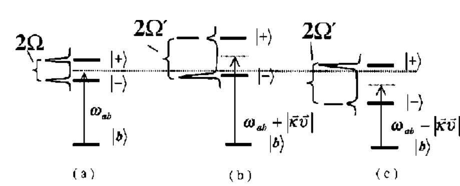

FIG. 4. Dressed-state representation of the coherently driven atom: (a) for the stationary atom; (b) and (c) for the moving one at velocities $\pm v$, respectively.

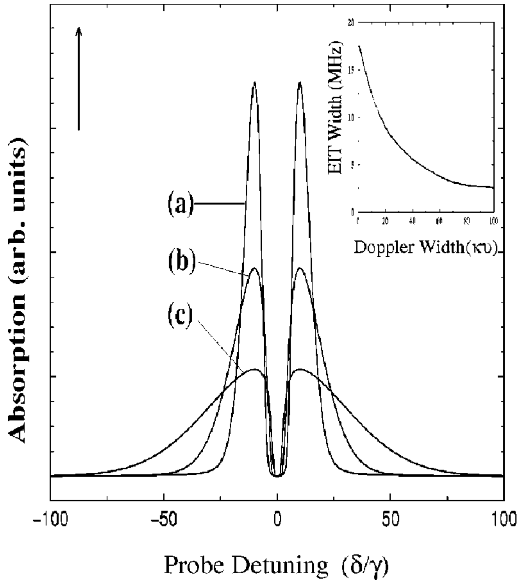

FIG. 5. Numerically calculated EIT line shape vs probe detuning at various Doppler broadenings, $\Omega=10\gamma$, $\alpha=0.01\gamma$, $\gamma_{bc}=0.001\gamma$, and $\Delta=0$. (a) $\Delta_D=5\gamma$, (b) $\Delta_D=10\gamma$, and (c) $\Delta_D=15\gamma$. The inset presents the EIT width vs Doppler broadening.

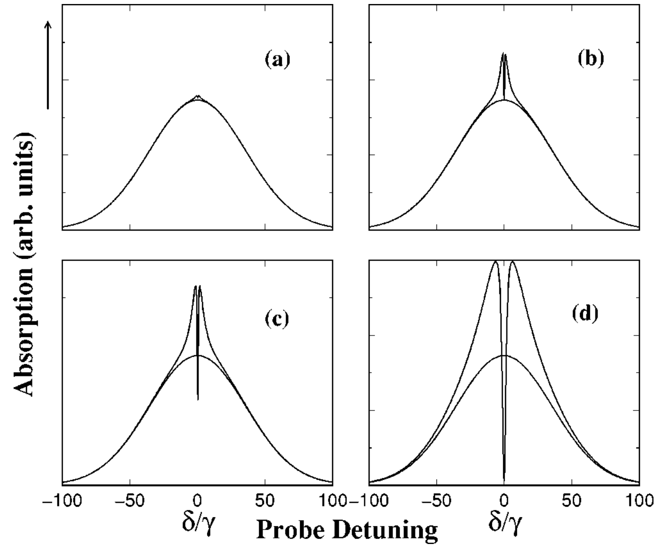

FIG. 6. Numerically calculated EIT spectra vs Doppler shift at various drive-field intensities $\alpha=0.01\gamma$, $\gamma_{bc}=0.1\gamma$, and $\Delta=0$. (a) $\Omega=0.1\gamma$, (b) $\Omega=0.5\gamma$, (c) $\Omega=\gamma$, and (d) $\Omega=5\gamma$. The identical spectra from (a)–(d) are Doppler-broadened absorption lines with no drive field.

for certain intensity, at which the optical pumping exceeds the relaxation rate in the coupled states $\Omega^2/\gamma \geqslant \gamma_{bc}$ [23]. Physically, it means that the medium begins to demonstrate nonlinear response. Optical pumping rate for moving atoms is less than that of immobile atoms and equals $\Omega^2(\gamma/(\gamma^2+|\vec{\kappa}\cdot\vec{v}|^2))$, that means not all the atoms (only with velocities $|\vec{\kappa}\cdot\vec{v}|\approx\Omega\sqrt{\gamma/\gamma_{bc}}$) of hot gas are coherently prepared. Substituting $|\vec{\kappa}\cdot\vec{v}|$ into the formula (3), we find that Γ obeys the square-root dependence on intensities $\Gamma=\Omega\sqrt{\gamma_{bc}/\gamma}$. This relation holds up to the intensity at which all moving atoms will be pumped into the state $|b\rangle$. After the moment $\Omega^2 \geqslant \Delta_D^2(2\gamma_{bc}/\gamma)$, EIT width has linear dependence as indicated in Eq. (2), that means the moving atoms do not greatly affect EIT width.

It is not easy to get the exact solution in such simple approach. Nevertheless, the impact of the moving atoms involved in optical pumping on the EIT are clarified. The probe-absorption spectra vs different drive intensities obtained by numerical calculation are shown in Fig. 6, from which one can see that only some part of the atoms participate in EIT resonance (that is developed in domelike resonance on the Doppler curve).

The experimentally observed EIT width is 3–10 times larger than the expected from the Eq. (1) (see, Fig. 3). At the same time the EIT width manifests nonlinear dependence that resembles the predicted square-root function. The comparison of the EIT width with the results of Ref. [18], also indicates the significant difference. Contrary to the results of the Ref. [15], we found that the divergence and parallelism of the beams might contribute to some broadening in EIT. This broadening equals $\Delta_D\varphi\approx 3$ MHz in experiment [24]. To test this, we perform the experiment on the magneto-optic rotation [25]. We create the coherence between ground-state sublevels of ^{87}Rb by parallel beam of linearly polarized drive light (probe laser is switched off and beam splitters are removed, beam diameter equals 3.5 mm) and detect the transmission when B_z magnetic field is swept. The EIT-width dependence is shown in Fig. 7. The conditions of this experiment obviously rule out any geometrical disturbance on the EIT width. The EIT dependence, similar to the previous case, exhibits nonlinear behavior, but with width is less than the theoretical prediction [see Fig. 7(b)]. The possible explanation of this may be related to the Ramsey effect, which narrows the EIT resonance due to diffusion motion [26]. However, Ramsey narrowing is effective at the presence of the buffer gas but in vacuum cell its influence is not so obvious.

In addition to the geometrical factors, there are some other reasons for disturbing the EIT-width dependence. Taichenachev, Tumaikin, and Yudin [18] treat the atoms with infinite lifetime of coherence γ_{bc}^{-1}. In our experiment, coherence lifetime is defined by the transit time of atom passing through the laser beam. This time changes from 3 μs in the

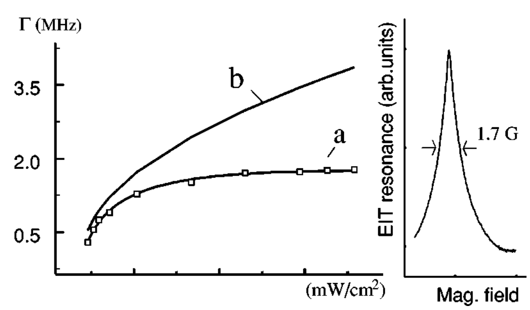

FIG. 7. Experimental (a) and theoretical (b) dependence of the EIT width on the intensities of the drive field in the case of nonlinear magneto-optic rotation. Curve (a) is drawn as a guide to the eye. Curve (b) is calculated according to Eq. (1).

middle of cell to 6 μs at the ends of the cell due to the variation of the diameter of the laser beams.

The effects of laser propagation and the variation of the drive laser intensity along the cell on the EIT width are not taken into account in Refs. [17,18]. They treat the medium with extremely small absorption in the center of the EIT resonance, whereas in experiment it is actually as high as ~45%. Our study shows that at 1 W/cm^2 drive intensity almost 80% of light reaches the end of cell. However, at low intensities, only 20–60 % of power is transmitted.

In the cited papers, very simple three-level atom are considered, though Rb atom has very rich structure with a lot of Zeeman sublevels. These sublevels may accumulate a lot of atoms due to optical pumping and induce some additional dark states because of Zeeman multicoherences [8].

In summary, we present an experimental dependence of the width of EIT resonance on the drive-field intensity. It is pointed out the nonlinear behavior of this dependence caused by Doppler broadening and optical pumping. Fulfilled numerical simulation gives adequate illustrations of their role in shaping the EIT resonance. Qualitative agreement with the theoretical predictions of [17,18] is observed. Our results may explain why in Refs. [11–13] are observed linear dependence of EIT width whereas in Refs. [15,16]—the nonlinear one. Namely, the buffer gas used in experiments [11–13] increases the homogeneous broadening (γ) up to 50–150 MHz. In this case Doppler narrowing and optical pumping do not disturb EIT width significantly. Besides this, studied intensity range in experiments is rather small $\Omega \leqslant \gamma$.

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APPENDIX: NUMERICAL CALCULATIONS

The density-matrix equations of motion for the Λ system shown in Fig. 8 in the rotating frame are given by

$$\dot{\rho}_{aa} = -(\gamma_b + \gamma_a + \gamma_0 + r)\rho_{aa} + r\rho_{bb} - i\Omega(\rho_{ac} - \rho_{ca}) - i\alpha(\rho_{ab} - \rho_{ba}), \tag{A1a}$$

$$\dot{\rho}_{bb} = r_b + \gamma_a\rho_{aa} - (\gamma_0 + r)\rho_{bb} + r\rho_{aa} + i\alpha\rho_{ab}, \tag{A1b}$$

$$\dot{\rho}_{cc} = r_c + \gamma_a\rho_{aa} - \gamma_0\rho_{cc} + i\alpha(\rho_{ac} - \rho_{ca}), \tag{A1c}$$

$$\dot{\rho}_{ab} = -\Gamma_{ab}\rho_{ab} - i\alpha(\dot{\rho}_{aa} - \rho_{bb}) + i\Omega\rho_{cb}, \tag{A1d}$$

$$\dot{\rho}_{cb} = -\Gamma_{cb}\rho_{cb} - i\alpha\rho_{ca} + i\Omega\rho_{ab}, \tag{A1e}$$

$$\dot{\rho}_{ca} = -\Gamma_{ca}\rho_{ca} + i\Omega(\rho_{aa} - \rho_{cc}) - i\alpha\rho_{cb}, \tag{A1f}$$

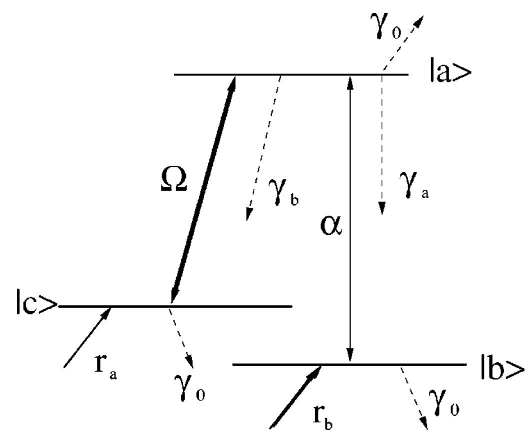

FIG. 8. Schematic of energy level. γ_a and γ_b are the population relaxation rates from state $|a\rangle$ to $|b\rangle$ and $|c\rangle$ states; α, Ω are the Rabi frequencies of probe and drive fields, respectively; γ_0 is the decay rate out of states $|b\rangle$ and $|c\rangle$: r_b and r_c are the pumping rate into respective states; r is the incoherent pump rate between $a \leftrightarrow b$.

where γ_a and γ_b are the population relaxation rates from state $|a\rangle$ to $|b\rangle$ and to $|c\rangle$ states, respectively; α and Ω are the Rabi frequencies of probe and drive fields respectively; γ_0 is the decay rate out of states $|a\rangle$, $|b\rangle$, and $|c\rangle$; r_b and r_c are the pumping rate into respective states $|b\rangle$ and $|c\rangle$; r is the incoherent pump rate between $a \leftrightarrow b$. Off-diagonal decay rates are $\Gamma_{ij} = \gamma_{ij} + i\Delta_{ij}$, γ_{ij} are the decay rates of level i and j, and Δ_{ij} are the detunings from atomic resonances for laser fields. We find absorption coefficient of the system for probe laser by solving the equations for ρ_{ab} in steady state.

In order to include the Doppler-broadening effect, we need to express the ρ_{ab} as a function of Doppler shift kv and integrate over the Doppler-velocity distribution $f(kv)$,

$$\chi = \int_{-\infty}^{\infty} d(kv)\, \eta \frac{\rho_{ab}}{\alpha} f(kv), \tag{A2}$$

where η is a constant, $f(kv)$ is the atomic Maxwellian distribution

$$f(kv) = \frac{1}{ku\sqrt{\pi}} e^{-(kv)^2/(ku)^2}, \tag{A3}$$

where $u = \sqrt{2k_B T/M}$ is the most probable speed of atom at given temperature T and atomic mass M, k_B is the Boltzmann constant, therefore, the full width at half maximum of the Doppler-broadened linewidth is given by $2\Delta_D = 2\sqrt{\ln 2}ku$. We numerically integrate Eq. (A2) and obtain the velocity-averaged absorption coefficient as a function of probe detuning.

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