Decoherence of electromagnetically induced transparency in atomic vapor

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We report characterization of electromagnetically induced transparency (EIT) resonances in the $D1$ line of ^{87}Rb under various experimental conditions. The dependence of the EIT linewidth on the power of the pump field was investigated at various temperatures for the ground states of the lambda system associated with different hyperfine levels of the atomic $5S_{1/2}$ state as well as magnetic sublevels of the same hyperfine level. Strictly linear behavior was observed in all cases. A theoretical analysis of our results shows that dephasing in the ground state is the main source of decoherence, with population exchange playing a minor role.

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Since the discovery of electromagnetically induced transparency (EIT) in the early 1990s,[1] the phenomenon has received significant attention due to its vast range of applications, in particular in quantum computation and quantum communication.[2] They include storage of quantum states of light,[3–5] logic gates, magnetometry, and routing of optical information.[6] Many of these applications require thorough understanding of the phenomena responsible for the width of EIT resonances, in particular of the relevant decoherence processes.

Decoherence in EIT is caused by several mechanisms, such as flight-through broadening, population exchange, and atom–atom and atom–wall collisions, but it is still not clear which is the most significant. Insight into this question can be gained by measuring the width of the EIT resonance as a function of the pump field intensity. Most existing experiments in atomic vapors[7–10] showed this dependence to be linear, with the exception of the work by Ye and Zibrov[11] (which was performed in unusual conditions, without buffer gas and with a very small beam diameter). On the other hand, an existing theoretical treatment of EIT in Doppler-broadened gases,[12,13] assuming the population exchange between the ground levels $|b\rangle$ and $|c\rangle$ [Fig. 1(a)] to be the main source of decoherence, predicts a nonlinear dependence for weak pump powers.

In the present work, we measured the width of the EIT resonances on the $D1$ transition in rubidium vapor in a variety of settings. Our results show a consistent linear behavior with y-axis intercepts of the order of a few kilohertz. Based on these findings, we propose an alternative theory relying upon pure dephasing (i.e., decay of the off-diagonal density matrix elements) as the dominant decoherence mechanism. This treatment does predict linear dependencies, yielding much better fits to our data. From the fits, we also obtain the ground states' decoherence rates on the scale of a few kHz, which are consistent with independent verifications.

The experiments were performed in atomic ^{87}Rb vapor at temperatures 60–100 °C, which correspond to the optically thick regime, using two Λ energy level configurations. In the first configuration (hereafter referred to as Zeeman), the pump and signal fields of wavelength $\lambda = 795$ nm couple pairs of Zeeman sublevels of the atomic ground state ($5S_{1/2}$, $F=2$) via the excited state ($5P_{1/2}$, $F'=2$). The light source is a Coherent MBR-110 Ti:saphire laser with a narrow spectral width (~100 kHz) and high long-term stability. The relative frequencies of both fields were precisely controlled by acousto-optic modulators. After modulation, both fields were recombined in a polarizing beam splitter, converted to orthogonal, circular polarizations by means of a quarter-wave plate, and directed into a 5 cm long rubidium vapor cell with 1 Torr Ne buffer gas, located inside a magnetically shielded oven. The measured diameter of the pump beam just before the cell was ~10 mm. After the cell, the linear polarization was recovered and the signal field detected.

In the second configuration (referred to as hyperfine), the signal field was obtained from an additional diode laser that was phase locked at 6.8 GHz to the Ti:sapphire laser to ensure a two-photon resonance

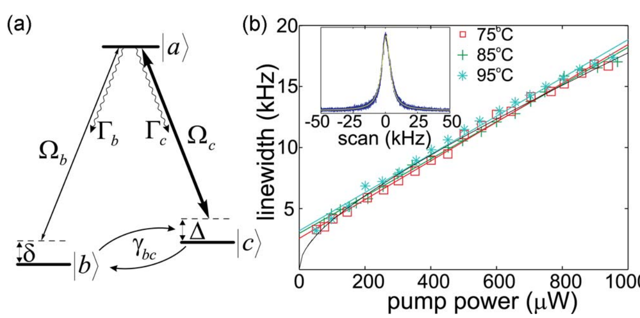

Fig. 1. (Color online) (a) Relevant atomic level structure. The excited state $|a\rangle$ is coupled to two ground states $|b\rangle$, $|c\rangle$ by a weak signal field with Rabi frequency Ω_b and detuning δ and a stronger pump field with Rabi frequency Ω_c and detuning Δ. Γ_b and Γ_c are rates of spontaneous emission into the respective states $|b\rangle$ and $|c\rangle$, whereas γ_{bc} describes the decay rate of coherences between these states. (b) Measured width of the Zeeman configuration EIT resonances as a function of the pump laser power for different temperatures together with linear fits and a fit to the theory of Javan and co-workers,[12,13] assuming a decoherence rate of $\gamma_{bc} = 110$ Hz. The inset shows an example of the measured EIT resonance.

with the ($|b\rangle = |5S_{1/2}, F=1\rangle$, $|a\rangle = |5P_{1/2}, F=2\rangle$, $|c\rangle = |5S_{1/2}, F=2\rangle$) transition. In this setting, we used linear polarizations for the pump and signal.

To measure the full width at half-maximum (FWHM) of the EIT resonance, we swept the signal field frequency. A typical scan is depicted in the inset of Fig. 1(b) and approximates a Lorentzian distribution. The pump power was varied from 100 μW to 1.2 mW ($\Omega_c \sim 1-3$ MHz), while the signal power was kept constant at about 20 μW ($\Omega_b \sim 500$ kHz). To compensate for the effect of transparency line narrowing with increasing optical thickness,[14] in calculating the FWHM the logarithm of the transmitted signal (which is proportional to the optical susceptibility) was used.

Figure 1(b) shows the results of the Zeeman measurement. The behavior is linear, temperature independent, and shows a y-intercept of 3 kHz. Similar measurements were performed in a 10 Torr Ne buffer gas cell (not depicted), showing a similar behavior and a slightly different y-intercept.

Experimental results for the hyperfine configuration are presented in Fig. 2. Similarly, a linear behavior is observed, but the slope and the zero crossing depend on temperature. Additional measurements were done with a 0.1 Torr cell (not depicted), showing a similar trend.

Both figures also display the best fit obtained using the theory of Javan and co-workers.[12,13] These fits do not follow well the experimentally observed data and yield unrealistically low values for the decoherence rates: 110 and 117 Hz, respectively. We conclude that the population exchange cannot be the dominant decoherence mechanism in our cells.

Motivated by this discrepancy, we developed an alternative theory assuming that the ground-state decoherence is dominated by dephasing, i.e., decay of the off-diagonal density matrix element ρ_{bc} at a rate γ_{bc}. With this assumption, and the notation defined in Fig. 1(a), we first determine the stationary density matrix of a single atom. Without the signal field, it is just $\rho^{(0)} = |b\rangle\langle b|$: all atoms are pumped into the $|b\rangle$ state. In the first (linear) order with respect to the signal, one obtains the well-known expression[15]

$$\rho_{ab}^{(1)} = \Omega_b \frac{1}{\Delta - \delta_2 + \dfrac{|\Omega_c|^2}{i\gamma_{bc} + \delta_2} - i\dfrac{\Gamma}{2}}, \tag{1}$$

where $\Gamma = \Gamma_b + \Gamma_c$ is the inverse lifetime of the excited state and $\delta_2 = \Delta - \delta$ is the two-photon detuning.

Atoms in a hot gas "see" the light fields Doppler shifted according to their momentary momentum. In our experiment, both fields are co-propagating, and because their optical frequencies are similar, they experience a similar Doppler shift: the two-photon detuning δ_2 is almost independent of the individual atoms' motion while the pump field detuning Δ is not.

Assuming that the pump field is tuned to the center of the Doppler-broadened line, averaging expression Eq. (1) over all atoms weighted by their velocity

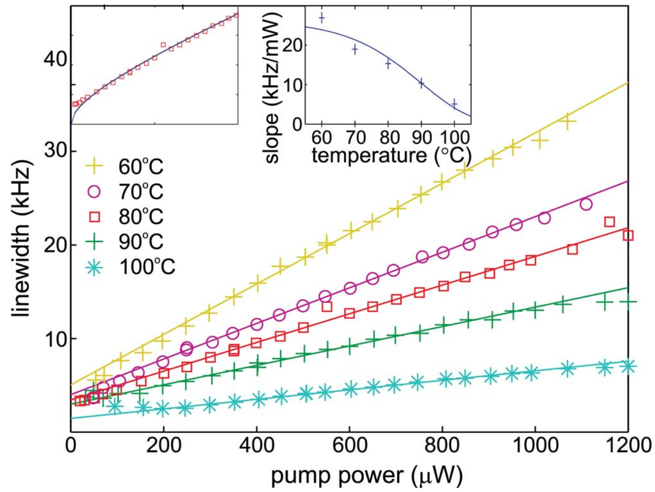

Fig. 2. (Color online) Dependence of the EIT linewidth on the pump power for the hyperfine configuration. The lines are linear fits to the measured data. The left inset shows a fit of the population exchange theory[12] to the data acquired at $T = 80°$C, with $\gamma_{bc} = 117$ Hz. The right inset shows the temperature dependence of the experimentally determined slopes together with a theoretical estimate.

distribution $p(\Delta) = p(2\pi v/\lambda)$ gives the susceptibility for the signal field:

$$\chi_b = \frac{\wp}{\Omega_b} \int p(\Delta)\rho_{ab} d\Delta, \quad \wp = \frac{N\mathcal{D}_{ab}^2}{\hbar\epsilon_0}, \tag{2}$$

where N is the atomic density, $\mathcal{D}_{ab}$ the dipole moment of the $|a\rangle \leftrightarrow |b\rangle$ transition, ϵ_0 is the free-space permittivity and

$$p(\Delta) = \frac{\sqrt{\ln 2}}{W_d\sqrt{\pi}} e^{-\ln 2(\Delta/W_d)^2}, \tag{3}$$

with $2W_d = (4\pi/\lambda)\sqrt{2 \ln 2 k_B T/m_{\mathrm{Rb}}} \approx 0.55$ GHz being the FWHM of the Doppler-broadened line.

To simplify our calculations, we approximate the Gaussian Boltzmann distribution by a Lorentzian of the same width and maximum:

$$p_{\mathrm{Lorentz}}(\Delta) = \frac{\sqrt{\ln 2}}{W_d\sqrt{\pi}} \frac{1}{1 + (\Delta/W_d)^2}. \tag{4}$$

Equation (4) approximates the Gaussian distribution reasonably well for small Δ; it differs only in its wings, which correspond to high-velocity atoms that see both fields far detuned and thus interact only weakly. Performing the integration (2) in this approximation, we find the average susceptibility for the Ω_b field:

$$\chi_b = 2\wp\sqrt{\pi \ln 2} \frac{i\gamma_{bc} + \delta_2}{(\gamma_{bc} - i\delta_2)(\Gamma + 2W_d - 2i\delta_2) + 2|\Omega_c|^2}. \tag{5}$$

The EIT linewidth is several orders of magnitude smaller than the Doppler width $2W_d$, hence we can drop the $2i\delta_2$ term in the denominator. The (intensity) absorption coefficient α is then of Lorentzian form:

$$\alpha(\delta_2) = \frac{\omega}{c}\,\mathrm{Im}(\chi_b) = \alpha_{\max} - \frac{\alpha_{\max} - \alpha_{\min}}{1 + \left(\dfrac{2\delta_2}{\mathrm{FWHM}}\right)^2}, \qquad (6)$$

with

$$\alpha_{\max} = 2\frac{\omega\wp}{c}\sqrt{\pi \ln 2}\,\frac{1}{2W_d + \Gamma}, \qquad (7)$$

$$\alpha_{\min} = 2\frac{\omega\wp}{c}\sqrt{\pi \ln 2}\,\frac{1}{2W_d + \Gamma + 2\dfrac{|\Omega_c|^2}{\gamma_{bc}}}, \qquad (8)$$

$$\mathrm{FWHM} = 2\gamma_{bc} + \frac{4|\Omega_c|^2}{2W_d + \Gamma}. \qquad (9)$$

Since $|\Omega_c|^2$ is proportional to the beam intensity, the EIT linewidth scales linearly with the pump power and intersects the y-axis at a minimum of $2\gamma_{bc}$.

As is evident from Figs. 1 and 2, linear dependence (9) provides an excellent fit to our experimental data. The y-axis intercept is twice the ground-state decoherence rate, which can be measured independently, e.g., by means of light storage.[3] With our setup, storage times of 100–250 μs were observed, which is in reasonable agreement with the measured intercepts.

We note that the population exchange theory[12,13] predicts that at high pump intensity the dependence of the EIT linewidth on the pump power approaches linear:

$$\mathrm{FWHM} \to 4\gamma_{\mathrm{pe}}\frac{W_d}{\Gamma} + \frac{2|\Omega_c|^2}{W_d}, \qquad (10)$$

γ_{pe} being the population exchange rate. This asymptotic dependence has a slope similar to that in Eq. (9), but its intercept, for $\gamma_{\mathrm{pe}} \sim \gamma_{bc}$, is higher by a factor of ~ 180. Therefore, even a minute population exchange rate would have a significant effect on the dependencies studied. Our experimental results thus show that the fraction of population exchange mechanism in the ground-state decoherence is indeed minor. We can estimate γ_{pe} to be below 50 Hz.

Decoherence in atomic vapor cells is known to be dominated by the flight-through mechanism. The atoms arrive into the interaction area in an arbitrary ground state, so one would expect at least some signature of population exchange decoherence, and it is surprising that none is present. We explain this by the geometry of the laser beams, which was close to Gaussian. Before entering the interaction area, an atom, initially in a random ground state, propagates through the "wings" of the Gaussian profile. In this region, the signal field is negligible, but the pump field is already sufficiently strong to pump the atom out of $|c\rangle$. When entering the central part of the interaction region, most of the atoms will be in state $|b\rangle$, albeit with a random phase.

The temperature dependence of the slope in the hyperfine configuration is due to the absorption of the pump field by the saturated rubidium vapor. To quantify this effect, we utilized the relation between the saturated rubidium vapor density and the temperature to calculate the expected pump absorption. We then integrated Eq. (5) over the length of the cell, assuming that the pump absorption is governed by Beer's law, and obtained the temperature dependence of the slope, which was then compared with the measured slopes (Fig. 2, right inset). The best fit was obtained for the 12 mm beam diameter, which agrees well with our measured diameter of 10 mm.

In the Zeeman configuration, the pump absorption was negligible at all temperatures due to optical pumping, which explains a constant slope. Also the zero crossing ($2\gamma_{bc}$) in this setting was temperature independent, which shows that the density-dependent Rb–Rb atomic collision rate is not a significant decoherence mechanism.[16] In the hyperfine configuration (Fig. 2), this behavior is not exactly reproduced in that the zero crossing of extrapolated linear fits *decreases* with temperature. This may be a consequence of a slight deviation of the experimental data from the ideal linear behavior predicted by Eq. (9) at high pump powers.

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