Investigation of dephasing rates in an interacting Rydberg gas

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Abstract. We experimentally and theoretically investigate the dephasing rates of the coherent evolution of a resonantly driven pseudo spin emersed in a reservoir of pseudo spins. The pseudo spin is realized by optically exciting $^{87}\text{Rb}$ atoms into a Rydberg state. Hence, the upper spin states are coupled via the strong van der Waals interaction. Two different experimental techniques to measure the dephasing rates are shown: the ‘rotary echo’ technique, known from nuclear magnetic resonance physics, and electromagnetically induced transparency. The experiments are performed in a dense frozen Rydberg gas, either confined in a magnetic trap or in an optical dipole trap. Additionally, a numerical simulation is used to analyse the dephasing in the rotary echo experiments.

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1. Introduction

A detailed understanding of dephasing processes in a quantum system is essential if one desires to obtain coherent control over quantum matter. Spin-$\frac{1}{2}$ systems imbedded into a strongly coupled environment are the paradigm for decoherence theory. Most noticeably, the spin-boson model has been studied in great depth [1, 2]. Ultracold atoms, that are coherently excited into a Rydberg state, can serve as a test bed for strongly interacting quantum matter. The coherent dynamics of an individual spin is coupled strongly to the bath of the other simultaneously driven atoms leading to a dephasing of the quantum mechanical evolution. In this paper, the dephasing mechanism is investigated experimentally by two complementary methods and compared with a numerical model calculation. This work is also related to possible quantum computational schemes such as frozen Rydberg gases, i.e. the centre of mass motion of the atoms is negligible on the timescale of the experiment, and can be used to build fast quantum gates [3, 4]. The latter article suggests using the strong interaction among Rydberg atoms, namely the dipole–dipole interaction, to create a mesoscopic collective state, which has the advantage that not every single atom must be controlled separately on a microscopic scale. The blockade of the excitation into the Rydberg state due to the dipole–dipole and van der Waals interaction has been observed in laser-cooled atomic clouds in various experiments [5]–[8] and in magnetically trapped atomic samples at much higher number densities [9].

The coherent excitation of interacting Rydberg systems is investigated in [9]–[12]. In order to successfully build a quantum gate operation, it is crucial to know the mechanisms and the timescales on which the dephasing of the states happen in the system. This knowledge can be incorporated into the technical realization of quantum gate operation in future experiments. The gate operation in a quantum system must take place on a timescale faster than the dephasing rate $\gamma_d$ in order to avoid a decoherence of the system. In this paper, the dephasing rates are investigated by two complementary methods, namely using the rotary echo technique [13] and electromagnetically induced transparency (EIT) [14, 15].

Further insight into the processes causing the dephasing is obtained by numerical simulations of the rotary echo experiment. The rotary echo technique is known from the research field of nuclear magnetic resonance and provides a tool to overcome disturbing effects due
Figure 1. Schematic view of the experimental setup used for the rotary echo experiment in the magnetic trap (a) and the dipole trap (ODT) (b). Panel (c) shows the setup for the EIT experiment in an ODT. The ODT beam is depicted in black and is propagating along the $z$-direction. The laser is linearly polarized along the $x$-axis and far red detuned (826 nm) with respect to transitions from the ground state. For the rotary echo sequence in the magnetic trap (dashed ellipse), both excitation lasers are collinear and propagating along the $-z$-direction. The laser for the lower transition (see figure 3(a)) of the two-photon excitation via the $5P_{3/2}$ into the $43S_{1/2}$, shown in red, has a wavelength of 780 nm and a $1/e^2$-diameter of 1 mm. The laser for the excitation into the $43S_{1/2}$ Rydberg state is shown in blue and has a wavelength of 480 nm and a waist of 42 $\mu$m. The lasers are $\sigma^+$ (780 nm) and $\sigma^-$ (480 nm) polarized with respect to the quantization axis (green arrow) given by a small magnetic field along the $z$-direction. In the ODT, the quantization is chosen along the $x$-axis, while the 780 nm beam is travelling along the $-x$-axis. For the EIT sequence, a quantization axis along the $y$-direction is chosen. The 780 nm beam is again $\sigma^+$ polarized and imaged with a charge-coupled device (CCD) camera, which is also used in the rotary echo sequence to image the ground state atoms after the Rydberg excitation has taken place. The Rydberg atoms are field ionized and detected by a multichannel plate (MCP).

to inhomogeneously distributed Rabi frequencies. An echo proves the coherence of a system directly by exciting and de-exciting the atomic sample by a time reversal in the driving field.

Another way to investigate the coherence properties of the Rydberg sample is the observation of EIT. Here, the coherence of the system is proven by the existence of a coherent superposition between two states, namely the ground state and the Rydberg state.

2. Setup

In this paper, two different experimental schemes are presented for the investigation of the dephasing of the driven Rydberg system excited from a cloud of ultracold $^{87}$Rb atoms, trapped either in a magnetic trap or in an optical dipole trap (ODT) in the $5S_{1/2}(f = 2, m_f = 2)$ state. The first experiments discussed here are rotary echo experiments [11, 13] in a magnetic trap and in an ODT. The second experiment uses EIT of an atomic sample trapped in an ODT. Details of the experimental setup are extensively discussed in [16]. Figure 1 shows specific information about the configurations used in different experiments.
2.1. Magnetic trap

Each experimental sequence described in this paper starts with an evaporatively cooled cloud of $^{87}$Rb atoms in a magnetic trap. The trapping frequencies are radially $\omega_r = 2\pi \times 322$ Hz and axially $\omega_z = 2\pi \times 18$ Hz. In order to study the density dependence of the dephasing, the density of ground state atoms in the trap is varied by means of a Landau–Zener sweep \[17\]. Using this technique atoms are transferred from the initial $5S_{1/2}(f = 2, m_f = 2)$ of $^{87}$Rb via a microwave photon with a frequency of 6.8 GHz into the magnetically untrapped $5S_{1/2}(f = 1, m_f = 1)$ state. This is the starting point for the rotary echo measurement done with magnetically trapped atoms described in section 3.1.

2.2. ODT

An ODT was set up to measure the dephasing rates using the rotary echo sequence (see section 3.1) and the EIT with the advantage that a homogenous magnetic field over the atomic sample can be applied in an arbitrary direction. Regarding the experiments on EIT (see section 3.2), it is important to switch off the trapping potential quickly. This is much easier to achieve in an ODT than in a magnetic trap. Furthermore, it is possible to tailor the geometry of the cloud in a wide range using optical potentials. The laser beam for the ODT is propagating collinearly to the long axis of the cigar-shaped cloud in the $z$-direction and is linearly polarized along the $x$-axis (see figure 1). The trapping laser has a wavelength of 826 nm, a waist of $w_0 = 21 \mu$m and a quality factor of the Gaussian beam of $M^2 = 1.5$. The power of the ODT is stabilized using a photodiode and an acousto-optical modulator to $P = 22$ mW. With a trap depth of $U_0 = 24 \mu$K one ends up with trap frequencies of $\omega_r = 2\pi \times 735$ Hz and $\omega_z = 2\pi \times 10$ Hz in the radial and axial directions, respectively.

After the evaporative cooling of the atomic sample and the Landau–Zener sweep in the magnetic trap, the ODT is ramped up within 50 ms. Afterwards the magnetic trap is switched off within 20 ms and the cloud is kept for 100 ms in the ODT for thermalization.

At a temperature of $(5.8 \pm 0.6) \mu$K the atomic cloud has a Gaussian width ($1/\sqrt{e}$ -radius) of $\sigma_z = (400 \pm 10) \mu$m in axial direction and of $\sigma_r \approx 5 \mu$m in radial direction. Due to the spatial resolution of the imaging system of 5.6 $\mu$m it is not possible to accurately resolve the radial size of the trapped atomic cloud.

The peak density of ground state atoms in the harmonic approximation is $n_g = N_0/((2\pi)^{3/2} \sigma_r^2 \sigma_z)$, with the trapped atom number $N_0$. In order to calculate $n_g$, the atom number $N_0$ and the size in axial direction $\sigma_z$ are taken from absorption images of the atomic cloud. The radial size of the atomic sample is given by $\sigma_r = \omega_r^{-1}(k_B T/m)^{1/2}$, where $k_B$ is the Boltzmann constant, $T$ is the temperature and $m$ is the mass.

3. Dephasing measurements

3.1. The rotary echo experiment

In the rotary echo experiment, the atoms are excited for a time $200 \leq \tau \leq 800$ ns into the $43S_{1/2}$ Rydberg state, by means of a two-photon excitation (see figure 3(a)). After a time $\tau_p \leq \tau$, the phase of the upper excitation laser is flipped by $\pi$ and thereby also the sign of the effective two-level excitation is inverted from $\Omega$ to $-\Omega$. In contrast, the sign of the interaction among the Rydberg atoms is not changed.
Visibility of the rotary echo signal in the ODT as a function of the pulse duration $\tau$ for different atomic peak densities $n_g$. In order to obtain the dephasing rates $\gamma_d = -\ln(V)/\tau$, an exponential decay of the visibility with increasing pulse length is assumed. Two fits are shown exemplarily in red. The inset shows the Rydberg atom number $N_R$ as a function of $\tau_p$, resulting in a typical rotary echo signal. This particular measurement is done for $\tau = 200$ ns and $n_g = 4.4 \times 10^{17}$ m$^{-3}$. A visibility of 75% is obtained from a parabolic fit.

Assuming undamped Rabi oscillations without spontaneous decay, interaction or finite linewidth of the excitation laser, the population is completely reversed from the Rydberg state into the ground state when $\tau_p = \tau/2$ (see inset of figure 2).

The data presented in this paper are taken for two different trap types. Firstly, the rotary echo experiment in the magnetic trap is done at a temperature of $T = (3.8 \pm 0.6) \mu$K with a peak density of ground state atoms of $n_g = 5.2 \times 10^{19}$ m$^{-3}$ [11]. Here, the lasers propagate along the $-z$-axis with $\sigma^+$ (780 nm) and $\sigma^-$ (480 nm) polarization resulting in a selective excitation of the $43S_{1/2}(j = 1/2, m_j = 1/2)$ state. The lasers are detuned with respect to the $5P_{3/2}$ state by $2\pi \times 470$ MHz, but tuned to resonance with respect to the two-photon excitation. As the detuning to the radiative state is large, the system reduces to an effective two-level atom. The two-photon excitation to the $43S_{1/2}(j = 1/2, m_j = 1/2)$ Rydberg state is insensitive to magnetic field inhomogeneities since the magnetic moments of this state and the ground state are equal. The waist of the 780 nm laser is 1 mm and 42 $\mu$m for the 480 nm laser. A small electric field of 200 V m$^{-1}$ is applied during the excitation in order to remove ions from the trap volume [9].

Secondly, the rotary echo is performed in an optical dipole trap with a slightly different setup compared with the rotary echo experiment of magnetically trapped atoms described in [11]. The 780 nm laser for the lower transition has a waist of 1 mm and is travelling along the $-x$-direction (see figure 1(b)). In this scheme, the atoms are quantized along $x$. The atoms are radially irradiated to reduce effects due to absorption, i.e. inhomogeneous intensities. The experiment is now performed for densities $n_g$ of ground state atoms between $3 \times 10^{17}$ and $2 \times 10^{18}$ m$^{-3}$.

The number of atoms in the Rydberg state is measured by field ionization of the Rydberg atoms. Due to the applied electric field, the ions are additionally pushed towards an MCP.
The atoms remaining in the ground state are imaged using resonant light on the 5S\textsubscript{1/2}(f = 2, \(m_f = 2\)) to 5P\textsubscript{3/2}(f = 3, \(m_f = 3\)) transition in the \(y\)-direction by a CCD camera.

Figure 2 shows the visibility of the rotary echo signal in the ODT for various pulse durations \(\tau\) and peak densities of ground state atoms \(n_g\). The inset of figure 2 shows a typical rotary echo measurement for a pulse duration of 200 ns. The visibility

\[
\mathcal{V} = \frac{N_R(\tau_p = 0) - N_R(\tau_p = \tau/2)}{N_R(\tau_p = 0) + N_R(\tau_p = \tau/2)}
\]

is obtained using a parabolic fit. Here \(N_R\) is the Rydberg atom number. The dephasing rate \(\gamma_d = -\ln(\mathcal{V})/\tau\) is obtained by assuming an exponential decay of the visibility with increasing pulse duration \(\tau\). This assumption is the simplest model of dephasing and is used due to the nescience of the exact functional behaviour.

In order to study the dephasing of the Rydberg state, theoretically numerical simulations of the rotary echo are performed. Therefore, the time evolution of a system consisting of \(N\) atoms in a box of volume \(V\), according to the Hamiltonian

\[
\mathcal{H} = \frac{\hbar \Omega(t)}{2} \sum_{i=1}^{N} \sigma_x^{(i)} + C_b \sum_{j<i} \hat{P}_{rr}^{(i)} \hat{P}_{rr}^{(j)},
\]

is computed. Here \(\sigma_x^{(i)}\) are the Pauli matrices \((\alpha \in \{x, y, z\})\), \(\hat{P}_{rr}^{(i)} = |r_i\rangle\langle r_i| = (1 + \sigma_x^{(i)})/2\) is the projector onto the excited Rydberg state, \(C_b\) denotes the strength of the van der Waals interaction, and the positions of the atoms \(r_i\) are randomly distributed but fixed. The van der Waals coefficient has a value of \(C_b = 1.6 \times 10^{-60} \text{ J m}^6\) for the 43S\textsubscript{1/2} Rydberg state [18], which is solely used for the measurements presented in this paper. We neglect resonant dipole–dipole interaction, i.e. Förster resonance, since the interatomic distances that are required for such a process are much smaller than the typical length scale in the experiment, i.e. the blockade radius [9, 19].

In the simulation, as in the experiment, the Rabi frequency \(\Omega(t)\) changes from \(\Omega\) to \(-\Omega\) at \(\tau_p\). The dipole blockade is used to drastically reduce the size of the Hilbert space as described in detail in [20]. The simulations are performed for \(N = 44, \ldots , 54\) and the interaction strength is \(C_b/(V^2 \hbar \Omega) = 0.01\). The Rabi frequency and the excitation time are chosen such that \(\Omega \tau = 0.32, \ldots , 0.74\). The resulting echo curves have been fitted the same way as the experimental data with a parabolic function. The dephasing rate is again obtained by \(\gamma_d = -\ln(\mathcal{V})/\tau\).

3.2. EIT in an interacting Rydberg gas

EIT occurs in systems with at least three levels coupled by two laser modes [14]. If two of the three levels have a long lifetime compared with the third level, the system evolves on the timescale of the third level into a dark state. This dark state is a superposition of the two long living states without a contribution of the ‘radiative’ intermediate state. Thus, probing with a weak laser on the transition between one of the long living states and the radiative state results in a narrow transparency window around the resonance of this transition.

Such a three-level system can be realized within the two-photon excitation scheme into long living Rydberg states. Again the 43S\textsubscript{1/2}(\(j = 1/2, m_j = 1/2\)) Rydberg state, which has a lifetime of \(\sim 100 \mu\text{s}\), is used. In the following, this state is referred to as \(|r\rangle\). The ground state \(|g\rangle\)
is the $5S_{1/2}(f = 2, m_f = 2)$ state, which does not decay on the timescale of our experiments. Finally, the radiative state $|e\rangle$ is given by the intermediate $5P_{3/2}(f = 3, m_f = 3)$ state, which has a decay rate of $\Gamma_{eg} = 2\pi \times 6 \text{ MHz}$. EIT in such a ladder system involving a Rydberg state has been investigated previously in a thermal vapour of $^{85}\text{Rb}$ [10] and in a weakly interacting Rydberg gas [12].

The system can be described using the Lindblad equation of motion for the density matrix $\rho$

$$\dot{\rho} = -\frac{i}{\hbar}[\mathcal{H}, \rho] + \mathcal{L}. \tag{3}$$

With the rotating wave approximation, the Hamilton operator reads in the basis $|g\rangle = (1, 0, 0)^t$, $|e\rangle = (0, 1, 0)^t$ and $|r\rangle = (0, 0, 1)^t$

$$\mathcal{H} = \frac{\hbar}{2} \begin{pmatrix} 0 & \Omega_p & 0 \\ \Omega_p^* & -2\delta_p & \Omega_c \\ 0 & \Omega_c^* & -2\delta \end{pmatrix}, \tag{4}$$

with the Rabi frequency defined as $\hbar \Omega_{c/p} = -E_0 d_{c/p}$, where $E_0$ is the electric field amplitude and $d_c$ and $d_p$ are the matrix elements for the dipole transition of the coupling and probe lasers, respectively. The detunings are (see figure 3(a)) $\delta_p = \omega_p - \omega_{eg}$ and $\delta_c = \omega_c - \omega_{re}$ for the probe and the coupling laser, respectively. The two-photon detuning is $\delta = \delta_p + \delta_c$.

The Liouville operator

$$\mathcal{L} = \begin{pmatrix} \Gamma_{ge}\rho_{ee} & -\frac{1}{2}\Gamma_e\rho_{ge} & -\frac{1}{2}\Gamma_{re}\rho_{re} \\ -\frac{1}{2}\Gamma_e\rho_{eg} & -\Gamma_{eg}\rho_{ee} + \Gamma_{re}\rho_{rr} & -\frac{1}{2}(\Gamma_e + \Gamma_r)\rho_{ee} \\ -\frac{1}{2}\Gamma_{re}\rho_{eg} & -\frac{1}{2}(\Gamma_e + \Gamma_r)\rho_{re} & -\Gamma_{re}\rho_{rr} \end{pmatrix}, \tag{5}$$

in (3) refers to the dissipation and dephasing of the three-level atom. The decay rates are $\Gamma_e = \gamma_{ed} + \Gamma_{eg}$ and $\Gamma_r = \gamma_{rd} + \Gamma_{re}$, with $\Gamma_{eg}$ and $\Gamma_{re}$ being the natural linewidths of $|e\rangle$ and $|r\rangle$, respectively. In addition, analogous to [22], the dephasing rates $\gamma_{ed}$ and $\gamma_{rd}$ of these states are taken into account.

Since the interaction energy among the Rydberg atoms exceeds all other energy scales, the dominant contribution to the dephasing is due to the interaction among Rydberg atoms, which can be assumed to be much larger than any other dephasing. Thus, in the following the dephasing rate of the intermediate state $\gamma_{ed}$ is neglected.

The imaginary part of the coherence between the ground state and the intermediate state $\text{Im}(\rho_{ge})$ can be calculated by solving (3) in steady state, i.e. $\dot{\rho} = 0$. On the other hand, the coherence $\rho_{ge}$ can be obtained from absorption images in steady state in the following way. The scattering cross-section is $\sigma = \sigma_0 \Gamma_{eg} \text{Im}(\rho_{ge})/\Omega_p$, with the on-resonance scattering cross-section $\sigma_0 = \sigma(\delta_p + \delta_c = 0)$. The atom number obtained from an off-resonant absorption image is $N(\delta) \sim \text{OD}(\delta)/\sigma_0$, where $\text{OD}(\delta)$ is the optical density depending on the two-photon detuning. Since $\sigma/\sigma_0 = N(\delta)/N_0$, the imaginary part of the coherence $\rho_{ge}$ in steady state is found to be

$$\text{Im}(\rho_{ge}) = \frac{\Omega_p}{\Gamma_{eg}} \frac{N(\delta)}{N_0}. \tag{6}$$
In the case of EIT, a certain number of the ground state atoms are excited to the Rydberg state and do not contribute to the coherence $\rho_{ge}$. Equation (6) must be corrected by this atom number and reads

$$\text{Im}(\rho_{ge}) = \frac{\Omega_p}{\Gamma_{eg}} \frac{N(\delta)}{N_0 - \max(N_R)},$$

where $\max(N_R)$ is the Rydberg atom number on two-photon resonance.

The two-photon excitation scheme to measure the dephasing rate $\gamma_d$ is mostly the same as for the rotary echo in the ODT. However, the direction for the probe laser driving the $5S_{1/2}(f = 2, m_f = 2)$ to $5P_{3/2}(f = 3, m_f = 3)$ excitation is now the $y$-axis (see figure 1(c)), which is also the quantization axis. This beam has a $1/e^2$-diameter of 13.5 mm and is detected by a CCD camera. The 480 nm coupling laser is tuned to the $5P_{3/2}(f = 3, m_f = 3)$ to $43S_{1/2}$ $(j = 1/2, m_j = 1/2)$ resonance and linearly polarized along the $x$-direction, while still travelling along the $-z$-axis. The Rabi frequency is $\Omega_p = 2\pi \times 2$ MHz and $\Omega_c = 2\pi \times 8$ MHz.
for the 780 and 480 nm lasers, respectively. The Rabi frequency for the coupling transition is obtained by calculating the dipole matrix element between the $5\text{P}_{3/2}$ and $43\text{S}_{1/2}$ states. The calculations are in good agreement with [23].

In order to decrease the optical density of the atomic sample, a Landau–Zener sweep and a time of flight of 100 $\mu$s are used. This lowers the peak density to values between $n_{\text{g}} \sim 2 \times 10^{17} \text{ m}^{-3}$ and $5 \times 10^{17} \text{ m}^{-3}$ at a temperature of 6.2 $\mu$K. The cloud is excited for 100 $\mu$s after which the Rydberg atoms are field ionized and detected by the MCP. Note that during illumination the density of the cloud is only decreasing by $\sim 50\%$.

Figure 3(b) shows three measurements of $\text{Im}(\rho_{ge})$ with $\Omega_\text{c} = 0$ (red) and $\Omega_\text{c} = 2\pi \times 4.5 \text{ MHz}$ (green and blue lines). In the presence of the coupling laser, a clear signature for the population of the dark state is visible, resulting in a decrease of the absorption for $\delta_\text{g} \to 0$.

Without dephasing, i.e. $\gamma_{\text{ed}} = \gamma_{\text{rd}} = 0$, the absorption would tend to zero, shown as a dashed blue line in figure 3(b). However, in the presence of interaction among Rydberg atoms, the dephasing $\gamma_{\text{rd}}$ of the Rydberg level is finite. The curves in figure 3(b) are calculated with $\gamma_{\text{ed}} = 2\pi \times 3.0 \text{ MHz}$ and $2\pi \times 3.5 \text{ MHz}$ for the blue and the green data sets, respectively. The measurements were taken with a detuning of $\delta_\text{c} = 2\pi \times 0.75 \text{ MHz}$ and $-2\pi \times 2.0 \text{ MHz}$ for the blue and the green data sets, respectively. The feature has a linewidth smaller than the natural linewidth of the $|e\rangle$ to $|g\rangle$ transition and, hence, is a direct proof of the coherent population of the dark state.

The dephasing rates can be measured by scanning the coupling instead of scanning the probe laser with the advantage of removing the absorption line shape from the data. The measurements of $\text{Im}(\rho_{ge})$ as a function of $\delta_\text{g}$ are shown in figure 4 for different peak densities of ground state atoms $n_{\text{g}}$ and different coupling Rabi frequencies $\Omega_\text{c}$. Again, approaching the two-photon resonance, a reduction of $\text{Im}(\rho_{ge})$ is found. A fit function to obtain the dephasing rates from the data shown in figure 4 can be found by solving (3) in the steady state. In the case of a perturbative probe laser, it is sufficient to take only the first-order expansion in the probe Rabi frequency $[22]$

$$\text{Im}(\rho_{ge}) \propto \frac{4\delta^2 \Gamma_\text{c} + \Gamma_\text{f} (|\Omega_\text{c}|^2 + \Gamma_\text{e} \Gamma_\text{f})}{(|\Omega_\text{c}|^2 + (\Gamma_\text{c} - 2i\delta_\text{p}) (\Gamma_\text{c} - 2i\delta))^2}. \quad (8)$$

In a first step, this function is used to fit the data shown in figure 4. The only free fit parameters in this function are the two-photon detuning $\delta$ and the rate $\Gamma_\text{f} = \Gamma_\text{ec} + \gamma_{\text{rd}} \approx \gamma_{\text{ed}}$. The value for the maximal coupling Rabi frequency is $\Omega_\text{c} = 2\pi \times 6.5 \text{ MHz}$ and the rate $\Gamma_\text{e} = \Gamma_\text{eg} + \gamma_{\text{rd}} \approx \Gamma_\text{eg}$. The parameters $\Gamma_\text{e} = \Gamma_\text{eg}$ and $\Omega_\text{eg}$ were kept fixed. In order to account for the finite probe Rabi frequency of $\Omega_\text{p} = 2\pi \times 800 \text{ kHz}$, the imaginary part of the coherence is later calculated in a second step by numerically solving (3). The red curves in figure 4 show the result of these calculations.

4. Results and discussion

The dephasing rates gathered from different experiments are shown in figure 5(a) against the measured maximal Rydberg atom number. In the case of the rotary echo data, for the experiment as well as for the simulation, this number is $\text{max}(N_\text{R}) = N_\text{R}(\tau_\text{p} = 0)$. In the EIT experiments $\text{max}(N_\text{R})$ is the Rydberg atom number on two-photon resonance, i.e. $\text{max}(N_\text{R}) = N_\text{R}(\delta = 0)$. The vertical dashed line separates the rotary echo data (left side) from the EIT data (right side). Note that both experiments are conducted on completely different timescales and that
Figure 4. Imaginary part of $\rho_{ge}$ as a function of the detuning $\delta_c$ of the coupling laser for different peak densities $n_g$ of ground state atoms and Rabi frequencies $\Omega_c$. The data are normalized for $|\delta_c| \gg 0$ to the steady-state value of $\text{Im}(\rho_{eg})$ for $\Omega_c = 0$. The red lines show the numerically obtained solution of (3) with a probe Rabi frequency $\Omega_p = 2\pi \times 800$ kHz. The coupling Rabi frequency in (a) and (b) is $\Omega_c = 2\pi \times 6.5$ MHz, in (c) $\Omega_c = 2\pi \times 4.6$ MHz and in (d) $\Omega_c = 2\pi \times 3.3$ MHz. The insets in (b) show two typical absorption pictures of the atomic cloud for $\delta_c \neq 0$ (left) and $\delta_c = 0$ (right). The red and grey colours in these pictures correspond to an optical density of 0.4 and 0, respectively.

on a longer timescale additional dephasings or decoherences could contribute to the dephasing rate $\gamma_d$.

The data show a tendency that for higher Rydberg atom numbers the dephasing is increasing. This assertion is supported by the results of the numerical calculation shown in figure 5(b). The only dephasing in the simulation is due to the interaction among Rydberg atoms and is not affected by any experimental uncertainties, e.g. laser linewidth.

In order to identify a power law dependence of the dephasing on $\text{max}(N_R)$ the data are plotted double logarithmic. A power law of the form $\gamma_d = a_c \text{max}(N_R)^b$ is shown in figure 5. The fit to the calculated rotary echo data yields $a_c = (0.03 \pm 0.01) \Omega$ and $b = 2.16 \pm 0.12$. The obtained exponent $b$ reflects the nature of the dephasing of the ultracold sample of Rydberg atoms, namely the van der Waals interaction $V_{vdW} = C_6 N_R^2$. Power laws with a constant exponent $b = 2.16$ are fitted subsequently to the data shown in figure 5(a). The data for both experiments, i.e. rotary echo and EIT, are fitted independently. The resulting coefficients are
Figure 5. Collection of the dephasing rates $\gamma_d$ as a function of the measured maximal Rydberg atom number $N_R$. All figures are plotted double logarithmic. Panel (a) shows the results from the rotary echo measurements in the magnetic trap (green squares) and in the ODT (brown squares) on the left side of the dashed vertical line. The right side of (a) shows the dephasing rates $\gamma_d = \gamma_{rd}$ obtained from the EIT measurements presented in figures 3 (red circles) and 4 (blue circles). Theoretically investigated dephasing rates for numerically calculated rotary echo experiments are shown in (b). The dashed line in (b) is a fit to the theoretical data, while the dashed dotted and solid line in (a) has the same slope as in (b), but fitted offsets (see the text).

$\alpha_{RE} \simeq \alpha_{EIT} = (0.09 \pm 0.02) \text{ s}^{-1}$. Note that the definitions of the dephasing rates are different for both experimental sequences and, thus, they might differ by a constant factor. However, although both dephasing rates are based on entirely different measurements the results are described by the same power law, which is furthermore in good agreement with the theoretical prediction of the rotary echo experiment. Thus, the additional energy scales, given by the decoherences $\Gamma_{eg}$ and $\Gamma_{re}$, due to the longer timescale of the EIT experiments do not affect the dephasing mechanism. It remains to be investigated if the obtained power law for the dephasing rate is a result of an underlying universality similar to the observed evidence for a universal scaling in ultracold Rydberg gases [24].

The dephasing rate due to a frequency uncertainty of the excitation laser is estimated to be $\simeq 1.5 \text{ MHz}$ on the minute timescale. An upper bound for the laser linewidth on the 100 ns timescale of $\sim 200 \text{ kHz}$ can be found from the measurement shown in the inset of figure 2. The laser linewidth causes the same effect on the data as the dephasing due to the interaction does, namely it decreases the visibility of the signal. This constant dephasing is observed as a kink on the left side of figure 5(a) for Rydberg atom numbers $\max(N_R) < 2 \times 10^3$. However, the ‘dephasing’ caused by the instrumentation is much smaller than the observed additional dephasing due to the interaction among the Rydberg atoms. This argument is again strengthened by the numerical simulation shown in figure 5(b), since these calculations do not take any laser linewidth into account.

In future experiments, the dephasing of the Rydberg state could be tailored by increasing the confinement of the atomic sample in the radial direction in order to study the dephasing rate more quantitatively in future experiments. If the radial width becomes smaller than the...
blockade radius, i.e. the situation becomes purely one dimensional, the reduced number of next neighbours reduces the dephasing due to the interaction. Conducting the experiments in an ODT comes with the advantage that the dimensionality is already reduced in comparison with the confinement in the magnetic trap. A disadvantage, which will be addressed in future experiments, is the adjustability of the atomic density in the ground state. The Landau–Zener sweep technique to decrease the density in the ODT comprises an additional heating due to the changed number of atoms in the magnetic trap, which are loaded into the ODT. The lowered thermalization rate in the axial direction might be an explanation for this effect.

5. Conclusion

To conclude, two different types of experimental techniques, which allow one to investigate the dephasing rates in a system of ultracold Rydberg atoms, are presented. The dephasing rates due to the interaction among Rydberg atoms were found to exceed all other dephasing rates, e.g. caused by instrumentation. The tendency of an increase of the dephasing rate with increasing Rydberg atom number is confirmed by conducting numerical calculations for the rotary echo experiment. Both the measurement and the numerical simulations are well described by the same exponent for a power law dependence of the dephasing rate on the maximal Rydberg atom number and is close to the naively expected exponent for a dephasing due to the van der Waals interaction. Furthermore, it is shown that the dephasing coefficients of the power law dependence are equal within the accuracy of the measurement. This result is remarkable in the sense that the two experiments that are used to obtain the dephasing rates are conducted on completely different timescales. In future experiments, strongly interacting frozen Rydberg gases will serve as a test ground for decoherence theories, where the dimensionality, the strength and the nature of the interaction can be tailored.

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