Rydberg–Stark deceleration and trapping of helium in magnetic fields

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Abstract

Triplet (S = 1) He Rydberg atoms in supersonic beams with an initial velocity of 350 m s$^{-1}$ have been decelerated to zero velocity and loaded into an off-axis electric trap in the presence and absence of magnetic fields. Comparing the deceleration efficiencies and the radiative decay of the population of trapped He Rydberg atoms to the (1s)$^3$(2s)$^1$ $^3$S$^1$ metastable level in the two sets of deceleration and trapping experiments revealed that the effects of magnetic fields up to 30 mT are negligible provided that a background dc electric field is maintained in the decelerator. A magnetic quadrupole trap of 30 mT depth corresponds to a He temperature of about 40 mK. The results thus represent an important step towards achieving high densities of cold paramagnetic samples following successive cycles of Rydberg–Stark deceleration, trapping, and radiative decay in overlaid electric and magnetic traps.

Keywords: Rydberg states, Rydberg–Stark deceleration, magnetic fields

(Some figures may appear in colour only in the online journal)

1. Introduction

Following the proposals of Breeden and Metcalf [1], and Wing [2], who suggested the exploitation of the large electric-dipole moments of Rydberg states to manipulate the motion of neutral atoms and molecules, Softley and co-workers developed the first devices which used these principles. With them, they deflected beams ofKr Rydberg atoms [3] and decelerated beams ofH$_2$ Rydberg molecules [4, 5]. Similar experiments revealed the influence of avoided crossings on the deceleration of Ar Rydberg atoms [6].

The deflection and deceleration of Rydberg atoms and molecules initiated in those studies is now referred to as Rydberg–Stark deceleration [7] and is related to the methods of multistage Stark [8–13] and multistage Zeeman [14–16] deceleration that are used to generate cold-molecule samples. Rydberg–Stark deceleration was improved using time-dependent potentials and more complex electrode structures, and soon allowed for excited atomic and molecular species to be trapped in three-dimensional space [17–23]. The dynamics of the processes causing loss of particles from the traps were then studied in detail in the case of H Rydberg atoms [24] and H$_2$ Rydberg molecules [25].

Compared to decelerating ground-state or metastable particles, Rydberg–Stark deceleration has the disadvantage that the atoms and molecules must first be excited to long-lived Rydberg–Stark states, but offers the advantage of very efficient deceleration because of the large electric-dipole moments of Rydberg states. Complete deceleration of atoms and molecules initially propagating at velocities between 500 and 1000 m s$^{-1}$ in supersonic beams can be achieved over short distances of several mm using electric fields that are only of the order of several hundreds of V cm$^{-1}$. In contrast, Stark deceleration and trapping of polar ground-state molecules requires the use of metre-long multistage devices and of electric fields about 100 times stronger [8]. A limitation of Rydberg–Stark deceleration is that the trapped particles are in excited states and can decay to states which do not have an electric-dipole moment sufficiently large for trapping. The decay typically takes place...
in a cascade of radiative transitions through which the excited atoms or molecules are brought back to the ground state or a metastable state [21, 24]. If this state has a magnetic-dipole moment, the particles might still be contained after radiative decay if a magnetic trap is superimposed on the electric trap. Many cycles of Rydberg–Stark deceleration and trapping could therefore be performed, allowing the accumulation of ground-state or metastable particles and the increase of phase-space density of cold atoms or molecules in the magnetic trap.

A similar strategy has been followed by Riedel et al [26] who decelerated ground-state NH radicals to standstill by multistage Stark deceleration and held them confined in a superimposed on-axis static magnetic trap (see also reference [13]). The magnetic-trap-loading process was repeated several times and an increase of the phase-space density was observed. However, collisions of the trapped molecules with the unaccelerated molecules in the supersonic beams eventually limited the phase-space density of cold NH molecules in the magnetic trap. Carrying out a similar experiment with Rydberg particles has the potential of a more efficient trap-loading process because of the larger efficiency of Rydberg–Stark deceleration. Moreover, electrode configurations for Rydberg–Stark deceleration can be designed to deflect the beams during deceleration, enabling off-axis trapping [21]. As a first step towards realizing such experiments, the effect of an inhomogeneous magnetic field on the Rydberg–Stark deceleration, deflection, and trapping process needs to be characterized because Rydberg states with non-zero angular momentum are sensitive to magnetic fields and exhibit energy splittings which may influence their behaviour during Rydberg–Stark deceleration.

In a previous study, we measured high-resolution spectra of the Rydberg states of triplet He in the presence of electric fields in the range 0–250 V cm\(^{-1}\) and magnetic fields in the range 0–50 mT needed for magnetic trapping of triplet He following Rydberg–Stark deceleration [27]. In those field ranges and for the Rydberg states with the quantum number \(n\) between 20 and 30 relevant for these experiments, the diamagnetic term of the Zeeman interaction is negligible and both the paramagnetic Zeeman and the Stark effects are linear. The study of the Stark and Zeeman effects for a broad range of angles between electric- and magnetic-field vectors revealed the ubiquity of avoided crossings. Such crossings would, in general, prevent efficient Rydberg–Stark deceleration by changing the magnitude and, in some cases, even the sign of the electric dipole moment. The results indicated that most avoided crossings take place at low electric-field strengths, below 20 V cm\(^{-1}\), suggesting that the application of a background dc electric-field component should alleviate adverse effects of the magnetic field on Rydberg–Stark deceleration and trapping.

Understanding the role of magnetic fields in Rydberg–Stark deceleration is also relevant in the context of experiments involving the manipulation of antihydrogen [28, 29] and positronium [30, 31] using inhomogeneous electric fields. We present here experiments which characterize Rydberg–Stark deceleration of beams of triplet He Rydberg atoms in the presence of inhomogeneous magnetic fields and show that the effects of the magnetic fields on the deceleration and trapping process are minimal provided that a background dc electric field is maintained in the decelerator, as expected from the results obtained in [27].

2. Experimental setup

The experimental setup is presented schematically in figure 1(a). It is a modified version of a setup used to decelerate, deflect, and load H Rydberg atoms to study their trap-decay processes [21, 24] and consists of three distinct regions labelled (i)–(iii). In region (i), a cold supersonic beam of metastable \((1s)^1(2s)^1{^1S}_1\) He (He\(^+\) hereafter) atoms is generated in an electric discharge through a supersonic beam of pure He immediately after the orifice of a pulsed valve. After collimation, it enters the second region, region (ii), enclosed in two heat shields and connected to a two-stage pulse-tube cooler, where the atoms are excited to a selected Rydberg–Stark state, decelerated, deflected, and trapped electrostatically. The heat shields and pulsed-tube cooler enable the control of the blackbody-radiation temperature [24], which has been known for a long time to induce transitions between neighbouring Rydberg states [32]. After a selected time, the He Rydberg atoms contained in the off-axis trap are pulsed field ionized and the He\(^+\) ions are extracted towards a microchannel-plate (MCP) detector located at the end of a field-free region (region (iii)).

The supersonic beam of pure He atoms in the ground state is generated using a home-built pulsed valve operated at a repetition rate of 25 Hz. It produces gas pulses with a duration of about 20 \(\mu\)s. The valve body is cooled with a two-stage pulse-tube cooler to temperatures in the range between 10 and 20 K, resulting in supersonic beams with mean forward longitudinal velocities in the range between 350 and 450 m s\(^{-1}\). These atoms are excited to the metastable \((1s)^1(2s)^1{^3S}_1\) state in an electric discharge at the exit of the valve. The low starting velocities facilitate the deceleration and trapping procedures.

The He\(^+\) beam is skimmed with a 2 mm-diameter skimmer located \(\sim\)20 cm downstream from the valve before it enters the decelerator (region (ii) on figure 1), where it is photoexcited to Rydberg–Stark states with principal quantum number \(n\) around 20 using a pulse-amplified diode-laser system described in [27]. The laser beam intersects the He\(^+\) beam at right angles between two parallel copper electrodes labelled 1 and 2 in figure 1(b), of dimensions 20 mm \(\times\) 7 mm separated by 3 mm in the \(y\) direction. These electrodes are used to generate a static electric field \(F_{\text{exc}}\) at excitation of sufficient strength and spatial homogeneity for individual Stark states to be excited selectively. The laser polarization can be set to be parallel (or perpendicular) with respect to \(F_{\text{exc}}\) by using a Berek compensator placed in front of the entrance window of the vacuum chamber. In the former (latter) case, the He\(^+\) atoms are photoexcited to Rydberg–Stark states with \(n\ell = 0\) (\(\pm 1\)). In the experiments presented here, \(F_{\text{exc}}\) was chosen to be 500 V cm\(^{-1}\), which is just below the field, called Inglis–Teller (IT) field, at which the \(n = 20\) and 21 Stark manifolds start to overlap spectrally [33].
Immediately after photoexcitation, the atoms enter the deceleration and trapping region, which consists of 10 copper electrodes (3–12) depicted on an enlarged scale in figure 1(b). Electrodes 9 and 11 are located below, and electrodes 10 and 12 above the plane of the figure. They generate the constant offset electric field required to suppress avoided crossings in the Stark map when a magnetic field is present [27]. Electrodes 3–8 are cylinders with a diameter of 1 mm. The out-of-plane electrodes are cylinders with a diameter of 2 mm separated by 10 mm when no magnets are present, or 5 mm when magnets are present.

The operational principle of the decelerator is inspired from the decelerator used to decelerate and deflect H and D Rydberg atoms, and H_2 Rydberg molecules [21, 24, 34].
Figure 2. (a) Electric-potential pulses applied to electrodes 3, 4, 5, 6, 7 and 8 used to decelerate, deflect and trap the He Rydberg atoms. (i) . . . (iv) denote the times for which the electric-field distributions are shown on panel (b). (b) Electric-field distributions in the decelerator at the times indicated in panel (a). The lines of constant electric field range from 0 to 200 V cm$^{-1}$ in steps of 20 V cm$^{-1}$. The black dot indicates the position of the centre of the Rydberg-atom cloud assuming an initial velocity of 350 m s$^{-1}$ in the $z$ direction and Rydberg–Stark states with $n = 20$, $k = 12$. 


Time-dependent exponentially-decaying potentials

\[ V_i(t) = V_{\text{off}} + \theta \left( t - t_{0,i} \right) \left( V_{\text{amp},i} - V_{\text{off}} \right) \exp \left( -\frac{t - t_{0,i}}{\tau_i} \right) \]  

are applied to each of the electrodes numbered 3–8 to create inhomogeneous electric fields with strong gradients required for the deceleration of He atoms in Rydberg–Stark states. In equation (1), \( i \in \{3, 4, 5, 6, 7, 8\} \) represents the electrode index, \( V_{\text{off}} \) is the permanent potential applied on electrode \( i \) to generate a constant quadrupole trapping field, \( \theta \left( t - t_{0,i} \right) \) is the Heaviside step function, \( V_{\text{amp},i} \) is the amplitude of the potential pulse applied to the electrode \( i \), \( t_{0,i} \) is the time at which the potential is switched on, and \( \tau_i \) is the time constant of the \( i \)th exponential decay. The potential pulses are generated by transmitting square-shaped pulses through RC circuits with the desired time constant \( \tau_i \). The amplitude \( V_{\text{amp},i} \) is chosen in the range between 400 and 800 V, depending on the selected \( n \) and \( k \) values of the excited Rydberg–Stark states, to ensure the strong decelerating and deflecting forces needed to steer the atoms into the electric-quadrupole trap between electrodes 4, 5, 7, and 8.

The parameters \( V_{\text{amp},i} \), \( t_{0,i} \), and \( \tau_i \) are chosen to maximize the number of Rydberg atoms loaded into the off-axis trap. The optimal sequence and exemplary electric-field distributions are depicted in figure 2.

The forces acting on a Rydberg atom and resulting from the inhomogeneous electric-field distribution generated by the decelerator electrodes are given to first order by [33]:

\[ -\nabla E_{\text{Stark}} = -\frac{3}{2} nke_0 \nabla F, \]

where \( \nabla F \) is the gradient of the electric-field strength at the position of the atom.

Figure 2 presents the pulses applied to the electrodes and the resulting electric-field distributions. Electrodes 1, 2, and 13 on figure 1(b) are kept at constant potentials in the range between 10 and 30 V. The out-of-plane electrodes [larger circles in figure 1(b)] are also kept at constant potentials in the range between 10 and 50 V to maintain the offset electric field necessary to mitigate the negative effects of the magnetic field on the deceleration process [27]. These electrodes also close the trap in the \( x \) direction.

The magnetic field in the setup is generated by permanent magnets attached to the out-of-plane electrodes 9 and 10 [see figure 1(c)]. Six cylindrical magnets (2 mm diameter × 2 mm height) in two stacks of three separated by 5 mm above and below the plane depicted in figure 1(b) were used to create a quadrupole magnetic-field distribution. The magnets are made of compressed NdFeB powder encased in a stainless-steel shell, with each magnet having a magnetization in the range 860–995 kA m\(^{-1}\). To estimate the magnetic field in the experiment, a finite-element-method simulation was performed using the commercial software COMSOL Multiphysics 5.4 (comsol.ch/release/5.4). The contour plot of the magnetic-field distribution is depicted in figure 3. It is designed so that the magnetic field is maximal (∼30 mT) in the region the atoms traverse during deceleration and deflection. The depth of the magnetic trap corresponds, for the (1s)\(^3\)(2s)\(^1\) 3S\(_1\) metastable state of He, to a temperature of ∼40 mK.

After deceleration, the He Rydberg atoms are detected by pulsed field ionization (PFI). PFI is achieved by applying pulsed potentials \( V_{\text{ion},1} = +4 \) kV on electrodes 1 and 2 and \( V_{\text{ion},2} = +2.5 \) kV on electrodes 3 and 4, both with a rise time of 50 ns and a duration of about 2 \( \mu \)s. Electrodes 5 and 6 remain grounded if detection is to be carried out along the axis of the He\(^+\) beam. When detection of the deflected and trapped sample is performed, +4 kV is applied on electrodes 2 and 13, +2.5 kV on electrodes 4 and 8, and electrodes 5 and 7 remain grounded. The He\(^+\) ions produced by PFI are accelerated towards an MCP detector with a diameter of 40 mm located ∼80 mm beyond the PFI region, where their times-offlight (TOF) distribution is monitored. An Einzel lens is used to focus the ions onto the MCP detector regardless of whether they were generated on the original beam-propagation axis or in the off-axis trap. Typical results following PFI by pulsing electrodes 1, 2, 3, and 4 without a magnetic field present are displayed in figure 4. All three traces were recorded by PFI at the same time delay after photoexcitation. The pulse strengths used are indicated in the caption. When the He Rydberg atoms are not decelerated, PFI gives rise to the He\(^+\) TOF trace depicted in black in figure 4, with a maximum at about 2.5 \( \mu \)s. The TOF traces obtained after deceleration (the blue and red traces in figure 4) have their maxima at shorter times because, at the time of PFI, the Rydberg atoms are located closer to electrodes 3 and 4 and are therefore produced at a higher electric potential, which accelerates the He\(^+\) ions produced by PFI more strongly. The He\(^+\) TOF distributions thus directly reflect the efficiency of the deceleration, which can
be used to quantify the effects of Rydberg–Stark deceleration pulse sequences and optimize them.

The Lorentz force deflects the trajectories of the \( \text{He}^+ \) ions if they are not moving exactly parallel to the magnetic field. The inhomogeneous magnetic-field distribution generated by the permanent magnets creates regions in which this is the case, an effect we could not eliminate in our setup and which reduced the detection efficiency of the trapped Rydberg atoms.

The \(^{3}\text{He}-\text{Rydberg}\)-atom beam is first decelerated from \( \sim350 \text{ m s}^{-1} \) to zero velocity component in the direction of the \( \text{He}^+ \) beam while being deflected by 90° and loaded into the off-axis trap. This process lasts for \( \sim70 \mu\text{s} \). After this time, the atoms are kept in the trap for a defined duration before they are pulsed-field ionized. By monitoring the \( \text{He}^+ \) ion signal at the MCP as a function of time after trap loading, the decay of the \( \text{He} \) Rydberg atoms in the trap can be monitored. This experiment was first conducted without the permanent magnets at different temperatures of the decelerator environment to study the effects of the black-body-radiation-induced transitions. Once these experiments revealed that black-body-radiation-induced transitions do not play a significant role, the trapping experiment was then repeated at room temperature with the permanent magnets to assess the effects of the magnetic field.

3. Experimental results

3.1. Off-axis trapping at different black-body-radiation temperatures without magnetic fields

The first experiments were performed to study the dynamics inside the trap when the temperature of the black-body radiation \( T_{BB} \) was varied. Temperature variation in the range between 4.7 and 298 K was achieved through the use of a two-stage pulsed-tube cooler attached to the decelerator and the two heat shields (see figure 1). The heat shields have small openings to allow passage for the helium and laser beams. These openings, in principle, allow external black-body radiation to access the trap and prevent accurate temperature control, however, this contribution was reduced to less than 1%, as explained in reference [24].

The experiments were carried out after the initial preparation of the \( \text{He} \) atoms in the \( |n, k, m_{\ell} \rangle = |20, 12, \pm 1 \rangle \) Rydberg–Stark states. The states are labelled with the magnetic quantum number \( m_{\ell} \) because the spin–orbit coupling in the \(^3\text{P}\) term at \( n = 20 \) (\( \sim30 \text{ MHz} \)) is much less than both the Stark splittings below the IT field and the laser bandwidth (\( \sim150 \text{ MHz} \)).

Rydberg states around \( n = 20 \) were chosen for these experiments because their IT fields are much larger than the fields present in the decelerator (the IT field scales as \( n^{-5} \) [33]) while their dipole moment is still considerable—the electric-dipole moment scales as \( n^2 \) and is \( \sim600 \text{ D} \) for the \( |20, 12, \pm 1 \rangle \) state [33]. Consequently, an overall more efficient deceleration could be achieved than when using states of higher \( n \) values. The reduced dimensions of the decelerator compared to that used to decelerate, deflect, and trap \( \text{H} \) atoms in states with \( n \) values around 30 [21] increased the gradients and thus the forces applied so that the effects of the heavier mass and lower dipole moments could be compensated.

The \( \text{He} \) atoms enter the decelerator with initial velocity of 350 m s\(^{-1}\) and undergo deceleration between electrodes 3 and 4 (\( \pm500 \text{ V} \) with decay constants 3.2 and 2.7 \( \mu\text{s} \), respectively). The electric-potential pulses applied on electrodes 5 and 6 (\( \pm500 \text{ V} \) with decay constants 9.8 and 5.3 \( \mu\text{s} \), respectively) serve as deflecting potentials and deviate the atom beam around electrode 4 into the off-axis trap between electrodes 4, 5, 7 and 8. The final deceleration is achieved by applying electric-potential pulses on electrodes 7 and 8 (\( \pm500 \text{ V} \) with decay constants 2.4 \( \mu\text{s} \) for both electrodes). The whole process from laser excitation to trap loading takes \( \sim70 \mu\text{s} \) (see section 2 and figure 2).

After trap loading, the decay of the atoms from the trap is observed by progressively delaying the PFI pulse, as depicted in figures 5(a) and (b). The decay curves were fitted with an exponential function and the trapping lifetimes were determined. The fitted function has the following form:

\[
f(t) = 0.95 \cdot \exp \left( \frac{t - t_{\text{loading}}}{\tau} \right) + 0.05,
\]

where \( t \) is the time measured from photoexcitation, \( t_{\text{loading}} \) is the time it takes to load the trap (about 70 \( \mu\text{s} \)), and \( \tau \) is the trapping lifetime. \( \tau \) is the only adjustable parameter in equation (4). A 5% background signal in figure 5 is attributed to the PFI of He Rydberg atoms that have undergone transitions to very long-lived Rydberg–Stark states with \( |m_{\ell}| > 3 \) or \( \text{He}^+ \) ions produced from \( \text{He}^+ \) through collisions in the background gas. Because the data, the fits, and the calculations are normalized to 1, the pre-factor of the exponential is adjusted so that the decay curve starts at 1 at \( t = t_{\text{loading}} \).

The trapping times for different \( T_{BB} \) under otherwise identical conditions are summarized in table 1. The statistical uncertainties represent one standard deviation of the fitted
values and mainly originate from the uncertainty of the background level at very long trapping times, beyond 160 μs.

Within the experimental uncertainties, these values cannot be distinguished from each other. We thus conclude that varying \( T_{BB} \) does not significantly influence the trapping experiments carried out with He Rydberg atoms originally prepared in the \( |20, 12, ±1\rangle \) triplet Rydberg–Stark states.

The observed and calculated trapping times (see section 4 below for the description of the calculation) obtained at \( T_{BB} = 4.7 \) K agree within the experimental uncertainties. The experimental data recorded at higher values of \( T_{BB} \) do not reflect the trend towards shorter lifetimes predicted by the calculations. The differences in these cases can be explained by transitions to longer-lived Rydberg–Stark states during the 70 μs-long trap-loading phase of the experiment. Indeed, only 0.3% of the initially prepared Rydberg atoms is predicted to remain in the original state at 298 K, whereas this fraction increases to ∼5% at 4.7 K.

3.2. Off-axis trapping at room temperature with and without magnetic fields

Because the black-body-radiation temperature of the decelerator and its environment does not play a significant role in the trapping experiments (see previous subsection), we studied how inhomogeneous magnetic fields influence the deceleration and trapping process at room temperature. The hypothesis, based on the conclusions from reference [27], is that inhomogeneous magnetic fields do not significantly perturb the deceleration and trapping process of Rydberg He as long as precautions are taken to ensure a sufficiently strong background electric field in the decelerator.

The approach chosen to test this hypothesis was to perform off-axis trapping experiments with and without permanent magnets on electrodes 9 and 10 under otherwise identical conditions. The initial velocity of the He* beam was 350 m s\(^{-1}\), the polarization of the laser radiation was perpendicular to the electric field \( F_{\text{exc}} \) in the photoexcitation region, and the He atoms were initially prepared in the \( |20, 12, ±1\rangle \) Rydberg–Stark states.

After trap loading, which was performed as described in section 3.1, the decay of the atoms from the trap was observed by progressively delaying the PFI pulse, as depicted on figures 6(a) and (b). The decay curves were fitted with exponential functions to determine the trapping lifetimes. The fitted function has the following form:

\[
f(t) = 0.98 \cdot \exp \left( -\frac{t - t_{\text{loading}}}{\tau} \right) + 0.02,
\]

where \( t \) is the time measured from photoexcitation, \( t_{\text{loading}} \) is the time it takes to load the trap (∼70 μs), and \( \tau \) is the trapping lifetime. \( \tau \) is the only adjustable parameter in equation (4). The mean observed normalised background level (0.02) in these measurements is smaller than in the measurements described in section 3.1, which is attributed to the slightly different background pressure in the chamber and the density of He* generated in the discharge.

The trapping lifetime obtained from the fit of the decay curve recorded following deceleration in the absence and in the presence of magnetic fields are:

\[
\tau(a) = (24 \pm 2_{\text{stat}} \pm 3_{\text{syst}}) \mu s,
\]

and

\[
\tau(b) = (26 \pm 2_{\text{stat}} \pm 3_{\text{syst}}) \mu s,
\]

Figure 5. Integrated PFI signal of the He Rydberg atoms (\( |n, k, m_k\rangle = |20, 12, ±1\rangle \)) in the off-axis trap as a function of the time between the loading of the trap and the application of the PFI pulse. \( T_{BB} \) of the decelerator and its environment have been varied in the range 4.7–298 K. This decay is fitted with an exponential-decay curve. (a) Experimental data showing the total integrated PFI signal. (b) Fits of the experimental data points (full lines) and the trap-decay curves with a trapping lifetime corresponding to the theoretically calculated value which takes into account the temperature (\( T_{BB} \)) of the environment (dashed lines).

| \( T_{BB}/K \) | \( \tau_{\text{meas}} \) (μs) | \( \tau_{\text{calc}} \) (μs) |
|-----|-----|-----|
| 298 | (40 ± 4_{\text{stat}} ± 5_{\text{syst}}) | 13.7 |
| 210 | (33 ± 5_{\text{stat}} ± 6_{\text{syst}}) | 16.2 |
| 120 | (41 ± 6_{\text{stat}} ± 5_{\text{syst}}) | 20.1 |
| 30  | (34 ± 4_{\text{stat}} ± 5_{\text{syst}}) | 25.3 |
| 4.7 | (31 ± 4_{\text{stat}} ± 5_{\text{syst}}) | 26.8 |

Table 1. Comparison of calculated lifetimes (\( \tau_{\text{calc}} \)) and observed trap-decay times (\( \tau_{\text{meas}} \)) of the \( |20, 12, ±1\rangle \) triplet Rydberg–Stark states of He at \( T_{BB} \) in the range 4.7–298 K.
The calculated radiative lifetimes in an electric field of 50 V cm\(^{-1}\), which corresponds to the average electric-field strength in the off-axis trap, for \(n = 20\), \(|m_\ell| = 1\) Rydberg–Stark states of He at temperatures of 4.7 K (red), 100 K (blue), and 298 K (black) are presented in figure 7. They do not change significantly as the electric-field strength is raised or lowered.

The calculated \(k\)-dependence of the natural lifetimes becomes more pronounced at lower black-body temperatures.
At 4.7 K, for example, the \( k = 0 \) state has a radiative lifetime of 49.6 \( \mu s \), while the outer-most states with \( |k| = 18 \) have radiative lifetimes under 20 \( \mu s \). This effect can be explained by the fact that the states located at the edges of the manifold have the largest \( \ell \) = 1 character and thus their decay to the metastable state is more probable. At 298 K, the calculated lifetimes all lie between 12 and 17 \( \mu s \), which indicates the black-body-radiation-induced transitions play a more pronounced role. For the \( n = 20, k = 12, |m| = 1 \) Rydberg–Stark states and at the environment temperature of 298 K used in the experiments, the radiative lifetime is predicted to be 13.7 \( \mu s \).

The data presented in figures 5 and 6 are well described with a single exponential curve. Influence from collisions between the Rydberg atoms and the background gas would lead to a multi-exponential decay [24]. We therefore conclude that these collisions do not contribute significantly to the observed decay.

The lack of a strong dependence of the trapping time on \( T_{BB} \) indicates that black-body-radiation-induced transitions to states which are not confined in the trap do not contribute significantly to the trapping time. The measured trapping times, however, are significantly different from the times expected for the temperatures presented in figure 5 with only the values obtained at 30 K and 4.7 K being within one standard deviation of the calculated values. This is attributed to the fact that the atoms have time to undergo collisions and black-body-radiation-induced processes during the relatively long time it takes to load the trap, as discussed in section 3.1. Both types of processes can increase \( |m| \) and prolong the expected trapping times.

The experimental results obtained in the trapping experiments carried out at 298 K with and without magnetic fields yield trapping efficiencies and trapping times that are identical within the experimental uncertainties, but larger than the calculated radiative lifetimes. The observed lifetimes are also consistent with those observed in the experiments carried out without magnetic fields but for different values of \( T_{BB} \). We therefore conclude that the trapping is independent of the magnetic field and that the observed trapping times at \( T_{BB} = 298 \) K are the results of black-body-radiation-induced transitions during the 80 \( \mu s \)-long deceleration period preceding trap loading. The calculated rates and the propensity rules for \( n \) - and \( m \) -changing transitions suggest that the states contributing to the decay have values around 20 and \( |m| \) values up to \( |m| = 2 \). Under these conditions, the trapping times are prolonged but the non-exponential nature of the decay is not observable within the experimental uncertainties.

This study showed that varying \( T_{BB} \) during the deceleration and off-axis trapping of He Rydberg atoms prepared in the \( 20,12,\pm1 \) Stark states does not significantly influence the efficiency of the deceleration and trap-loading process nor the trapping lifetimes, and that the deceleration and off-axis trapping of Rydberg He in \( 20,12,\pm1 \) states can be successfully performed in the presence of magnetic fields of up to \( \sim 30 \) mT, corresponding to magnetic-trap depths of \( \sim 40 \) mK for metastable He. The results confirm the prediction drawn from the results of previous spectroscopic experiments (see, e.g. figure 17 in reference [27]), that a linear Stark-effect regime, which enables Rydberg–Stark deceleration, can be established to mitigate magnetic-field effects by applying a sufficiently strong background dc electric field, i.e., \( \sim 20 \) V cm\(^{-1}\) for the experiments presented here.

Experiments in which the magnets were attached to electrodes 10 and 12 were conducted with the goal of detecting the off-axis-trapped He Rydberg atoms from the superimposed traps; however, these efforts remained unsuccessful because the magnetic fields of the magnetic trap deflect the ion trajectories too strongly for them to be detected in our experiment. This problem can be overcome, in future work, by a different design of the trap and the ion-extraction electrodes.

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References

[1] Breeden T and Metcalf H 1981 Phys. Rev. Lett. 47 1726–9
[2] Wing W H 1980 Phys. Rev. Lett. 45 631–4
[3] Townsend D, Goodgame A L, Procter S R, Mackenzie S R and Softley T P 2003 J. Phys. B: At. Mol. Opt. Phys. 34 439–50
[4] Procter S R, Yamakita Y, Merkt F and Softley T P 2003 Chem. Phys. Lett. 374 667–75
[5] Yamakita Y, Procter S R, Goodgame A L, Softley T P and Merkt F 2004 J. Chem. Phys. 121 1419–31
[6] Vliegen E, Wörner H J, Softley T P and Merkt F 2004 Phys. Rev. Lett. 92 33005
[7] Hogan S D 2016 EPJ Tech. Instrum. 3
[8] Bethlem H L, Berden G, Crompoets F M H, Jongma R T, van Roij A J A and Meijer G 2000 Nature 406 491–4

[9] Hoekstra S, Metsälä M, Zieger P C, Scharfenberg L, Gilijamse J J, Meijer G and van de Meerakker S Y T 2007 Phys. Rev. A 76 063408

[10] Sawyer B C, Lev B L, Hudson E R, Stuhl B K, Lara M, Bohn J L and Ye J 2007 Phys. Rev. Lett. 98 253002

[11] van den Berg J E, Turkesteen S H, Prinsen E B and Hoekstra S 2012 Eur. Phys. J. D 66 235

[12] Jansen P, Quintero-Pérez M, Wall T E, van den Berg J E, Hoekstra S and Bethlem H L 2013 Phys. Rev. A 88 043424

[13] Haas D, von Planta C, Kierspel T, Zhang D and Willitsch S 2019 Commun. Phys. 2 101

[14] Wiederkehr A W, Schmutz H, Motsch M and Merkt F 2012 Mol. Phys. 110 1807–14

[15] Akerman N, Karpov M, Segev Y, Bibelnik N, Narevicius J and Narevicius E 2017 Phys. Rev. Lett. 119 073204

[16] Plomp V, Gao Z, Cremers T and van de Meerakker S Y T 2019 Phys. Rev. A 99 033417

[17] Vanhaecke N, Comparat D and Pillet P 2005 J. Phys. B: At. Mol. Opt. Phys. 38 S409–19

[18] Vliegen E, Hogan S D, Schmutz H and Merkt F 2007 Phys. Rev. A 76 023405

[19] Hogan S D and Merkt F 2008 Phys. Rev. Lett. 100 043001

[20] Hogan S D, Seiler C and Merkt F 2009 Phys. Rev. Lett. 103 123001

[21] Seiler C, Hogan S D, Schmutz H, Agner J A and Merkt F 2011 Phys. Rev. Lett. 106 073003

[22] Lancuba P and Hogan S D 2014 Phys. Rev. A 90 053420

[23] Lancuba P and Hogan S D 2016 J. Phys. B: At. Mol. Opt. Phys. 49 074006

[24] Seiler C, Agner J A, Pillet P and Merkt F 2016 J. Phys. B: At. Mol. Opt. Phys. 49 094006

[25] Seiler C, Hogan S D and Merkt F 2011 Phys. Chem. Chem. Phys. 13 19000–12

[26] Riedel J, Hoekstra S, Jaeger W, Gilijamse J J, van de Meerakker S Y T and Meijer G 2011 Eur. Phys. J. D 65 161–6

[27] Tkáč O, Žeško M, Agner J A, Schmutz H and Merkt F 2016 J. Phys. B: At. Mol. Opt. Phys. 49 094006

[28] The AEGIS Collaboration Kellerbauer A, Allkofer Y et al 2012 Hyperfine Interact. 209 43–9

[29] The AEGIS Collaboration Doser M, Amsler C et al 2012 Class. Quantum Grav. 29 184009

[30] Deller A, Alonso A M, Cooper B S, Hogan S D and Cassidy D B 2016 Phys. Rev. Lett. 117 073202

[31] Rayment M H, Gurung L, Sheldon R E, Hogan S D and Cassidy D B 2019 Phys. Rev. A 100 013410

[32] Spencer W P, Vaidyanathan A G, Kleppner D and Ducis T W 1982 Phys. Rev. A 25 380–4

[33] Gallagher T F 1994 Rydberg Atoms (Cambridge: Cambridge University Press)

[34] Hogan S D, Seiler C and Merkt F 2013 J. Phys. B: At. Mol. Opt. Phys. 46 045303

[35] Zimmerman M L, Littman M G, Kash M M and Kleppner D 1979 Phys. Rev. A 20 2251–75

[36] Drake G W F and Swainson R A 1991 Phys. Rev. A 44 5448–59

[37] Zhelyazkova V, Žeško M, Schmutz H, Agner J A and Merkt F 2019 Mol. Phys. 117 2980–9