Single-atom detection on a chip: from realization to application

A Stibor, H Bender, S Kühnhold, J Fortágh, C Zimmermann and A Günther

CQ Center for Collective Quantum Phenomena and their Applications, Eberhard-Karls-Universität Tübingen, Auf der Morgenstelle 14, D-72076 Tübingen, Germany
E-mail: aguenth@pit.physik.uni-tuebingen.de

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Abstract. In this paper, we describe the preparation and detection of ultracold atoms on a microchip with single-atom sensitivity. The detection scheme is based on multi-photon ionization of atoms and the subsequent guiding of the generated ions by ion optics to a channel electron multiplier. We resolve single atoms with a detection efficiency above 60%. The detector is suitable for real-time observations of static and dynamic processes in ultracold quantum gases. Although the ionization is destructive, sampling a small subset of the atomic distribution is sufficient for the determination of the desired information. We take full high-resolution spectra of ultracold atoms by ionizing only 5% of the atoms. Using an additional microwave near 6.8 GHz, the detection scheme becomes energy, position and state selective. This can be used for in situ determination of the energy distribution and temperature of atom clouds inside the trap and applied for future correlation measurements.

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1 Author to whom any correspondence should be addressed.
1. Introduction

Single-atom manipulation and detection is moving to the forefront of cold atom research. The aim to investigate many particle physics beyond the mean-field approach [1–3] and to measure quantum correlations in ultracold gases [4–6] has been the driving force for a number of new developments. Recently demonstrated single-atom detectors are based on optical cavities [4, 7–9], fluorescence imaging [10], laser ionization [11], electron impact ionization [12] and field ionization [13]. Metastable atoms have been directly counted using microchannel plates [14]. The integration of such single-atom detectors on atom chips outlines even new applications in the areas of engineered entanglement, quantum information processing, real-time atom interferometry and surface diagnostics [15, 16].

In this work, we describe the operation of an integrated single-atom detector based on photo-ionization of ultracold rubidium atoms and efficient ion counting. Due to single-atom sensitivity, we measure the energy distribution and temperature of atoms inside the trap, ionizing only a small fraction (∼5%) of them. This is in contrast to the well-established destructive measurements method of absorption imaging. Thus our detection scheme solves a long-sought goal of in situ measurements on ultracold atom clouds without destroying the whole sample. Nonetheless, even an ideal detector introduces unavoidable losses in the sense that it removes particles from the clouds quantum state and projects it onto the eigenstates of the detector.

In optical spectroscopy, electronically excited atoms can be measured with unchallenged sensitivity by photo-ionization and subsequent ion detection [17]. Here we apply this approach to two-photon spectroscopy of ultracold rubidium atoms on an atom chip. The trapped rubidium atoms are excited to the five-dimensional (5D) state by two-photon absorption of light from
Figure 1. Experimental setup of the single-atom detection scheme on an atom chip (not to scale). Below the microchip a cloud of ultracold rubidium atoms is captured in a magnetic microtrap. The atoms get ionized by two overlapping laser beams with wavelengths of 778 and 1080 nm, which are focused through a 200 \( \mu \)m wide hole in the chip. The ions are extracted by ion optics, consisting of three electrodes, and guided to a CEM detector.

A diode laser near 778 nm. An additional fiber laser near 1080 nm subsequently ionizes the excited atoms [18]. The ions are collected by ion optics and guided to a channel electron multiplier (CEM), where they are detected with a probability above 60%. With this setup, we record decay curves of the atom number and complete two-photon excitation spectra with a single cloud of about 10\(^6\) atoms. Collective properties, such as temperature, density distribution and elementary excitations, can be monitored in situ and with high temporal, spatial and energetic resolution. The scheme is particular well suited to the large class of experiments where atoms are trapped in an optical dipole trap with a laser wavelength shorter than 1250 nm such that the trapping laser simultaneously provides the light for the ionization.

The real-time character of our measurement allows us to monitor the density distribution of the atoms in the trap in situ and to determine the trap lifetime, which, however, is limited by off-resonant scattering of the ionization lasers. Applying an additional microwave yields a high energetic and spatial resolution that can be used to measure the energy distribution of the atoms across the trap and therefore the temperature of the cloud.

2. Experimental setup

An illustration of the experimental setup is given in figure 1. The procedure for the operation of the single-atom detector can be separated into three parts. The first part concerns the photoionization process, the corresponding laser system and the interaction with the rubidium atoms. Secondly, ion optics are needed to guide the ionized atoms to a sensitive single-ion detector.
The last part concerns the preparation and manipulation of ultracold atoms at the chip surface, within the detector’s sensitivity volume.

2.1. Photo-ionization of rubidium

Starting with rubidium atoms in the 5S\(\frac{1}{2}\) ground state, a total energy of \(>4.18\) eV [20] per atom is required for ionization. This energy can be provided not only by photo-ionization, but also via field ionization [13] or electron impact ionization [12]. Here, we focus on photo-ionization since this process can be driven efficiently with commercial laser systems.

Figure 2 shows the relevant electronic level scheme of \(^{87}\text{Rb}\) for the multi-photon ionization process. Initially, the atoms are prepared in the 5S\(\frac{1}{2}\), \(F = 2\) hyperfine ground state. Using a laser at 778.1065 nm, the atoms are first excited via a two-photon transition to the 5D\(\frac{5}{2}\) state. The transition rate for this process is resonantly enhanced due to the existence of the 5P\(\frac{3}{2}\) state, which is only 2 nm detuned from the laser wavelength. Once the atoms are in the 5D\(\frac{5}{2}\) state, an arbitrary laser with wavelength below 1250 nm (\(\hbar \omega > 0.99\) eV) is required to ionize the atoms.
Figure 3. Laser setup for rubidium photo-ionization. A master/slave laser system is used for two-photon spectroscopy of the rubidium 5D state. The fluorescence signal is used to stabilize the frequency of the lasers on an appropriate $5S_{1/2} \rightarrow 5D_{5/2}$ transition. To allow a free tuning of the two-photon laser between all $5D_{5/2}$ final states, a second laser near 778 nm is locked onto its optical beating signal with the fluorescence-locked master/slave system. This beam is then combined with the 1080 nm laser beam and guided into the vacuum chamber.

to ionize the atoms. In principle, this final ionization step could be done using the two-photon 778 nm laser beam. However, the $5D_{5/2}$ state has a natural lifetime of 230 ns corresponding to a natural decay rate of $\Gamma = 2\pi \times 0.69 \text{ MHz}$. This rate has to be overcome by the ionization laser, since the natural decay causes heating and atom losses and thus a reduction in the detection efficiency. A high-power fiber laser at 1080 nm is thus used to achieve the desired transition rates to the continuum.

A schematic drawing of the ionization laser system is shown in figure 3. A grating stabilized diode laser in Littrow configuration [22] is used as the reference laser for the 778 nm laser beam. The laser injects a slave laser to provide enough power for the two-photon fluorescence spectroscopy. Therefore, the slave beam is focused and retroreflected through a rubidium glass cell of 10 cm length. To compensate the low two-photon transition rates, the cell is heated to 100 °C, which increases the rubidium partial pressure in the glass cell by three orders of magnitude to about $1 \times 10^{-4} \text{ mbar}$. An additional mu-metal shielding around the cell prevents the spectrum from being influenced by magnetic stray fields. For spectroscopy we detect the blue fluorescence light at 420 nm emitted during the decay cascade $5D \Rightarrow 6P \Rightarrow 5S$ (cf figure 2). Using an optical filter, this light can easily be separated from the exciting 778 nm background light. Besides filtering, the fluorescence light from the spectroscopic cell is collected with a lens and focused onto a photo-multiplier, which allows for small signal detection. A laser power of 40 mW is sufficient to observe reasonable two-photon spectra of the $^{85}\text{Rb}$ and $^{87}\text{Rb}$
5S\(_{1/2}\) ⇒ 5D\(_{5/2}\) absorption lines. We stabilize the reference laser directly to a strong absorption line of the spectrum. To allow a state selective two-photon excitation of the ultracold atoms, we want the 778 nm work beam to be freely tunable within the two-photon spectra. A second grating stabilized laser is thus stabilized via an optical beating signal with the reference laser and used as two-photon laser in the experiment. For beam shaping and timing, the 778 nm beam is sent through a single-mode optical fiber, followed by a programmable shutter.

For the ionization at a wavelength of 1080 nm we use a fiber laser with a maximum power of 4 W (FiboTec, Fiber Laser 1080 nm). The laser with a Gaussian TEM\(_{00}\) beam profile and a spectral width of 1 nm is sent for timing purposes through a mechanical shutter before it is made to overlap with the 778 nm beam. Both beams are then sent into the experiment chamber and focused by an \( f = 40 \text{ mm} \) lens through the hole in the atom chip. Before the beams leave the chamber again, they are re-collimated by an \( f = 80 \text{ mm} \) lens inside the ultrahigh vacuum (UHV) chamber. Measuring the beam diameters at different positions and taking into account the focal length and position of the in-vacuum lenses, we determine the beam parameters of both lasers at the position close to the chip surface where experiments will take place. The beam waist (1/e\(^2\) radius) of the 778 nm laser is 34 \( \mu \text{m} \) with a corresponding Rayleigh length of 4.7 mm. The beam waist of the 1080 nm laser is 26 \( \mu \text{m} \) and the Rayleigh length is 2 mm. Due to the large Rayleigh lengths of both beams, we assume that the atomic cloud is always positioned in the beam waist.

We calculate the laser power required for each of the ionization steps using the low-intensity, single-atom transition rates. For the two-photon transition a general formalism based on a three-state model is given in [23]. In the case of an overall resonant transition and a large detuning \( \Delta \) of the single-photon energy with respect to the intermediate state, the transition rate can be described by the natural decay rate \( \Gamma \) of the final state and the involved Rabi frequencies:

\[
R_1 = \frac{\Omega_{\text{eff}}^2}{4\Delta^2} \cdot \frac{1}{\Gamma} = \frac{\Omega_{\text{eff}}^2}{\Gamma} \sim I_{778}^2.
\]

Here, \( \Omega_{\text{eff}} \) describes the Rabi frequency between the 5S ground state and the 5P intermediate state, and \( \Omega_{\text{mp}} \) the Rabi frequency between the 5P intermediate and the 5D final state. Usually, the single-photon Rabi frequencies are combined to an effective two-photon Rabi frequency \( \Omega_{\text{eff}} = \Omega_{\text{sp}}\Omega_{\text{pd}}/(2\Delta) \), such that the rate in equation (1) has the same form as in the case of a single-photon transition (see the right side of equation (1)). In contrast, the corresponding transition rate depends quadratically on the laser intensity \( I_{778} \). Since the transition rate decreases with increasing detuning to the intermediate state, we restrict our further analysis to the nearest intermediate state, which is the 5P\(_{3/2}\) state at 2 nm detuning. With \( \Gamma = 2\pi \times 0.69 \text{ MHz} \) and \( \Omega_{\text{sp}} \) and \( \Omega_{\text{pd}} \) calculated from the corresponding dipole transition matrix elements, the single-atom, two-photon transition rate at the center of the 778 nm beam becomes \( R_1 \sim 10.07 \text{ kHz} \) (P mW\(^{-1}\))^2 with the laser power \( P \).

For the final ionization step, the transition rate can be described in terms of the ionization cross section \( \sigma_i \) and the photon flux density \( \Phi = I/(\hbar \omega) \):

\[
R_2 = \sigma_i \Phi = \frac{\sigma_i I}{\hbar \omega}.
\]

with \( \omega \) being the laser frequency and \( I \) the intensity. The ionization cross section depends on the laser wavelength and is \( \sigma_i \sim 17.5 \times 10^{-22} \text{ m}^2 \) for the 1080 nm laser light used in the experiment [24]. For the beam waist of 26 \( \mu \text{m} \), the ionization rate at the center of the beam is thus \( R_2 = 8.96 \text{ MHz} \) P/W, corresponding to a transition time of \( \tau \sim 112 \text{ ns} \) (P/W)\(^{-1}\).

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To overcome the spontaneous decay rate of the $5D_{5/2}$ state with $\Gamma = 1/230$ ns, a 1080 nm beam power of at least 0.5 W is required. For optimized detection efficiencies we typically use beam powers above 1 W.

As both lasers off-resonantly couple the $5S$ and the $5P$ states, they cause an ac-Stark shift of the $5S$ ground state and a scattering limited lifetime of the atomic cloud. The ac-Stark shift, also known as light shift or dipole potential, depends on the laser power and reaches its maximum value at the center of the beam. For the blue detuned 778 nm laser, the potential becomes repulsive and the $5S$ energy shift amounts to $\delta \nu_{778} = (-97.6 P) \text{kHz mW}^{-1}$ [25]. The fiber laser in contrast is red detuned, causing an attractive dipole potential with a maximum light shift of $\delta \nu_{1080} = (2.9 P) \text{MHz W}^{-1}$. As the power of the 1080 nm laser exceeds the power of the 778 nm laser by three orders of magnitude, the dipole potential is largely dominated by the 1080 nm laser. This creates an additional dipole trap on top of the magnetic trap with corresponding trap frequencies of $\omega_{\perp} = 2\pi \times 1.45 \text{kHz} \sqrt{P} \text{W}^{-1/2}$ and $\omega_{\parallel} = 2\pi \times 13.6 \text{Hz} \sqrt{P} \text{W}^{-1/2}$ in the directions perpendicular and parallel to the linear polarized beam path. The scattering rates also depend on the beam powers and can be calculated to $\Gamma_{778} = (3.2 P) \text{s}^{-1} \text{mW}^{-1}$ and $\Gamma_{1080} = (0.47 P) \text{s}^{-1} \text{W}^{-1}$. For typically used laser powers of $P_{778} = 300 \mu\text{W}$ and $P_{1080} = 1 \text{W}$, they limit the atomic lifetime to about 1 s. To achieve longer detection times, either the laser powers may be reduced, or the lasers may be chopped with a small duty cycle. The latter method keeps the high laser power during detection and thus allows for a high detection efficiency and a long detection time.

### 2.2. Ion optics and detection

Once the ions are generated they are guided with ion optics to the ion detector. Figure 4 shows a schematic view of the electrode arrangement with respect to the chip surface and the ionization beams together with an image of the final realization in the UHV chamber. A detailed and scaled technical drawing of the ion detector can be seen in figure 5.

The ion optics consist of three electrodes: the extraction electrode, the tube electrode and the deflection electrode. They are used to guide the ions to a CEM (Burle channeltron, model CEM 4502 hc mounted), which is positioned about 55 mm away from the ionization region. The extraction electrode is made of a copper disc with an outer diameter of 8 mm and a central bore of radius 0.5 mm. It is positioned directly under the chip surface, at a distance of 1.6 mm. Setting the electrode on a small negative voltage ($>-100$ V), the ions are attracted and guided through the central bore directly into the tube electrode of 5 mm diameter and 33.5 mm length. Although the inside of the negatively charged tube is field-free, it forms together with the extraction electrode an electrostatic lens for the ions. Following the tube, a deflection electrode is mounted beside the optical beam path at a distance of about 40 mm to the extraction electrode. The deflector is set on a high negative voltage ($\sim -1000$ V) and redirects the charged particle beam out of the optical path onto the ion detector, more precisely onto the mesh in front of the detector, which is set on the tubes voltage. The mesh prevents the influence of the strong electric field of the channeltron on the ion trajectories through the ion optics. Once the ionized atoms have passed the fine grid (transmission: 87%), they are accelerated towards the front side of the channeltron ($\sim -2.3$ kV), where they are counted with single-ion precision. Figure 6 shows a numerical simulation of the charged particle beam for different initial positions of the generated ions. Due to the electrostatic lens effect, the ions are efficiently guided to the channeltron in a compact particle beam allowing for proper illumination of the ion detector.
Figure 4. Electrode arrangement for the ion optics made up of the extraction electrode, the tube electrode and the deflection electrode. The ion optics allow us to directly extract ions from the ionization area close to the chip surface and to guide them through the magnetic coil (around the tube electrode) onto the CEM. To reduce optical stray light, the ionization beams are also guided through the electrodes and take the same path as the ionized particle beam. However, the charged particle beam is bent out of the optical path by using the deflection electrode. The channeltron is used to detect the ions with single-atom resolution.

A simulation of this effect in our setup is visualized in figure 7. If the electrostatic lens is on the same potential as the tube electrode (see figure 7(a)), the lens effect vanishes and the ion beam strongly diverges. Depending on the initial position of the ions, a large number of trajectories end up on the tube electrode and do not reach the channeltron. This would thus result in a low detection efficiency. If, however, the tube voltage is set to $-240 \, \text{V}$, the divergent particle beam becomes collimated and can efficiently be guided to the channeltron (see figure 7(b)). Typically a few microseconds after an atom has been ionized, it reaches the CEM detector, where it causes an electron avalanche inside the CEM tube. Using a combined discriminator and amplifier (Scientific Instruments, model F-100TD), the resulting charge pulse is converted to a TTL-level signal and can be counted with standard detection electronics (Stanford Research Systems, model SR430 Multi-Channel Scaler). The dead time between two successive pulses is limited by the pulse width of the CEM pulses and the pulse pair resolution of the discriminator/amplifier, both below 20 ns. This is true only if the detector is operated in the non-saturated regime below a continuous count rate of $10^7$–$10^8 \, \text{counts} \, \text{s}^{-1}$. All measurements reported here are well below this value.

2.3. Preparation and manipulation of ultracold atoms

The procedure to efficiently load ultracold atoms at an atom chip is described in detail elsewhere [26]. We usually start with a cloud of $3 \times 10^8 \, ^{87}\text{Rb}$ atoms, transferred from a six-beam
Figure 5. Side-view of the ion detector, including the atom chip, the ion optics and the CEM detector. All parts are shown true to scale (units in mm), and the dimensions and relative positions of each individual element are illustrated. The length of the tube electrode is chosen such that the ion beam can be guided through a magnetic coil (not shown). To reduce background counts, the channeltron is placed in a separate housing, allowing ion access only from the direction of the grid.

magneto-optical trap to a conventional magnetic quadrupole trap. After changing the trapping potential to an Ioffe-type trap, the thermal atoms are evaporatively cooled to temperatures below 15 µK. Thereafter, the atomic cloud is moved towards the atom chip and loaded to a chip-based magnetic trap. For this purpose, the chip comprises several micro-patterned wires that allow for trapping of the spin-polarized rubidium atoms at distances up to a few hundred micrometers to the chip surface (see figure 8(a)). At the same time, the complex wire geometry features a three-dimensional positioning system, which allows for sub-µm-precise positioning and momentum control of the trapped atomic cloud. We station the cloud directly above the 200-µm-wide bore in the chip surface where ionization will take place. To avoid the accumulation of laser-generated ions on the surface around the hole, we covered the chip with a thin layer of pure graphite (Acheson, Aquadag). This shortcuts the wires with a resistance of several tens of kiloohms. This is enough to prevent the surface from charging, but does not influence the magnetic potential generated by the currents through the wires. To adjust the position of the cloud with respect to the ionization lasers, we make use of the repulsive/attractive dipole potential of the 778 nm/1080 nm laser. Therefore, we turn on the lasers and take absorption images of the atomic cloud in the combined magnetic and optical trap. Figures 8(b) and (c) show these absorption images for the ionization lasers turned on individually. The images clearly exhibit the position of the repulsive/attractive dipole potentials of the two lasers and allow us...
**Figure 6.** Simulation of the particle trajectories for different starting positions at the chip surface. The ion beam is first accelerated and collimated through an electrostatic lens made up of the extraction electrode and the tube electrode. After passing the tube, the particle beam is redirected by the deflection electrode, to separate the ions from the optical beam path. Subsequently, the ions hit the channeltron.

**Figure 7.** Simulation of the ion beam with and without the electrostatic lens effect. (a) If the extraction electrode is on the same potential as the tube electrode, the ion beam diverges strongly and only a few ions make their way through the tube. (b) For proper voltage at the tube electrode, the particle beam can be collimated and guided without losses through the tube.

to adjust the clouds center to this position. Decreasing the confining magnetic potential to zero in the longitudinal direction removes all atoms that are not trapped by the dipole potential from the 1080 nm laser. The remaining cloud can be seen in the absorption image in figure 8(d).
3. Characterization and calibration of the detector

3.1. Detector calibration

Due to the pulsed source of ultracold atoms on the chip with a duty cycle of about 70 s and relatively small ion numbers, it would have been a tedious process to optimize the ion optics voltages. For that reason and to calibrate the efficiency of the ion detector, a separate experiment was set up in a low-vacuum chamber ($2 \times 10^{-6}$ mbar) [19]. To determine the ion detection efficiency, we counted the ions and the electrons emerging from the ionization process simultaneously in different CEMs and recorded their coincidence signal.

We started by evaporating atoms from a continuous $^{87}$Rb dispenser source and photoionized them by a three-photon process within a standing wave optical resonator. The ions were collected by the same ion optics as those used in the ultracold atom experiment. The cavity enhanced laser beam is positioned at a distance of a few millimeters to the ion optics and aligned parallel to the extraction electrode. Opposite to the electrode, a grounded aperture with a diameter (0.5 mm) smaller than the one of the extraction electrode (1 mm) is installed. Due to
this circumstance, only electrons emerging from the ion detection volume can pass through the aperture. They impact on an aluminum plate behind the aperture, creating secondary electrons that are counted by a second CEM detector. Comparing the delay time between the ion and electron CEM signals, a clear Gaussian distributed coincidence curve was measured with a mean ion delay time of 3.45 µs and a standard deviation of 50 ns.

The number of detected ions/electrons in the CEMs is given by 
\[ N_i, e = \eta_i, e N \]
with \( N \) being the total number of ionized atoms and \( \eta_i, e \) the ion/electron CEM detection efficiency. The number of ion–electron coincidences is \( N_c = \eta_i \eta_e N \). Substituting for \( \eta_e \) yields \( \eta_i = N_c / N_e \). Therefore, after the correction of background electrons and false coincidences, the ion detection efficiency can be evaluated by dividing the measured coincidence and electron counts \[ 28 \]. The advantage of this technique is that there is no need to know \( \eta_e \). Using this method the total efficiency of our ion detection system was determined to be 47.8±2.6%.

We were also able to measure the detection efficiency with a different method in the main experiment where the ultracold atoms are ionized in the UHV chamber. The total number of ions counted in the CEM was compared with the number of atoms deduced from absorption imaging \[ 11 \] (see section 3.2). A detection efficiency of 67±12% was determined, which is, obviously, significantly higher compared to the value in the calibration, low-vacuum setup. We assign the deviation to the different starting velocities of the ions in the detection region. Whereas the rubidium atoms in the cavity setup directly emerge from a dispenser with a temperature of some 100°C before they get ionized, the ultracold atoms in the main UHV experiment have a temperature of the order of 10 µK. A higher temperature leads to a larger velocity spread and to a broader range of trajectories through the ion optics. This spreading results in a broader channeltron illumination. According to \[ 29 \], the efficiency of a channeltron depends on the impact region of the ions. On the outer horn, it is usually between 50 and 60% at a voltage of −2900 V. In the channel at the center, the ions have a lower impact angle relative to the surface, leading to a higher rate of secondary electrons and an efficiency of up to 80%. From ion optics simulations we conclude that almost all ions within the detection volume are guided to the CEM. This cylindrical volume is defined by the direct projection of the extraction electrode aperture on the chip surface. Since we collect all the ions in this region, the detection efficiency is mainly limited by the efficiency of the CEM. In the case of ultracold ions, we optimized our ion lenses to focus on a small spot on the area with the highest efficiency in the middle of the funnel. However, in the cavity experiment with hot atoms, the CEM is illuminated on a larger area due to the velocity spread, causing a decrease in detection efficiency.

A complete scan of the normalized ion detection rate of the hot cavity ions as a function of the tube electrode voltage \( U_t/U_e \) and the deflection electrode voltage \( U_d/U_e \) is given in figure 9. The ion trajectories and therefore the detection rate depend only on the ratios of the electrode voltages and are shown with a fixed value for the extraction electrode voltage \( U_e = −40 \) V. However, we noticed that this value should not be lower than −40 V to overcome stray fields in the ionization region \[ 19 \]. We have also included the optimized voltage ratios chosen in the ultracold atom chip experiment in figure 9 (blue dot: \( U_t = −126 \) V, \( U_d = −760 \) V, \( U_e = −56 \) V).

According to \[ 28 \], atom counting with sub-Poissonian precision is possible if the efficiency of the detector is well above 50%. This requirement is already achieved in the current work. By the implementation of a conversion dynode with high secondary electron emission, a detection efficiency close to 100% is possible and would be an interesting improvement for atom number squeezing experiments.
Figure 9. Normalized ion CEM signal measured in the low-vacuum calibration setup. The signal is plotted as a function of the tube electrode voltage $U_t/U_e$ versus the deflection electrode voltage $U_d/U_e$. The data are recorded at $U_e = -40$ V, although essentially the same dependence was observed for other extraction electrode voltages. Numerical simulations predict the impact of the ions on the CEM in the area between the dashed white lines. The blue dot represents the voltage ratio used in all measurements described in this paper with ultracold atoms on the chip.

3.2. Detection efficiency

The detection efficiency of the ion detector in the ultracold atom experiment can be determined by comparing the ionization signal of a well-defined cloud with an atom number measured from absorption imaging. Therefore, we need to prepare a defined number of atoms in the ionization volume, which can then be ionized on a timescale faster than other loss mechanisms. For that reason, as a first step, a cold thermal cloud is prepared in the microtrap. The 1080 nm laser is then linearly ramped up within 1 s to a power above 1 W increasing the atomic density at the center of the trap due to the dipole potential (see section 3.3). The surrounding atoms still trapped in the magnetic potential in the longitudinal direction are removed by opening the axial confinement ($x$-axis in figure 1) of the magnetic trap within 100 ms. Enclosure in the radial direction is still provided by the quadrupole microtrap. The number of atoms, now all trapped within the laser beam and therefore in the ionization volume (see figure 8(d)), is of the order of a few ten thousands and can be well determined by absorption imaging after 100 $\mu$s time-of-flight. In order to count the number of trapped atoms with our single-atom detector, the 778 nm laser is turned on with a power of 316 $\mu$W immediately after the axial confinement of the magnetic trap is opened. While still trapped in the dipole potential, the cloud is ionized within 300 ms.
Comparison of the total number of counts with the number deduced from absorption imaging results in a detection efficiency of 67±12%.

3.3. Efficient loading into the detection volume

We demonstrate the single-atom detector by preparing a cloud of several $10^5$ rubidium atoms below the hole in the chip. To ensure high detection efficiency the laser powers are chosen to be 440 µW and 1.6 W for the 778 nm laser and the 1080 nm laser, respectively. Once the 778 nm laser is tuned resonantly to the $5S_{1/2} \rightarrow 5D_{5/2}$ transition and both lasers are activated, the single-atom detector starts counting individual ionization events. Summing up the events over a given time window, we end up with typical ionization curves as shown in figure 10. However, due to the large dipole potential of the 1080 nm laser, the ionization signal strongly depends on the power up process of the fiber laser. Figure 10(a) shows the ionization signal for three different start up schemes of the 1080 nm laser. If both lasers are turned on instantaneously, the ionization rate shortly increases to its maximum value and subsequently decreases exponentially on the timescale of the atomic lifetime, which is $\sim 750$ ms (see blue curve (i) in figure 10(a)). An efficient loading of the attractive dipole dimple can be achieved by turning on the 1080 nm laser adiabatically, i.e. within a 1 s linear ramp. This increases the maximal counting rate by more than a factor of two, as compared to the non-adiabatic case (see red curve (ii) in figure 10(a)). The adiabatic switching thus strongly increases the atomic density inside the ionization volume. After the activation of the 778 nm laser, the atoms loaded on the dipole trap are ionized on a fast timescale of $\sim 25$ ms, corresponding to the single-atom ionization rate. Thereafter, we observe a second decay on a much longer timescale of again $\sim 750$ ms, which is the same as in the non-adiabatic case. This long time ionization signal results from magnetically trapped atoms passing the ionization volume. The corresponding timescale is given by the cloud lifetime in the combined magnetic and optical trap, mainly limited by ionization losses and off-resonant scattering due to the ionization lasers. To demonstrate off-resonant scattering losses, the green curve (iii) in figure 10(a) shows the ionization rate for non-adiabatic laser switching, whereas the 1080 nm laser was activated 1 s before the diode laser. As compared to the simultaneous laser activation (blue curve (i)), the measurement shows the same behavior, but with an ionization rate reduced by about one third. These losses are purely due to the fiber laser off-resonantly coupling the rubidium 5S and 5P states.

We repeated the experiment for different temperatures of the thermal cloud and observed that the relative gain in density for the adiabatic loading process compared with the non-adiabatic loading is reduced for increasing cloud temperature. This is obvious, since the adiabatic loading efficiency decreases with the atoms temperature for a dipole trap of given depth. For illustration, figure 10 shows the same measurements for two different cloud temperatures of 6 µK (a) and 12 µK (b). While in the first case, the adiabatic loading increases the atomic density in the ionization volume by more than a factor of two as compared to the non-adiabatic case (cf figure 10(a)), the hotter cloud shows a remaining gain of only 1.3.

3.4. Transition rates of multi-photon excitations

To achieve broad usability of our detector for various applications, the dependence of the ionization rates on the laser powers needs to be studied. Experiments for which only the total atom counts are relevant require e.g. transition rates much higher than the overall loss rate.
Figure 10. Ionization rates for different activation schemes of the lasers (i–iii) and cloud temperatures of (a) 6 µK and (b) 12 µK. If the 1080 and 778 nm lasers are activated instantaneously (blue curve and inset (i)), the counting rate shows a single exponential decay on the timescale of the atomic lifetime in the combined magnetic and dipole trap ($\tau \sim 750$ ms). A slow activation of the fiber laser in the form of a linear ramp (red curve and inset (ii)) results in an adiabatic loading of the attractive dipole potential and thus increases the maximum counting rate. Subsequently, the rate drops on a double exponential scheme: firstly, the atoms adiabatically loaded into the dipole trap are ionized on a fast timescale ($\tau \sim 25$ ms). Secondly, the remaining part of the cloud is again ionized on the long timescale of the trap lifetime ($\tau \sim 750$ ms). If the 1080 nm laser is switched on instantaneously but 1 s before the 778 nm laser (green curve and inset (iii)), we observe the same behavior as in the case of simultaneous switching (blue curve (i)) but with a strongly reduced counting rate.

from the trap. However, in order to make use of our detector’s big advantage, the possibility to measure atomic properties \textit{in situ} by sampling only a small subset of the whole system, small transition rates are crucial. The excitation rates of the 778 and 1080 nm lasers are thus important quantities to measure.

To determine the excitation rates of the 778 nm laser, a cold ensemble of atoms is loaded into the dipole trap, in the same manner as stated in section 3.2. The cloud is now completely within the ionization volume. For sufficient power of the 1080 nm laser, the loss rate after activating the 778 nm laser is thus given by the excitation rate to the 5D state. From the exponential decay of the ionization signal, this rate is deduced for different powers of the 778 nm
Figure 11. Ionization rates for different 778 nm laser powers. The atoms are initially in a dipole trap created by the 1080 nm laser with a power of 0.8 W. Sudden switch-on of the 778 nm laser leads to an exponentially decaying signal of the ion count rate. The losses are mainly caused by the excitation rate to the 5D state and depend strongly on the power of the 778 nm laser. Typical decay curves with respective decay times $\tau$ are shown: for a power of 100 $\mu$W in (a) and for 300 $\mu$W in (b). The corresponding decay times differ by a factor of 8.6. (c) Since the $5S_{1/2} \Rightarrow 5D_{5/2}$ transition is a two-photon process, the decay rate depends quadratically on the 778 nm laser power. The measured decay rates are thus well fitted by a quadratic power dependence.

laser and plotted in figures 11(a) and (b). As expected for a two-photon transition, the rate increases quadratically with the intensity of the exciting laser (cf equation (1) and figure 11(c)). The decay rates are basically the same, as have been observed in section 3.3 after adiabatically loading the dipole trap (red curve (ii) in figure 10). The fast decay time is identical to the decay time observed in the pure dipole trap.

The natural lifetime of the 5D state is 230 ns, limited by the decay to the 6P and 5P states. The rate at which the 1080 nm laser excites the atoms from the 5D state has to overcome this natural decay in order to achieve a good detection efficiency. To determine this condition, a cold cloud of atoms is constantly irradiated by the 778 nm laser at a power of 90 $\mu$W, while the power of the 1080 nm laser is linearly ramped within 200 ms to the maximum value of 1.6 W. In order to compensate for the ac-Stark shift of the 1080 nm laser, the frequency of the 778 nm laser is simultaneously swept by an amount of 2.18 MHz. The ionization signal is recorded for a total
Figure 12. Recorded ion signal and loss corrected ion signal for varying fiber laser power. (a) The 1080 nm laser is ramped within 200 ms to 1.6 W and kept constant thereafter. The red dashed line shows an exponential decay fit from which the trap lifetime and therefore loss rates are deduced. (b) The trap loss corrected curve shows saturation characteristics. A power above 1 W seems sufficient to saturate the ionization process.

of 330 ms, whereas the last 130 ms are measured at constant powers and frequencies of both lasers. This additional 130 ms gives access to the overall loss rates from the trap. The ionization signal is plotted in figure 12(a). Figure 12(b) shows the count rate dependence from the 1080 nm laser power, which is already corrected for the atomic losses. A 1080 nm laser power above 1 W thus seems to be sufficient to overcome the spontaneous decay and to saturate the ionization process. This is in good agreement with theoretical calculations in section 2.1, which showed that at 0.5 W the transition rate to the continuum equals the natural decay rate.

4. Single-atom detector applications

4.1. Monitoring the atomic density distributions

As a first application of our single-atom detector, we measure the atomic density distribution of a trapped atomic cloud on the chip. Therefore, we move an ultracold ensemble through the ionization lasers and monitor the corresponding ion count rate. We start with a cloud of $\sim 1 \times 10^6$ atoms at a temperature of $10 \, \mu$K prepared in a microtrap next to the ionization region. Subsequently, the 778 nm laser is tuned to resonance at a power of 90 $\mu$W, while the 1080 nm laser is set to 0.5 W. While monitoring the ion signal, we apply a homogeneous magnetic field to move the atoms in the radial ($y$) direction (see figure 1) within 200 ms through the ionization beams. Figure 13 shows the corresponding ion counts in the detector. The shift of the trap as a function of the applied field can be simulated numerically and provides the position calibration in the plot. The signal is thus a scan of the atomic density distribution along the radial direction of the cloud. Knowing the temperature, this measurement allows for the determination of the radial trap frequency and vice versa. For the 10 $\mu$K cold cloud in figure 13, we deduced a trap frequency of $\omega_r = 2\pi \times 155$ Hz. However, the drawbacks of this method are possible excitation of oscillations and the change of the trap geometry during the movement.
Figure 13. Scan of the atomic density via laser ionization of a moving cloud. Within 200 ms the cloud is shifted in the axial direction through the ionization volume. The signal is fitted by a Gaussian distribution (red line). From the width of the Gaussian, the radial trapping frequency can be deduced to be $\omega_r = 2\pi \times 155$ Hz for a thermal cloud with a temperature of 10 $\mu$K. The two-peak structure within the data can possibly be related to excitations caused by the movement of the cloud.

to simulations, the radial trap frequencies in the demonstrated measurement vary by less than 30%. This limits the precision in the temperature determination.

4.2. Energy selective detection and determination of the clouds temperature

An elegant and precise way to measure the density distribution and even the energy distribution of the atoms across the trap could be by adding an additional microwave field to the ionization scheme [11]. This is done by irradiating a tuneable microwave near 6.8 GHz, driving the transition $5S_{1/2}, F = 2, m_F = 2 \Rightarrow 5S_{1/2}, F = 1, m_F = 1$ and simultaneously ionizing the atoms from the $F = 1$ state. The process is illustrated in figure 14(a). The microwave is coupled into the chamber via a helix antenna providing a Rabi frequency of $\Omega = 2\pi \times 23.7$ kHz at the position of the trap. Due to the Zeeman shift of the magnetic sub-levels within a hyperfine manifold, the transition frequency is dependent on the position in the trapping potential. The resonance condition for the microwave is given by

$$\mu_B B(x, y, z) = \frac{1}{3}\hbar \Delta \omega,$$

with $\Delta \omega$ being the microwave frequency detuning (with respect to the resonance frequency at the bottom of the trap), $\mu_B$ the Bohr magneton and $B$ the local magnetic field. This defines an ellipsoid around the center of the trap from which the atoms are excited. Ionization, however, occurs only at the areas where the resonance shell intersects with the laser beams. For small detunings $\Delta \omega < 2\pi \times 1.3$ kHz the ellipsoid is completely covered by the laser beams. For large detunings $\Delta \omega > 2\pi \times 0.25$ MHz, however, the resonance shell leaves the ionization volume in two dimensions and reduces the intersection area to two single planar sheets located at $z = \pm \sqrt{4\hbar \Delta \omega / (3m\omega_t^2)}$. If we neglect the axial confinement of the dipole potential caused by the
Figure 14. Energy selective ionization for temperature determination of the ultracold atomic ensemble. (a) Sketch of the $^{87}$Rb $5S_{1/2}$ hyperfine ground state with Zeeman splitting in the harmonic magnetic trap. The blue arrow indicates the microwave coupling between the $m_F = 2$ and $m_F = 1$ states. From the $m_F = 1$ level the atoms are excited to the 5D state by the 778 nm laser (red arrow). (b) Ionization signal of the atomic energy distribution in the microtrap. The atoms were initially prepared in the $F = 2$, $m_F = 2$ state. Tuning the microwave frequency allows us to monitor the atomic energy distribution across the trap. Starting from high-frequency detunings, the ion count rate increases exponentially while moving towards the bottom of the trap. A sudden drop is observed at this point and defines the magnetic trap offset field $B_0$. The red line shows a fit with the Boltzmann factor in equation (4), revealing a cloud temperature of 9.6 $\mu$K. Due to residual atoms in the $F = 2$, $m_F = 1$ state, we additionally observe corresponding hyperfine transitions, as indicated in the plot. The frequency axis origin is set to the $F = 2$, $m_F = 0 \rightarrow F = 1$, $m_F = 0$ transition.

1080 nm laser, the potential energy at the position of the sheets is proportional to the microwave transition energy. The atomic density distribution function imaged in the ionization signal can be expressed in the form [11]:

$$n(\Delta \omega) = \frac{1}{\sqrt{\Delta \omega}} \exp \left( -\frac{2\hbar \Delta \omega}{3k_B T} \right).$$  \hspace{1cm} (4)

For large detunings, the behavior of this function is purely exponential and allows for an *in situ* temperature determination.

We demonstrate this method by tuning the 778 nm laser to the transition $5S_{1/2}$, $F = 1 \rightarrow 5D_{5/2}$, $F' = 1$. After ramping the 1080 nm laser within 500 ms to a power of 0.2 W, we wait 500 ms and turn on the 778 nm laser with a power of 300 $\mu$W. Simultaneously, the microwave is switched on and tuned within 320 ms from 6.841 GHz to 6.835 GHz. The resulting ion signal is exhibited in figure 14(b). The data can be fitted by the exponential part in equation (4), yielding a cloud temperature of 9.6 $\mu$K. This is in good agreement with time-of-flight measurements.
The advantage of this method is the possibility to measure the temperature and energy distribution of a trapped atomic cloud by sampling only a small subset of the original ensemble. This can be achieved by adjusting the lasers to low powers. The remaining cloud can still be used for further experiments, which is not possible in time-of-flight measurements via absorption imaging.

4.3. Spectroscopy of ultracold atoms

With the single-atom detector on the microchip it is possible to perform optical spectroscopy directly and in situ on the ultracold rubidium cloud [11]. Initially a cloud of $1.8 \times 10^6$ $^{87}$Rb atoms was prepared at a temperature of $18 \mu$K in the $5S_{1/2}$, $F = 2$ state. We turned on simultaneously the 778 nm excitation and 1080 nm ionization lasers and waited 500 ms for the atoms to thermalize into the suddenly generated strong attractive dipole potential from the ionization laser. Then we scanned the 778 nm laser over a frequency range of 65 MHz, at a power of 270 $\mu$W and a rate of $d\nu/dt = -45$ MHz s$^{-1}$. Due to the two-photon transition, a frequency detuning $\Delta \nu$ of the 778 nm diode laser corresponds to a detuning of $2\Delta \nu$ with respect to the $5S$–$5D$ level spacing. For that reason the actual scan covers a frequency range of 130 MHz, including all hyperfine resonances in the $5D_{5/2}$ state. The atoms are excited each time the frequency reaches resonance (see figure 15(a)), revealing clear spectral line profiles. By scanning over all resonances and ionizing from the $5D$ states with the 1080 nm laser at a power of 0.6 W, the spectrum in figure 15(b) was recorded. It includes a total number of $6 \times 10^4$ ion counts. Taking into account a detection efficiency of 67%, this number corresponds to only 5% of the initial atoms. This opens up the possibility to perform further experiments on the same cloud.

In principle, the ionization from the $5D$ state can also be performed by a third photon from the 778 nm laser without the use of the 1080 nm laser. We demonstrate this with a 778 nm laser power of 28.6 mW and the resulting three-photon spectrum is exhibited in figure 15(c). Since the transition from the $5S$ to the $5D$ state is already saturated at low 778 nm intensities in the mW regime, the ionization signal is strongly dependent on the power of the ionizing laser, which should be of the order of 1 W. For that reason, the count rate is significantly decreased due to the low 778 nm diode laser power available for the excitation to the continuum, making this method rather inefficient. Nevertheless, the two-photon peaks in the spectrum are still clearly resolved. As compared to the measurement in figure 15(b), the ion count rates are reduced by a factor of about 20. This corresponds well to the 20 times smaller ionization power applied to record the spectra shown in figure 15(c).

By irradiating a microwave with a frequency of 6.8373 GHz, the atoms are transferred from the $5S_{1/2}$, $F = 2$ to the $5S_{1/2}$, $F = 1$ ground state. From this level they are excited to the $5D$ hyperfine levels again by the absorption of two photons. We scanned the 778 nm diode laser frequency for this process within a range of $\Delta \nu = 40$ MHz to occupy the $5D_{5/2}$, $F = 2$, 3 and 4 levels before the atoms are again ionized by the 1080 nm laser. The resulting ionization spectrum is illustrated in figure 15(d). Due to the selection rule for the two-photon process $\Delta F = 0, \pm 1$ and $\pm 2$, only three spectral peaks are visible in this excitation, starting from the $F = 1$ ground state, in comparison to the four peaks in figures 15(b) and (c), starting from $F = 2$.

To determine the influence of the fiber laser power, we recorded two-photon spectra, such as the one illustrated in figure 15(b), for ten different powers for the 1080 nm laser between 80 mW and 1.6 W, whereas the 778 nm laser was kept constant at a power of 270 $\mu$W. For each
Figure 15. (a) Sketch of the three methods to record an ionization spectrum of the \(^{87}\)Rb \(5S \Rightarrow 5D\) hyperfine transitions ionizing only a fraction of the ultracold cloud. The atoms are prepared in the \(5S_{1/2}, F = 2\) state and get excited by a tunable 778 nm laser via a two-photon process into the \(5D_{5/2}\), \(F' = 1, 2, 3\) and 4 states (blue arrows). From there they get ionized by a 1080 nm laser (gray arrow). The resulting (blue) spectra are shown in (b) together with four Gaussian fits (red) to determine the linewidths and positions of these transitions. (c) It was possible to ionize the atoms in the \(5D_{5/2}\) states after the two-photon excitation by a third 778 nm laser photon without irradiating the 1080 nm laser (red arrows in (a)). Due to the low diode laser power, this process is less efficient, but the corresponding (red) spectrum can still be resolved well. (d) Irradiating a 6.8373 GHz microwave transfers the atoms from the \(5S_{1/2}, F = 2\) to the \(5S_{1/2}, F = 1\) ground state. By adjusting the two-photon transition wavelength of the 778 nm laser, the \(5D_{5/2}, F' = 1, 2\) and 3 states could be excited and ionized again by the 1080 nm laser (green and gray arrows in (a)), as can be seen from the three (green) ionization peaks. The origin of the frequency detuning is chosen arbitrarily and without relation between each spectral image.
Figure 16. (a) Frequency shift of the peak positions in the $^{87}\text{Rb} \, 5S_{1/2} \Rightarrow 5D_{5/2}$ ionization spectra for different 1080 nm fiber laser powers due to the ac-Stark effect. (b) Power broadening (FWHM) of the spectral linewidths due to an increased ionization rate, thus lowering the 5D lifetime. An additional contribution arises from an inhomogeneous light shift due to the Gaussian beam profile of the fiber laser. The experimental data (blue dots) were fitted with a linear fit (red line).

illustrates the 1080 nm laser power broadening of the spectral linewidths. The full-width at half-maximum (FWHM) increases by $3.0 \pm 0.4 \text{ MHz W}^{-1}$. This is partly due to an increased ionization rate, thus lowering the effective 5D lifetime ($\sim 0.7 \text{ MHz W}^{-1}$), and to the Gaussian beam profile of the fiber laser, leading to an inhomogeneous light shift and therefore a power-dependent line broadening. A zero power linewidth of $2.6 \pm 0.4 \text{ MHz}$ is extrapolated from the data. There are three contributions to this value: the natural transition linewidth ($\sim 0.33 \text{ MHz}$), the linewidth of the diode laser ($\sim 1.5 \text{ MHz}$) and the inhomogeneous Zeeman broadening ($\sim 0.5 \text{ MHz}$) arising from the thermal distribution of the atoms in the magnetic trap.

As discussed previously, the photo-ionization scheme can easily be made energy and space selective by coupling an additional microwave to the ionization process. Changing the microwave frequency then allows one to adjust the area of particle detection within the atomic ensemble. This spatial sensitivity in combination with the nanosecond temporal resolution of the CEM detector opens the door to new correlation measurements within a cloud of ultracold bosonic or fermionic particles.

5. Outlook and conclusion

We have demonstrated single-atom detection by means of directly loading a cloud of ultracold $^{87}\text{Rb}$ atoms on an atom chip and ionizing the single atoms. This was achieved by a combination of two laser beams, the first one stimulating a two-photon process to an excited level and the second one ionizing the atoms. We described the design of optimized ion optics to efficiently extract the ions from the cloud and guide them into a channeltron detector. These were tested in a separate setup, where the efficiency of this mechanism was additionally obtained by an ion–electron correlation measurement. Comparing the data with absorption images of the
ultracold cloud, a detection efficiency up to $67 \pm 12\%$ was extracted. The cloud was loaded from the magnetic trap below the surface of the microchip into a dipole trap created by the ionizing fiber laser. We studied experimental decay curves due to ionization for different laser powers, proving the single-atom resolution in our measurements.

Another important application of this setup is in performing high-resolution optical spectroscopy on the cloud \textit{in situ} and in real time, ionizing only about 5\% of the atoms. This was realized by scanning the two-photon excitation wavelength over the hyperfine levels in the excited state. With increasing laser power, the spectral lines get shifted linearly to higher frequencies by the ac-Stark effect and show power broadening. Introducing an external microwave frequency, the energy distribution of the atoms can be determined with high precision. This allows for the measurement of the temperature and the local density of the atomic ensembles with only little ionization losses.

Quantum correlations and atom number statistics are important entities for describing and understanding quantum degenerate gases [31–33] and controlled quantum phase transitions [1, 3]. From that perspective, applications of our detection technique in such experiments could be crucial in various fields of research. Although second-order correlation measurements have already been realized in different systems [4–6], our setup allows for \textit{in situ} observations of trapped quantum gases and phase transitions and, more important, for a local correlation measurement within the atomic ensemble. Having a Bose–Einstein condensate, this enables us to measure particle correlations not only within the condensed part and the surrounding thermal cloud but also in the intermediate region. This becomes especially interesting, since just recently different theoretical approaches emerged for describing combined systems of condensed and thermal fractions [34–36].

Such second-order correlation measurements require us to reduce the detection volume of the ionization-based detector to the lengthscale of the atomic coherence length. Within a Bose–Einstein condensate, this coherence length is relatively large and is usually given by the size of the condensate. Moving towards the surrounding thermal cloud, however, this coherence length rapidly shrinks to the size of the corresponding de-Broglie wavelength, which in the case of a 100 nK cold cloud decreases to $\sim 0.6 \mu$m. Proper local correlation measurements thus require one to minimize the detection volume below $1 \mu$m$^3$. The minimum detection volume of our microwave enhanced detector is given by the microwave Rabi frequency and the magnetic Zeeman shift across the atomic cloud. By choosing appropriate field gradients of typically $1 \text{T m}^{-1}$ and Rabi frequencies of about $2\pi \times 1 \text{kHz}$, the spatial detection volume along one dimension could be reduced to $\sim 50 \text{nm}$. The other directions might be limited by reducing the beam waist of the 778 nm laser. For a waist of $\sim 2 \mu$m, the spatial detection volume would hence be limited to about $2 \mu\text{m} \times 2 \mu\text{m} \times 50 \text{nm} = 0.2 \mu\text{m}^3$, fulfilling the mentioned requirements. Concerning coherence times, our calculations also show that, even for a 100 nK cold cloud, the coherence time is at least two orders of magnitude larger than our single-atom detection rate. The microwave enhanced detector would thus allow us, in a unique way, to locally measure particle statistics within a given coherence length and time.

The dynamical properties of our detector offer a promising approach to studying the behavior of atomic and molecular quantum gases in real time. This could also be useful in the investigation of the recently experimentally verified Efimov states [37] by recording three-body decay rates. Our single-atom detector could be an important constituent for the possible construction of a high-resolution quantum sensor for gravity gradients. Imagine a periodic surface potential on the microchip reflecting or transmitting a single-atom wave running along a
magnetic waveguide with a certain velocity and therefore de-Broglie wavelength. Extremely small tilts of the chip could be noticed when measuring the transmitted single-atom signal behind the periodic barrier. Modification of such a setup could also be adopted for quantum computation schemes on atom chips [38]. A still very young scientific topic is the study of atom–ion mixtures [39]. With our setup it is possible to ionize extremely small quantities of atoms within the cloud. The influences of their polarizing electric fields on the surrounding neutral atoms and the possible change in the interaction energy of the whole quantum gas are fascinating topics for future investigations.

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