

Continuous-time ultrahigh-frequency sensing using cold Rydberg atoms

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We present a technique for detecting ultrahigh-frequency (UHF) radio fields using a three-photon Rydberg excitation scheme in a continuously laser-cooled sample of ^{87}Rb atoms. By measuring Autler-Townes splitting, we demonstrate resonant detection of UHF fields with frequency range around 500–900 MHz through probing $nF \rightarrow nG$ transitions, achieving a lowest minimum detectable field of 2.5(6) mV/cm at 899 MHz. We also demonstrate continuous-time detection of a modulated rf signal, with a 3-dB bandwidth of 5.7(4) kHz at a carrier frequency of 899 MHz. Our approach employs atom loss spectroscopy rather than electromagnetically induced transparency, which enables detection while the atomic sample is simultaneously laser-cooled. We investigate this operating regime to determine the feasibility of combining the benefits of reduced thermal dephasing (and subsequent increased sensitivity) of cold atoms with the continuous operation associated with thermal atoms. Our continuous-time detection scheme provides an advantage over existing pulsed cold atom systems as we avoid the slow experimental duty cycles typically associated with replenishing the cold atom ensemble. We characterize the excitation scheme by varying laser detunings and analyze the impact and limitations due to various broadening mechanisms on the detection sensitivity.

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I. INTRODUCTION

Rydberg atoms are of great interest for quantum sensing applications due to the large transition dipole moments between nearby Rydberg states, exaggerated interaction strengths, and large polarizabilities [1,2]. Rydberg electrometry systems have the potential to provide better sensitivity [3] and smaller form factors than traditional dipole antennas [4,5]. This is of particular interest in the ultrahigh-frequency (300–3000-MHz) regime, where the radio frequency (rf) wavelength becomes of the order of 1 m, limiting the minimum size of a dipole antenna if high performance is required. Applications employing frequencies within this band are extensive, including radio astronomy [6] and broadcast; satellite navigation [7] and sensing for Earth observation [8,9]; and aeronautical communications [10,11] and mobile telemetry [12]—where detection technology could especially benefit from miniaturization [13].

One typical Rydberg electric field sensing modality [3, 14] employs an electromagnetically induced transparency (EIT) Autler-Townes (AT) technique, whereby two laser fields are used to couple the ground state of an atom to an nS/nD Rydberg state, generating an EIT signal [15]. The signal is detected via the transmission of the ground-state transition laser (probe) through the atomic sample. An rf field resonant with a nearby Rydberg state dresses the initial Rydberg state, resulting in an AT splitting of the EIT signal, of magnitude proportional to the strength of the electric field. This electric field strength is related to the measured splitting only via atomic constants, which can be calculated precisely [16], and hence Rydberg systems are excellent candidates for “calibration-free” electric field metrology.

To construct an rf sensor with this method, a Rydberg state is chosen such that there is a neighboring state whose transition frequency matches the desired radio frequency. Commonly, radio frequencies around 10–100 GHz are detected [3,14,17–20]. In order to detect lower frequencies, various resonant [21] and nonresonant methods [22,23] have been demonstrated. For resonant excitation, one can excite to higher n to couple to high-sensitivity Rydberg transitions with lower radio frequencies [21]. This is possible as the energy-level spacing between pairs of neighboring Rydberg states scales as n^{-3} . Unfortunately,

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exciting to high n comes with a loss of coupling strength, due to the transition dipole moments coupling from the lower to the Rydberg states scaling as $n^{-3/2}$ [1]. In order to compensate for this, the Rydberg laser intensity must be increased as n is increased, which becomes unfeasible at the n required to access UHF transitions. At high n , the exaggerated scaling laws can also become unfavorable for rf detection—for example, the increased polarizability (proportional to n^7) can lead to unwanted perturbations due to stray electric fields, and the increased collisional cross section can introduce additional dephasing due to interactions with other atoms [24–26].

Nonresonant rf detection methods typically require large E-field strengths that mix the involved states [22,27,28] and/or application of additional electric fields for homodyne/heterodyne measurement [19,23,29]. These methods introduce additional complexity into the system. To model such systems, in order to extract physical parameters, one must go beyond the rotating-wave approximation using Floquet analysis [28]. Demodulation of the detected signal may also be required to extract information, for example, in communications protocols [26,30–32]. The use of additional electric fields also moves the system away from SI traceability, as the additional fields must be calibrated themselves.

An alternative method to allow for resonant detection of low-frequency rf fields is to excite to higher angular momentum states by introducing either additional rf fields [33,34] or additional laser fields [35,36]. The energy-level spacing is smaller for larger l so the n required to reach nearby low-frequency transitions is reduced. In this work, we employ this method by use of an additional (coupling) laser field to reach UHF transitions, in a cold atomic sample.

The majority of previous work has been performed in thermal vapor systems, where the sensitivity can be limited by the residual Doppler effect. Doppler broadening can be minimized through careful k -vector matching [37], but remains a limitation. On the other hand, cold atomic systems do not suffer from this effect to first order as the velocity profile is reduced by laser cooling. This reduced velocity profile also enables the probing of a larger proportion of atoms in the sample, as they are largely in the same velocity class and so fewer atoms are dark due to the Doppler shift. This is in contrast to thermal vapor systems where the velocity-selected atomic excitation fraction is typically between 10^{-3} and 10^{-2} [35].

Some work has been done using cold atoms for sensing, such as Refs. [38] and [39], with minimum detectable fields of the order of $100 \mu\text{V cm}^{-1}$ at around 10 GHz. These systems suffer, however, from slow experimental duty cycles between sample preparation and detection. They cool and prepare the atomic sample over many seconds, and then perform rf detection in a single “shot” before repeating the process. This limitation excludes

this technique from many real-world applications where detection must be performed continuously or where data are temporally encoded in signals.

Here, we introduce a detection technique that employs a three-photon excitation scheme to access the UHF regime in a continuously cooled ^{87}Rb sample. The rf field is applied to the $nF \rightarrow nG$ transitions, such that frequencies below 3000 MHz can be detected at $n > 30$. We perform UHF detection using atom loss spectroscopy, whereby the splitting of the Rydberg energy levels is observed via loss of atoms from the cold sample. Crucially, this scheme allows for high bandwidth detection compared to previous cold detection schemes, demonstrating a potential modality for Rydberg rf sensors.

The structure of the paper is as follows. In Sec. II, we describe the experimental setup used to prepare the cold atomic sample. We also describe the Rydberg excitation scheme and the atom loss spectroscopy method. In Sec. III, we demonstrate UHF detection below 1 GHz, first of a static carrier frequency, and then of an amplitude-modulated signal. In Sec. IV, we discuss the limiting factors for the sensitivity of detection, namely through an analysis of the various broadening mechanisms involved. Finally, in Sec. V, we conclude and highlight possible future directions of work.

II. METHODS

A. Experimental Setup

A cold atomic cloud of ^{87}Rb atoms is prepared using a two-stage vacuum system and a single cooling laser. The vacuum chamber is a commercially available double-MOT system from Inflection with antireflective coatings suitable for Rydberg excitation lasers. In the lower vacuum chamber, a two-dimensional (2D) magneto-optical trap (MOT) is prepared before the atoms are pushed vertically upwards by an additional laser beam into a separate chamber at 10^{-11} mbar pressure, where they are further cooled in a three-dimensional (3D) MOT with a resultant atom number around 10^8 . The temperature of the cloud is estimated to be 145–200 μK . A pair of magnetic coils in anti-Helmholtz configuration provide the magnetic field for trapping. The cooling laser consists of a distributed-feedback laser at 780 nm that is amplified by a tapered amplifier and then split across four double pass acousto-optic modulators (AOMs) to select for specific offset frequencies. A portion of the seed laser is used for laser locking via a modulation transfer spectroscopy lock [40]. Using this spectral reference, the laser is locked to the crossover feature arising from the $|5S_{1/2}, F = 2\rangle \rightarrow |5P_{3/2}, F' = 1, 3\rangle$ transitions. After an AOM double pass, the 3D cooling light is approximately -20 MHz detuned from the $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition. An additional electro-optic modulator (EOM) adds repump sidebands onto the cooling light at 6.8 GHz, to drive the

$|F = 1\rangle \rightarrow |F' = 2\rangle$ transition. Of the four AOMs, three are used for the 2D MOT, push, and 3D MOT beams respectively, and the last is tuned to be resonant with the $F' = 3$ transition for absorption imaging, which is used only to estimate the approximate atom number.

Figure 1 shows the Rydberg excitation scheme and an optical beam path diagram. Each laser beam is prepared away from the vacuum chamber and central optical breadboard, and is then fiber-coupled into the setup. The probe laser at 780 nm, resonant with the $|5P_{3/2}\rangle$ state, propagates perpendicular to the horizontal cooling beam and magnetic coil axis, before being detected on a high-gain photodetector. This beam is focused by a lens to a $1/e^2$ radius beam waist of 85 μm at the center of the vacuum cell, and defines the interaction region inside the atom cloud. Its intensity is chosen to remain in the weak probe regime [41], approximately 1 nW ($\Omega_p/2\pi \approx 0.2$ MHz). The coupling beam is generated by a fiber-coupled RIO external cavity laser package operating at 1529 nm. This is overlapped with the probe beam on a short-pass dichroic mirror and blue detuned slightly from the $|4D_{5/2}\rangle$ state. The intensity at the atoms is approximately 0.5 mW ($\Omega_c/2\pi \approx 35$ MHz). The Rydberg light is generated by an MSquared Solstis laser, at approximately 700 nm, and overlaps in the counter-propagating direction via a long-pass dichroic mirror, with a maximum power of approximately 200 mW ($\Omega_r/2\pi \approx 10$ –35 MHz).

The rf fields are transmitted by an Ettus LP0410 log-periodic printed circuit board antenna, oriented in the xy plane, mounted approximately 40 cm above the cell. The antenna generates a linearly polarized field in the xy -plane, while the laser beams are horizontally polarized in the xz - plane. The rf signal is generated with a Windfreak SynthHDpro microwave (MW) generator.

Experimental control and data acquisition and analysis are performed using Python, with an I/O National Instruments PXI-7852R card, mounted in a PXIe-1073 controller chassis. This is in turn connected to a control computer via a NI PCIe-8361 card.

B. Rydberg excitation and detection scheme

We probe the Rydberg state by monitoring the absorption of a weak beam locked on the $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition. Upon excitation outside of the cooling cycle, atoms are lost from the trap, reducing the optical depth (OD) of the sample and increasing the overall transmission of the probe beam. This provides a larger signal-to-noise ratio than an equivalent EIT signal, where the change in transmission is of the order of a few percent, as our dynamic range is set between zero and full transmission. This optical readout also provides a much faster response time than the similar work described in Ref. [42], where fluorescence images are taken and scans take tens of seconds. Here, a single full scan can reliably be performed in

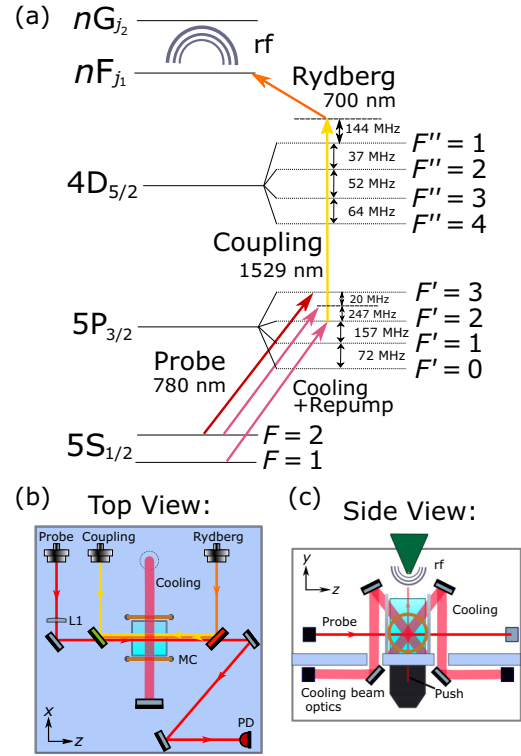


FIG. 1. (a) Rydberg excitation scheme diagram. (b) MOT breadboard setup from top-down view, showing excitation laser beam paths. L1, lens 1 ($f = 315$ mm); MC, (anti-Helmholtz) magnetic coils; PD, photodetector. (c) MOT breadboard setup from side view, showing the cross-beam cooling beam setup. Note that the beam path diagrams are schematic only, and not to scale.

about 500 ms or the on-resonant probe transmission can be continuously monitored to perform real-time rf detection, as in Ref. [43].

The continuously cooled regime we operate in is such that the effect of the excitation beams is to perturb the cooling cycle. For example, resonantly driving the $|5P_{3/2}, F' = 3\rangle \rightarrow |4D_{5/2}, F'' = 4\rangle$ transition with the coupling beam both scatters atoms out of the MOT, and induces AT splitting of the $5P_{3/2}$ levels (see the Appendix). This splitting modifies the effective detuning of the cooling light from the closed transition, reducing the effectiveness of the MOT. This, combined with the loss of atoms due to scattering, means that detection of a Rydberg signal is very difficult using resonantly driven transitions. Instead, by detuning the coupling laser away from resonance, and performing a multiphoton transition, these perturbative effects on the intermediate levels are reduced, while enabling the Rydberg state to be probed. The remainder of this section details spectroscopy that explains the choice of this multiphoton excitation scheme, that is ultimately shown in Fig. 1(a).

Figure 2 shows the probe transmission through the atomic cloud while varying the coupling and Rydberg laser

detunings, with the probe locked on the $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition, the coupling laser centered on the $|F' = 3\rangle \rightarrow |F'' = 4\rangle$ transition, and the Rydberg laser centered around the $45F_{7/2}$ transition (at 700 nm). Note that for all measurements in this paper, regions of larger transmission represent a loss of atoms from the MOT, leading to a reduced OD. The probe transmission is also converted from a voltage to a transmission percentage. The three horizontal features marked by the red dotted lines arise from the resonant driving of the $|5P_{3/2}, F' = 3\rangle \rightarrow |4D_{5/2}, F'' = 4, 3, 2\rangle$ transitions, from bottom to top, respectively [44]. At this moderate coupling power ($P_c \approx 1$ mW), the transitions are power-broadened, and there is in excess of 99% transmission along these resonances. The diagonal feature that passes through zero Rydberg detuning, marked by the purple dot-dashed line, arises from the two-photon resonance condition for Rydberg excitation, $\Delta_p + \Delta_c + \Delta_r = 0$, with a resonant probe ($\Delta_p = 0$). This highlights the ability to excite to the Rydberg state when far detuned from the two single resonance conditions at $\Delta_c = \Delta_r = 0$. The other feature that can be seen along the diagonal is a result of the three-photon resonance, $\Delta_{\text{cool}} + \Delta_c + \Delta_r = 0$ via the detuned cooling light, where $\Delta_{\text{cool}}/2\pi = -20$ MHz. The features highlighted by the three additional white/blue dotted horizontal lines around $(\Delta_r, \Delta_c)/2\pi = (-32, 297)$ MHz are a result of the three-photon resonance ($|5S_{1/2}, F = 1\rangle \rightarrow |5P_{3/2}, F' = 2\rangle \rightarrow |4D_{5/2}, F'' = 3, 2, 1\rangle \rightarrow |45F_{7/2}\rangle$). Here, the excitation resonance condition is fulfilled by the repump on the cooling light, with $\Delta_{\text{repump}} + \Delta_c + \Delta_r = 0$, and $\Delta_{\text{repump}}/2\pi = -267$ MHz detuned from the $F' = 3$ level. In other words, we can either off-resonantly excite to the Rydberg state by coupling via the $F' = 3$ state with the probe or the cooling light, or via the $F' = 2$ state with the repump sideband.

When obtaining a Rydberg signal, we lock the coupling laser at $\Delta_c/2\pi = +297$ MHz (from the $|5P_{3/2}, F' = 3\rangle \rightarrow |4D_{5/2}, F'' = 4\rangle$ transition, indicated by the blue dotted line in Fig. 2) and scan the Rydberg laser near the three-photon resonance. The benefits of doing so can be made clear by examining Fig. 3. In this figure, with the coupling laser fixed, the Rydberg laser is scanned around the $45F_{7/2}$ transition, while simultaneously scanning the cooling light, centered on the $|F = 2\rangle \rightarrow |F' = 3\rangle$ cooling transition. Note that the zero Rydberg detuning point has been shifted with respect to Fig. 2—it is now centered on the three-photon resonance fulfilled by the probe light (indicated by the red star in Fig. 2). The broad horizontal banding visible in the background of the figure shows the trend of atom number as a function of cooling detuning, with an optimum around $\Delta_{\text{cool}}/2\pi = -20$ MHz. On top of this background, there are three distinct features. The vertical feature centered at $\Delta_r = 0$ (indicated by the dotted red line) arises from the three-photon transition to the Rydberg state, with the resonant probe fulfilling the resonance condition

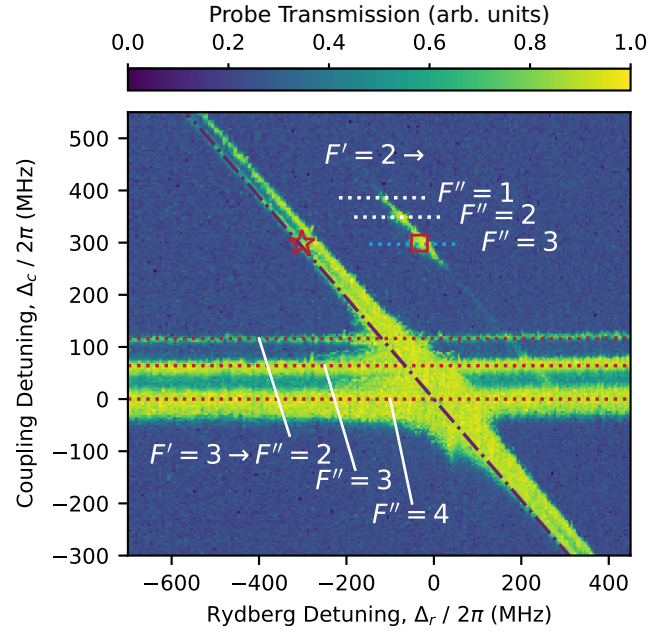


FIG. 2. Probe transmission through the atomic sample, as a function of Rydberg laser and coupling laser detuning, centered on the $|4D_{5/2}, F'' = 4\rangle$ and unresolved $|45F_{7/2}\rangle$ transitions, respectively. The red horizontal dotted lines indicate where the coupling laser couples to the $|4D_{5/2}, F'' = 4, 3, 2\rangle$, hyperfine levels, from bottom to top. The purple dot-dashed line indicates the two-photon resonance condition, for Rydberg excitation via the resonant probe pathway, $\Delta_c = -\Delta_r$. The three horizontal dotted white/blue lines indicate the positions of the three-photon resonance condition for Rydberg excitation via the cooling repump pathway, $\Delta_{\text{repump}} + \Delta_c + \Delta_r = 0$. They correspond to excitation via the $|4D_{5/2}, F'' = 3, 2, 1\rangle$ hyperfine levels, from bottom to top. The red star indicates the zero detuning point of Fig. 3, while the red square indicates the zero detuning point of Fig. 4, and subsequent measurements. The blue dotted line represents this detuning axis.

($\Delta_p + \Delta_c + \Delta_r = 0$) via the $|5P_{3/2}, F' = 3\rangle$ state. The other two diagonal features, that change as a function of the cooling detuning and are indicated by purple dashed lines, arise from the three-photon transition, with the cooling light and repump light fulfilling the resonance condition ($\Delta_{\text{cool(repump)}} + \Delta_c + \Delta_r = 0$), respectively. The diagonal feature nearest to $\Delta_r = 0$ does so via the $|5P_{3/2}, F' = 3\rangle$ state, whereas the feature at larger detuning does so via the $|5P_{3/2}, F' = 2\rangle$ state. At the optimal cooling detuning used for detection ($\Delta_{\text{cool}}/2\pi = -20$ MHz), the two Rydberg features around $\Delta_r = 0$ are almost indistinguishable. This is evident in Fig. 2, where the two transition pathways along the diagonal are barely resolvable until they are far from resonance. When an rf field is applied, the simultaneous splitting of these two components effectively leads to a much broader lineshape, which ultimately reduces the sensitivity of the detection. In order to observe a much cleaner and more narrow feature, we perform rf detection on the

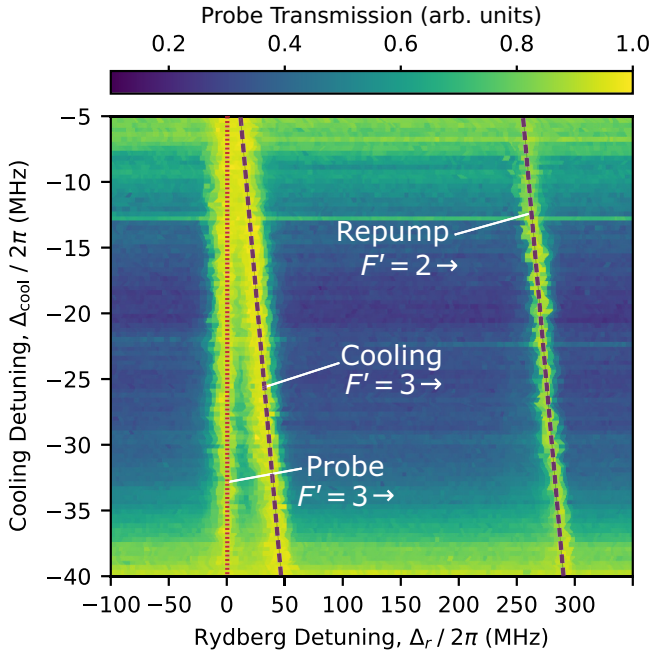


FIG. 3. Probe transmission as a function of the Rydberg and cooling detuning, with the coupling detuning fixed at $\Delta_c/2\pi = +297$ MHz. The changing background reflects the changing trapped atom number with cooling detuning. The vertical dotted line indicates the three-photon resonance condition via the resonant probe. The two dashed diagonal lines indicate the three-photon resonances via the cooling and repump light, exciting via the $|5P_{3/2}, F' = 3, 2\rangle$ from left to right, respectively. The excitation via the $F' = 2$ produces a “cleaner” and narrower signal and is therefore more suitable for rf detection.

feature from excitation via the $|5P_{3/2}, F' = 2\rangle$, at the slight expense of a smaller signal to noise, due to the weaker coupling through this pathway.

When performing a scan of the Rydberg detuning, due to the dynamics introduced by the MOT population cycling, the lineshape also depends on the chosen scan speed and width. Under typical experimental conditions, the MOT loads with a $1/e$ rate of 62 ms, which sets the approximate timescale for loading from an empty MOT. For scan speeds approaching this timescale, the signal becomes asymmetric, as atoms are lost more quickly than population can be cycled back up on the return sweep. Therefore, to ensure we observed the desired symmetric spectral lineshape when performing spectroscopy, all scans were performed at an arbitrarily larger rate of 2.5 s, to allow for differences in scan width.

III. UHF SENSING

When an rf field couples to a Rydberg level, it undergoes an AT splitting that results in two effective dressed

Rydberg transitions as

$$\omega_{\pm} = \omega_0 - \frac{\Delta_{\text{rf}}}{2} \pm \frac{1}{2} \sqrt{\Omega_{\text{rf}}^2 + \Delta_{\text{rf}}^2}, \quad (1)$$

where ω_0 is the bare angular transition frequency and $\Omega_{\text{rf}} = \vec{E} \cdot \vec{\mu}_d / \hbar$, Δ_{rf} are the rf Rabi frequency and angular detuning respectively, $\vec{\mu}_d$ is the rf transition dipole moment, and $\vec{E}$ is the electric field vector. If the applied rf field is perfectly resonant with the nearby Rydberg transition ($\Delta_{\text{rf}} = 0$), then the splitting, $\Delta\omega = \omega_+ - \omega_-$, is equal to the Rabi frequency of the driving field. The electric field, averaged over the atomic density and position, can then be calculated from the determined Rabi frequency and the calculated transition dipole moment for each transition. Here, we have calculated μ_d using $\sigma^{\pm}$ transitions. We have taken the quantization axis to be the B field, which we assume to be purely axial along the propagation direction of the beams (z), as the beam is centered between the anti-Helmholtz coils. The rf field is linearly polarized in the xy -plane so, in this geometry, will drive $\sigma^{\pm}$ transitions.

Figure 4 shows the evolution of the $|45F\rangle \rightarrow |45G\rangle$ transition, as the power of the driving rf field at 899 MHz is increased. Note that the Rydberg detuning zero has again been shifted, to center around the $F' = 2$ pathway signal (indicated by the red square in Fig. 2). Compared to this signal, the signal from the $F' = 3$ pathway (beginning at $-2\pi \cdot 247$ MHz) displays a broader overall lineshape, with a more complex spectrum as it splits. This is due to the probe and cooling lights’ nearly overlapping resonances. This results in a reduced sensitivity compared to the zero detuning signal. In contrast, the zero detuning signal is sufficiently detuned from other resonance to clearly display the expected AT splitting. The observed “wobble” in the colormap is simply a result of small mean frequency drift of the Rydberg laser frequency during data acquisition. Taking horizontal slices of this map and fitting two equal Gaussians to each lineshape, we can extract the splitting as a function of applied rf power, as the difference in the center points of each Gaussian.

In order to determine the $F \rightarrow G$ transition frequencies, the ARC (Alkali-Rydberg-Calculator) Python package [16] is used to calculate the approximate resonance. Then the rf field is scanned near the expected transition frequency, and the resonant frequency is determined by matching the amplitude of the two AT features [34,45]. In this regime (where the measured splittings are $\Delta\omega/2\pi > 6$ MHz), the error on the final electric field measurement due to the error in the rf detuning is much less than the error in the measured splitting, so the transition frequency need only be accurate to several megahertz to attain an accurate field measurement.

Figure 5 shows the calculated electric field strengths against the square root of the applied rf power for $n = 45, 50$, and 55 , corresponding to detected radio frequencies of

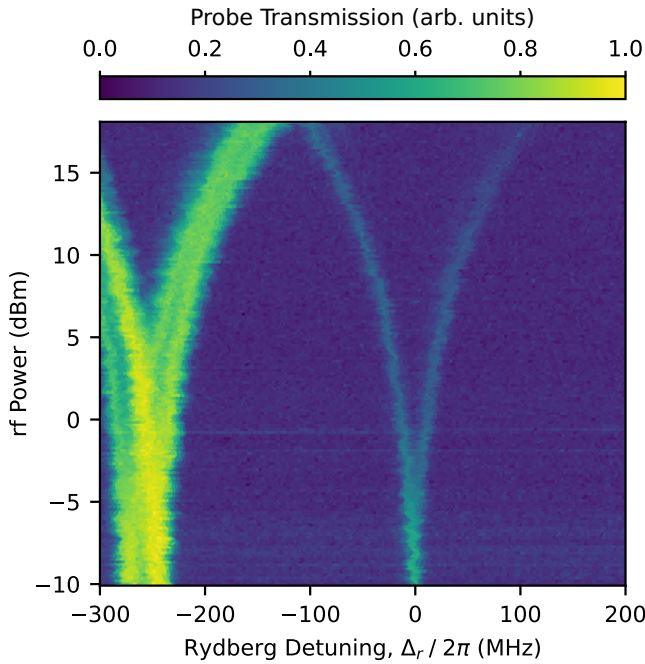


FIG. 4. Evolution of the $|45F\rangle \rightarrow |45G\rangle$ transition, driven by an 899-MHz field, with increasing applied rf power. The cooling light detuning is fixed at $\Delta_{\text{cool}}/2\pi = -20$ MHz. The signal centered around zero detuning is the excitation via the $F' = 2$ level, while the signal at negative detuning is via the $F' = 3$ level. Both undergo clear Autler-Townes splitting, but the zero detuning signal is resolvable at lower rf powers.

899, 659, and 493 MHz, respectively. The error bars are derived and propagated from the uncertainty in the fitted peak positions [46]. The rf power quoted is the rf power applied from the Windfreak to the antenna, so does not account for the radiation efficiency nor the directivity of the antenna or cable loss. Nevertheless, the expected proportional relation is shown by the linear fits to each dataset, along with an error region, representing a calibration factor for the antenna, between the applied and detected field, scaled by the atomic linestrength factor. For large electric field strengths there is good agreement, but at lower electric field strengths there is some evidence that the splitting becomes nonlinear with the applied E field. This nonlinear region has been studied and observed in Refs. [47,48] and is understood to arise when the Rabi frequency of the driving rf field is of similar order to that of the excitation lasers. At a certain lower limit of applied rf power, the peaks are no longer resolvable, such that the difference between the fitted peak positions saturates and flattens out. To make the figure easier to read, most of these datapoints have not been displayed, but the onset of this point has been defined as the minimum detectable field at each frequency. The minimum detectable field for descending frequencies was measured as 2.5(6), 4.4(6), and 5.6(6) mV/cm, respectively. This is a somewhat surprising result, as the transition dipole moment for the rf transitions increases

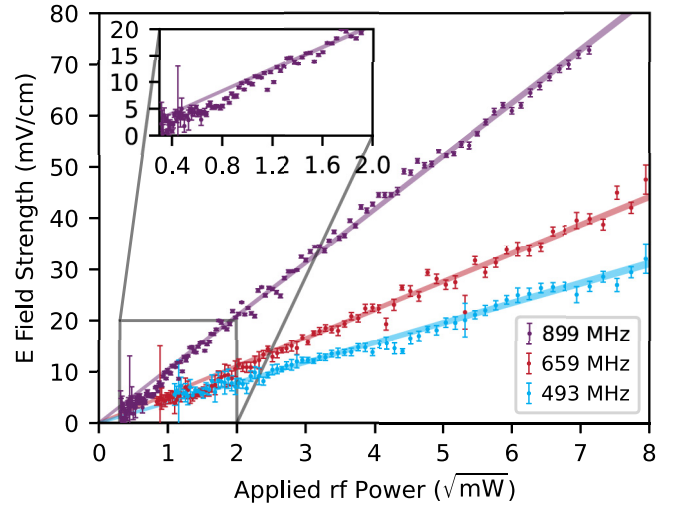


FIG. 5. Measured electric field as a function of applied rf power, for radio frequencies $f = 899, 659$, and 493 MHz, corresponding to measurements made at $n = 45, 50$, and 55 , respectively. The minimum detectable fields are 2.5(6), 4.4(6), and 5.6(6) mV/cm, respectively. The linear fits represent a calibration curve between the output from the antenna, and the detected E field, whose shaded areas represent the error region of the fit. The inset shows a zoomed region at low power for the 899 MHz data only, to demonstrate the deviation from the fit at low power.

with n , meaning one might naively expect the sensitivity to increase with n . However, other atomic factors negate from this extra sensitivity—for example, the increasing linewidth at increasing n [see Fig. 7(c)]. The decreasing gradient with decreasing frequency could also partially be attributed to a change in rf-related factors, such as a frequency dependence in the antenna's rf efficiency. It could also be due to the longer wavelengths interacting with the environment and creating a more complex field distribution in the cell [25,49]. For example, the distance between the roof of the metal optical bench and the breadboard is approximately 0.5 m while the wavelength at 499 MHz is 0.6 m. As these are of similar order, the emitting modes of the antenna could be altered [50], and the rf intensity distribution around the cell could become nontrivial due to interference effects.

Figure 6 demonstrates the ability for the system to detect continuously, by amplitude modulation of an rf signal at a carrier frequency of 899 MHz, and varying modulation frequencies. Here, the Rydberg laser is tuned to the center of the resonance and the average rf power applied to the antenna is 10 dBm. The rf power is amplitude modulated by approximately 40%. The modulation is generated externally and combined at the rf generator. A bias-tee is used to pick off an rf signal before the antenna, to use for a reference of the driving field. The raw data are extracted directly from the scope via its analog-to-digital converter, and the fast Fourier transform (FFT) is calculated in Python. Figures 6(a) and 6(b) show

the data for an example dataset at 60 Hz and 1300 Hz modulation frequency, respectively. In the inset of Fig. 6(a), the driver signal (red) shows the raw bias-tee pickoff. The atomic response (purple) is simultaneously observed, via the photodetector transmission on another channel, with the same modulation frequency and a small phase offset. Some discretization noise can be seen on the driver signal in this example due to the small pickoff voltage. In the main panel of Fig. 6(a), the FFT of the driving signal contains the largest component at the modulation frequency, but also contains overtones at multiples of this, due to imperfect MW generation. The atomic response captures the modulation frequency and overtones well, as can be seen in the purple trace. In Fig. 6(b), the atomic response has begun to reduce in amplitude compared to the driving signal, as can be seen by the reduced amplitude of the raw data in the inset. By extracting the peak prominence (signal minus background) of the FFT peak at the modulation frequency for varying modulation frequencies, we generate the Bode plot seen in Fig. 6(c). Here, the error bars represent the standard error of three repeat measurements. By fitting a two-variable Lorentzian to this data, we extract a 3-dB bandwidth response of 5.7(4) kHz. This represents a lower limit on the possible bandwidth, due to unaccounted uncertainties in this demonstration. For example, the modulation depth and initial power were left unoptimized, and the Rydberg laser frequency was left unlocked during the measurements. Use of a more sophisticated spectrum analyzer with a higher bit depth could capture smaller modulations, and sampling a larger number of points would ensure no potential aliasing occurs. A detailed analysis of the frequency response for the various components used was also not considered.

IV. DISCUSSION

In the AT splitting regime in which we operate, a primary limiting factor in the minimum detectable field for a given Rydberg electrometer will be the linewidth of the measured Rydberg signal without rf present, Γ , as this defines the minimum resolvable difference in signal used to measure the electric field. In the shot-noise-limited regime, this minimum resolvable change in transmission can be the limiting factor for the sensitivity of the system [51]. There are a number of broadening mechanisms that contribute to the overall linewidth of an observable signal. Some fundamental factors are the lifetime and dephasing rates of the involved states, blackbody broadening [52], and shot noise. Dephasing can arise from collisional interactions between Rydberg atoms and between ground-state and other Rydberg atoms [3]. Other factors include laser dephasing due to the laser linewidth, power broadening, transit time broadening, Doppler broadening, and Stark and Zeeman broadening.

In this system, the collisional interactions are likely to be negligible, as Rydberg atoms are lost from the trap

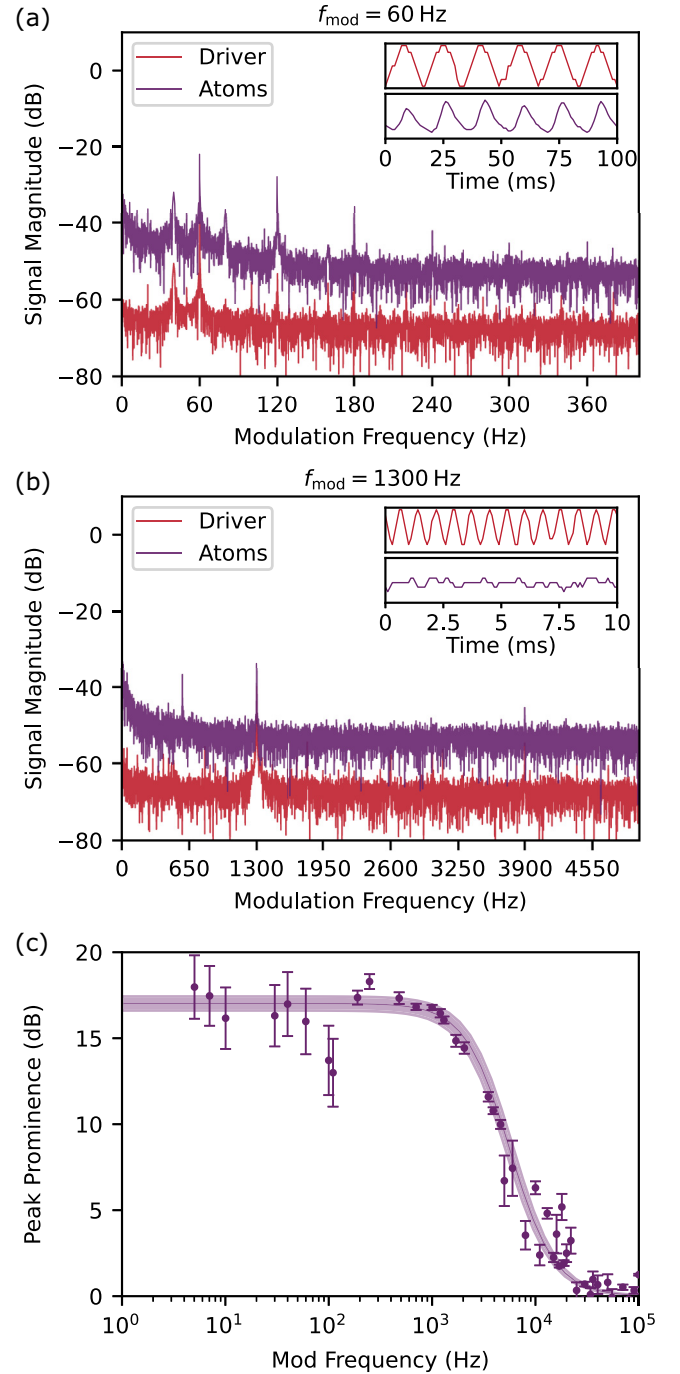


FIG. 6. Continuous-time detection of the amplitude modulated rf, with a carrier frequency of 899 MHz. Panels (a),(b) show example fast Fourier transforms of the probe transmission signals for $f_{\text{mod}} = 60$ Hz and 1300 Hz modulation, respectively. Shown in each inset is the driver signal (red, top)—a pickoff of the rf used to drive the antenna for reference—while the atom signal (purple, bottom) shows the photodetector transmission. Panel (c) shows a Bode plot of the response at various modulation frequencies, where the error bars are derived from the standard error of three repeats. The Lorentzian fit is shown, with a width of 5.7(4) kHz, and the shaded area represents a 1σ error region.

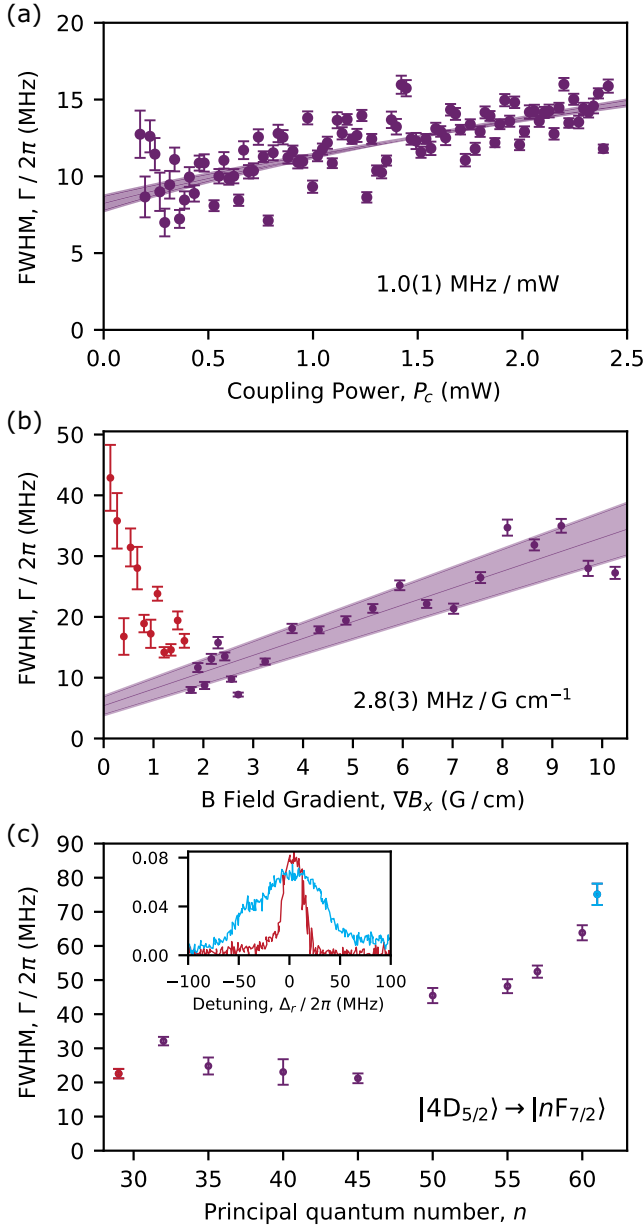


FIG. 7. Various broadening mechanisms that affect the linewidth of the measured Rydberg signal. (a) The power broadening arising from increasing coupling power, at $n = 45$ with B field gradient, $\nabla B_x = 3 \text{ G/cm}$. A fit of Eq. (2) extracts an initial linewidth, $\Gamma_0/2\pi = 8.2(5) \text{ MHz}$. A linearized estimate of the power broadening in the region of interest is $2\pi \cdot 1.0(1) \text{ MHz/mW}$. (b) Zeeman broadening due to the anti-Helmholtz B field, at $n = 45$ with $P_c = 1.2 \text{ mW}$. After the point of optimal linewidth, the Zeeman broadening term can be estimated with a linear fit as $2\pi \cdot 2.8(3) \text{ MHz/Gcm}^{-1}$. The datapoints used in the fit have been colored in purple. (c) The linewidth of the signal increases with n . This is likely due to increased dc Stark broadening from stray electric fields, due to the increased polarizability with n . The inset shows two example spectra at $n = 29$ and $n = 61$, in red and blue, respectively. Their corresponding full width, half maximum (FWHM) is highlighted in the same color. The shaded areas in (a),(b) represent the error regions associated with the fits.

upon excitation. Any Doppler broadening and transit time broadening is reduced as the velocity profile of the atoms is significantly reduced by the MOT cooling. This also means that one does not have to account for the usual Doppler mismatch factor [53] when operating in a thermal cell where the velocity profile is large. On the other hand, Zeeman and Stark broadening are likely much larger issues, due to the active inhomogeneous anti-Helmholtz magnetic field used for trapping, and the ability for stray electric fields to form inside the high-pressure vacuum cell, which can create an electric field gradient across the atoms.

In order to minimize the Rydberg signal linewidth and maximize electric field sensitivity, we perform an analysis of the broadening effect of several experimental factors, the results of which are shown in Fig. 7. For all measurements, the full width half maximum is extracted from a Gaussian fit to the Rydberg signal without an rf field present. The error bars in Figs. 7(a) and 7(b) represent the uncertainty on the fitted widths, while in Fig. 7(c) the error bars are calculated from both the uncertainty on the fit and the standard error from three repeat measurements.

Figure 7(a) shows the effect of varying the coupling power on the linewidth of the $|45F\rangle$ feature, at a magnetic field gradient of 3 G/cm . The probe beam power is minimized such that the weak probe condition is maintained, while maximizing signal to noise. The Rydberg laser power is also maximized, in order to maximize the weak coupling to the Rydberg states, leaving the coupling beam power as a free parameter. It can be seen from Fig. 7(a) that there is an increase in the linewidth that follows the form [54]

$$\Gamma = \Gamma_0 \left(1 + \frac{I}{I_{\text{sat}}} \right)^{1/2}, \quad (2)$$

where Γ_0 is the non-power-broadened linewidth, and I_{sat} is the saturation intensity. Fitting to this equation, the non-power-broadened linewidth can be extracted as $\Gamma_0/2\pi = 8.2(5) \text{ MHz}$. In the power range of interest, we can approximate a linear broadening term of $2\pi \cdot 1.0(1) \text{ MHz/mW}$.

In Fig. 7(b), we vary the current supplied to the anti-Helmholtz coils in order to vary the magnetic field gradient, with a coupling power of $P_c = 2.4 \text{ mW}$. At low magnetic field gradients, the atomic cloud is loosely trapped across a larger spatial region in the cell, meaning that the probe beam (Rayleigh range $z_R = 45.8(1) \text{ mm}$) samples atoms experiencing a larger spread of magnetic field strengths, leading to increased inhomogeneous broadening. The signal to noise is also small in this region, meaning the fitted Gaussians are shallower and wider. As the magnetic field gradient is increased, an optimal linewidth is reached at approximately 3 G/cm , where the maximal OD is reached, as the cloud has become localized to a high enough density for near full absorption

of the probe beam off-resonance. After this point, the linewidth increases, due to the inhomogeneous sampling of the Zeeman-shifted degenerate energy levels within the linewidth of the original feature. There is a rich energy manifold within the observable linewidth: for example, for the $|35F\rangle$ state, the fine structure j -splitting between $j = \frac{5}{2}, \frac{7}{2}$ has been measured as 3.6 MHz [55]. This represents the largest splitting in the bare atomic Hamiltonian, so the hyperfine and Zeeman manifold is unresolvable. Under the effect of a B field, these unresolved levels split simultaneously, culminating in a perceived overall broadening. This broadening could be estimated with a linear fit to the datapoints after the OD saturation point (illustrated in purple), to be $2\pi \cdot 2.8(3) \text{ MHz/Gcm}^{-1}$. Based on a calculation of the Zeeman shifts associated with the ground and Rydberg states, the level of broadening in this regime is consistent with a cloud of constant size around 3 mm.

Finally, in Fig. 7(c) we observe the linewidth of the feature as a function of principal quantum number, by tuning the Rydberg laser to match the three-photon resonance condition for each $|4D\rangle \rightarrow |nF\rangle$ transition. The Rydberg laser power was left unchanged between measurements, so the three-photon Rabi frequency reduces as n is increased, as a result of the reducing transition dipole moment. The linewidth broadens significantly as a function of increasing n , which could be attributed to increased dc-Stark-induced state mixing and inhomogeneous sampling of the stray electric field [56,57]. As the dc polarizability scales as n^7 , the higher n states are significantly more sensitive to shifts and splittings due to dc electric fields. It is possible for patch charges to form within the cell due to atoms adsorbing to the walls, forming stray electric fields [58]. Given that the higher n states are closer to the ionization limit, these atoms may be more likely to undergo blackbody-induced ionization, which in turn could generate disruptive dc electric field distributions. The combination of these factors may result in Stark broadening becoming dominant over the reduction in linewidth due to the increased lifetimes and reduction in three-photon Rabi frequency. At $n > 60$, the signal feature was too broad and no longer resembled a single Gaussian, so no meaningful linewidth could be extracted. The substructure appearing at high n [shown in the inset to Figure 7(c)] further supports the Stark broadening hypothesis, however.

Minimizing these broadening factors simultaneously, the data suggest a minimum linewidth around $2\pi \cdot 8(1) \text{ MHz}$ for the $n = 45$ state. The likely absolute limiting factor is the linewidth of the $|6P_{3/2}\rangle$ transition, $2\pi \cdot 6.06 \text{ MHz}$, as the continuous scattering from the strong cooling light introduces decoherence by mixing of this state. The calculated minimum likely includes this factor plus some residual Stark and Zeeman broadening.

V. CONCLUSION

We have demonstrated UHF detection down to 493 MHz using continuously cooled atom loss spectroscopy. This is performed using a three-photon Rydberg excitation using the cooling light, while the atom number is sampled with a weak probe. This three-photon excitation is comprehensively characterized as a function of the laser detunings. A modest minimum detectable field is achieved using a basic AT measurement regime, limited by the linewidth of the atom loss spectrum, likely arising from the decoherence introduced by the cooling light. Whilst the minimum detectable field is an order of magnitude less sensitive than other cold atom detectors [38], this work highlights the possibility of continuous detection in an MOT, which is previously undemonstrated. Further work could implement more complex techniques to improve rf sensitivity, such as implementing an additional local oscillator field [19,23,30,35], or could improve existing experimental parameters, such as by more thoroughly frequency stabilizing the coupling and Rydberg lasers.

There is little literature available to compare the performance of this detection technique at these radio frequencies, in this modality of operation. Therefore, a direct comparative study between this system and a similar thermal vapor system using EIT could help to illuminate the differing effects in this system that could limit sensitive detection. Similarly, this system could be easily adapted to a pulsed excitation scheme in the UHF regime, whereby the cooling light is pulsed off while detection is performed, so that cold EIT measurements could be made, as in Refs. [38,59]. Investigation of this mode of operation would isolate the broadening and perturbing effects of the cooling light and B field from the other broadening mechanisms, enabling a more thorough characterization of the limits of this detection, and a more direct comparison to previous cold and thermal studies using traditional EIT techniques.

These investigations could serve to inform the development of this system to take advantage of the benefits of cold atoms for sensing. For example, off-resonant transitions/excitations are made more feasible by the reduced Doppler-matching requirements compared to a thermal vapor. This could enable further unexplored routes to different frequency bands. The reduced collision rates should also reduce the collisional ionization that can cause issues with stray electric fields in thermal cells when exciting to high- n Rydberg states.

Finally, this system could provide a new method for the investigation of the energy levels and quantum defects of Rydberg F and G states, similarly to Ref. [53]. Larger angular momentum states could also be investigated through the application of additional rf fields, with the method used in Refs. [34,60].

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DATA AVAILABILITY

The data that support the findings of this article are openly available [61].

APPENDIX: COUPLING LASER-INDUCED AT SPLITTING OF $|5P_{3/2}\rangle$ STATE

To illustrate the necessity of exciting to the Rydberg state off-resonantly, here we demonstrate the effect of resonantly driving the $|5P_{3/2}, F' = 3\rangle \rightarrow |4D_{5/2}, F'' = 4\rangle$ coupling transition under continuous cooling. In a system with only the probe, cooling, and coupling lasers active, resonant driving of this transition leads to AT splitting of the $|5P_{3/2}\rangle$ state. This alters the effective detuning of the cooling light, which alters the effective damping coefficient of the cooling. This, in turn, modifies the effectiveness of the trap, culminating in a reduced trapped atom

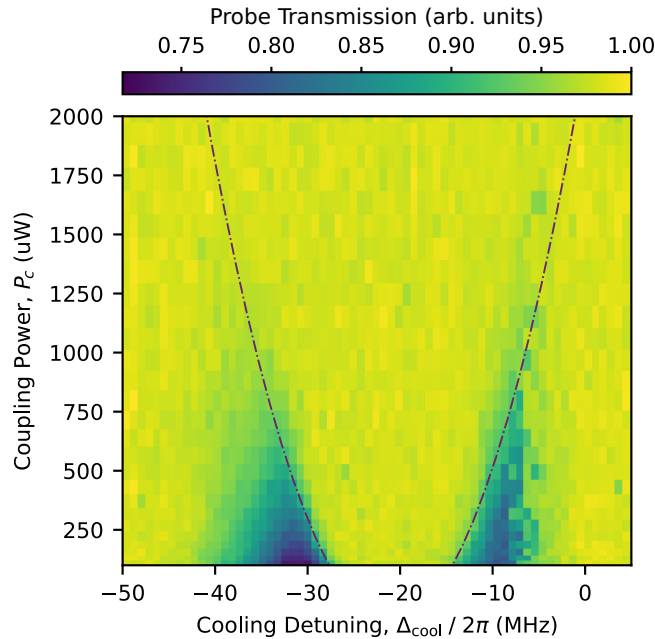


FIG. 8. AT splitting of the $|5P_{3/2}\rangle$ state induced by the coupling light. The probe absorption indicates the OD of the trapped atom cloud. The optimal trapping detuning becomes split as a function of the coupling power. At large coupling powers, the atoms are fully scattered out of the trap. The dot-dashed line indicates the Rabi frequency as a function of coupling power, calculated using the measured beam waist of the coupling beam.

number, which can be clearly seen in Fig. 8. Here, the power in the coupling beam is varied for varying cooling detunings, and the probe beam used to monitor the OD. For the unperturbed state ($P_c = 0$), the optimal cooling light detuning, with highest OD, is $\Delta_{\text{cool}} \approx -3\Gamma_{5P}$, where $\Gamma_{5P}/2\pi = 6.06$ MHz is the natural linewidth of the $5P$ transition (see Fig. 3). When the coupling laser is applied, the optimum cooling detuning shifts and splits, so that the most absorption continues to occur at $-3\Gamma_{5P}$ from the split levels. This results in two regions of increased absorption in the colormap, separated by the Rabi frequency of the coupling field. The trendline indicates the calculated coupling Rabi frequency as a function of power, with standard experimental parameters, that is, a $1/e^2$ beam waist of $400\ \mu\text{m}$. There is good agreement with the splitting, until at larger powers, the increased scattering rate causes all of the atoms to be lost from the trap. In order to compensate for this splitting, we must either detune the cooling light in step with the Rabi frequency of the coupling light, or detune the coupling light away from resonance. For experimental simplicity, we detune away from resonance and instead perform a three-photon excitation to the Rydberg state.

- [1] C. S. Adams, J. D. Pritchard, and J. P. Shaffer, Rydberg atom quantum technologies, *J. Phys. B* **53**, 012002 (2019).
- [2] T. F. Gallagher, *Rydberg Atoms* (Cambridge University Press, Cambridge, 1994).
- [3] H. Fan, S. Kumar, J. Sedlacek, H. Kübler, S. Karimkashi, and J. P. Shaffer, Atom based RF electric field sensing, *J. Phys. B* **48**, 202001 (2015).
- [4] D. H. Meyer, Z. A. Castillo, K. C. Cox, and P. D. Kunz, Assessment of Rydberg atoms for wideband electric field sensing, *J. Phys. B* **53**, 034001 (2020).
- [5] K. C. Cox, D. H. Meyer, F. K. Fatemi, and P. D. Kunz, Quantum-limited atomic receiver in the electrically small regime, *Phys. Rev. Lett.* **121**, 110502 (2018).
- [6] P. Padovani, The faint radio sky: radio astronomy becomes mainstream, *Astron. Astrophys. Rev.* **24**, 13 (2016).
- [7] J. Ingerski and A. Sapp, in *MILCOM 2002. Proceedings* (IEEE, Anaheim, CA, USA, 2002), Vol. 1, pp. 642–646.
- [8] H. Carreno-Luengo, C. S. Ruf, S. Gleason, and A. Russel, Detection of inland water bodies under dense biomass by CYGNSS, *Remote Sens. Environ.* **301**, 113896 (2024).
- [9] M. M. Farhad, A. M. Alam, S. Biswas, M. A. S. Rafi, A. C. Gurbuz, and M. Kurum, SDR-based dual polarized l-band microwave radiometer operating from small UAS platforms, *IEEE J. Sel. Top. Appl. Earth Obs. Remote Sens.* **17**, 6617 (2024).
- [10] N. Neji, R. de Lacerda, A. Azoulay, T. Letertre, and O. Outtier, Survey on the future aeronautical communication system and its development for continental communications, *IEEE Trans. Veh. Technol.* **62**, 83 (2013).
- [11] S. Orasch, H. Flühr, C. Rihacek, and B. Haindl, in *2024 47th MIPRO ICT and Electronics Convention (MIPRO)* (IEEE, Opatija, Croatia, 2024), pp. 195–200.

- [12] J. R. Rigley, Aeronautical mobile satellite services: The launching of a new era in mobile communications, *Can. J. Electr. Comput. Eng.* **17**, 175 (1992).
- [13] F. Raab, R. Caverly, R. Campbell, M. Eron, J. Hecht, A. Mediano, D. Myer, and J. Walker, HF, VHF, and UHF systems and technology, *IEEE Trans. Microwave Theory Tech.* **50**, 614 (2002).
- [14] C. L. Holloway, J. A. Gordon, S. Jefferts, A. Schwarzkopf, D. A. Anderson, S. A. Miller, N. Thaicharoen, and G. Raithel, Broadband Rydberg atom-based electric-field probe for SI-traceable, self-calibrated measurements, *IEEE Trans. Antennas Propag.* **62**, 6169 (2014).
- [15] A. K. Mohapatra, T. R. Jackson, and C. S. Adams, Coherent optical detection of highly excited Rydberg states using electromagnetically induced transparency, *Phys. Rev. Lett.* **98**, 113003 (2007).
- [16] N. Šibalić, J. Pritchard, C. Adams, and K. Weatherill, ARC: An open-source library for calculating properties of alkali Rydberg atoms, *Comput. Phys. Commun.* **220**, 319 (2017).
- [17] D. H. Meyer, K. C. Cox, F. K. Fatemi, and P. D. Kunz, Digital communication with Rydberg atoms and amplitude-modulated microwave fields, *Appl. Phys. Lett.* **112**, 211108 (2018).
- [18] J. S. Otto, M. Chilcott, A. B. Deb, and N. Kjaergaard, Distant RF field sensing with a passive Rydberg-atomic transducer, *Appl. Phys. Lett.* **123**, 144003 (2023).
- [19] M. Jing, Y. Hu, J. Ma, H. Zhang, L. Zhang, L. Xiao, and S. Jia, Atomic superheterodyne receiver based on microwave-dressed Rydberg spectroscopy, *Nat. Phys.* **16**, 915 (2020).
- [20] L. Ma, M. A. Viray, D. A. Anderson, and G. Raithel, Measurement of dc and ac electric fields inside an atomic vapor cell with wall-integrated electrodes, *Phys. Rev. Appl.* **18**, 024001 (2022).
- [21] C. L. Holloway, M. T. Simons, J. A. Gordon, P. F. Wilson, C. M. Cooke, D. A. Anderson, and G. Raithel, Atom-based RF electric field metrology: From self-calibrated measurements to subwavelength and near-field imaging, *IEEE Trans. Electromagn. Compat.* **59**, 717 (2017).
- [22] E. Paradis, G. Raithel, and D. A. Anderson, Atomic measurements of high-intensity VHF-band radio-frequency fields with a Rydberg vapor-cell detector, *Phys. Rev. A* **100**, 013420 (2019).
- [23] D. H. Meyer, P. D. Kunz, and K. C. Cox, Waveguide-coupled Rydberg spectrum analyzer from 0 to 20 GHz, *Phys. Rev. Appl.* **15**, 014053 (2021).
- [24] S. M. Bohaichuk, D. Booth, K. Nickerson, H. Tai, and J. P. Shaffer, Origins of Rydberg-atom electrometer transient response and its impact on radio-frequency pulse sensing, *Phys. Rev. Appl.* **18**, 034030 (2022).
- [25] T. Minhas, in *Quantum Sensing, Imaging, and Precision Metrology II* (SPIE, San Francisco, CA, USA, 2024), Vol. 12912, p. 129120B.
- [26] Z. Song, H. Liu, X. Liu, W. Zhang, H. Zou, J. Zhang, and J. Qu, Rydberg-atom-based digital communication using a continuously tunable radio-frequency carrier, *Opt. Express* **27**, 8848 (2019).
- [27] A. K. Mohapatra, M. G. Bason, B. Butscher, K. J. Weatherill, and C. S. Adams, A giant electro-optic effect using polarizable dark states, *Nat. Phys.* **4**, 890 (2008).
- [28] A. P. Rotunno, S. Berweger, N. Prajapati, M. T. Simons, A. B. Artusio-Glimpse, C. L. Holloway, M. Jayaseelan, R. M. Potvliege, and C. S. Adams, Detection of 3–300 MHz electric fields using Floquet sideband gaps by “Rabi matching” dressed Rydberg atoms, *J. Appl. Phys.* **134**, 134502 (2023).
- [29] S. Kumar, H. Fan, H. Kübler, J. Sheng, and J. P. Shaffer, Atom-based sensing of weak radio frequency electric fields using homodyne readout, *Sci. Rep.* **7**, 42981 (2017).
- [30] M. T. Simons, A. H. Haddab, J. A. Gordon, and C. L. Holloway, A Rydberg atom-based mixer: Measuring the phase of a radio frequency wave, *Appl. Phys. Lett.* **114**, 111105 (2019).
- [31] D. H. Meyer, J. C. Hill, P. D. Kunz, and K. C. Cox, Simultaneous multiband demodulation using a Rydberg atomic sensor, *Phys. Rev. Appl.* **19**, 014025 (2023).
- [32] A. B. Deb and N. Kjaergaard, Radio-over-fiber using an optical antenna based on Rydberg states of atoms, *Appl. Phys. Lett.* **112**, 211102 (2018).
- [33] P. K. Elgee, J. C. Hill, K.-J. E. LeBlanc, G. D. Ko, P. D. Kunz, D. H. Meyer, and K. C. Cox, Satellite radio detection via dual-microwave Rydberg spectroscopy, *Appl. Phys. Lett.* **123**, 084102 (2023).
- [34] G. Allinson, M. J. Jamieson, A. R. Mackellar, L. Downes, C. S. Adams, and K. J. Weatherill, Simultaneous multiband radio-frequency detection using high-orbital-angular-momentum states in a Rydberg-atom receiver, *Phys. Rev. Res.* **6**, 023317 (2024).
- [35] R. C. Brown, B. Kayim, M. A. Viray, A. R. Perry, B. C. Sawyer, and R. Wyllie, Very-high- and ultrahigh-frequency electric-field detection using high angular momentum Rydberg states, *Phys. Rev. A* **107**, 052605 (2023).
- [36] A. Duspayev, R. Cardman, D. A. Anderson, and G. Raithel, High-angular-momentum Rydberg states in a room-temperature vapor cell for dc electric-field sensing, *Phys. Rev. Res.* **6**, 023138 (2024).
- [37] S. M. Bohaichuk, F. Ripka, V. Venu, F. Christaller, C. Liu, M. Schmidt, H. Kübler, and J. P. Shaffer, Three-photon Rydberg-atom-based radio-frequency sensing scheme with narrow linewidth, *Phys. Rev. Appl.* **20**, L061004 (2023).
- [38] K.-Y. Liao, H.-T. Tu, S.-Z. Yang, C.-J. Chen, X.-H. Liu, J. Liang, X.-D. Zhang, H. Yan, and S.-L. Zhu, Microwave electrometry via electromagnetically induced absorption in cold Rydberg atoms, *Phys. Rev. A* **101**, 053432 (2020).
- [39] F. Zhou, F. Jia, X. Liu, Y. Yu, J. Mei, J. Zhang, F. Xie, and Z. Zhong, Improving the spectral resolution and measurement range of quantum microwave electrometry by cold Rydberg atoms, *J. Phys. B* **56**, 025001 (2023).
- [40] D. J. McCarron, S. A. King, and S. L. Cornish, Modulation transfer spectroscopy in atomic rubidium, *Meas. Sci. Technol.* **19**, 105601 (2008).
- [41] B. E. Sherlock and I. G. Hughes, How weak is a weak probe in laser spectroscopy? *Am. J. Phys.* **77**, 133 (2009).
- [42] R. Duverger, A. Bonnin, R. Granier, Q. Marolleau, C. Blanchard, N. Zahzam, Y. Bidet, M. Cadoret, A. Bresson, and S. Schwartz, Metrology of microwave fields based on trap-loss spectroscopy with cold Rydberg atoms, *Phys. Rev. Appl.* **22**, 044039 (2024).
- [43] C. L. Holloway, M. T. Simons, A. H. Haddab, C. J. Williams, and M. W. Holloway, A “real-time” guitar recording using Rydberg atoms and electromagnetically

- induced transparency: Quantum physics meets music, *AIP Adv.* **9**, 065213 (2019).
- [44] M. Allegrini, E. Arimondo, and L. A. Orozco, Survey of hyperfine structure measurements in alkali atoms, *J. Phys. Chem. Ref. Data* **51**, 043101 (2022).
- [45] M. T. Simons, J. A. Gordon, C. L. Holloway, D. A. Anderson, S. A. Miller, and G. Raithel, Using frequency detuning to improve the sensitivity of electric field measurements via electromagnetically induced transparency and Autler-Townes splitting in Rydberg atoms, *Appl. Phys. Lett.* **108**, 171101 (2016).
- [46] I. Hughes and T. P. A. Hase, *Measurements and Their Uncertainties: A Practical Guide to Modern Error Analysis* (Oxford University Press, New York, 2010).
- [47] C. L. Holloway, M. T. Simons, J. A. Gordon, A. Dienstfrey, D. A. Anderson, and G. Raithel, Electric field metrology for SI traceability: Systematic measurement uncertainties in electromagnetically induced transparency in atomic vapor, *J. Appl. Phys.* **121**, 233101 (2017).
- [48] S. Chen, D. J. Reed, A. R. MacKellar, L. A. Downes, N. F. A. Almuhawish, M. J. Jamieson, C. S. Adams, and K. J. Weatherill, Terahertz electrometry via infrared spectroscopy of atomic vapor, *Optica* **9**, 485 (2022).
- [49] H. Fan, S. Kumar, J. Sheng, J. P. Shaffer, C. L. Holloway, and J. A. Gordon, Effect of vapor-cell geometry on Rydberg-atom-based measurements of radio-frequency electric fields, *Phys. Rev. Appl.* **4**, 044015 (2015).
- [50] W. L. Barnes, S. A. R. Horsley, and W. L. Vos, Classical antennas, quantum emitters, and densities of optical states, *J. Opt.* **22**, 073001 (2020).
- [51] N. Prajapati, N. Bhusal, A. P. Rotunno, S. Berweger, M. T. Simons, A. B. Artusio-Glimpse, Y. Ju Wang, E. Bottomley, H. Fan, and C. L. Holloway, Sensitivity comparison of two-photon vs three-photon Rydberg electrometry, *J. Appl. Phys.* **134**, 024504 (2023).
- [52] I. I. Beterov, I. I. Ryabtsev, D. B. Tretyakov, and V. M. Entin, Quasiclassical calculations of blackbody-radiation-induced depopulation rates and effective lifetimes of Rydberg nS , nP , and nD alkali-metal atoms with $n \leq 80$, *Phys. Rev. A* **79**, 052504 (2009).
- [53] C. L. Holloway, J. A. Gordon, A. Schwarzkopf, D. A. Anderson, S. A. Miller, N. Thaicharoen, and G. Raithel, Sub-wavelength imaging and field mapping via electromagnetically induced transparency and Autler-Townes splitting in Rydberg atoms, *Appl. Phys. Lett.* **104**, 241108 (2014).
- [54] C. Foot, *Atomic Physics* (Oxford University Press, New York, 2005).
- [55] J. Han, Y. Jamil, D. V. L. Norum, P. J. Tanner, and T. F. Gallagher, Rb nf quantum defects from millimeter-wave spectroscopy of cold ^{85}Rb Rydberg atoms, *Phys. Rev. A* **74**, 054502 (2006).
- [56] A. Osterwalder and F. Merkt, Using high Rydberg states as electric field sensors, *Phys. Rev. Lett.* **82**, 1831 (1999).
- [57] J. D. Carter, O. Cherry, and J. D. D. Martin, Electric-field sensing near the surface microstructure of an atom chip using cold Rydberg atoms, *Phys. Rev. A* **86**, 053401 (2012).
- [58] B. Liu, L. H. Zhang, Z. K. Liu, Z. A. Deng, D. S. Ding, B. S. Shi, and G. C. Guo, Electric field measurement and application based on Rydberg atoms, *arXiv:2305.16696*.
- [59] K. J. Weatherill, J. D. Pritchard, R. P. Abel, M. G. Bason, A. K. Mohapatra, and C. S. Adams, Electromagnetically induced transparency of an interacting cold Rydberg ensemble, *J. Phys. B* **41**, 201002 (2008).
- [60] G. Allinson, L. A. Downes, C. S. Adams, and K. J. Weatherill, Determination of quantum defects and core polarizability of atomic cesium via terahertz and radio-frequency spectroscopy in thermal vapor, *Phys. Rev. A* **111**, 062802 (2025).
- [61] <https://doi.org/doi:10.15128/r28w32r568j>