

Low-frequency electric field sensing with a Rydberg beam

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We present a method for performing low-frequency electric-field sensing via ionization detection of Rydberg atoms in a collimated atomic beam. A collimated beam avoids many of the electric-field screening effects that are common in warm vapor cells due to the accumulation of alkali-metal atoms on glass surfaces. Further, a beam facilitates a spatially separated region for high signal-to-noise readout via ionization detection. Using this approach, we measure dc Stark shifts from external fields with frequencies as low as 1 Hz. The sensor demonstrates a sensitivity of better than 1 mV/m $\sqrt{\text{Hz}}$ for frequencies above 20 Hz and 0.14(4) mV/m $\sqrt{\text{Hz}}$ above 500 Hz with a linear dynamic range of over 50 dB.

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I. INTRODUCTION

Electrometry with Rydberg atoms has been well demonstrated over wide frequency ranges from kilohertz to terahertz [1–4]. Rydberg states have various properties that scale advantageously with principle quantum number n including the dipole moment ($\propto n^2$) and polarizability ($\propto n^7$) making them highly sensitive to electric fields. The properties of Rydberg atoms and response to electric fields are traceable to fundamental constants making them ideal for use as calibration tools and in precision measurement applications [5–7]. Rydberg sensors can simultaneously respond to disparate frequency bands [8], perform polarimetry [9], image terahertz frequency fields [10,11], reach the standard quantum limit for individual particles [12,13], and in principle reach the Heisenberg limit when leveraging entanglement [14,15].

Current Rydberg sensors commonly use warm vapor cells containing alkali-metal atoms, which pose challenges for sensing in the extremely low frequency (ELF) band. Sensing in the ELF band has applications in geophysics [16,17], submarine communication [18], and instances where the detection of weak voltage signals is required. A primary challenge to ELF sensing with Rydberg vapor cells is that alkali-metal atoms can accumulate on the glass, increasing the conductivity and leading to shielding of low-frequency fields [19–22]. Field plates can be placed within cells to circumvent this issue. These “port-coupled” devices require an external aperture (i.e. antenna) to collect the free-space field and couple it through a vacuum port to internal field plates [23,24]. Vapor cells constructed of sapphire, or with internal paraffin coatings, have shown a lower timescale of screening enabling sensitive detection at 1 kHz [25] and down to single Hertz [22,26,27].

Here, rather than a warm vapor, we create a collimated beam of Rydberg atoms inside the vacuum cell, and this offers two major benefits for ELF sensing. First, the atoms can be directed away from the glass surfaces that serve as apertures for electromagnetic fields, thus reducing issues with alkali-metal buildup and shielding effects. Second, the collimated beam readily enables high-efficiency ionization detection of the Rydberg atoms. This yields large signal-to-noise (SNR) ratio ion spectra, and high sensitivity measurements of external fields. Field ionization of Rydberg atoms has a long history, such as in studies of Rydberg-Rydberg interactions [28–30], Stark shifts of Rydberg states [31,32], and state selective ionization of Rydberg states [33–35]. Rydberg ionization techniques have also been used for field sensing of internal charge distributions within vacuum systems [36,37].

In this report we measure the dc Stark shift of Rydberg atoms in $101P_{3/2}$ to detect electric fields external to our vacuum chamber. External electric fields with frequencies down to 1 Hz are observable, and we achieve a sensitivity better than 1 mV/m $\sqrt{\text{Hz}}$ for frequencies above 20 Hz and 0.14(4) mV/m $\sqrt{\text{Hz}}$ above 500 Hz.

II. METHOD

A. Design overview

An overview of the experimental setup is shown in Fig. 1. A reservoir containing rubidium is heated, producing a vapor that flows through a nozzle and creates a collimated effusive atomic beam. This beam travels approximately 4.5 cm and passes through a laser-excitation region where a portion of the atoms are excited to a Rydberg state. The atoms then travel about 1.7 cm where they enter between two parallel mesh plates, which produce an electric field sufficient to ionize the Rydberg atoms. The positive ions then pass through the mesh and are deflected towards a channel electron multiplier (CEM) located about 6 cm from the ionization region. A preamplifier along with a data acquisition system (DAQ) and oscilloscope are used to read out the ion signal. A blue LED external to

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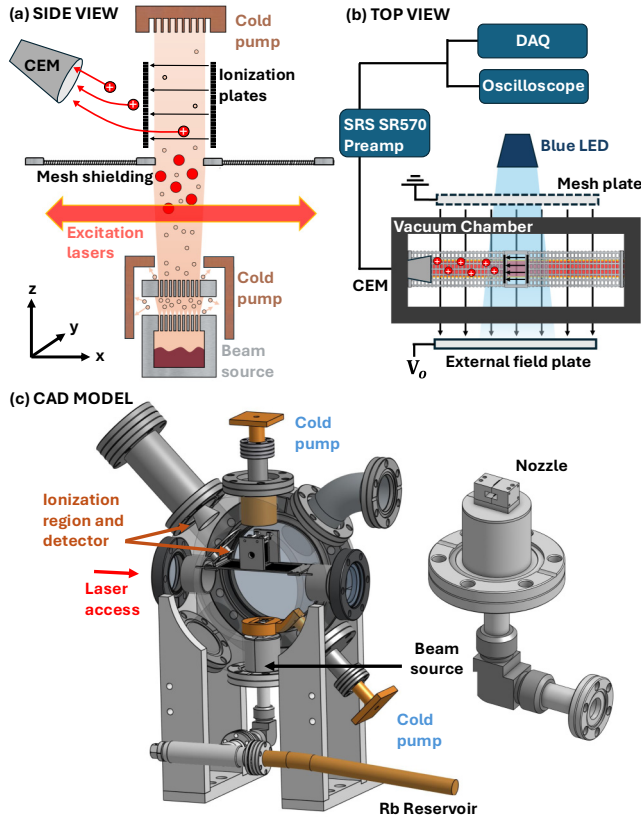


FIG. 1. Overview of the experimental apparatus; (a) and (b) are not drawn to scale. (a) A collimated ^{85}Rb beam passes through lasers exciting a portion of the atoms to a Rydberg state. These atoms pass through a set of mesh plates that ionize the Rydberg atoms and deflect the ions towards a CEM detector. (b) Two plates external to the vacuum system produce a low-frequency electric field at the excitation region in the chamber. A blue LED is used to reduce time-varying perturbations on the ion signal. (c) CAD model of the vacuum system and atomic beam source. Some components have been made transparent to aid in visual clarity.

the vacuum chamber is shone onto the windows to reduce unwanted time-varying perturbations in the experiment. In order to characterize the sensitivity, we place plates external to the vacuum system to produce electric fields at the excitation region. This induces a dc Stark shift resulting in a change in the measured ion signal.

B. Atomic beam source

We have a linear array of collimated atomic beams, which together create a quasi-two-dimensional (2D) sheet of vapor. This is produced using a cascaded collimator nozzle that is similar to Ref. [38]. In our system, a bottom reservoir is filled with natural abundance rubidium, and is heated to around 390 K. A top portion of the oven including the nozzle is heated separately to around 475 K to reduce the migration of liquid rubidium that could otherwise clog the nozzle. The nozzle, machined from aluminum, is a two-stage collimator shown in Figs. 1 and 2(c) with ten channels that have a total length of 9.3 mm and width and depth of (0.5×0.5) mm. The channels have a spacing of 0.127 mm between them and a gap that is

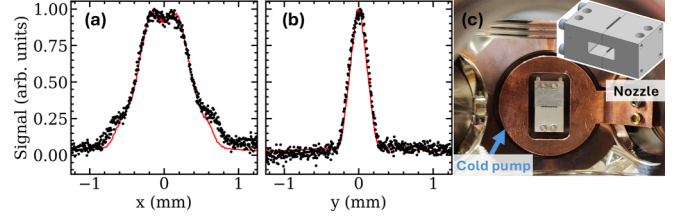


FIG. 2. Spatial characteristics of the atomic beam at the location of the laser excitation region parallel (a) and perpendicular (b) to the laser propagation directions. Red curves are simulated spatial profiles while black markers are experimental values. (c) Image of the collimating nozzle and surrounding cold pump.

cut out in the center that is 3.2 mm in length. Atoms with high transverse velocity have higher probability of leaving through the gap resulting in a collimated effusive atomic beam.

The spatial characteristics of the beam at the location of the excitation lasers, measured via fluorescence spectroscopy, is shown in Figs. 2(a) and 2(b), along with simulated spatial characteristics determined using MOLFLOW+ software [39]. The full width at half maximum of the atomic beam is approximately 8 mm along the x direction and 0.3 mm along the y direction. The small size in the y direction ensures spatial overlap of all the atoms with the laser beams. Fluorescence and absorption spectroscopy measurements show that the peak density in the beam is $\approx 10^8 \text{ cm}^{-3}$. While this density is lower than that found in warm vapor cells ($\approx 10^{10} \text{ cm}^{-3}$ for a room-temperature rubidium cell), the narrow transverse velocity spread results in reduced Doppler shifts helping improve excitation efficiency to the Rydberg state.

Although the atomic beam is well collimated, atoms can strike a surface of the vacuum chamber, scatter, and stick to the vacuum windows. To reduce the accumulation of rubidium on the windows, two copper cold pumps are employed. One surrounds the nozzle shown in Figs. 1 and 2(c) while the other is placed above the beam source as shown in Fig. 1. These are cooled to around 238 K using two-stage thermoelectric coolers mounted externally to the vacuum system. The copper feedthroughs used for the cold pumps along with the nozzle are mounted on thin walled stainless steel tubing to reduce thermal contact with the rest of the vacuum chamber.

C. Laser systems

The rubidium atoms are excited to a Rydberg state via a three photon ladder scheme shown in Fig. 3. A “probe” laser at 795 nm excites the $5S_{1/2}(F=3) \rightarrow 5P_{1/2}(F'=3)$ transition. A electro-optic modulator adds 3.036-GHz frequency sidebands to the probe laser to also excite the $5S_{1/2}(F=2) \rightarrow 5P_{1/2}(F'=3)$ transition, improving the total excitation efficiency to the final Rydberg state. The atoms are then excited from $5P_{1/2}(F=3) \rightarrow 6S_{1/2}(F'=3)$ with a 1324-nm “dressing” laser and finally from $6S_{1/2}(F=3) \rightarrow nP_{3/2}$ with a titanium-sapphire 740-nm “Rydberg” laser. In this work we tune the Rydberg laser to target principle quantum numbers $n = 55, 60$, and 101.

To maximize the excitation efficiency the lasers are arranged in a counterpropagating orientation with the probe and dressing laser wave vectors being $\vec{k}_p = k_p \hat{x}$ and $\vec{k}_d = k_d \hat{x}$ while

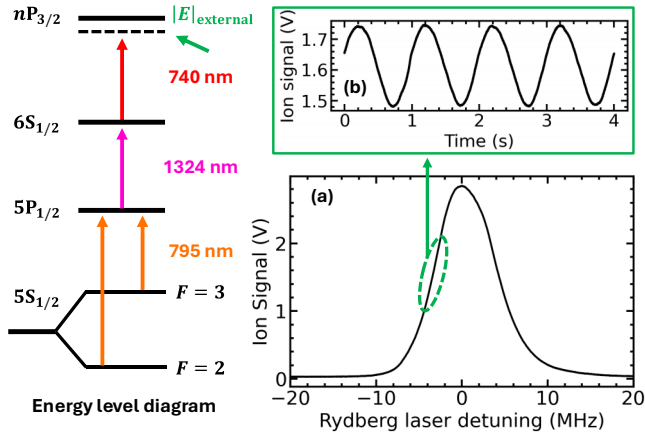


FIG. 3. Energy level and laser scheme for excitation to the Rydberg state. (a) Example ion spectra for $60P_{3/2}$. (b) Observed oscillation of the ion signal when a 1-Hz signal is applied externally to the system while the Rydberg laser is locked to the side of the ion spectra.

the Rydberg laser's wave vector is $\vec{k}_R = -k_R \hat{x}$ (see Fig. 1). All lasers are linearly polarized with Gaussian beam waists ($1/e^2$ value) of 1 mm. The typical probe and dressing laser powers are about 1 mW while the Rydberg laser power is about 300 mW. The probe and dressing lasers are power stabilized using acousto-optic modulators, and frequency stabilized by locking to an ultrastable optical cavity.

D. Field ionization and detection

Due to the excitation to a highly energetic state, the Rydberg atoms can be state selectively ionized with only modest electric fields. A simple picture in which the valence electron experiences the Coulomb potential from the nucleus summed with a one-dimensional homogeneous ionizing potential gives reasonably accurate values for the classical ionization threshold. This occurs when the binding energy exceeds the saddle-point energy requiring an electric-field strength of

$$|\vec{E}| \geq \frac{\pi \epsilon_0 R_y^2}{e^3 n^{*4}}, \quad (1)$$

where R_y is the Rydberg constant and $n^* = n - \delta_l$. Here δ_l is the quantum defect for the corresponding n state. A more complete picture requires including dc Stark shifts of the state energy. Assuming an adiabatic evolution of the initial state, ionization occurs when the state energy crosses the classical ionization threshold. For $n = 60$ the required field strength is 3000 V/m while for $n = 100$ it is only 350 V/m.

We achieve the necessary field strengths for ionization using two mesh ionization plates as shown in Fig. 1. For electrical insulation, the plates are suspended on ceramic rods, separated by 9.5 mm, and have an open area of 67%. The ceramic rods are mounted on stainless steel plates connected to a mesh shielding structure shown in Fig. 1. The holes shown in the side of the steel plates are used for separate experiments of laser excitation within the ionization region and are not necessary for this study. Upon passing between the plates,

the Rydberg atoms are state ionized with the positive ions being deflected through the mesh and towards a CEM detector that is biased to -1800 V. To reduce stray fields from the ionization plates and CEM detector from reaching the laser excitation region, additional mesh shielding is placed between the laser excitation region and ionization region as shown in Fig. 1. This reduces the dc Stark shift from fringe fields of the ionization plates, and allows for a continuous operation of the system without necessitating a pulsed measurement scheme.

To observe the ion spectra, the Rydberg laser frequency is scanned, and an example spectrum for $60P_{3/2}$ is shown in Fig. 3(a). The full width at half maximum is about 8 MHz with the dominant linewidth broadening mechanism being power broadening. Detection of external electric fields is performed by locking the Rydberg laser to the side of the ion spectra where relatively small frequency shifts $\Delta\nu$ result in a linear change in the ion signal voltage ΔV_{ion} . An example of this is shown in Fig. 3(b) for a 1-Hz frequency. In this example, the ion signal from the CEM is processed with a 10-Hz low-pass filter. Frequency shifts occur due to the quadratic Stark shift of the energy level with a magnitude of [40]

$$\Omega_{\text{quad}} = -\frac{1}{2}\alpha E^2, \quad (2)$$

where α is the polarizability of the Rydberg state, scaling as $\alpha \propto n^7$.

In the linear region of the ion spectra, the detectable external electric field E_{ext} for a quadratic Stark shift will be proportional to

$$E_{\text{ext}} \propto \sqrt{\frac{\Delta V_{\text{ion}}}{\beta n^7}}, \quad (3)$$

where β is the slope of the ion spectra. The minimal detectable change in the ion signal voltage, ΔV_{ion} , will be proportional to the SNR, which in the shot-noise regime is proportional to $\sqrt{N}$, where N is the number of Rydberg atoms detected. As $E_{\text{ext}} \propto N^{1/4}$, we can see that a 16 times increase in the number of Rydberg atoms only improves the sensitivity by a factor of 2.

To improve sensitivity in this quadratic Stark shift regime, we use a bias electric field E_b where $E_b \gg E_{\text{ext}}$, as has been done in prior work [22]. The total electric-field amplitude at the atoms is now $E = E_b + E_{\text{ext}}$ resulting in a Stark shift of

$$\Omega_{\text{quad}}^{\text{bias}} = -\frac{1}{2}\alpha(E_b^2 + 2E_b E_{\text{ext}} + E_{\text{ext}}^2). \quad (4)$$

This introduces a component of the Stark shift that is linear with respect to the external field. The minimum detectable external electric field $E_{\text{ext}}^{\text{bias}}$ is proportional to

$$E_{\text{ext}}^{\text{bias}} \propto \frac{\Delta V_{\text{ion}}}{E_b \beta n^7}, \quad (5)$$

where the quadratic dependence in SNR and the ion spectra slope have been removed. While it is clearly advantageous to operate with a high bias field, the scaling in Eq. (5) does not persist indefinitely. At sufficiently high electric-field strengths, state mixing occurs leading to broadening of the ion spectra which decreases β . In Fig. 4, the theoretical Stark map for $101P_{3/2}$ is presented where state mixing can be observed beginning at around 12 V/m due to linear Stark shifts of

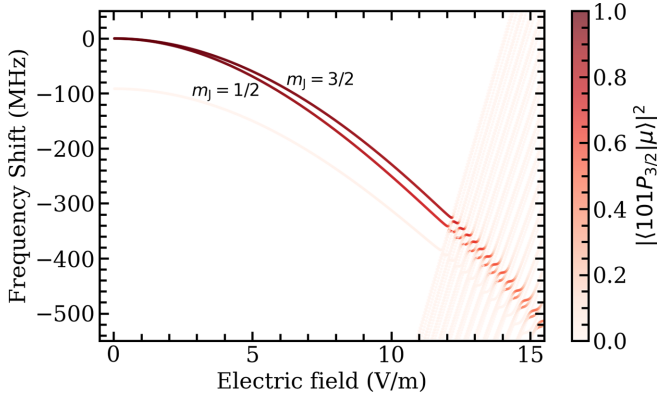


FIG. 4. Theoretical Stark map of $|101P_{3/2}, m_J = 1/2\rangle$ and $|101P_{3/2}, m_J = 3/2\rangle$ showing energy shifts and state mixing as a function of electric-field amplitude. The colorbar indicates the fraction of the initial state in each atom eigenstate $|\mu\rangle$.

$n = 98$ with $l = 3$ to 97. The Stark map is produced using the ARC (Alkali Rydberg Calculator) PYTHON package [41]. Since the energy spacing between Rydberg states decreases with increasing n , the onset of state mixing and linewidth broadening will occur at lower field strengths as n increases. As such, excitation to arbitrarily high- n states will also not result in an indefinite improvement to sensitivity.

We perform sensitivity measurements at $101P_{3/2}$ and various calibration measurements at lower- n states of around $n = 60$. States around $n = 60$ are preferable for calibrating external electric fields because splitting of the m_J states can be easily resolved, and they are relatively less susceptible to environmental noise. In fact, at $n = 101$ we observe unintentional bias fields, including primarily from ionization plates, and also the CEM and other sources. We empirically set our controlled bias field for good sensitivity, and the unintentional fields contribute to this beneficial Stark shift, though we expect the inhomogeneity of these fields could be broadening the Rydberg ion spectra. Further improvements to the shielding and uniformity of the bias field could reduce the linewidth and improve sensitivity.

We empirically found that shining a blue (400 nm) LED uniformly across the windows reduced unwanted slow time-varying perturbations on the experiment. We believe the LED increases the mobility of charges on the glass, allowing shielding of slow external field noise [42]. The shielding does affect the signal strength as well, and we found that for frequencies below ≈ 10 Hz, the signal amplitude is about half of its value without the blue LED. Further, power fluctuation of the blue LED adds noise to the ion signal. Eliminating this blue LED power noise would reduce the overall ion signal noise floor by a factor of 2. Nonetheless, the benefits of increased stability over longer timescales outweighed the signal reduction and additional noise. A more detailed study of the underlying physical dynamics of the LED-induced surface physics will be the subject of future work.

III. RESULTS

In order to determine the minimum detectable electric field, it is necessary to know the electric-field strength at the

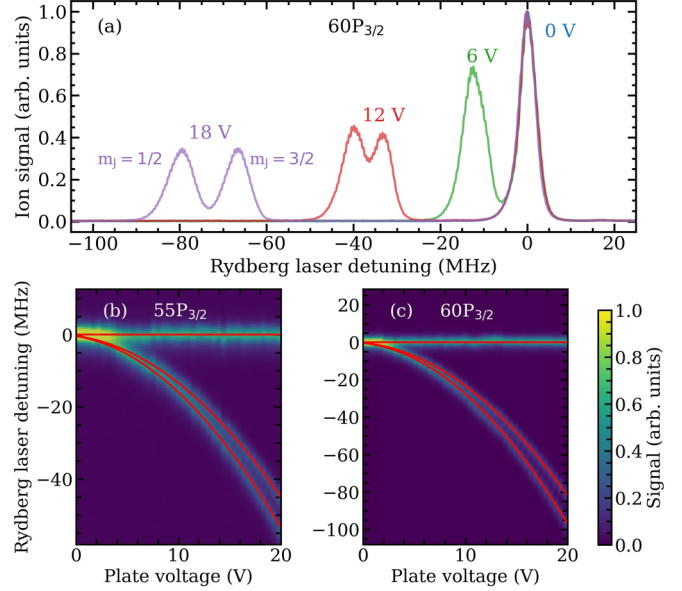


FIG. 5. (a) Example dc Stark shifts and splitting of the m_J states for various voltages applied to the external field plates. (b, c) Measured 2D Stark maps with theoretical fits (red curves) for $55P_{3/2}$ (b) and $60P_{3/2}$ (c). The theoretical fits are used to determine the conversion factor between the voltage applied to the external plates and electric-field magnitude at the laser excitation region in the system.

location of the atoms for a given voltage applied to the external field plate. Due to shielding effects from the metal chamber and internal components this cannot be easily computed from the plate separation and applied voltage. Instead we apply a 600-Hz square wave signal to the external field plate and measure the time averaged ion signal as a function of Rydberg laser detuning for various plate voltages. Measuring the time averaged signal results in an effective measurement of a dc Stark shift and at sufficiently high plate voltages splitting of the m_J levels also occurs as shown in Fig. 5(a). From these results a 2D Stark map can be constructed and fit to theoretical curves where the fitting parameters provide a conversion factor from plate voltage to electric-field amplitude as well as an initial Stark shift to account for residual background fields present in the chamber.

Two-dimensional Stark maps and theoretical curves are given in Figs. 5(b) and 5(c) for $55P_{3/2}$ and $60P_{3/2}$. From the fit, the conversion parameters are such that 1 V on the external field plate corresponds to an electric-field amplitude at the atoms of $1.75(5)$ and $1.78(5)$ V/m for $55P_{3/2}$ and $60P_{3/2}$ respectively. As a second means of comparison, the vacuum chamber, relevant internal components, and external field plates are modeled with the COMSOL MULTIPHYSICS software. Figure 6(a) shows the simulated electric field with shielding effects from the chamber and 1 V applied to the external plate. A field strength of ≈ 2.4 V/m at the location of the atoms is predicted. We do not include possible shielding effects from surface charges on the glass windows in this model. In Fig. 6(b) are simulations of the field strength in which the chamber and internal metal are modeled as vacuum. This together with Fig. 6(a) indicates that about 55% of the field strength from the plates is being shielded by the chamber.

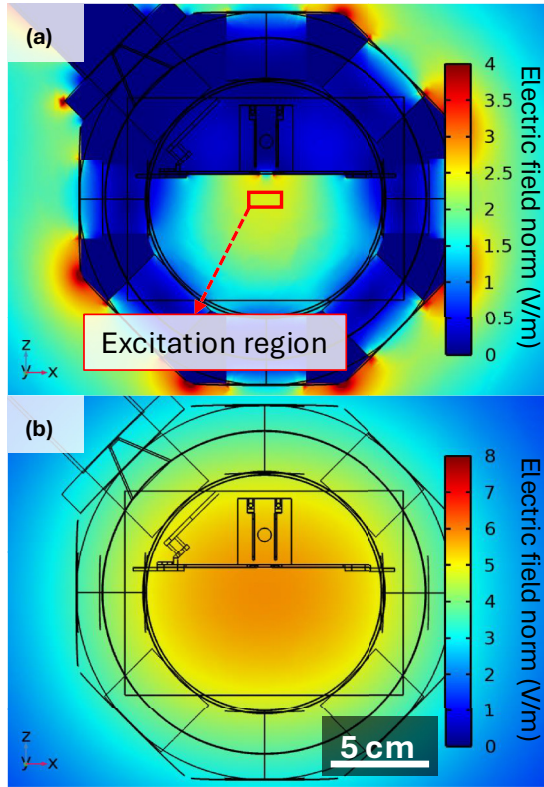


FIG. 6. COMSOL simulations of the magnitude of the electric field at the center of the excitation region ($y = 0$) for 1 V applied to the external field plate. (a) Results including shielding effects from the metal chamber and internal components. (b) The metal chamber along with all internal components are treated as vacuum providing no shielding of the applied external field. The bottommost bar indicates the length scale.

To determine the sensitivity of the system, we excite atoms to $101P_{3/2}$ and the ion signal is recorded for 1 s. The fast Fourier transform (FFT) is then computed with a software FFT with a 1-Hz bandwidth. Over the linear-response region of the ion spectra, the FFT amplitude will scale linearly with the applied electric-field amplitude. Figure 7(a) shows this for a 5-Hz signal where the linear portion of the data is fit to a line. The sensitivity is given by the point where this line crosses the noise floor, or equivalently where the SNR is unity for a 1-s measurement. The noise floor is measured with no applied voltage to the external field plate. Figure 7(b) is zoomed in, showing the FFT amplitude rolling off as it approaches the noise floor.

The sensitivity of our sensor over a frequency range of 1 to 1200 Hz is shown in Fig. 8. The sensitivity is $0.14(4)$ mV/m $\sqrt{\text{Hz}}$ for frequencies above 500 Hz and we see the sensitivity is better than 1 mV/m $\sqrt{\text{Hz}}$ for frequencies above 20 Hz. The two noticeable peaks at 60 and 180 Hz are from environmental electronic noise. We note that the sensitivity in Fig. 8 is the field strength measured at the location of the atoms and does not correct for the shielding effects from the chamber. If these chamber shielding effects were not present the sensitivity to external fields would be better by approximately a factor of 2, based on our COMSOL modeling. Further, for frequencies below approximately 10 Hz, an additional

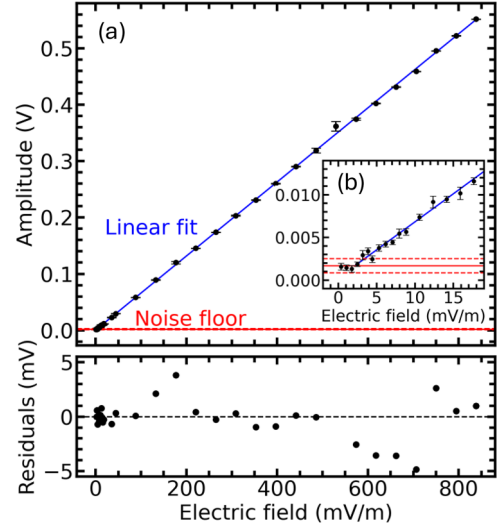


FIG. 7. (a) FFT amplitude of the ion signal at 5 Hz for a 1-s measurement time vs the applied electric field. The linear portion of the data is fit to a line (blue curve), and the point where this line crosses the 5-Hz noise floor (red solid line) gives the sensitivity. The dashed red lines indicate the uncertainty in the noise floor. The inset (b) shows the FFT amplitude rolling off at lower electric-field values as it approaches the noise floor.

factor of 2 improvement in sensitivity would be achieved if shielding effects from the blue LED were not present. With ionization detection we are able to achieve high single shot SNR. In Fig. 9, the single shot SNR versus Rydberg laser power is presented along with an example single shot spectra for 250 mW of Rydberg laser power. This spectra showing an SNR of 1500 is obtained with a Rydberg laser frequency scan rate of 20 Hz and is only processed with our preamplifier with a low-pass 10-kHz filter. As the SNR versus Rydberg laser power follows a linear trend, we expect that with 1 W of Rydberg laser power the sensitivity would be better by a factor of 3, or around 0.05 mV/m $\sqrt{\text{Hz}}$. It is important to note that this assumes that the sensitivity is not limited first by environmental noise. We can observe environmental noise when there is human activity close to the apparatus, and the best sensitivity is obtained during periods when there is minimal to no activity in adjacent laboratory spaces.

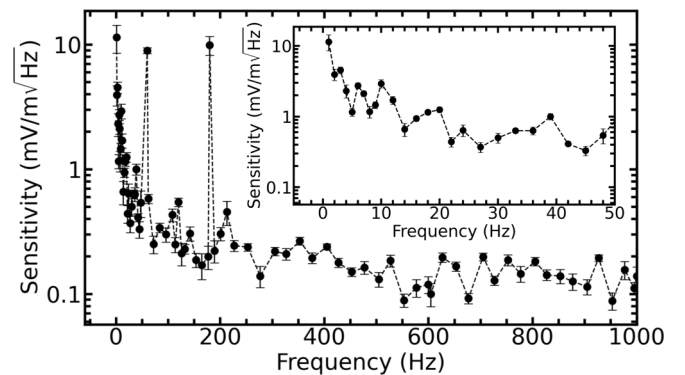


FIG. 8. Measured electric-field sensitivity vs electric-field frequency. The inset shows the sensitivity for frequencies below 50 Hz.

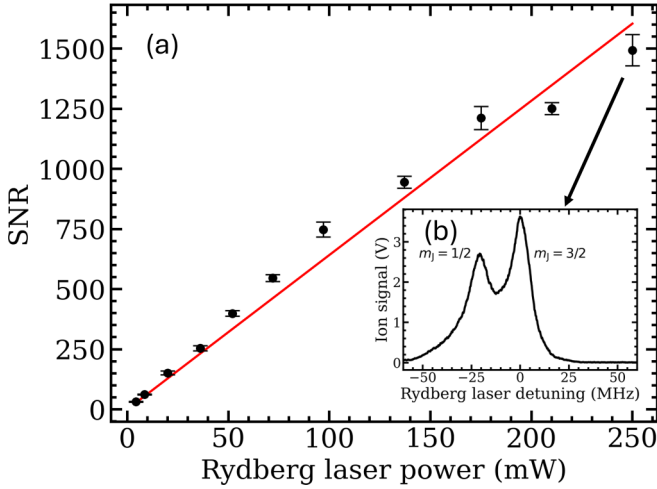


FIG. 9. (a) Single shot SNR of the $101P_{3/2}$, $m_J = 3/2$ Rydberg ion signal vs Rydberg laser power. Data are black points fit to a red line. Error bars represent the standard deviation of the SNR from ten shots. (b) Example single shot ion spectra for $101P_{3/2}$ with 250 mW of Rydberg laser power.

The linear dynamic range of the system is determined by measuring the FFT amplitude for 1-s integration time and observing the point in which the linear fit exceeds the experimentally measured amplitude by 1 dB. Figure 10 shows the dynamic range for a 20- and 108-Hz signal applied to the external field plate. The linear dynamic range is over 50 dB for both frequencies.

IV. CONCLUSION

We have presented an approach for performing low-frequency electric-field sensing using ionization detection on a collimated beam of Rydberg atoms. This reduces challenges associated with electric-field screening effects that are present in warm alkali-metal vapor cells. With a sensitivity of $0.14(4)$ mV/m $\sqrt{\text{Hz}}$ above 500 Hz, and ability to measure down to single Hertz frequencies, the use of a Rydberg beam and ionization detection offers a promising alternative to typical Rydberg electrometry schemes using warm vapor cells.

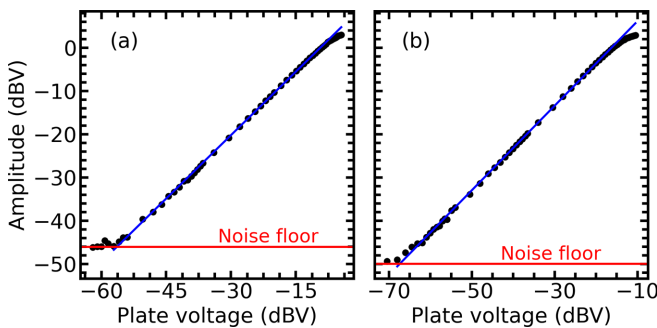


FIG. 10. FFT amplitude in 1-s integration time for 20 Hz (a) and 108 Hz (b) as a function of voltage applied to the external plate. Data (black points) are fit to a line (blue) to determine the linear dynamic range, which is over 50 dB for both frequencies.

There are avenues for further improving the sensitivity of the current system. While we have restricted ourselves to a “warm” atomic beam, a compact laser cooled beam source such as one from Ref. [43] would further reduce the transverse velocity spread, improving the excitation efficiency to the Rydberg state while also reducing Doppler broadening. A simple comparison can be performed by numerically solving the Lindblad master equation for the laser excitation scheme discussed above to determine the steady-state Rydberg population. Using the RYDIQULE PYTHON library [44,45], we find that for an atomic beam with a transverse temperature of 10 mK, the excitation efficiency for comparable laser powers is 15 times higher, and with a linewidth that is three times narrower than our warm atomic beam. This would result in an order-of-magnitude improvement in the sensitivity, assuming the dominant linewidth broadening mechanism is not from electric-field gradients. We note that at sufficiently high Rydberg densities, Rydberg-Rydberg interactions would play an important role limiting the Rydberg excitation efficiency and lifetime, and could lead to line broadening.

Shielding effects from the metal vacuum chamber can be reduced by using a mostly quartz vacuum chamber which has been utilized in a number of magneto-optical trap experiments [46–48]. Effects of charge accumulation on the glass surfaces can be reduced by using sapphire or sapphire coated glass windows [19,22]. While the interaction volume of the laser excitation region is comparable to centimeter scale vapor cells, the physical footprint of our vacuum chamber and internal components is considerably larger. The use of a miniature glass vacuum chamber in combination with a compact beam source [43,49,50] would allow a significant reduction in the overall system size. Particularly in small systems, a pulsed ionization scheme could be beneficial to allow for a miniature source while avoiding unwanted Stark shifts from ionization plates in close proximity to the laser excitation region.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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