

Resolving magnetic-sublevel structure in Rydberg Autler-Townes spectra with arbitrary rf polarization

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We investigate the role of magnetic sublevels in Autler-Townes spectra of Rydberg atoms driven by radio-frequency (rf) fields with arbitrary polarization. While conventional treatments predict two symmetric sidebands from independent m_J transitions, experiments have reported additional unexplained spectral features. We show that these arise from elliptical rf polarization, which coherently couples multiple magnetic sublevels and requires a full multilevel treatment. We develop and diagonalize a Hamiltonian including all coupled m_J sublevels, predicting polarization-dependent degeneracies that produce varying numbers of resolved peaks. Using long-wavelength transitions and an anechoic environment we realize homogeneous rf fields that enable complete resolution of the m_J -dependent dressed states. We observe excellent agreement with theory as the rf ellipticity is varied. These results demonstrate that rf polarization fundamentally modifies Autler-Townes spectra and provide a consistent framework for interpreting magnetic-sublevel structures, with implications for Rydberg-based rf electrometry and polarimetry.

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I. INTRODUCTION

Autler-Townes splitting in Rydberg states of alkali atoms has been demonstrated as a sensitive, SI-traceable form of radio-frequency (rf) electrometry [1,2]. It has been shown that angular momentum plays a role in the shape of these spectra, and this effect has been leveraged to use Rydberg sensors as field polarimeters [3–9]. The interpretation of these angular momentum effects has important implications for the traceability of these sensors [10]. Theoretical treatments [8] of the angular momenta predict that a $D \rightarrow P$ transition should split into two peaks on either side, in addition to the central unsplit peak. These peaks correspond to individual m_J levels. However, in previous measurements aiming to resolve angular-momentum-dependent behavior [11,12], field inhomogeneity broadened the dressed states beyond their energy separation. Experiments that applied enough field strength to observe structure related to magnetic sublevels [13–15] showed more than the two peaks predicted, and their relative locations did not agree with theory. Here, we show that the observed additional peaks are due to elliptical polarization of the rf field (the presence of scatterers prevents pure rf polarization in many environments). We derive theoretical Autler-Townes spectra for arbitrary polarizations, and measure the spectra experimentally. To allow for resolution of the m_J peaks, we apply rf with a wavelength long compared to the vapor cell

to ensure spatial homogeneity [16] and place the vapor cell in an anechoic chamber to allow for pure rf polarization. With this, we apply a large enough field to separate the m_J structure and study the resulting spectra as a function of rf polarization. While our study focuses on a $D_{5/2} \rightarrow P_{3/2}$ transition, the analysis can easily be extended to other fine structure states by following the theoretical treatment laid out in Sec. II.

II. THEORY

The system is shown in Fig. 1. We consider rf radiation resonant between two Rydberg states $|n_D D_{5/2}\rangle$ and $|n_P P_{3/2}\rangle$. The resonant rf radiation dresses the energies of these states [17]. We then read out the energy of these dressed states using two-photon electromagnetically induced transparency (EIT), following the experimental scheme of Ref. [18].

Pure linear [Fig. 1(b)] and circular [Fig. 1(d)] rf polarizations independently couple m_J levels, and thus can be thought of as independent systems, whereby each m_J is dressed and probed by the lasers separately with no coherent interactions. However, for elliptical polarization [Fig. 1(c)], the effects of each m_J are no longer independent as the rf coherently couples nearly all the m_J states together. Because of this, the spectral peaks cannot be treated as resulting from individual transitions with separate dipole moments. Instead, we must consider all the coherent couplings in a Hamiltonian which accounts for all magnetic sublevels and couplings, with the spectral peaks corresponding to eigenvalues of this Hamiltonian.

We write the rf electric field in terms of a complex phasor amplitude

$$\mathbf{E}_{\text{rf}}(t) = \text{Re}[\tilde{\mathbf{E}}_{\text{rf}} e^{-i\omega t}]. \quad (1)$$

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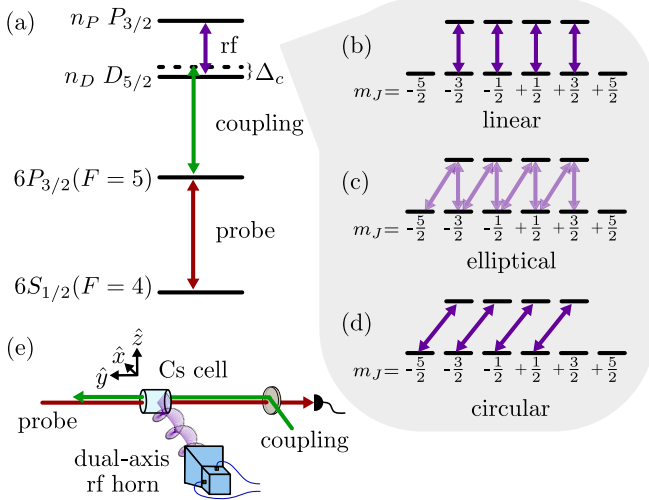


FIG. 1. (a) The energy level diagram for the rf field sensing. The coupling of the two Rydberg states' m_J subspaces are shown for (b) linear, (c) elliptical, and (d) circular polarizations of the rf field. (e) The experimental layout and definition of the coordinate system.

Consider the Hamiltonian in the basis of all m_J sublevels of the $D_{5/2}$ and $P_{3/2}$ Rydberg states coupled by the rf field. If we choose the quantization axis z to be aligned with linear polarization of an rf field [consistent with the coordinate system defined in Fig. 1(e)], we can write the complex rf field as

$$\tilde{E}_{\text{rf}} = E_0(\cos \chi \hat{z} + i \sin \chi \hat{y}), \quad (2)$$

where χ parametrizes the ellipticity of the rf polarization. Here, $\chi = 0$ corresponds to linear polarization along z , which will drive π transitions. In this basis, the spherical components of the field are

$$\epsilon_0 = \cos \chi, \quad \epsilon_+ = -\frac{1}{\sqrt{2}} \sin \chi, \quad \epsilon_- = -\frac{1}{\sqrt{2}} \sin \chi. \quad (3)$$

In this convention ϵ_0 drives π transitions ($\Delta m_J = 0$), while $\epsilon_{\pm}$ drive $\sigma^{\pm}$ transitions ($\Delta m_J = \pm 1$). In the ordered $|J, m_J\rangle$ basis,

$$\begin{aligned} & \{|D, -\frac{5}{2}\rangle, |D, -\frac{3}{2}\rangle, |D, -\frac{1}{2}\rangle, |D, +\frac{1}{2}\rangle, |D, +\frac{3}{2}\rangle, \\ & |D, +\frac{5}{2}\rangle, |P, -\frac{3}{2}\rangle, |P, -\frac{1}{2}\rangle, |P, +\frac{1}{2}\rangle, |P, +\frac{3}{2}\rangle\}, \end{aligned} \quad (4)$$

using the appropriate Wigner 3- j symbols [19], the rotating-wave Hamiltonian (for zero detuning) takes the form

$$H_{\text{RWA}} = \frac{\hbar \Omega_0}{2} \begin{pmatrix} 0 & 0 & 0 & 0 & 0 & 0 & \frac{1}{\sqrt{6}}\epsilon_+ & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & -\frac{1}{\sqrt{15}}\epsilon_0 & \frac{1}{\sqrt{10}}\epsilon_+ & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & \frac{1}{\sqrt{60}}\epsilon_- & -\frac{1}{\sqrt{10}}\epsilon_0 & \frac{1}{\sqrt{20}}\epsilon_+ & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & \frac{1}{\sqrt{20}}\epsilon_- & -\frac{1}{\sqrt{10}}\epsilon_0 & \frac{1}{\sqrt{60}}\epsilon_+ \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & \frac{1}{\sqrt{10}}\epsilon_- & -\frac{1}{\sqrt{15}}\epsilon_0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & \frac{1}{\sqrt{6}}\epsilon_- \\ \frac{1}{\sqrt{6}}\epsilon_+^* & -\frac{1}{\sqrt{15}}\epsilon_0^* & \frac{1}{\sqrt{60}}\epsilon_-^* & 0 & 0 & 0 & -\frac{2\Delta_{\text{rf}}}{\Omega_0} & 0 & 0 & 0 \\ 0 & \frac{1}{\sqrt{10}}\epsilon_+^* & -\frac{1}{\sqrt{10}}\epsilon_0^* & \frac{1}{\sqrt{20}}\epsilon_-^* & 0 & 0 & 0 & -\frac{2\Delta_{\text{rf}}}{\Omega_0} & 0 & 0 \\ 0 & 0 & \frac{1}{\sqrt{20}}\epsilon_+^* & -\frac{1}{\sqrt{10}}\epsilon_0^* & \frac{1}{\sqrt{10}}\epsilon_-^* & 0 & 0 & 0 & -\frac{2\Delta_{\text{rf}}}{\Omega_0} & 0 \\ 0 & 0 & 0 & \frac{1}{\sqrt{60}}\epsilon_+^* & -\frac{1}{\sqrt{15}}\epsilon_0^* & \frac{1}{\sqrt{6}}\epsilon_-^* & 0 & 0 & 0 & -\frac{2\Delta_{\text{rf}}}{\Omega_0} \end{pmatrix}, \quad (5)$$

where Ω_0 is the Rabi rate associated with the reduced dipole matrix element μ_0 ,

$$\Omega_0 = \frac{\mu_0 E_0}{\hbar} = \frac{\langle n_P P_{3/2} \| \mu \| n_D D_{5/2} \rangle E_0}{\hbar}, \quad (6)$$

and Δ_{rf} is the angular frequency detuning of the rf radiation. Throughout this work we consider the resonant case, $\Delta_{\text{rf}} = 0$, corresponding to the experimental conditions. The eigenvalues of this Hamiltonian at zero detuning are shown as a function of the ellipticity χ in Fig. 2(a). Because there are ten states involved in the coherent coupling, there will be ten dressed states with some degeneracy. The distribution of these states are shown for linearly polarized rf in Fig. 2(b) and for circularly polarized rf in Fig. 2(c). For simplicity, the eigenstates are given in the quantization axis for which m_J is a good quantum number.

We can note a few things from the theoretical result. The degeneracy of the dressed eigenstates changes for different χ , meaning the polarization will affect the number of observed

spectroscopic peaks. Furthermore, for pure linear and circular polarizations, each eigenstate is a simple superposition of the two states corresponding to each allowed coupling of the rf.

III. MEASUREMENT

We confirm the energy spectra predicted by Eq. (5) with spectroscopic measurements using the EIT readout scheme from Fig. 1(e). To achieve resolution of the dressed eigenstates, we must ensure field homogeneity. We pick a set of Rydberg states $n_D = 65$, $n_P = 66$ corresponding to an rf transition wavelength ($\lambda = 120$ cm) which is long compared to the vapor cell (2.2 cm diameter by 2.2 cm length with 1 mm wall thickness) in order to minimize internal reflections. In addition, we perform the measurement in an anechoic chamber lined with rf-absorbing foam and mount the vapor cell with masking tape on low-density closed cell foam to minimize scattering. With this and an rf amplifier to apply fields on the

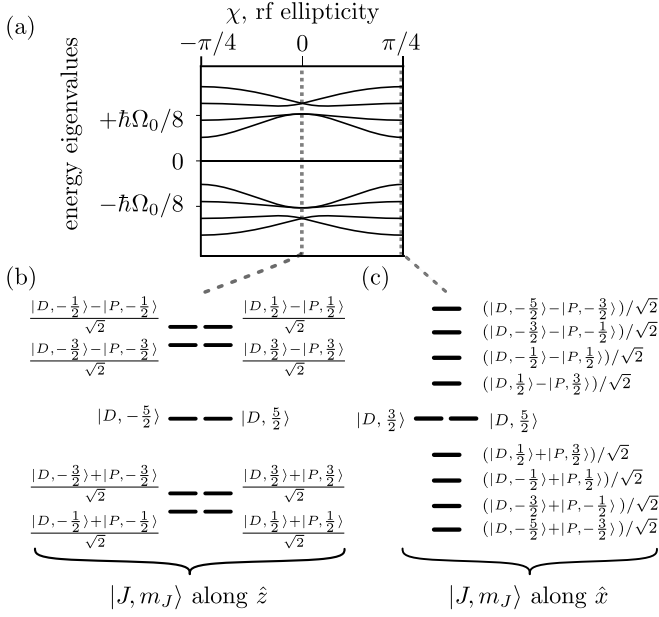


FIG. 2. (a) The energy eigenvalues as a function of the ellipticity χ . (b) An energy level diagram of the dressed eigenstates for linearly polarized rf ($\chi = 0$) in the basis along $\hat{z}$. (c) The dressed eigenstates for circularly polarized rf ($\chi = \pi/4$) in the basis along $\hat{x}$.

order of V/m, it becomes possible to separate the individual dressed eigenstates by more than their linewidth.

We lock the probe on resonance and scan the coupling laser over an angular-frequency detuning Δ_c from resonance [see Fig. 1(a)]. The measured spectra are shown in Fig. 3. The probe laser has a power of 1.6 μ W in a 300- μ m full width at half maximum (FWHM) Gaussian beam, while the coupling

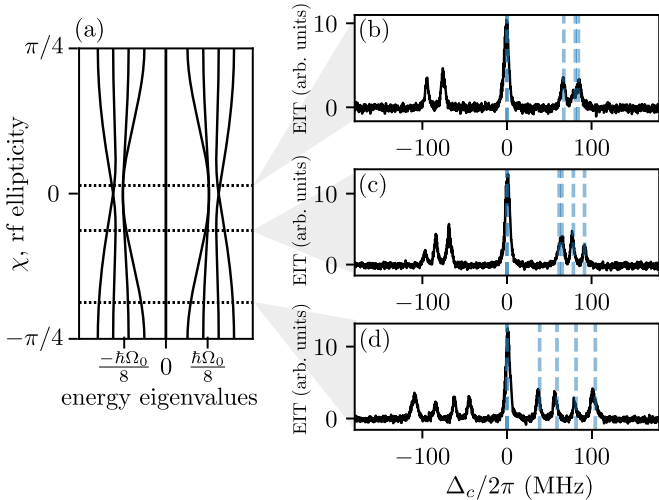


FIG. 3. (a) Eigenvalues of the Hamiltonian from Eq. (5) as a function of the rf ellipticity χ . (b)–(d) Experimental spectroscopy of the Autler-Townes splitting at various ellipticities giving rise to (b) two, (c) three, and (d) four sets of dressed peaks on either side of the central unsplit peak at a field strength of $E_0 = 5$ V/m. The predicted peak locations from the eigenvalues of Eq. (5) are overlaid as dashed lines. The lasers are linearly polarized along $\hat{x}$ in (b) and (c) and along $\hat{z}$ in (d) to populate all m_J .

laser has a power of 70 mW in a 300- μ m FWHM Gaussian beam. To increase the signal-to-noise ratio, the coupling laser power is modulated at 10 kHz with a chopper wheel and the EIT signal is read out from the probe transmission using a lock-in amplifier. The frequency of the probe laser is locked to an ultralow expansion cavity while the coupling laser is scanned. The frequency of the coupling laser is calibrated to the spacing of the $D_{3/2}$ and $D_{5/2}$ states. Drift of the coupling laser during data acquisition is accounted for by aligning each trace to the unsplit central peak. The background magnetic field is shimmed away with permanent magnets placed on the optics table near the vapor cell.

At $\chi = 0$, where the rf polarization is purely linear, the eigenstates are maximally degenerate, leading to only two peaks on either side of the spectrum. However, the degeneracy of these peaks is lifted at impure polarizations. Spectra at various ellipticities are measured in Figs. 3(b)–3(d). The apparent number of peaks varies from two to four with the ellipticity of the rf field.

To control the ellipticity with our dual-axis horn, we set the $\hat{y}$ component to be 90° out of phase with the $\hat{z}$ component. We then scan the relative powers P_z and P_y of the two rf channels to realize the field defined in Eq. (2),

$$P_z = P_0 \cos^2(\chi), \quad P_y = P_0 \sin^2(\chi), \quad (7)$$

where P_0 is the total power sent to the rf horn. The measured spectra are shown as we scan χ continuously in Fig. 4(a).

By repeating the measurement at two laser polarizations and comparing the spectra, we can perform a crude method of quantum state tomography to measure the $|m_J|$ of each peak. While the eigenenergies are independent of the basis/coordinate system, the basis in which m_J is a good quantum number can rotate with the rf polarization. However, because the linearly polarized lasers have a coordinate system that does not change, they will measure m_J along a consistent direction. Consider $\hat{z}$ to be the quantization axis [that in which Eq. (5) is defined]. The polarization of the lasers can be used to preferentially populate different m_J states along $\hat{z}$. When the lasers are linearly polarized in $\hat{z}$, they preferentially populate and measure low $|m_J|$ states because they drive π transitions whose transition dipole moments are stronger going to low $|m_J|$. When the lasers are pointed in $\hat{x}$, they drive σ^+ and σ^- transitions in $\hat{z}$, which have larger dipole moments going to larger $|m_J|$. Therefore, the relative difference in the amplitudes of the spectra is a good qualitative estimation of the angular momentum magnitude $|m_J|$. In Fig. 4, we measure the spectra as a function of the rf ellipticity, which we set via the relative amplitudes of out-of-phase signals on the dual-axis horn [consistent with Eq. (2)]. We then take the difference of the measured spectra to approximate the angular momentum along $\hat{z}$. We then compare this to the expected $|m_J|$ in the D state for each eigenstate [Fig. 4(b)], calculated as

$$\langle |m_J| \rangle_D = \frac{\sum_{m_J} |m_J| \langle D, m_J | \psi \rangle}{\sum_{m_J} \langle D, m_J | \psi \rangle}, \quad (8)$$

where ψ is the rf-dressed eigenstate.

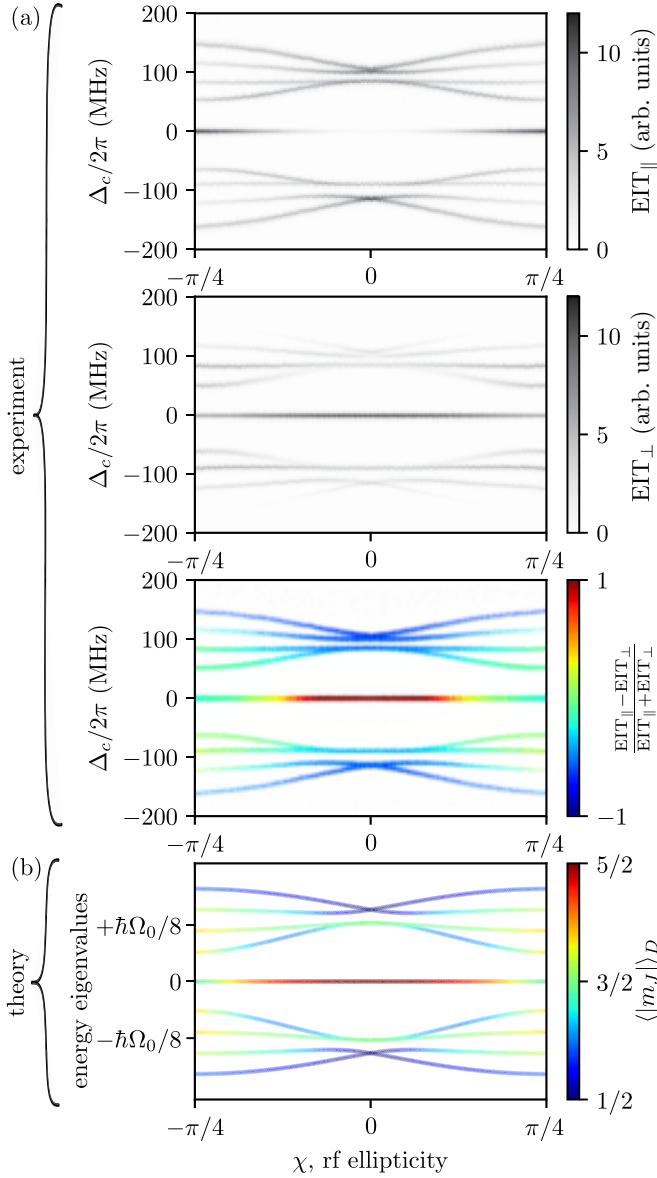


FIG. 4. (a) Measured spectra of the $65D_{5/2} \rightarrow 66P_{3/2}$ transition at $E_0 = 7$ V/m scanning the ellipticity with the dual-axis horn. The EIT spectra is measured with the laser polarizations parallel to $\hat{z}$ ($EIT_{\parallel}$) (top), and perpendicular to $\hat{z}$ ($EIT_{\perp}$) (middle). We then take the normalized difference of the two spectra ($EIT_{\parallel} + EIT_{\perp}$) to qualitatively approximate the angular momentum magnitude $|m_J|$ along $\hat{z}$. An α mask is applied based off of the sum of the spectra $EIT_{\parallel} + EIT_{\perp}$ to show the color map only where an EIT signal is present. (b) The eigenvalues of Eq. (5) are plotted with a color map calculated from the expectation value of $|m_J|$ along $\hat{z}$ in the D state (that which the lasers probe).

IV. PURE rf POLARIZATION

In general, linearly polarized rf radiation is difficult to obtain in the laboratory without special treatment due to scattering of the rf. However, theory and interpretation of spectra can be much simpler for pure linear polarization. Because it does not couple disparate m_J states together, we can treat each m_J ladder as an independent two-level subsystem. Each $|m_J|$ will have a different transition dipole matrix element for the

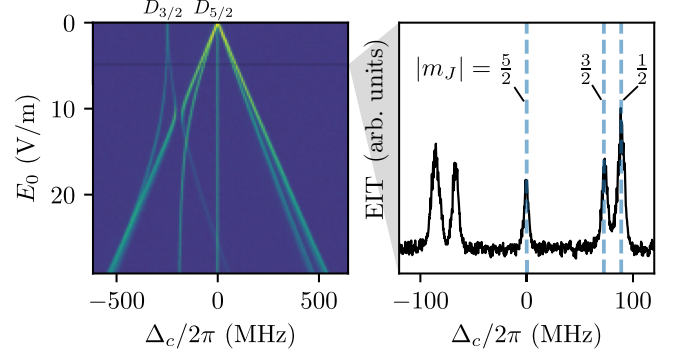


FIG. 5. Measured Autler-Townes spectra with nearly pure linear rf polarization. Left: experimental Autler-Townes spectra varying the rf field strength. Right: an individual spectrum at an E_0 of 4.9 V/m. The predicted peak locations from Eq. (9) are overlaid as dashed lines. In the plot on the right, we reduce the Δ_c -axis scale to only highlight the $D_{5/2}$ states.

same rf coupling and will therefore split a different amount. Looking at the couplings in Fig. 1(b), we can see that there are three distinct values of $|m_J|$. The highest value $|m_J|$ will not split, as there is no $|m_J| = 5/2$ in the P state to couple to. The other two couplings are given by angular momentum algebra via the appropriate Wigner 3- j symbol [19]:

$$\Omega = \frac{E_0}{\hbar} \cdot \begin{cases} 0, & |m_J| = 5/2, \\ \frac{1}{\sqrt{15}}\mu_0, & |m_J| = 3/2, \\ \frac{1}{\sqrt{10}}\mu_0, & |m_J| = 1/2. \end{cases} \quad (9)$$

We map out Autler-Townes spectra for a π -polarized rf field in Fig. 5. Here, the rf and probe lasers are polarized in $\hat{z}$, and we rotate the coupling laser polarization slightly off from $\hat{z}$ towards $\hat{x}$ to populate all m_J levels. We find excellent agreement between the measured peak locations and the splittings predicted by Eq. (9), with the peak locations matching theory to within one percent. This line shape is consistent with that predicted by Ref. [8]. Note that at large fields, we observe an avoided crossing between the dressed $D_{5/2}$ and $D_{3/2}$ states. This interaction perturbs the negative detuning side of the spectrum in all of the provided measurements.

While rf circular polarization is less common, it also can be simplified into individual splittings, going as

$$\Omega = \frac{E_0}{\hbar} \cdot \begin{cases} \frac{1}{\sqrt{6}}\mu_0, & m_J = -5/2 (\rightarrow -3/2), \\ \frac{1}{\sqrt{10}}\mu_0, & m_J = -3/2 (\rightarrow -1/2), \\ \frac{1}{\sqrt{20}}\mu_0, & m_J = -1/2 (\rightarrow 1/2), \\ \frac{1}{\sqrt{60}}\mu_0, & m_J = 1/2 (\rightarrow 3/2), \\ 0, & m_J = 3/2, \\ 0, & m_J = 5/2. \end{cases} \quad (10)$$

V. CONCLUSION

We have shown that the structure of Autler-Townes spectra in Rydberg atoms is governed by the polarization of the driving rf field through its coupling of magnetic sublevels. In contrast to the conventional picture of independent m_J transi-

tions, elliptical polarization induces coherent coupling across the full manifold of sublevels, requiring a complete multilevel treatment. By constructing and diagonalizing the corresponding Hamiltonian, we predict polarization-dependent degeneracies that give rise to spectra with two, three, or four resolved peaks on either side of the central unsplit peak.

Using a controlled experimental environment that minimizes rf inhomogeneity and polarization impurities, we directly resolve the magnetic-sublevel structure of the rf-dressed Rydberg states. This constitutes a clear experimental resolution of individual m_l -dependent features in resonant Rydberg Autler-Townes spectra. The measured peak locations and their evolution with rf ellipticity show excellent agreement with the eigenvalues of the full Hamiltonian. Furthermore, by comparing spectra taken with orthogonal laser polarizations, we extract qualitative information about the angular momentum composition of the dressed states.

These results resolve long-standing discrepancies between theory and experiment in Rydberg Autler-Townes measurements and demonstrate that rf polarization is a critical parameter in interpreting spectroscopic features. More broadly, this work establishes a framework for incorporating magnetic sublevel structure into Rydberg-based sensing, with

direct implications for improving the accuracy and traceability of rf electrometry and enabling polarization-resolved measurements in atomic rf sensors.

Note added. Recently, we became aware of another work [20] demonstrating similar results, with great experimental and theoretical agreement to the findings presented in this work.

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The authors declare no conflict of interest.

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DATA AVAILABILITY

All of the data presented in this paper and used to support the conclusions of this article is openly available [21].

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- Correction:* An explanatory note in Ref. [12] was omitted during the proof process and has been inserted.