



Autler–Townes doublet in single-photon Rydberg spectra of cesium atomic vapor with a 319 nm UV laser

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Abstract

We demonstrate the single-photon excitation spectra of cesium Rydberg atoms by means of a Doppler-free purely all-optical detection with a room-temperature vapor cell and a 319 nm ultra-violet (UV) laser. We excite atoms directly from $6S_{1/2}$ ground state to $71P_{3/2}$ Rydberg state with a narrow-linewidth 319 nm UV laser. The detection of Rydberg states is performed by monitoring the absorption of an 852 nm probe beam in a V-type three-level system. With a strong coupling light, we observe the Autler–Townes doublet and investigate experimentally the dependence of the separation and linewidth on the coupling intensity, which is consistent with the prediction based on the dressed state theory. We further investigate the Rydberg spectra with an external magnetic field. The existence of non-degenerate Zeeman sub-levels results in the broadening and shift of the spectra. It has potential application in sensing magnetic field.

1 Introduction

Rydberg atoms, highly-excited atoms with principal quantum numbers $n > 10$, have shown great promise in terms of their unique structure. These physical properties scale with principal quantum number n , such as long radiation lifetime ($\sim n^3$), strong dipole–dipole interaction ($\sim n^4$) and large polarizability ($\sim n^7$) [1]. It makes Rydberg atoms promising for applications such as quantum information processing [2], quantum metrology [3, 4] and simulations of quantum many-body physics [5]. To exploit coherent quantum dynamics, these experiments are performed on timescales shorter than the lifetime of Rydberg state. However, for the realization of supersolids [6], frustrated quantum magnetism [7] or spin squeezing for enhanced metrology [8], it is necessary to extend the investigation time of Rydberg atoms. To take advantage of the long coherence time of ground-state atoms and the strong long-range interaction between

Rydberg atoms, we can create a wave function that is mostly ground state with an adjustable Rydberg component by off-resonantly coupling the ground state to Rydberg state [9–13], which is called Rydberg dressing approach. Therefore, experimental research of Autler–Townes (AT) splitting will open a route to produce the dressed-state atoms with the long coherence time and controllable long-range interaction by adjusting the coupling intensity and detuning between the ground state and Rydberg state. Besides, AT splitting [14] can be applied to study various physical phenomena, such as four-wave mixing [15], slow light [16], dephasing rates of Rydberg states [17], enhancement of spin–orbit interaction [18], and the measurement of the interval of atomic hyperfine levels. Moreover, the AT splitting can be also used to measure the Rabi frequency [19].

Most studies about AT splitting are commonly performed on the Λ - and ladder-type three-level system [20, 21], which consists of the ground state and the low excited state. For Rydberg state, the relevant research is rare, especially in V-type three-level system. Most experiments about AT splitting of Rydberg spectra are demonstrated by a two-photon excitation process in a ladder-type three-level system [22, 23]. However, due to the very short lifetime of the intermediate state in the two-photon scheme, the Rydberg-dressed experiment is generally implemented by a single-photon excitation [10–13]. Bounds et al. [19] demonstrated Rydberg-dressed magneto-optical trap of strontium atoms by off-resonantly coupling a long-lifetime metastable state

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to a high-lying Rydberg state with a 319 nm UV laser. The trapping dressed-state atoms had enhanced electric polarizability. Recently, Arias et al. [13] realized a Rydberg-dressed potassium atomic clock with sub-kilohertz accuracy in a far detuned optical dipole trap by a 286 nm UV laser via the single-photon transition. Therefore, to implement the follow-up Rydberg-dressed experiment of cesium atoms, it is important to perform the AT splitting of Rydberg spectra via single-photon excitation. Thoumany et al. [24] demonstrated the single-photon Rydberg excitation spectroscopy with a 35 mW output of 297 nm UV dye laser in a room-temperature Rb atomic vapor cell. However, it is difficult to observe the AT splitting of Rydberg spectra with high signal-to-noise ratio and then perform Rydberg-dressed experiment by utilizing the low-power UV light.

In addition, combining the sensitivity of Rydberg atoms to the external field and a three-level system, high-precision Rydberg spectra related to the Zeeman effect play an important role in quantum information processing [25]. Bao et al. [26] demonstrated the importance of the laser-field polarizations in Rydberg EIT, and modeled qualitatively the Zeeman spectra using standard density matrix equations. Level populations and coherences in ladder-type Rydberg EIT systems were mainly affected by the Zeeman shifts. Recently, Cheng et al. [27] observed Rydberg spectra of ladder-type EIT in a magnetic field at the room-temperature ^{87}Rb vapor. These experiments were performed by a two-photon excitation process via an intermediate state, resulting in complexity of Rydberg Zeeman spectra. Moreover, compared with single-photon Rydberg excitation, the two-photon scheme has following disadvantages: atomic decoherence via photon scattering of the upper and lower transitions, and AC Stark shifts due to the upper and lower excitation laser beams. For alkali metal atoms, direct excitation from the ground state to the desired Rydberg state usually requires high-power UV laser which is now not commercially available, therefore, the single-photon Rydberg excitation experiments are rare. In recent years, with the quickly-developed fiber lasers, fiber amplifiers, and the efficient nonlinear laser frequency conversion technique based on quasi-phase matching materials, high-power narrow-linewidth continuous-wave UV laser [28, 29] has been implemented to meet actual demands of Rydberg excitation of neutral atoms by single-photon transition in the field of atomic physics.

In this paper, using the Doppler-free purely all-optical detection method and the home-made narrow-linewidth Watt-level 319 nm UV laser, we perform the single-photon Rydberg excitation for $6S_{1/2} (F=4) \rightarrow 71P_{3/2}$ transition in a room-temperature Cs atomic vapor cell. With a strong coupling laser beam, we observe the AT splitting of Rydberg spectra and investigate experimentally the dependence of the separation and linewidth upon the coupling intensity, which is explained by the dressed state theory. Furthermore,

we investigate Rydberg spectra with an external magnetic field, and analyse the dependence of the frequency shift and linewidth upon magnetic field.

2 Experimental arrangement

The lifetime of Rydberg state increases with principal quantum number. Considering vacuum-induced radiative decay and blackbody radiation effects [28], the lifetime of Cs $71P_{3/2}$ state is $185 \mu\text{s}$, resulting in a natural linewidth of $\sim 2\pi \times 860 \text{ Hz}$. The oscillator strength of $6S_{1/2} \rightarrow 71P_{3/2}$ transition is six orders of magnitude smaller than that of $6S_{1/2} \rightarrow 6P_{3/2}$ transition. Compared to $6P_{3/2}$ excited state, the lower transition probability and the smaller light-scattering cross sections lead to very weak absorption signals for $71P_{3/2}$ Rydberg state. To enhance the signal, the Rydberg transition is detected by monitoring the reduced absorption of the 852 nm probe beam. A schematic diagram of the experimental setup is shown in Fig. 1a. The narrow-linewidth UV laser is produced by the cavity-enhanced second harmonic generation following sum-frequency generation of two infrared lasers at 1560 nm and 1077 nm, leading to more than 2 W output power at

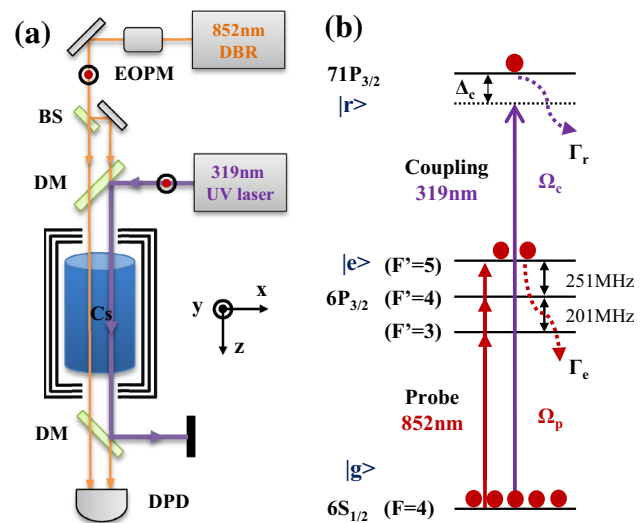


Fig. 1 **a** Schematic of the experimental arrangement. The 319 nm coupling beam co-propagates with the 852 nm probe beam in a 10-cm-long Cs vapor cell placed in a μ -metal tank. DBR distributed-Bragg-reflector diode laser, EOPM electro-optic phase modulator, BS beam splitter, DM dichroic mirror, DPD differential photodiode. **b** Schematic diagram of a V-type three-level system for Cs atomic single-photon Rydberg excitation. The coupling and probe beams have Rabi frequency Ω_c and Ω_p , respectively. The Γ_i ($i=r$ or e) denote the decay rates from the respective levels. The 852 nm probe beam is resonant on the transition from $6S_{1/2} (F=4)$ to $6P_{3/2} (F'=3, 4, \text{ or } 5)$, and the 319 nm coupling beam is scanned over $6S_{1/2} (F=4) \rightarrow 71P_{3/2}$ Rydberg transition

319 nm [29, 30]. The 852 nm probe beam is provided by a distributed-Bragg-reflector (DBR) diode laser which is frequency stabilized to Cs $6S_{1/2}$ ($F=4$) \rightarrow $6P_{3/2}$ ($F'=5$, or 4, or 3) hyperfine transition by means of polarization spectroscopy [31]. A waveguide-type electro-optic phase modulator (EOPM) (Photline, NIR-MX800) at 852 nm is used to calibrate the frequency interval. The probe beam has a $1/e^2$ radius of ~ 350 μm while the coupling beam is slightly larger with a $1/e^2$ radius of ~ 400 μm . The probe light is split into two beams of equal intensity via a beam splitter. One of them is superposed with the 319 nm coupling beam using a dichroic mirror, and co-propagate through a 10-cm-long fused-quartz Cs cell placed in a μ -metal tank. This tank reduces the magnetic field to less than 0.2 mG. To eliminate the Doppler broadened background, the transmission of the probe beams are detected with a differential photodiode (DPD) (New Focus, Model 2107) after passing through another dichroic mirror. When the UV light is scanned across a specific single-photon Rydberg transition, the atoms on ground state are partially transferred to the long-lifetime Rydberg state, leading to the reduced absorption of the probe beam overlapped with the coupling beam.

Figure 1b shows a V-type three-level system interacting with two light fields composed of the ground state $|g\rangle$ [$6S_{1/2}$ ($F=4$)], the excited state $|e\rangle$ [$6P_{3/2}$ ($F'=3, 4, 5$)], and a Rydberg state $|r\rangle$ [$71P_{3/2}$]. For the V-type system, the common ground $|g\rangle$ is coupled to two excited states: $|g\rangle \leftrightarrow |e\rangle$ transition by the probe light and $|g\rangle \leftrightarrow |r\rangle$ transition by the UV coupling light. The coupling beam has Rabi frequency of Ω_c and detuning of Δ_c from resonance. The probe beam has Rabi frequency Ω_p and detuning Δ_p . The decay rates from the excited state $|e\rangle$ and Rydberg state $|r\rangle$ are Γ_e and Γ_r , respectively. Here Γ_e is $2\pi \times 5.2$ MHz.

If the probe light is locked on $6S_{1/2}$ ($F=4$) \rightarrow $6P_{3/2}$ ($F'=5$) cycling transition, the atoms with zero-velocity component along the propagation direction of the probe beam populate in the $6P_{3/2}$ ($F'=5$) state, as shown in Fig. 1b. Considering the Doppler effect in a thermal ensemble, the $6P_{3/2}$ ($F'=4$ and 3) states are also populated atoms with a certain velocity component in the same direction. For atoms with a velocity component of v_z , the detuning of probe and coupling fields in the same direction are expressed as:

$$\begin{aligned} \Delta_p &= k_p v_z = -\frac{\omega_0}{c} v_z \\ \Delta_c &= k_c v_z = -\frac{\lambda_p}{\lambda_c} \Delta_p, \end{aligned} \quad (1)$$

where ω_0 is the reference frequency; c is the speed of light in vacuum; λ_p , λ_c are the wavelengths of the probe and coupling beams, respectively.

3 Results and discussion

3.1 Single-photon Rydberg excitation velocity-selective spectra

In general, the interaction between the atoms and the coupling light field is studied by the absorption of probe beam. In this way, we can obtain the coherent spectra without Doppler background, and also improve the spectral resolution. Since the frequency interval between $6P_{3/2}$ ($F'=3$) and ($F'=4$) is 201 MHz which is much larger than the laser linewidth of 1 MHz, and the probe light is locked to the $6S_{1/2}$ ($F=4$) \rightarrow $6P_{3/2}$ ($F'=5$) cycling transition, the absorption of the probe beam attributes mainly to the atoms with velocity of $v_z=0$ in the direction of beam propagation. When the probe light with a certain intensity passes through Cs vapor cell, it will immediately reach a population balance between the ground state and $6P_{3/2}$ excited state, and a stable transmission signal is obtained simultaneously. Only when some atoms are transferred from the ground state to the $71P_{3/2}$ Rydberg state, the balance between the ground state and $6P_{3/2}$ excited state is broken, and the absorption of the probe light is reduced. Therefore, the absorption spectra of single-photon Rydberg excitation is observed by the enhanced transmission of the probe beam while scanning the frequency of the coupling beam.

Figure 2 shows the single-photon excitation spectra for the $71P_{3/2}$ Rydberg state when the probe light is locked to different hyperfine transitions from Cs $6S_{1/2}$ ($F=4$) to $6P_{3/2}$. The probe and coupling beams have identical

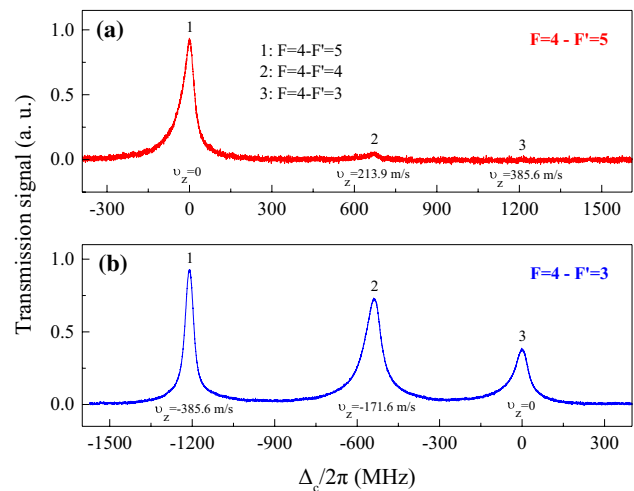


Fig. 2 The single-photon Rydberg excitation velocity-selective spectra for the $6S_{1/2}$ ($F=4$) \rightarrow $71P_{3/2}$ transition when the probe light is locked to different transitions: **a** Cs $6S_{1/2}$ ($F=4$) \rightarrow $6P_{3/2}$ ($F'=5$) hyperfine transition, **b** Cs $6S_{1/2}$ ($F=4$) \rightarrow $6P_{3/2}$ ($F'=3$) hyperfine transition

linear polarization. Figure 2a, b correspond to the probe light are resonant with $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=5)$ and $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$ transitions. When the UV light is scanned over the $71P_{3/2}$ state, three transmission peaks (marked as 1, 2, 3) appear and correspond to $(F=4) \rightarrow (F'=5)$, $(F=4) \rightarrow (F'=4)$, and $(F=4) \rightarrow (F'=3)$ hyperfine transitions. Here, we use an EOPM to produce two sidebands to calibrate the frequency interval of the transmission spectra. Setting $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=5)$ cycling transition as the reference frequency, for atoms with velocity of $v_z=0$, the detuning of the probe light is $\Delta_p=0$; for atoms with velocity groups of $v_z=213.9$ and 385.6 m/s, Δ_p are -251 and -452 MHz given by Eq. (1), which will result in resonance with $(F=4) \rightarrow (F'=4)$ and $(F=4) \rightarrow (F'=3)$ hyperfine transitions, respectively. Considering the wavelength mismatch of $\lambda_p/\lambda_c=2.675$ between the probe and coupling beams, when the UV beam matches these atomic groups to the $71P_{3/2}$ state, the corresponding detuning Δ_c are 671 and 1210 MHz, which are consistent with the measurement values with a high-precision wavelength meter (HighFinesse, WS-7), as shown in Fig. 2a. For Fig. 2b, the interacting process is also similar and it does not vary with different $nP_{3/2}$ Rydberg states ($n=70-100$).

Furthermore, we have observed that the relative intensity of the three transmission peaks (marked as 1–3) is different when the probe light is resonant with different atomic transitions. There are two main reasons: the strength of the interaction between atoms and nearly-resonant laser, and the atoms with a certain velocity interacting with lasers. The strength of the interaction between two Zeeman sub-levels is characterized by the dipole matrix elements [32]:

$$\langle JIFm_F | e r_q | J'I'F'm_{F'} \rangle = \left\langle J \left\| e \vec{r} \right\| J' \right\rangle (-1)^{2F'+m_F+J+I} \sqrt{(2F+1)(2F'+1)(2J+1)} \begin{pmatrix} F' & 1 & F \\ m_{F'} & q & -m_F \end{pmatrix} \begin{Bmatrix} J & J' & 1 \\ F' & F & I \end{Bmatrix}. \quad (2)$$

Here, $q = m_{F'} - m_F$ is polarization dependent, labeling the component of \mathbf{r} in the spherical basis, and the doubled bars indicate that the matrix element is reduced. I is the nuclear spin momentum, J is the angular momentum, F is the total angular momentum, and m_F is the Zeeman sub-levels. The values in the parentheses and curly brackets denote the Wigner 3- j symbol and Wigner 6- j symbol, respectively. Because the two beams have identical linear polarization, we only calculate the relative transition strengths between two Zeeman sub-levels for π transitions $[(F=4, m_F) \rightarrow (F'=5, 4, 3, m_{F'}=m_F)]$, the corresponding ratio of the three transmission peaks should be 44:21:7. In fact, due to laser interaction with the atoms of different velocities, the relative strength of these transmission peaks does not perfectly satisfy this condition.

When the physical system is in thermal equilibrium, most of the atoms are in the ground state, and its velocity distribution is described by the Maxwell–Boltzmann distribution. Taking the transmission peak (marked as 3) as an example, when the detuning of the UV beam is relatively small, the transmission enhancement of the signal becomes more obvious with the increase of the number of atoms interacting with the laser in the direction of beam propagation. On the contrary, for the peak-1 in the red detuning, the relative strength of the transmission signal will gradually decrease with the increase of the detuning, as shown in Fig. 2.

3.2 Autler–Townes splitting

We further investigate the transmission spectra for different coupling intensities with the probe light of ~ 43 μW , as shown in Fig. 3. The 852 nm probe light is locked to the hyperfine transition of $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$. The transmission peak (marked as 3) corresponds to Cs $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$ hyperfine transitions, as depicted in Fig. 2b. With a moderate intensity of the coupling light, we can see that the AT doublet appears in the transmission spectra. The separation of the doublet increases with the enhancement of the coupling light, and its linewidth is also gradually widened.

For the above physical process, we first consider a simple V-type three-level system, as shown in Fig. 1b. When the strong coupling light interacts with a two-level system ($|g\rangle \leftrightarrow |r\rangle$ transition), the atomic energy level is split into two which is induced by dressing splitting, resulting in a new resonance. The absorption of the probe beam is proportional to $\text{Im}(\rho_{eg})$, where ρ_{eg} is the induced polarizability

on the $|g\rangle \leftrightarrow |g\rangle$ transition coupled by the probe light. From the density-matrix equations, the steady-state value of ρ_{eg} is given by [33]:

$$\rho_{eg} = i\Omega_p \frac{(\gamma_{rg}\Gamma_{rg} + 2\Omega_c^2)(i\Delta_p + \gamma_{rg})}{[(i\Delta_p + \Gamma_{eg})(i\Delta_p + \Gamma_{re}) + \Omega_c^2](\gamma_{rg}\Gamma_{rg} + 4\Omega_c^2)}, \quad (3)$$

where γ_{ij} and Γ_{ij} denote the spontaneous emission rates and the decay rates from $|i\rangle$ to $|j\rangle$, respectively. In our case, the $|r\rangle \leftrightarrow |e\rangle$ transition is dipole-forbidden, and Γ_{re} represents higher-order transition. The theoretical analysis process is similar to that of the Λ -system [20], showing that the absorption spectra of probe beams are separated into two peaks located at the position of two dressed states, named the AT doublet. Their locations are given by [34]:

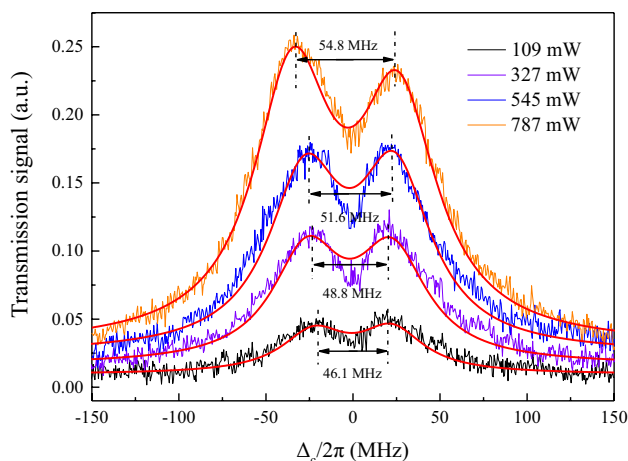


Fig. 3 The AT splitting spectra at different coupling intensities for the certain intensity of probe light. The probe light is locked to the hyperfine transition of $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$. The transmission peak corresponds to Cs $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$ hyperfine transition. The red solid lines are the results of the multiple-peaks fitting of Lorentz function

$$\Delta_{\pm} = \frac{A_c}{2} \pm \frac{1}{2} \sqrt{4\Omega_c^2 + 4\Omega_c^2}, \tag{4}$$

with the corresponding linewidths:

$$\Gamma_{\pm} = \frac{\Gamma_{rg} + D}{2} \left(1 \mp \frac{A_c}{\sqrt{4\Omega_c^2 + 4\Omega_c^2}} \right) + W. \tag{5}$$

Here, D is the Doppler width and W is the other broadenings analyzed in the following section. It is clear from the above expression that, if $\Delta_c \approx 0$, the location of the two peaks have $\pm \Omega_c$ symmetrically shifted away from the atomic hyperfine transition and they have identical linewidth of $(\Gamma_{rg} + D)/2$. In this situation, the separation between the AT doublet uniquely depends on the coupling Rabi frequency. In our experiment, the probe light is locked to Cs $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$ transition, therefore, the transmission signal obtained with DPD is mainly caused by zero velocity atoms which have a certain probability to absorb the coupling and probe beams.

Figure 4 shows that the separation and linewidth of the AT doublet as a function of the coupling beam’s power. The relation between the effective Rabi frequency and the power depends on many factors such as the intensity profile of the coupling beam, absorption through the cell, scattering loss at the surface of the cell. Moreover, the measured lifetime of the Rydberg state is much shorter than the calculated value due to atom loss induced by an external decoherence mechanism [1, 28], such as laser irradiation of surfaces. It is difficult to determine experimentally these factors precisely, therefore,

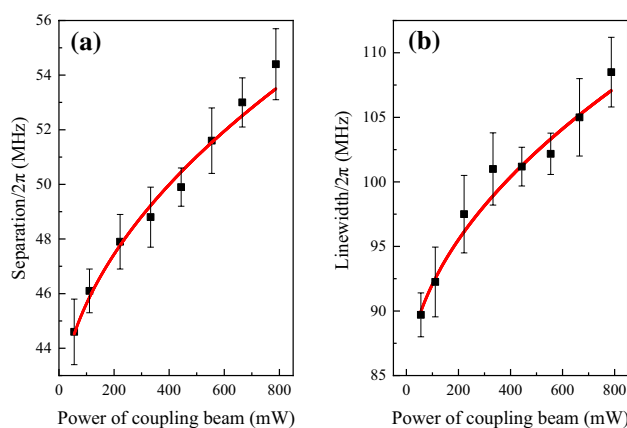


Fig. 4 The separation and linewidth of the peak-3 in the AT doublet of single-photon Rydberg spectra as a function of the coupling beam’s power, while the probe light is resonant with Cs $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=3)$ transition. The solid lines are the expected variation from the theoretical analysis, as explained in the text. The error bar is obtained by fitting the AT doublet, including the fluctuation of the laser beams and the asymmetry of the spectrum due to the nonlinear sweep of the triangular wave for the UV laser

we have left it as an overall fit parameter to obtain the solid lines in Fig. 4. It is important to note that this parameter does not change the trend of the data and its shape, but only makes the fitted curve to move up or down. The best fitting results show that the coupling beam has an ideal Gauss radius of $430 \mu\text{m}$, which is slightly larger than the measured value of $400 \mu\text{m}$. Considering the experimental condition of $\Delta_c \approx 0$ and the existing broadening mechanisms described below in the text, the solid line in Fig. 4a is only the variation predicted by Eq. (4). The trend of experimental data is in good agreement with the theoretical prediction. In Fig. 4b, with the fixed intensity of the probe beam, the linewidth of the AT doublet increases with the coupling beam’s power, which is mainly caused by the power broadening. The power-broadened linewidth varies as $w = \Gamma_{\text{Nat}}(1 + s)^{1/2}$, where $s = III_s$ is the saturation parameter. Taking into account various complex broadening mechanisms, the expression can be written as:

$$w = \Gamma_{\text{Nat}} \sqrt{1 + s} + W. \tag{6}$$

From the fitting results shown in Fig. 4b, we estimate that the measured narrowest linewidth is about 83.8 MHz (not zero) without the coupling beam. We first attribute the remaining linewidth to power broadening by the probe beam. Considering the wavelength mismatch of $\lambda_p/\lambda_c = 2.675$ between the probe and coupling beams, we calculate that the power broadened linewidth of probe beam for one of the doublets is 32.2 MHz , which is smaller than the measured value. Therefore, there exist other broadening mechanisms in the V-type Rydberg system, not only the power broadening. It also contributes from the following broadening mechanisms:

natural linewidth of ~ 14 MHz ($\sim 2.675 \times 5.2$ MHz) [24], transit-time broadening of ~ 450 kHz ($\sim 2.675 \times 168$ kHz), collisional broadening [35], misalignment of the probe and coupling beams in Cs atomic vapor cell [20], the long-range interactions between Rydberg atoms, and the broadening caused by stray electric fields. The fact that an applied dc electric field is not effective on the EIT spectra regardless of the probe and coupling laser polarizations because Rydberg atoms are being screened in the vapor cell [36]. The shielding effect is attributed to ions and electrons produced by photodesorption induced by the coupling beam at the cell inner surfaces. In the presence of an additional radio-frequency field, the shielding effect persists and is reduced by an external high-frequency field (higher than dozens of MHz) [37]. Therefore, the spectral lines broadening from the Stark shifts with the ambient static and low-frequency alternating electric field is negligible. To reduce the interference of high-frequency electric field, we will use the shielding material with good conductivity and being grounded in the future. Considering the above broadening mechanisms, the observed linewidth is well fitted. However, it is important to note that the intensity of the probe light is relatively strong to observe the Rydberg spectra with high signal-to-noise ratio. Subsequently, we need to improve the signal-to-noise ratio

$$E_{|Jm_J; Im_I\rangle} \approx A_{\text{hfs}} m_I m_J + B_{\text{hfs}} \frac{9(m_I m_J)^2 - 3J(J+1)m_I^2 - 3I(I+1)m_J^2 + I(I+1)J(J+1)}{4J(2J-1)I(2I-1)} + \mu_B (g_J m_J + g_I m_I) B. \tag{10}$$

of the signal and suppress spectral broadening of the weak excitation signal which is critical for precision spectroscopy.

3.3 Rydberg spectra in an external magnetic field

Hyperfine structure is a result of the electron–nucleus interactions. Each of the hyperfine energy levels contains $2F + 1$ magnetic sub-levels that determines the angular distribution of the electron wave function. Without an external magnetic field, these Zeeman sub-levels are degenerate. However, when the magnetic field is applied, their degeneracy is broken. The total Hamiltonian of the system can be expressed as [32]:

$$H = H_{\text{hfs}} + H_B. \tag{7}$$

Using the first-order perturbation theory, the Hamiltonian of hyperfine structure and that of describing the interaction of atoms with magnetic fields are expressed as, respectively:

$$H_{\text{hfs}} = A_{\text{hfs}} \mathbf{I} \cdot \mathbf{J} + B_{\text{hfs}} \frac{3(\mathbf{I} \cdot \mathbf{J})^2 + \frac{3}{2}(\mathbf{I} \cdot \mathbf{J}) - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}$$

$$H_B = \frac{\mu_B}{h} (g_S \mathbf{S} + g_L \mathbf{L} + g_I \mathbf{I}) \cdot \mathbf{B}. \tag{8}$$

Here, A_{hfs} is the magnetic dipole constant, and B_{hfs} is the electric quadrupole constant. \mathbf{I} , \mathbf{J} are the total nuclear and electronic angular momentum, respectively. For ^{133}Cs , $\mathbf{I} = 7/2$, $\mu_B/h \approx 1.4$ MHz/Gauss is Bohr magneton, g_S , g_L , g_I are the electron spin, electron orbital and nuclear g -factor, respectively.

According to the interaction strength between the atoms and magnetic field, for our V-type three-level system, the splitting of the ground state $|g\rangle$ ($6S_{1/2}$) is linear Zeeman effect in the weak field, but that of the Rydberg state $|r\rangle$ ($71P_{3/2}$) is in the strong Paschen–Back region. For the Rydberg state, the decoupled basis ($|Jm_J; Im_I\rangle$) is used as the eigenbasis while the coupled basis ($|JIFm_F\rangle$) is the eigenbasis for $|g\rangle$. The dipole matrix elements in the coupled basis have been given in Eq. (2). The decoupled and coupled basis are related by [36]:

$$|Jm_J; Im_I\rangle = (-1)^{J-I+m_F\sqrt{2F+1}} \times \sum_{Fm_F} \begin{pmatrix} J & I & F \\ m_J & m_I & -m_F \end{pmatrix} |JIFm_F\rangle. \tag{9}$$

Using the first-order perturbation theory, the Zeeman splitting of the Rydberg state describing the interaction of atoms with magnetic fields is given by [32, 38]:

From the above Eqs. (8, 10), the shifts of three levels ($|g\rangle$, $|e\rangle$, and $|r\rangle$) are different because of different magnetic interactions, just as described in reference [39]. The absolute position of the spectrum varies with the detuning which is related to Zeeman splitting of different states in an external magnetic field. The Zeeman splitting of the ground state varies linearly with the magnetic field. For Rydberg states in the strong Paschen–Back region, the splitting also linearly depends on the magnetic fields in $|m_J, m_I\rangle$ basis; for the atoms in $6P_{3/2}$ excited state, the magnetic interaction is linear and quadratic Zeeman effect. When the magnetic field is small ($B \lesssim 10$ Gauss), it is approximately in the linear range of Zeeman effect. In the magnetic field, for the ground state (Δ_g), first excited state (Δ_e) and Rydberg state (Δ_r), the interval of adjacent Zeeman sub-levels can be expressed as:

$$\Delta_g = \mu_B B g_F m_F = \frac{1}{4} \mu_B B m_F$$

$$\Delta_e = \mu_B B g_{F'} m_{F'} = \frac{2}{5} \mu_B B m_{F'}$$

$$\Delta_r = \mu_B B g_J m_J = \frac{4}{3} \mu_B B m_J. \tag{11}$$

The detuning of probe and coupling beams in the magnetic field relative to the transition of $|g\rangle \leftrightarrow |e\rangle$ and $|g\rangle \leftrightarrow |r\rangle$ are $\Delta_p = \Delta_e - \Delta_g$ and $\Delta_c = \Delta_r - \Delta_g$, respectively. Considering the wavelength mismatch induced by Doppler effect, when the probe beam is resonant with the field-free $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=5)$ transition, the resonance condition satisfies the relation:

$$\Delta_c = (\Delta_r - \Delta_g) + \frac{\lambda_p}{\lambda_c} (\Delta_e - \Delta_g). \tag{12}$$

In Eq. (10), the fine energy level of $71P_{3/2}$ Rydberg state is split into four Zeeman sub-levels of $m_j = \pm 1/2, \pm 3/2$ in an external magnetic field. The Zeeman splitting of the $71P_{3/2}$ state is shown in Fig. 5. As shown in Fig. 1b, both the probe and coupling beams are linearly polarized along the y direction, and co-propagate passing through Cs cell, which is contained in a cylindrical solenoid. The coupling and probe beams travel in the direction of the magnetic field (parallel to the z direction). Figure 6 shows that the transmission signals of the probe beam in different strengths of magnetic field B from 0 to 6.8 Gauss. With the increase of magnetic field, the single-photon Rydberg excitation spectra

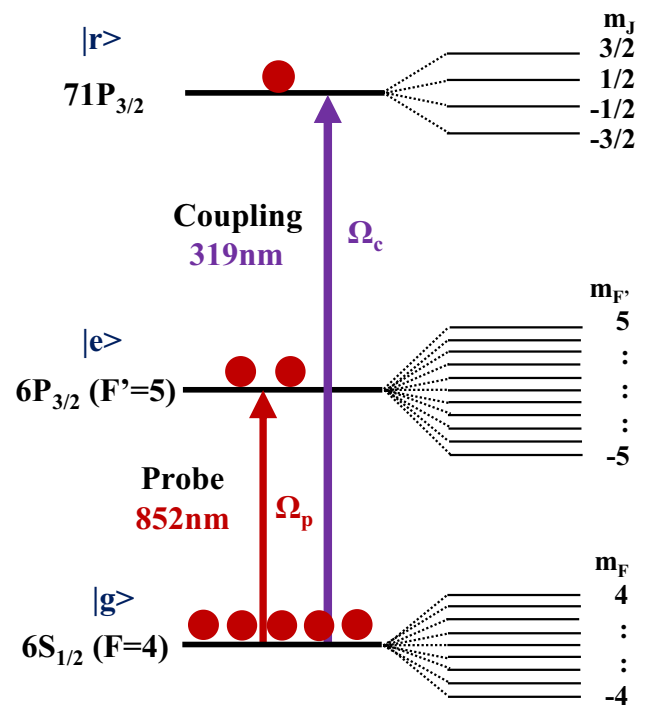


Fig. 5 The V-type scheme of the single-photon Rydberg excitation without (left) and with (right) magnetic field. The 852 nm probe beam is resonant with the field-free $6S_{1/2} (F=4) \rightarrow 6P_{3/2} (F'=5)$ transition, and the 319 nm coupling beam scans through the $6S_{1/2} (F=4) \rightarrow 71P_{3/2}$ Rydberg transition. The sketch on the right shows the Zeeman sub-levels of $|g\rangle$, $|e\rangle$, and $|r\rangle$ states when a magnetic field less than 10 Gauss is applied (separations not to scale)

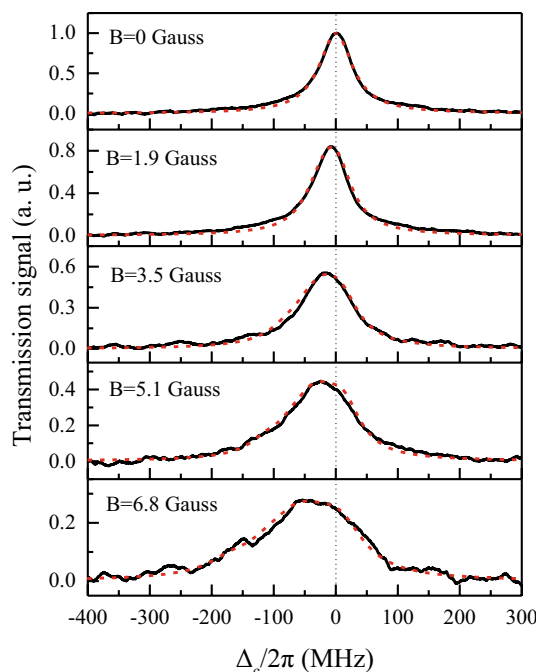


Fig. 6 The transmission signals of the probe beam in different strengths of magnetic field B from 0 to 6.8 Gauss. The solid lines (black) are the experimental data, the dotted lines (red) are the results of the theoretical fittings. With the increase of magnetic field, the single-photon Rydberg excitation spectra expand gradually, and the location of the peak is frequency shifted

expand gradually. The transitions between the degenerate Zeeman sub-levels are resonant, and the optical pump effect is perfect. But when an external magnetic field is applied, the degeneracy is broken. The Zeeman splitting is different for $|g\rangle$, $|e\rangle$ and $|r\rangle$ in the same magnetic field, therefore, the transitions between Zeeman sub-levels are no longer resonant. According to Eq. (11), the detuning increases with the magnetic field, and the corresponding optical pump effect [40] is also gradually weakened. Therefore, the transmission signal decreases with the magnetic field, and the signal is frequency shifted along the direction of red detuning of the coupling beam, as shown in Fig. 7. The solid lines are the expected variation from the theoretical prediction. In previous work [41], due to the influence of power broadening and other broadening mechanisms, the narrowest spectral linewidth is about 24 MHz in the absence of the magnetic field. Hence, in an external magnetic field, we do not observe clearly four Zeeman sub-peaks in the spectra because the broadened linewidth of the signal limits the spectral resolution. Even so, it has potential application in sensing magnetic field.

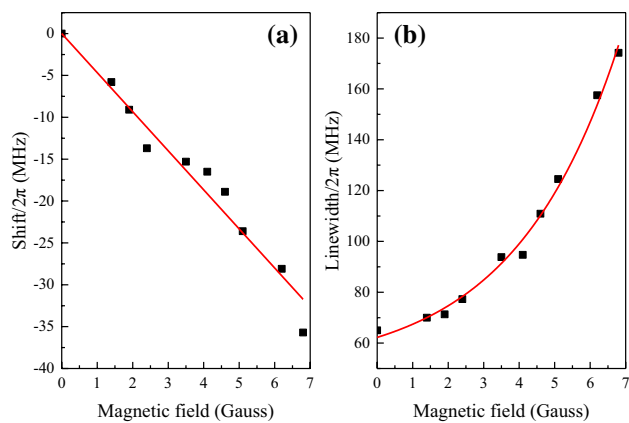


Fig. 7 The frequency shift and linewidth of the transmission signal as a function of magnetic field. The solid lines are the expected variation from the theoretical prediction

4 Conclusion

We have demonstrated the Doppler-free velocity-selective spectra of Cs Rydberg atoms with a single-photon excitation scheme in a room-temperature vapor cell. The weak atomic transition from the ground state ($6S_{1/2}$) to Rydberg state ($71P_{3/2}$) is detected by the reduced absorption of probe beam resonant on the D2 line of Cs in a V-type three-level system. With moderate intensity of the coupling light, the Autler–Townes doublet is clearly observed. The dependence of the separation and linewidth on the coupling intensity is experimentally investigated with dressed state model. The results show good agreement with the theoretical prediction. The implementation of this experiment opens the possibility to realize the Rydberg-dressed ground-state atoms by off-resonantly coupling the ground state to Rydberg state in future work. The Rydberg-dressed atoms combine the long coherence time of ground-state atoms with tunable, strong and long-range interactions between Rydberg atoms. It allows for increasing experimental tunability of these interactions and extending the investigation time of Rydberg atoms while the coherence between atomic states is kept, which is vital for quantum simulation and metrology [6–8]. Now, we have observed the dressed splitting of Rydberg spectra in a cold cesium ensemble [42]. Next, we will perform the Rydberg-dressed experiment by off-resonantly coupling a ground state (Cs $6S_{1/2}$) to a highly-excited Rydberg state ($nP_{3/2}$, $n = 70$ – 100) via a single-photon transition. Combining the high-power narrow-linewidth 319 nm UV laser with Cs cold atomic ensemble in a cross optical dipole trap, we will study some issues about Rydberg-dressed atoms, including the measurement and characterization of the many-body interaction and the ground-Rydberg quantum optical coherence.

Furthermore, we have observed and analyzed the spectra of non-degenerate Zeeman sub-levels when an external magnetic field is applied. Investigation of the single-photon Rydberg spectroscopy in room-temperature atomic vapor has promising application in precision spectroscopy, especially for studying the atomic quantum states and the interaction between Rydberg atoms.

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