Measurement of the Stark shift of the $6s^2S_{1/2} \rightarrow 7p^2P_J$ transitions in atomic cesium

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We report measurements of the Stark shift of the cesium $6s^2S_{3/2} \rightarrow 7p^2P_{3/2}$ and the $6s^2S_{1/2} \rightarrow 7p^2P_{1/2}$ transitions at $\lambda = 456$ nm and 459 nm, respectively, in an atomic beam. From these, we determine the static scalar polarizability for both $7P$ states, and the tensor polarizability for the $7P_{3/2}$ state. The fractional uncertainty of the scalar polarizabilities is $\sim 0.18\%$, while that of the tensor term is $0.66\%$. These measurements provide sensitive tests of theoretical models of the Cs atom, which has played a central role in parity nonconservation measurements.

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I. INTRODUCTION

The high precision attainable in measurements of the Stark shift of atomic transition frequencies makes them sensitive tests of theoretically determined radial matrix elements. Atomic cesium, which has played a central role in parity nonconservation measurements over the past forty years [1–6], is of particular interest in this regard, where accurate determinations of electric dipole matrix elements, experimental [7–11] and theoretical [12–16], are critical for precise determination of the weak charge. The high precision attainable in measurements of the tensor polarizability for the $7P$ states of atomic cesium, whose magnitudes de-
The general principle of the measurement is similar to that of several other recent works [25–28]. We use the output of a single, narrow-band tunable laser source, which we split into two separate beams, labeled the reference and Stark beams in Fig. 3. Using an electro-optic modulator (EOM) and an acousto-optic modulator (AOM) to offset the frequencies of these two beams, we concurrently bring the reference beam into resonance with the cesium transition in a field-free vapor cell (the reference cell), and the Stark beam into resonance with the transition in cesium atoms to which a uniform electric field has been applied. The difference between the frequency offsets of these two beams, which depends only on the rf frequencies driving the modulators, equals the Stark-shift of the resonance. This eliminates the requirement for calibration of the laser frequency scan, which can be problematic at the precision required in these measurements. The Doppler broadening of the resonances is largely suppressed in our measurements, allowing us to resolve the hyperfine structure of the transitions, and also allowing us to use relatively low dc electric field strengths in our measurements.

The laser for these measurements, which we operate at wavelengths of 455.7 nm (for the $6s^2S_{1/2} \rightarrow 7p^2P_{3/2}$ transition) or 459.4 nm (for the $6s^2S_{1/2} \rightarrow 7p^2P_{1/2}$ transition), is a home-made external cavity diode laser (ECDL) using an AR coated laser diode, which generates approximately 10 mW of optical power. We diffract the output beam in an AOM, and use the first-order diffracted beam, whose frequency is $f_l + f_{AO}$ (where $f_l$ is the frequency of the laser output and $f_{AO} = 110.0$ MHz is the AOM drive frequency), for the experiment (i.e. this is the Stark beam). The 110 MHz drive signal is produced by a synthesized signal generator and amplified by an rf amplifier. We direct the undiffracted beam, which we use as our reference beam, into a field-free cesium vapor cell, and frequency-lock the laser to one hyperfine component.
of the Doppler-free saturated absorption spectrum (the 
6s2S1/2, F = 4 → 7p2P1/2, F = 4 line at λ = 459 nm or
the 6s2S1/2, F = 4 → 7p2P3/2, F = 5 line at λ = 456
nm) of this spectrum. To obtain an error signal for lock-
ing to the peak of the hyperfine line, we dither the laser
frequency at 30 kHz. In either case, the laser fre-
cuency \( f_1 \) is resonant with and stabilized to the unshifted
atomic resonance, \( f_p \). The linewidth of the laser spec-
trum is \(< 1 \text{ MHz} \). Because the absorption strengths of
these transitions are relatively weak, we have to heat the
cesium vapor cell to a temperature in the range 80-110°C
to obtain sufficient Cs density within the cell.

We impose optical sidebands on the spectrum of the
Stark beam by modulating its phase in a traveling wave
EOM, driven by a separate signal generator and ampli-
fier at a frequency \( f_{EO} \), where we can adjust \( f_{EO} \) to
any frequency in the range from 110 to 1000 MHz. We
use the lower frequency sideband, whose frequency is
\( f_1 + f_{EO} - f_{EO} \), to excite the Stark-shifted absorption
resonance in the atomic beam. We carry out the mea-
surements in one of two different modes: we fix the fre-
cuency \( f_{EO} \) and vary the dc field amplitude \( E_0 \) to ‘tune’
the Stark-shifted absorption line into resonance with the
lower frequency sideband; or we fix the amplitude \( E_0 \)
and vary the frequency \( f_{EO} \) to match the Stark-shifted
resonance.

The cesium atom beam is formed inside an aluminum
vacuum chamber pumped with a turbomolecular pump
to a pressure of \( 5 \times 10^{-6} \) torr. We use an effusive cesium
oven with a nozzle consisting of an array of stainless steel
hypodermic needle tubes to form the atom beam. More
details are available in our earlier publications \([24,30]\).
This oven and nozzle generates a beam of dimension 12
mm × 8 mm near the nozzle. We insert an atomic beam
aperture (labeled skimmer in Fig. 3) before the inter-
action region to further reduce the width of the atomic
beam to \( \sim \frac{1}{3} \) the spacing of the field plates. This reduces
the accumulation of cesium on the electric field plates.
The spectral width of the \( 6s \rightarrow 7p \) absorption lines
in our beam geometry is \( \sim 6 \text{ MHz FWHM} \), largely due
to Doppler broadening in the slightly diverging atomic
beam. The natural linewidths of these transitions, cor-
responding to the 133 ns lifetime of \( 7P_{3/2} \) and the 155 ns
lifetime of the \( 7P_{1/2} \) state \([31,32]\), are 1.2 MHz and 1.0
MHz, respectively.

The uniformity of the static electric field, and the pre-
cision with which this field can be determined, depends
critically on the parallel conducting field plates used to
generate this field. We construct these field plates from a
pair of 76.2 × 25.4 mm (3 × 1 in) microscope glass slides,
coated on the inside surfaces with a thin conducting layer
of indium tin oxide (ITO). These field plates are spaced
by 4.928 (4) mm (0.19400 (15) in), and are mounted in-
side an aluminum framework with external ceramic posts
using a vacuum compatible epoxy. (The number enclosed
within parentheses following these parameters indicates
our estimate of the uncertainty.) We evaluated the non-
uniformity of the electric field within the interaction re-
gion due to fringing effects using a commercial software
package, and found that this variation is less than a part
in \( 10^5 \).

During assembly, we spaced the field plates with a set
carefully selected ceramic spacers to assure a high de-
gree of parallelism, then removed the spacers after the
epoxy had dried. (We observed drifts in some of our
early Stark shift measurements, which we attributed to
an accumulation of cesium on the internal spacers used
for those measurements. These drifts were absent af-
fer we removed the internal spacers.) We estimate the
0.00015 inch uncertainty in the spacing of the glass slides
based on the relative ease with which we can slip cal-
ibrated ceramic beads, whose lengths we measured to
\( \pm 0.00005 \) inch, at various locations near the central re-
gion of the field plates, similar to the technique described
in Refs. \([21,27]\). We also measured the parallelism of the
plates by reflecting a HeNe laser beam from the two sur-
faces, and observing the spacing of the fringes formed by
the interference of the two reflected beams. We estimate
that the angle between the two plates was less than 0.15
mrad. This high degree of parallelism between the field
plates is consistent with our estimate of the variation of
the plate spacing over the width of the plates.

We use a pair of stable high-voltage sources to bias the
field plates, one plate positively biased, the other
negative. Between sets of data, we switch the polarity of
the field plates. We measure the voltage applied to
each field plate using an Ohmcraft 1000:1 high voltage
resistive divider, which we have carefully checked and
calibrated for nonlinearity and stability. The fractional
uncertainty in the voltage measurement of each field plate
is \( \sim 2 \times 10^{-5} \).

Consistent with the treatment by Schmieder \([24]\), we
define the \( z \)-axis of the atomic system as the direction
of the applied field \( E_0 \). While the Stark beam for these
measurements propagates in a direction \( \mathbf{k} \) nearly parallel
to this \( z \)-axis, and its polarization state is linear, the ex-
periment is relatively insensitive to either of these con-
ditions, since the ground state components are degenerate,
and the various peaks in the Stark spectrum correspond
to different hyperfine components of the excited state
alone. Changes in polarization or imperfect alignment of
\( \mathbf{k} \) with the \( z \)-axis only change the relative height of the
peaks in the Stark spectrum, but not their frequency. By
contrast, it is important that the laser beam propagates
in a direction perpendicular to the atomic velocity to as-
sure narrow absorption linewidths and to minimize the
Doppler shift of the lines. Using an alignment laser, we
mount the parallel field plates inside the vacuum system
centered on and parallel to the atomic beam. In addi-
tion, we observe the reflection of the Stark beam from
the field plates and adjust this beam to normal incidence
on the field plates. After these alignment steps, only a
minor adjustment of the Stark beam direction \( \mathbf{k} \) is nec-
cessary to minimize the Doppler shift of the resonance in
the atomic beam, which we determine by zeroing the ap-
plied field and comparing the absorption resonance in the
atomic beam to that of the reference cell.

In order to detect the absorption resonances in the atomic beam, we use the detection system that we developed earlier \cite{29,31} for sensitive measurement of highly-forbidden optical transitions. We based this system on a technique reported earlier in Ref. \cite{2}. The population of the cesium atoms as they effuse from the oven is equally distributed among each of the F=3 and F=4 hyperfine components of the ground state. Before the atoms interact with the blue laser, we transfer all of the atoms into the F=4 level by optically pumping the population with the output of an 852 nm ECDL tuned to the 6s^2S_{1/2}, F = 3 \rightarrow 6p^2P_{3/2}, F' = 4 hyperfine transition of the D_2 resonance line. After interacting with the blue laser, the population in the F = 3 ground state is a measure of the excitation rate by the Stark beam to the 7P state, since these atoms decay spontaneously back to the ground state, where some end up in the initially empty F=3 hyperfine state. We detect this population using the output of a second ECDL tuned to the D_2 line at 852 nm (in this case resonant with the 6s^2S_{1/2}, F = 3 \rightarrow 6p^2P_{3/2}, F' = 2 cycling transition) and a large area photodiode to measure the scattered optical power in this region.

We use a lock-in amplifier for phase-sensitive detection of the photodiode current in order to improve the sensitivity of the measurement. We dither the frequency of the EO sideband at 145 Hz (with a 1 MHz amplitude), which modulates the rate of absorption by the atoms. The derivative signal produced by the lock-in amplifier, illustrated in Fig. 2(b), is a dispersion shaped resonance of width \sim 6 MHz. The zero-crossing is well suited for determination of line center.

During the course of our measurements, we found that the amplitude of the optical sideband of the Stark beam, as monitored with a scanning Fabry-Perot interferometer, varied across the 100-1000 MHz spectrum. To ensure that the optical power in the sideband used for the experiment is constant, we selected frequencies f_{EO} at which the sideband power is relatively uniform.

We must keep the laser intensity below the saturation intensity I_{sat} in order to minimize power broadening and light shifts. The laser intensity for these measurements is 4 mW/cm^2, of which only \sim \frac{4}{7} is in the lower sideband that interacts with the atoms. Using the reduced matrix dipole matrix elements for these transitions \cite{15}, we estimate that the saturation intensity I_{sat} of the 6s^2S_{1/2} \rightarrow 7p^2P_{3/2} transition is about 15 mW/cm^2, while that of the 6s^2S_{1/2} \rightarrow 7p^2P_{1/2} line is about 50 mW/cm^2. Therefore, the sideband intensity is well below I_{sat} in both cases.

### A. Scalar polarizability of 7p^2P_{1/2}

For our determination of the polarizability \alpha_0 of the 7p^2P_{1/2} state, we first set the EO modulation frequency f_{EO} to one of seven pre-determined values in the range between 110 and 1000 MHz. (The minimum of this range is the AO drive frequency, and corresponds to a zero Stark shift, while the maximum frequency is the maximum frequency of our signal generator.) At each value of f_{EO}, we adjust the voltage applied to the plates to shift the transition into resonance with the lower sideband of the Stark beam. In Fig. 4 we show one set of these data, plotted as f_{EO} vs. E_0^2. The solid line indicates the result of a linear least squares fit. The residual error of each data point is shown in (b).

The intercept of this fitted line is 109.8 (2) MHz, consistent with the 110.0 MHz frequency offset imposed by the AOM. The slope of this line is 3.6417 (12) MHz/(kV/cm)^2, and is equal to half the difference between the polarizabilities of the 7P_{1/2} state and the 6S_{1/2} state, \frac{1}{2}(\alpha_0(7P_{1/2}) - \alpha_0(6S_{1/2})). The uncertainty of 0.0012 MHz/(kV/cm)^2 is statistical and is determined from the scatter of the data points from the linear fit to the data. We show the difference between the data points and the linear fit, in Fig. 4(b). The rms residual for this set of data is 0.22 MHz.

We measured the Stark shift of the 7P_{1/2} state four times, reversing the direction of electric field between sets of data and observed no systematics due to electric field orientation. We show the slope \frac{1}{2}(\alpha_0(7P_{1/2}) - \alpha_0(6S_{1/2})) resulting from each of these measurements in Fig. 5. The error bars shown in the figure indicate the statistical uncertainty for each data point.

The reduced chi^2 for these measurements is 4.14, indi-
cating that the measurement uncertainty is larger than the statistical uncertainty. The weighted average of the four measurements yields $\frac{1}{4}\{\alpha_0(7P_{1/2}) - \alpha_0(6S_{1/2})\} = 3.6405 \pm 0.002 (6) \text{ MHz/(kV/cm)^2}$, as indicated by the diamond-shaped data point and horizontal line in Fig. 5. We have not scaled the statistical error despite the large $\chi^2$ factor. As we will discuss in Section IIII, the overall measurement uncertainty is dominated by the uncertainty in the field plate spacing, and scaling the statistical error has little impact on our final result.

Using the ground state polarizability $\alpha_0(6S_{1/2}) = 0.09978 (7) \text{ MHz/(kV/cm)^2}$ from Ref. [12], we find $\alpha_0 = 7.3808 (12) \text{ MHz/(kV/cm)^2}$, where the number in parenthesis denotes the statistical error only. In atomic units, this converts to 29,662 (5) $\alpha_o^3$.

### B. Scalar and tensor polarizability of $7p^2 P_{3/2}$

We base our determinations of $\alpha_0$ and $\alpha_2$ for the $7p^2 P_{3/2}$ line on two lines in the Stark-shifted spectrum, namely the ($F' = 2$, $m_F' = \pm 1$) line and the (4,2) line. We chose these particular $m_F$ peaks because they are well resolved from other peaks in the Stark spectrum, as shown in Fig. 2, and because their frequency difference due to the Stark shift is large, allowing for a more precise evaluation of $\alpha_2$. From Eq. (9), the polarizability of the $F' = 5$, $m_F = \pm 5$ components is $\alpha_0 + \alpha_2$. Our process for determination of the sum $\alpha_0 + \alpha_2$ for the $7p^2 P_{3/2}$ state then is similar to that of $\alpha_0$ for the $7p^2 P_{1/2}$, described in the last section. With the reference laser frequency $f_l$ tuned and locked to the $6s^2 S_{1/2}$, $F = 4 \rightarrow 7p^2 P_{3/2}$, $F' = 5$ resonance in the reference cell, we adjust the frequency $f_{EO}$ of the signal applied to the EOM to one of seven values in the range from 700 to 1000 MHz. (Below 700 MHz, the various peaks within the Stark spectrum partially overlap, introducing errors in the measurements of the line center.) Then we vary the voltage applied to the field plates to bring the (5,5) peak into resonance. We also take one measurement at zero electric field, varying $f_{EO}$ to find the line center.

We show an example of one data set in Fig. 6. The result of a linear least squares fit, represented by the straight line in this figure, yields an intercept of 110.7 MHz and a slope of $\frac{1}{4}\{\alpha_0(7P_{3/2}) + \alpha_2(7P_{3/2}) - \alpha_0(6S_{1/2})\} = 4.0386 (18) \text{ MHz/(kV/cm)^2}$. We show the deviation of each of the data points from the fitted line in Fig. 6(b). The rms residual is 0.3 MHz. We repeat this measurement with the electric field orientation reversed, and obtain a result which is in good agreement with our first measurement. Using the two measurements, we determine a weighted average slope of $\frac{1}{4}\{\alpha_0(7P_{3/2}) + \alpha_2(7P_{3/2}) - \alpha_0(6S_{1/2})\} = 4.0389 (13) \text{ MHz/(kV/cm)^2}$. Using $\alpha_0(6S_{1/2})$ from Ref. [12], we obtain $\alpha_0 + \alpha_2 = 8.1776 (26) \text{ MHz/(kV/cm)^2}$ for the $7p^2 P_{3/2}$ state. This uncertainty accounts for statistical effects only.

As we discussed earlier, the frequency difference between the hyperfine components of the Stark spectrum is quantified through the tensor polarizability $\alpha_2$, for which we base our determination on a measurement of the frequency difference between the (5,5) peak and the (4,2)
peak. At zero field, the frequency difference between these peaks is the hyperfine splitting $f_{4-5} = 83.0255 \text{ MHz}$ [22]. We measure this frequency difference by fixing the voltage applied to the field plates (and hence electric field) and adjusting the frequency $f_{EO}$ to bring the Stark laser sideband into resonance with each sublevel. For each of three voltage levels, we repeated the measurement twice. In order to determine the value of $\alpha_2$, we fit the diagonalized matrix $Q$ to the data points. In this case, we fixed the value of $\alpha_0 + \alpha_2$ to the value $8.1776 (26) \text{ MHz/(kV/cm)}^2$, as discussed above, and used Eqs. (40a) and (41) from Ref. [24] to generate curves for varying values of $\alpha_2$. The least rms deviation between the calculated and measured frequency difference between the $(F', |m'_F|) = (5, 5)$ and $(4, 2)$ Stark-shifted frequencies yields $\alpha_2 = -1.0981 (65) \text{ MHz/(kV/cm)}^2 = -4.413 (26) a_0^3$. In Fig. 7 we show the best fit curves for the Stark-shifted hyperfine peaks versus $E_0^2$, with the scalar component of the Stark shift suppressed. The circles denote our experimental data points. We also show measurements of the $F' = 5$, $m'_F = 4$ sublevels in this figure. We did not use these values in the determination of $\alpha_2$, since the smaller frequency difference between this peak and the $(5, 5)$ gave these values a larger uncertainty and larger error. Combining our results for $\alpha_0 + \alpha_2$ and $\alpha_2$ yields a value of $\alpha_0 = 9.2757 (70) \text{ MHz/(kV/cm)}^2$, or $37.277 (28) a_0^3$.

### III. MEASUREMENT ERRORS

The uncertainties in the polarizabilities that we presented in the previous section include only statistical effects derived from the scatter in the data points from the fitted lines. In addition, there are other experimental factors, as summarized in Table I that we must consider. In this section, we discuss these contributions and provide estimates of their magnitudes.

The largest uncertainty in our measurements is the systematic effect due to determination of the static electric field strength. These uncertainties derive from the uncertainty in the field plate spacing (including the uncertainty in the measurement of this spacing $d$, as well as any non-uniformity in $d$), the uncertainty in the measurement of the voltage applied to the plates, and edge effects that reach in to the center of the field plates. We have discussed the first of these in Section III where we estimate an uncertainty of the plate spacing of $0.08\%$. Since the Stark shift depends on $E_0^2$, the corresponding uncertainty in the polarizabilities is twice as large, or $0.16\%$. We also described the voltage dividers that we used to measure the voltage applied to the field plates in Section III. This fractional uncertainty of $5 \times 10^{-5}$ results in an uncertainty in the polarizabilities of $1 \times 10^{-4}$. We also list in Table I the measurement error of the voltameter as specified by the manufacturer.

We estimate that the precision with which we can measure the linecenter of each of the Stark shifted lineshapes is $\pm 0.2 \text{ MHz}$. This is primarily limited by signal asymmetry due to residual amplitude modulation of the Stark beam at 145 Hz. For instance, if the asymmetry of a dispersion-shaped resonance is $15\%$ of the maximum error signal, as was typical of our measurements, the zero crossing is shifted by $\sim 0.2 \text{ MHz}$, assuming a $6 \text{ MHz}$ linewidth of the absorption peak. Another limiting factor is dc offsets in the error signal, due to electronics and the overlap from adjacent peaks in the spectrum. We modeled the pulling of the line center due to adjacent peaks and found its effect on the polarizabilities to be less than $2 \times 10^{-4}$. We estimate that these limiting factors lead to a fractional uncertainty in the polarizabilities of $2 \times 10^{-4}$.

We also considered frequency shifts due to changes in the propagation direction of the laser beam. Such a change could introduce a Doppler shift in line center of the resonance. Heating effects in the EOM could deflect the beam, for example. We have projected the Stark beam onto a screen $10 \text{ m}$ beyond the EOM, and were unable to observe any such deflection. We place an up-

![Graph showing frequency shift vs. electric field strength](image)

**FIG. 7.** (Color on-line) Frequency $f_{EO}$ for the $(5,5)$ (blue), $(5,4)$ (gold), and $(4,2)$ (red) peaks vs $E_0^2$ for determination of the tensor polarizability $\alpha_2$. The solid lines are the calculated Stark shifts, less the $\alpha_0$ terms, with $\alpha_2$ adjusted to minimize the deviation with the data. The notation beside each line indicates $(F', |m'_F|)$ of the $7p^2P_{3/2}$ state.

| Source                  | Uncertainty % of $\alpha$ |
|-------------------------|----------------------------|
| Field plate spacing     | $4 \mu m$                  | 0.16 |
| Voltage divider ratio   | 0.005%                     | 0.01 |
| Voltage measurements    | 0.005%                     | 0.01 |
| Error signal line center| 0.2 MHz                    | 0.02 |
| AOM drive frequency     | 10 kHz                     | 0.01 |
| EOM drive frequency     | 10 kHz                     | 0.01 |
| Beam alignment into chamber | 0.05 mrad              | 0.01 |
| **Total systematic uncertainty** | 0.16 |

**TABLE I.** Sources of error, estimates of their uncertainties, and the resulting percentage uncertainty in $\alpha$ resulting from each source.
per limit of 0.05 mrad on any such shift. Estimating the Doppler shift to be about 0.7 MHz/mrad, this shift corresponds to an uncertainty of less than 0.05 MHz. This limit is consistent with our observed rms residuals of the measured peak positions in Figs. 4 and 6 of 0.2 MHz.

In order to determine $\alpha_2(7P_{3/2})$ from our measurement of the frequency difference between the $(5,5)$ and $(4,2)$ peaks of the stark spectrum, we used the hyperfine constants $a = 16.605 (6)$ MHz and $b = -0.15 (3)$ MHz from Ref. [22]. We consider here the effect of the uncertainty of these hyperfine constants on the uncertainties of the polarizabilities of the $7P_{3/2}$ state. By varying the values of the constants by one standard deviation and running the fitting function again, we can estimate their effect on our values of $\alpha_0(7P_{3/2})$ and $\alpha_2(7P_{3/2})$. This effect is estimated to be 0.21% for $\alpha_2(7P_{3/2})$ and $< 0.02\%$ for $\alpha_0(7P_{3/2})$.

For our measurements of the scalar polarizabilities $\alpha_0(7P_{1/2})$ and $\alpha_0(7P_{3/2})$, only the 0.08% variability in the field plate spacing is significant. These effects contribute a 0.16% uncertainty. For the tensor polarizability $\alpha_2(7P_{3/2})$, there is an additional 0.21% error due to the uncertainty in the hyperfine constants. We add these uncertainties in quadrature with the statistical uncertainty stated to obtain the total uncertainty. For the scalar polarizabilities $\alpha_0(7P_{1/2})$ and $\alpha_0(7P_{3/2})$, this results in an uncertainty in the final result of 0.17% and 0.18% respectively. For the tensor polarizability $\alpha_2(7P_{3/2})$, the statistical uncertainty is the primary contributor to the 0.66% uncertainty in our result. In the next section, we present our final results for each, and compare with prior experimental and theoretical determinations of these quantities.

### IV. DISCUSSION

Our results are in good agreement with, and of higher precision than, previous theoretical and experimental results. We present a summary of past theoretical and experimental results in Table II. Our measurement results for the scalar polarizabilities are

$$\alpha_0(7P_{1/2}) = 29,660 (50) a_0^3$$

and

$$\alpha_0(7P_{3/2}) = 37,280 (70) a_0^3.$$  (5)

For the tensor polarizability of the $7P_{3/2}$ state, we find

$$\alpha_2(7P_{3/2}) = -4.413 (29) a_0^3.$$  (6)

Each of these results agrees with past measurements reported in Refs. [19] and [21]. The measurement result for $\alpha_2$ for the $7P_{3/2}$ in Ref. [20] differs by $\sim 10\%$ from the others, including the present results. The precision of the present results is much higher than that of the previous reports, due to our use of narrow-band laser sources, Doppler-free resonances, and r.f. modulation techniques. The theoretical calculations of van Wijngaarden and Li [23] and of Iskrenova-Tchoukova et al. [17] are in good agreement with our results for all three polarizabilities as well. The former does not report uncertainties. Our results differ from those of Ref. [17] by typically less than 1%, while their stated uncertainties are about 2%.

### V. CONCLUSION

We have described our experimental determinations of the Stark shift of the $6s^2S_{3/2} \rightarrow 7p^2P_J$ transitions for $J = \frac{1}{2}$ and $\frac{3}{2}$ in atomic cesium. Through use of a narrow-band, frequency-stabilized diode laser and Doppler-free techniques, the precision of our measurements is higher than that of previous measurements. While our polarizability measurements do not yield radial matrix elements directly, the strong agreement between our polarizabilities and those of Ref. [17] do infer that the radial matrix elements calculated in that work are very accurate.

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