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Quantum defects and the $1/n$ dependence of Rydberg energies: Second-order polarization effects

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The principal result of this paper is a general expression for the second-order dipole polarization energy of a Rydberg electron resulting from the term $-\alpha_1/r^4$ in the asymptotic potential, where $\alpha_1$ is the core polarizability. It is shown that the second-order term contributes even as well as odd powers of $1/n$ in a $1/n$ expansion of the energies for Rydberg states. The results are used to extend the interpretation of the terms in a quantum-defect expansion. It is shown that the Ritz expansion for the quantum defect, which contains only odd inverse powers of the effective quantum number $n^*$, provides a powerful method for deducing the even-order terms in the second-order energy. Least-squares fits to high-precision variational calculations for the Rydberg states of helium, using $1/n$ and quantum-defect expansions, are presented. The results reveal well-defined "Ritz defects," which represent the degree to which the data cannot be represented by a Ritz expansion for the quantum defect. The implications for extrapolations of quantum defects are discussed. Finally, it is shown that the second-order polarization energy plays a significant role in understanding the quantum defects for the alkali metals.

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I. INTRODUCTION

The quantum-defect method [1] is now well established as the method of choice for the analysis of experimental data on the term energies of Rydberg sequences of states. A vast literature has grown up around its many varied applications [2]. However, advances over the past ten years in the precision of both experimental measurements [3–12] and theoretical calculations [13–17] raise new questions concerning the ultimate limits of the quantum-defect method, and in particular the Ritz expansion for the quantum defect, as a suitable functional form for the representation of data.

This paper has three main purposes. The first is to extend the physical identification that can be made for the terms in the Ritz expansion of the quantum defect, especially those arising from second-order polarization effects. Exact analytic results are obtained in Sec. III for the second-order dipole polarization term and compared with the predictions of quantum-defect theory. It is shown in Sec. IV that the second-order term is important in the analysis of data if the core polarizability is large. The second purpose is to apply these results in Sec. IV to recent high-precision calculations [13–15] for the Rydberg $P$, $D$, and $F$ states of helium, and to identify what extensions of the usual Ritz expansion might be necessary. The third is to identify the conditions under which the Ritz expansion is no longer capable of representing data in the broader context of quasihydrogenic spectra. Not considered are the more general problems of multiple Rydberg series and their perturbations [2].

II. QUANTUM-DEFECT THEORY AND $1/n$ EXPANSIONS

As first pointed out by Ritz [18], the term energies of a single Rydberg series of states for a quasihydrogenic atom are well represented by the formula

$$T_n = -2R_M/(2n^*)^2,$$

(1)

where $R_M$ is the reduced mass Rydberg constant, $2R_M$ is the reduced mass atomic unit of energy, and $n^*$ is an effective quantum number given by

$$n^* = n - \delta(n^*)$$

$$= n - \delta_0 - \frac{\delta_2}{(n - \delta)^2} - \frac{\delta_4}{(n - \delta)^4} - \cdots.$$ 

(2)

$\delta$ is called the "quantum defect." A theoretical justification of this result was given by Sommerfeld [19] based on the old quantum theory, and by Hartree [20] using wave mechanics. The key point proved by Hartree is that if the motion of an electron is describable by a Hamiltonian of the form

$$H = H_c + \lambda V,$$ 

(3)

where $H_c$ is the Hamiltonian for a purely Coulombic potential and $\lambda V$ is a short-range spherically symmetric correction potential, then the eigenvalues of $H$ are given exactly by Eq. (1) with only the even powers of $1/(n - \delta)$ appearing in Eq. (2). The proof is nonperturbative, and so applies for arbitrary values of the strength parameter $\lambda$. As will be shown in Sec. III, this fact can be used to advantage in calculating the general $n$ dependence of the terms in a perturbation expansion containing powers of $\lambda$. However, notice that $\delta$ appears in the denominators of Eq. (2), rather than $\delta_0$. Replacing $\delta$ by $\delta_0$ as is often done in fitting experimental data, may lead to an apparently adequate fit, but it spoils the theoretical significance of the coefficients. It will also ultimately limit the accuracy of the quantum-defect method. This point will be further discussed below and in Secs. III and IV.
The physical significance of the coefficients in the Ritz formula (2) is made evident by expanding Eq. (1) as a power series in $1/n$. Keeping terms up to quadratic in the coefficients, the result is

$$T_n = -2R'_M \left[ \frac{1}{2n^2} + \frac{\delta_0}{n^3} + \frac{\delta_2}{n^5} + \frac{\delta_4}{n^7} + \frac{\delta_6}{n^9} + \cdots \right]$$

$$+ \frac{3\delta_0^2}{2n^4} + \frac{5\delta_0\delta_2}{n^6} + \frac{7}{2n^8}(\delta_2^2 + 2\delta_0\delta_4)$$

$$+ \frac{9}{n^{10}}(\delta_0^2\delta_4 + \delta_2\delta_4) + \cdots . \ \ (4)$$

The influence of the $\delta$ expansion in the denominators of Eq. (2) first appears in the terms of order $1/n^8$. Replacing $\delta$ by $\delta_0$ reduces the coefficient of $\delta_2^2$ from $\frac{7}{2}$ to $\frac{3}{2}$ and the coefficient of $\delta_2\delta_4$ from 9 to 3. It will be shown in Sec. III that the value $\frac{3}{2}$ is in fact correct.

The physical significance of the terms in Eq. (4) now follows from a comparison with the asymptotic potential experienced by the Rydberg electron. In order for a single local potential as in Eq. (3) to be adequate, it is necessary that nonlocal exchange effects be negligible. If in addition the angular momentum $L$ is large enough so that there is little core penetration by the Rydberg electron, then the asymptotic potential has the well-known form [21]

$$V(r) = -\frac{1}{2} \left[ \frac{c_4}{r^4} + \frac{c_6}{r^6} + \frac{c_7}{r^7} + \frac{c_8}{r^8} + \cdots \right] , \ \ (5)$$

where

$$c_4 = \alpha_1 , \ \ (6)$$

$$c_6 = \alpha_2 - 6\beta_1 , \ \ (7)$$

$$c_7 = -8' - 16\gamma / 6 , \ \ (8)$$

$$c_8 = \alpha_3 - 15\beta_2 + e - \alpha_1\beta_1 + 72\gamma [1+L(L+1)/10] . \ \ (9)$$

$r$ is the radial coordinate of the Rydberg electron, $\alpha_L$ is the 2$^\lambda$-pole polarizability, $\beta_L$ is the corresponding nonadiabatic correction, and the other terms in $c_7$ and $c_8$ are higher-order corrections defined by Drachman [21]. The quantities $c_4, c_6, \ldots$ play the role of the strength parameter $\lambda$ in Eq. (3). The first-order correction to the energy is then

$$E^{(1)} = \langle V \rangle , \ \ (10)$$

where the expectation value is with respect to the Rydberg electron. With the notation

$$f_p(L) = (L+p)!/(L-p)! , \ \ (11)$$

$$G_p(L) = (2^p Z^2)^{p} / (2L+p+2)! \middle/ (2L+p+1)!$$

the general expressions for the expectation values of the powers of $1/r$ in Eq. (5) are [22]

$$\langle r^{-4} \rangle = G_4(L) \frac{1}{n^2} \left[ 3n^2 - f_4(L) \right] , \ \ (13)$$

$$\langle r^{-6} \rangle = G_6(L) \frac{1}{n^2} \left[ \frac{3}{2} n^4 - 10n^2 [f_4(L) - \frac{5}{8}] + f_6(L) \right] , \ \ (14)$$

$$\langle r^{-7} \rangle = \frac{3G_7(L)}{n^2} \left[ \frac{11}{2} n^5 - \frac{11}{2} n^3 [f_4(L) - \frac{5}{8}] + n [f_2(L) - \frac{2}{3} f_3(L) + \frac{7}{18}] \right] , \ \ (15)$$

$$\langle r^{-8} \rangle = \frac{10G_8(L)}{n^2} \left[ \frac{21}{2} n^6 - 63n^4 [f_4(L) - \frac{5}{8}] + 21n^2 [f_2(L) - \frac{3}{4} f_5(L) + \frac{15}{16}] + f_3(L) \right] . \ \ (16)$$

Comparing with Eq. (4) for the odd powers of $1/n$ leads immediately to the first-order identifications

$$\delta_0^{(1)} = \frac{1}{4} c_4 G_4(L) + \frac{15}{8} c_6 G_6(L) + \cdots , \ \ (17)$$

$$\delta_2^{(1)} = -\frac{1}{4} c_4 G_4(L) f_4(L) - 15 c_6 G_6(L) [f_4(L) - \frac{5}{8}] - 70 c_8 G_6(L) [f_4(L) - \frac{5}{8}] - 315 c_8 G_8(L) [f_4(L) - \frac{5}{8}] - \cdots , \ \ (18)$$

$$\delta_4^{(1)} = \frac{1}{4} c_4 G_4(L) f_4(L) + 15 c_6 G_6(L) [f_4(L) - \frac{5}{8}] + f_4(L) + \frac{15}{8} , \ \ (19)$$

$$\delta_6^{(1)} = -5 c_8 G_8(L) f_4(L) - \cdots , \ \ (20)$$

correct to first order in the $c_j$'s. The above are of course asymptotic expansions which must be terminated after a finite number of terms, depending on the value of $L$.

Except for the nonadiabatic and higher-order corrections contained in $c_6, c_7,$ and $c_8$, the above identifications coincide exactly with the standard ones discussed for example by Edlén [1] and Curtis [23]. What has not been adequately discussed before is the origin and significance of the even powers of $1/n$ in Eq. (4). The coefficients are all quadratic in the $\delta$'s. They are uniquely determined once the coefficients of the odd powers have been fixed because Eq. (4) contains twice as many terms as the original Ritz expansion for a given highest power of $1/n$.

In fitting data, it has sometimes been argued that $T_n$ should be written in the form

$$T_n = -2R'_M \left[ 1/2n^2 + a_3/n^3 + a_5/n^4 + \cdots \right] , \ \ (21)$$

where the even terms beyond $1/n^2$ have been dropped. One argument is based on the fact that no inverse power of $r$ (whether even or odd) in the asymptotic potential can produce even powers of $1/n$, and additional exchange and short-range effects not included in Eq. (5) decrease exponentially with $n$. For example, Chang and Poe [24] fitted their results of many-body perturbation theory (MBPT) calculations for He to a function of this form in
order to extrapolate to high \( n \). It has even been argued that the results of nonrelativistic MBPT can be represented exactly in the form of Eq. (21) [25]. However, if Hartree’s proof is to be taken seriously, then the even terms in Eq. (4) have a significance which cannot be ignored. As shown in Sec. III, they coincide exactly with the second-order perturbation corrections generated by the terms in the asymptotic potential. In fact, Eq. (4) provides a trivial way of calculating them. For example, if Eq. (21) is extended to include the even terms \( a_4/n^4 + a_6/n^6 + \cdots \), then Eq. (4) predicts that

\[
a_4 = \frac{3a_2}{2}, \tag{22}
\]

\[
a_6 = 5a_2a_5, \tag{23}
\]

and so on for the values implied by the assumed Ritz expansion. If a fit to experimental or theoretical energy levels yields some other value for \( a_4 \), say \( \bar{a}_4 \), then the difference

\[
\Delta a_4 = a_4 - \bar{a}_4 \tag{24}
\]

is the leading term in what might be called the Ritz defect. It represents the degree to which the data cannot be represented in the form of the Ritz expansion for the quantum defect. This result does not depend on the coefficients being small since Hartree’s proof is nonperturbative.

The even inverse powers also enter when relativistic corrections are included. This is easily seen from the well-known formula for the leading one-electron relativistic energy shift [26]

\[
\Delta T_{rel} = -a_4^{2}Z_{eff}^{2} \left[ \frac{1}{j+\frac{1}{2}} - \frac{3}{4n} \right]. \tag{25}
\]

A relativistic generalization of the quantum-defect method has been developed by Johnson and Cheng [27]. The presence of the \( 1/n^4 \) term was also noted in passing by Curtis [23] in his semiclassical analysis of the quantum-defect method. If experimental energies or theoretical calculations are analyzed in terms of a \( 1/n \) expansion, it is important to realize that the even inverse powers contain nonrelativistic as well as relativistic contributions.

### III. SECOND-ORDER DIPOLE POLARIZATION ENERGY

If the asymptotic expansion (5) is carried to terms of order \( \langle r^{-7} \rangle \) and \( \langle r^{-8} \rangle \), then one should also include the second-order dipole polarization correction (in atomic units)

\[
E^{(2)} = -\frac{1}{2} \langle \psi_1 | r^{-4} | \psi_0 \rangle , \tag{26}
\]

where \( \psi_0 \) is the unperturbed wave function for the \( nL \) Rydberg electron and \( \psi_1 \) satisfies the first-order perturbation equation

\[
(H_0 + Z_{eff}^2/2n^2)\psi_1 = -\frac{\alpha_1}{2r} \psi_0 = -\frac{1}{2} \langle \psi_0 | r^{-4} | \psi_0 \rangle \psi_0 . \tag{27}
\]

Numerical values for \( E^{(2)} \) have been calculated by Drachman [21] and shown to be important for the Rydberg states of helium. He succeeded in obtaining analytic expressions for \( E^{(2)} \) for sequences of states with constant \( n - L \), but the general \( n \) dependence of \( E^{(2)} \) for fixed \( L \) has not previously been obtained.

Introducing the scaled quantities

\[
x = Z_{eff}r/a_0, \quad \xi = E(a_0/e^2)/Z_{eff}^2 ,
\]

where \( Z_{eff} \) is the effective nuclear charge experienced by the Rydberg electron, and defining \( \bar{a}_1 = Z_{eff}^2 \alpha_1/a_0^3 \), the radial part of Eq. (27) becomes

\[
\left[ -\frac{1}{2} \frac{d^2}{dx^2} - \frac{d}{dx} + L(L+1)/2x^2 - \frac{1}{x} + \frac{1}{2n^2} \right] u_1(x) \]

\[
- \frac{\bar{a}_1}{2x^4} u_0(x) = \xi^{(1)} u_0(x) , \tag{28}
\]

where \( u_0(x) \) is the radial part of \( \psi_0 \) and

\[
\xi^{(1)} = -\left( \bar{a}_1/2 \right) \langle \psi_0 | x^{-4} \psi_0 \rangle . \tag{29}
\]

The assumed normalization condition \( \langle \psi_1 | \psi_0 \rangle = 0 \) can always be satisfied by adding an appropriate component of \( u_0 \) to \( u_1 \). The solution to (28) for a particular \( nL \) state \( (L \geq 2) \) is easily found by writing \( u_1(x) \) in the form

\[
u_1(x) = \frac{\bar{a}_1}{2} \sum_{j=L-1}^{n+1} g_j x^{j-1} e^{-x/n}
\]

\[
+ \frac{\bar{a}_1}{2} \sum_{j=L+1}^{n} h_j x^{j-1} \ln x e^{-x/n} . \tag{30}
\]

The hydrogenic \( u_0(x) \) is given by

\[
u_0(x) = \sum_{j=L+1}^{n} A_j x^{j-1} e^{-x/n} , \tag{31}
\]

with

\[
A_{L+1} = \frac{2L+1}{n L+2} \left[ \frac{(n+L-1)!}{(n-L+1)!} \right]^{1/2} \frac{1}{(2L+1)!} \tag{32}
\]

and

\[
A_{j+1} = \left[ \frac{2(j-n)}{n \left( j+1 \right) - L \left( L+1 \right)} \right] A_j ,
\]

\[
\quad j = L+1, \ldots, n-1 . \tag{33}
\]

Substituting into Eq. (28) yields the recursion relations

\[
h_{j+1} = \left[ \frac{2(j-n)}{n \left( j+1 \right) - L \left( L+1 \right)} \right] h_j \tag{34}
\]

and
\[ g_{j+1} = \frac{2}{[j(j+1)-L(L+1)]} \times \left( \frac{(j-n)}{n} - g_j - A_{j+1} + \left( x^{-4} \right) A_{j-1} \right) \]
\[ - \frac{(2j+1)}{2} h_{j+1} + \frac{1}{n} h_j \]  
(35)

for \( j = L+1, \ldots, n+1 \), with the starting values

\[ g_{L-1} = A_{L+1}/(2L-1), \quad h_{L-1} = 0; \]
\[ g_L = \frac{1}{L} \left[ A_{L+2} - \frac{(L-n-1)}{n} g_{L-1} \right], \quad h_L = 0; \]
\[ g_{L+1} = 0, \quad h_{L+1} = \frac{2}{(2L+1)} \left( \frac{(L-n)}{n} g_L - A_{L+3} \right) \]

for the unnormalized solution.

The recursion relations (34) and (35) always terminate after a finite number of terms, and so a closed form solution can be found for any particular values of \( n \) and \( L \) \((n \geq L+1)\). However, our aim here is to find the general functional dependence of \( E'(n) \) on \( n \) and \( L \). A direct derivation of the general solution for arbitrary \( n \) and \( L \) from the recursion relations appears to be a difficult task, even with the use of symbolic manipulation programs such as MACSYMA [21].

To circumvent this difficulty, we have followed the expedient of generating a sufficiently large array of values for \( g^{(2)}(n,L) \) in quadruple precision arithmetic (32 decimal digits) and deducing the general solution. To this end, values of \( g^{(2)}(n,L) \) were calculated for all \( 2 \leq L \leq 8 \) and \( L+1 \leq n \leq 20 \). To proceed from here a great deal can be learned about the form of the general solution from the particular cases

\[ \psi^{(2)}(n,n-1) = -\frac{2\sqrt{2}(128n^4 - 650n^3 + 848n^2 - 518n + 105)}{n^8(2n-5)(2n-1)(2n-1)(2n-3)^3} , \]
(36)
\[ \psi^{(2)}(n,n-2) = -\frac{2\sqrt{2}(128n^7 - 400n^6 - 2320n^5 + 12666n^4 - 19133n^3 + 4846n^2 + 10228n - 5880)}{n^9(2n-7)(2n-9)(2n-1)(2n-3)(2n-5)^3} , \]
(37)

derived by Drachman [21]. These results do not allow one to disentangle the dependence on \( n \) and \( L \), but one can extract the coefficients of \( 1/n^j \) in the expansion

\[ \psi^{(2)}(n,L) = \sum_j c_j(L)/n^j \]  
(38)

for fixed \( L \) by differencing our numerical values. The result is that only the powers \( 3 \leq j \leq 8 \) contribute. As a consequence, none of the factors in the denominators of (36) and (37) can be of the form \( (n \pm L + \cdots) \) because this would lead to infinite expansions in \( 1/n \). They must all be of the form \( (2L \pm p) \), where \( p \) is an integer. The denominators of (36) and (37) can therefore be written in the forms

\[ D(n,n-1) = n^8(2L-3) \]
\[ \times [(2L-1)(2L+1)(2L+2)]^3/8 \]  
(39)

and

\[ D(n,n-2) = n^9(2L-3) \]
\[ \times [(2L-1)(2L+1)(2L+2)]^3/64 . \]  
(40)

With this information in hand, it is relatively easy to identify the numerical coefficients \( c_j(L) \) as rational fractions. Observing how the prime factors in the denominators of the \( c_j(L) \) change with \( L \) for fixed \( j \) allows a unique identification with factors of \( (2L+p) \). The numerators are more complicated, but they can also be expressed as finite polynomials in \( L \) with integer coefficients. The final result is

\[ \psi^{(2)}(n,L) = -2\sqrt{2} \left[ \frac{(2L-3)!!}{(2L+3)!!} \right]^2 \left[ \frac{(2L-5)!!}{(2L+5)!!} \right]^2 \frac{\frac{9n^4}{[L(L+1)]^3} - \frac{6n^2}{[L(L+1)]^2}}{\left[ 45 + 7[9f_1(L) + 520f_2(L) + 80f_3(L)] \right] + \left[ 3 + 40[f_1(L) + 6f_2(L)] \right]} \]
\[ + \frac{1}{n^8} \left[ \frac{27n^4}{[L(L+1)]^2} - \frac{30n^2}{L(L+1)} + 7 \right] \]  
(41)
where
\[ f_1(L) = L(L + 1), \]
\[ f_2(L) = (L - 1)L(L + 1)(L + 2), \]
\[ f_3(L) = (L - 2)(L - 1)L(L + 1)(L + 2)(L + 3), \]
in accordance with definition (11). Finally \( E' \) is related to \( E' \) by
\[ E'(n) = Z_{\text{eff}}^2 E'(n). \] (42)

For the Rydberg states of helium, \( Z_{\text{eff}} = 1 \) and \( \alpha_1 = \alpha = a \). The numerical values obtained from the above general formula agree well with those tabulated by Drachman [21]. Explicit formulas for the first few \( L \) values are
\[
E^{(2)}(n, 2) = \frac{\alpha_1^2}{4[[13(5)(7)]^2} \left[ \frac{-10127}{3} \left( \frac{1}{n^3} - \frac{4}{n^5} \right) - \frac{5336}{n^7} \right] + \frac{\alpha_1^2}{2[[13(5)(7)]^2} \left[ - \frac{3}{n^4} + \frac{20}{n^6} - \frac{28}{n^8} \right],
\] (43)
\[
E^{(2)}(n, 3) = \frac{\alpha_1^2}{44[[5(7)(9)]^3} \left[ -94169 \left( \frac{1}{8n^3} - \frac{1}{n^5} \right) - \frac{78088}{n^7} \right] + \frac{\alpha_1^2}{4[[5(7)(9)]^2} \left[ - \frac{3}{2n^4} + \frac{20}{n^6} - \frac{56}{n^8} \right],
\] (44)
\[
E^{(2)}(n, 4) = \frac{\alpha_1^2}{260[[7(9)(11)]^3} \left[ -2487183 \left( \frac{3}{40n^3} - \frac{1}{n^5} \right) - \frac{697624}{n^7} \right] + \frac{\alpha_1^2}{2[[7(9)(11)]^2} \left[ L - \frac{27}{5n^4} + \frac{12}{n^6} - \frac{56}{n^8} \right],
\] (45)
\[
E^{(2)}(n, 5) = \frac{\alpha_1^2}{140[[9(11)(13)]^3} \left[ -319243 \left( \frac{1}{5n^3} - \frac{4}{n^5} \right) - \frac{540808}{n^7} \right] + \frac{\alpha_1^2}{2[[9(11)(13)]^2} \left[ - \frac{3}{25n^4} + \frac{4}{n^6} - \frac{28}{n^8} \right],
\] (46)
\[
E^{(2)}(n, 6) = \frac{\alpha_1^2}{68[[11(13)(15)]^3} \left[ -4445579 \left( \frac{1}{7n^3} - \frac{4}{n^5} \right) - \frac{359896}{n^7} \right] + \frac{\alpha_1^2}{2[[11(13)(15)]^2} \left[ - \frac{3}{49n^4} + \frac{20}{n^6} - \frac{28}{n^8} \right],
\] (47)
\[
E^{(2)}(n, 7) = \frac{\alpha_1^2}{836[[13(15)(17)]^3} \left[ -287142879 \left( \frac{3}{112n^3} - \frac{1}{n^5} \right) - \frac{5824024}{n^7} \right] + \frac{\alpha_1^2}{4[[13(15)(17)]^2} \left[ - \frac{27}{392n^4} + \frac{30}{7n^6} - \frac{56}{n^8} \right],
\] (48)
\[
E^{(2)}(n, 8) = \frac{\alpha_1^2}{364[[15(17)(19)]^3} \left[ -22740989 \left( \frac{1}{48n^3} - \frac{1}{n^5} \right) - \frac{3233288}{n^7} \right] + \frac{\alpha_1^2}{4[[15(17)(19)]^2} \left[ - \frac{1}{24n^4} + \frac{10}{3n^6} - \frac{56}{n^8} \right].
\] (49)

For the \( D \) states of He, the coefficient of the \( n^{-4} \) term corresponds to \(-70.81 \text{ MHz}\) and for the \( F \) states it corresponds to \(-1.967 \text{ MHz}\).

In view of the discussion in Sec. II, the even powers of \( n \) in Eq. (41) are of particular significance. Equation (4) predicts that the leading even terms are (in atomic units)
\[
T_n^{(\text{even})} = - \left[ \frac{380}{2n^4} + \frac{580\delta_2}{n^6} + \frac{7}{2n^8} (\delta_2^2 + 2\delta_0\delta_4 + \cdots) \right].
\] (50)
An examination of Eqs. (17)–(20) shows that only \( \delta_0 \) and \( \delta_2 \) depend in first order on \( c_4 = \alpha_1 \). Retaining only the \( \alpha_1 \) dependence, Eqs. (17) and (18) reduce to
\[
\delta_0^{(1)} = \frac{6\alpha_1}{L(L + 1)(2L - 1)(2L + 1)(2L + 3)},
\] (51)
\[
\delta_2^{(1)} = \frac{-2\alpha_1}{(2L - 1)(2L + 1)(2L + 3)}. \] (52)

Substituting (51) and (52) into (50) yields
\[ T_n(\text{even}) = \frac{-2a_i^2}{n^{8(2L-1)(2L+1)(2L+3)}} \times \left[ \frac{27n^4}{[L(L+1)]^2} - \frac{30n^2}{L(L+1)} + 7 \right], \]  
(53)

in exact agreement with the corresponding terms in Eq. (41) obtained from the second-order perturbation calculation. This provides a remarkable confirmation that Hartree's proof is an exact analytic result which can be used to advantage in deriving at least the even-\( n \) terms in second order perturbation energies. It also verifies the correctness of the coefficient of the \( n^{-8} \) term in Eq. (4). The numerical value \( \frac{1}{2} \) comes in part from the \( \delta \) expansion in the denominators of Eq. (2).

IV. COMPARISON WITH THEORETICAL AND EXPERIMENTAL TERM ENERGIES

A. Comparison with nonrelativistic energies for helium

High-precision variational calculations [13–15] and experimental data [12] are now available for the Rydberg states of helium. The purpose of this section is to construct \( 1/n \) and quantum-defect fits to the data in order to assess the validity of the Ritz expansion, and identify the contributions from \( E^{(2)} \).

As a preliminary, Table I shows the orders of magnitude to be expected for the coefficients \( a_i \) in the \( 1/n \) expansion

\[ T_n = -2R_M(1/2n^2 + a_3/n^3 + a_4/n^4 + a_5/n^5 + \ldots) \]  
(54)

for the term energy. The contributions are obtained from Eq. (4), using expressions (17)-(20) for the first-order quantum defects and Eq. (41) for \( E^{(2)} \). The totals are obtained from Drachman's [21] prescription for summing these asymptotic series; i.e., the total includes the contribution to \( a_i \) from \( \frac{1}{2}(c_{r}(r^{-7}) + c_{s}(r^{-8})) \) with \( \pm \frac{1}{2}(c_{r}(r^{-7}) + c_{s}(r^{-8})) \) being the uncertainty. For \( L=2 \) and 3, the total includes \( \frac{1}{2}c_{s}(r^{-6}) \) with \( \pm \frac{1}{2}c_{s}(r^{-6}) \) being the uncertainty. Previous comparisons of the total energies with high-precision variational calculations [15] show that these uncertainty estimates are in good accord with the actual errors resulting from truncating the asymptotic expansion. The numerical values of the core polarizabilities, etc., required to calculate the \( c_i \) coefficients in the asymptotic potential are

\[ a_1 = \frac{9}{32}, \quad a_2 = \frac{15}{64}, \quad a_3 = \frac{55}{1024}, \]

\[ \beta_1 = \frac{43}{512}, \quad \beta_2 = \frac{107}{2048}, \quad \gamma = \frac{319}{12288}, \]

\[ s_1' = \frac{213}{312}, \quad \epsilon = \frac{4329}{32768}. \]

The above values were checked independently and found to agree with those listed by Drachman [21].

Table I shows several interesting features. First, the dramatic increase in the rate of convergence and accur-
cy with increasing $L$ is clearly evident. Second, the coefficient $a_2$ is negative in magnitude and becomes progressively larger with increasing $L$ relative to the other terms. Since $a_3 \approx 6_0^{(1)}$ and $a_4 \approx 6_2^{(1)}$, the reason for this behavior is clear from the denominators in Eqs. (51) and (52). Third, $E^{(2)}$ makes a relatively small contribution to the odd-order terms, but it is solely responsible for the even-order terms. Since $E^{(2)}$ scales in proportion to $a_2^2$, it becomes relatively much larger in other hydrogenic atoms such as the alkali metals. For example, $\alpha_1$ (Cs$^+$) $\approx 56.2 \alpha_1$ (He$^+$) [28]. For $L = 2$, this makes the $E^{(2)}$ contribution to $a_3$ about the same size as the leading $c_4(r^{-4})$ contribution. Even for the $D$ states of helium, omitting the $E^{(2)}$ contribution would reduce the value of $\alpha_1$ deduced from the measured quantum defect by about 3%. Results for the alkali metals are further discussed below.

Turning now to the numerical fits, Table II lists nonrelativistic eigenvalues for the $P$, $D$, and $F$ states of helium, obtained by the double basis set variational calculations described previously [13–15]. A full account of the calculations is in preparation [29]. The exceptional accuracy of the results provides an ideal "experiment," free of relativistic and QED corrections, against which the Ritz expansion can be tested. The following three functional forms, all containing six adjustable parameters, will be compared:

\begin{equation}
T_n = -2R_M(1/2n^2 + a_3/n^3 + a_4/n^4 + a_5/n^5 + a_6/n^6 + a_7/n^7 + a_8/n^8),
\end{equation}

\begin{equation}
T_n = -R_M/(n - \delta_0 - \delta_1/n - \delta_2/n^2 - \delta_3/n^3 - \delta_4/n^4 - \delta_5/n^5)^2,
\end{equation}

\begin{equation}
T_n'' = -R_M/(n - \delta_0 - \delta_2/n^2 - \delta_4/n^4 - \delta_6/n^6 - \delta_8/n^8 - \delta_{10}/n^{10})^2.
\end{equation}

$T_n$ corresponds directly to Eq. (54) and the asymptotic numerical values for the coefficients shown in Table I. $T_n$ contains all inverse powers of $n^*$ in the quantum-defect expansion, while $T_n''$ contains only the even inverse powers (the Ritz expansion).

Beginning with $T_n$, Table III compares the first two or three fitting coefficients for the $P$, $D$, and $F$ states with the asymptotic values from Table I. In every case, the agreement is within the estimated accuracy of the asymptotic coefficients. The "Ritz" value for $a_4$ corresponds to $a_4 = \frac{3}{4}a_4$ variational [cf. Eq. (22)], and the Ritz defect is as defined by Eq. (24). In each case the Ritz defect is statistically significant, indicating that the variational eigenvalues cannot be represented by a Ritz expansion. However, the agreement is close enough to provide strong confirmation for the calculated $E^{(2)}$ contributions.

The $P$ states are a special case because the asymptotic

### Table II. Variational nonrelativistic eigenvalues for helium.

<table>
<thead>
<tr>
<th>State</th>
<th>$E$ (a.u.)</th>
<th>State</th>
<th>$E$ (a.u.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2^1P$</td>
<td>$-2.123 843 086 498 093(10)^a$</td>
<td>$2^1P$</td>
<td>$-2.133 164 190 779 274(9)$</td>
</tr>
<tr>
<td>$3^1P$</td>
<td>$-2.055 146 362 091 944(31)$</td>
<td>$3^1P$</td>
<td>$-2.058 081 084 274 274(34)$</td>
</tr>
<tr>
<td>$4^1P$</td>
<td>$-2.031 069 650 450 235(24)$</td>
<td>$4^1P$</td>
<td>$-2.032 324 354 296 619(16)$</td>
</tr>
<tr>
<td>$5^1P$</td>
<td>$-2.019 905 989 909 825(22)$</td>
<td>$5^1P$</td>
<td>$-2.020 551 187 256 247(23)$</td>
</tr>
<tr>
<td>$6^1P$</td>
<td>$-2.013 833 979 671 734(23)$</td>
<td>$6^1P$</td>
<td>$-2.014 207 958 773 734(12)$</td>
</tr>
<tr>
<td>$7^1P$</td>
<td>$-2.010 169 314 529 353(20)$</td>
<td>$7^1P$</td>
<td>$-2.010 404 960 007 936(21)$</td>
</tr>
<tr>
<td>$8^1P$</td>
<td>$-2.007 789 127 133 214(18)$</td>
<td>$8^1P$</td>
<td>$-2.007 947 013 771 117(13)$</td>
</tr>
<tr>
<td>$9^1P$</td>
<td>$-2.006 156 384 652 846(37)$</td>
<td>$9^1P$</td>
<td>$-2.006 267 267 366 410(42)$</td>
</tr>
<tr>
<td>$10^1P$</td>
<td>$-2.004 987 983 802 218(44)$</td>
<td>$10^1P$</td>
<td>$-2.005 068 805 497 766(99)$</td>
</tr>
<tr>
<td>$3^3D$</td>
<td>$-2.055 620 732 852 246(6)$</td>
<td>$3^3D$</td>
<td>$-2.055 636 309 453 261(3)$</td>
</tr>
<tr>
<td>$4^3D$</td>
<td>$-2.031 279 846 178 687(7)$</td>
<td>$4^3D$</td>
<td>$-2.031 288 847 501 795(3)$</td>
</tr>
<tr>
<td>$5^3D$</td>
<td>$-2.020 015 836 159 984(4)$</td>
<td>$5^3D$</td>
<td>$-2.020 021 027 464 911(6)$</td>
</tr>
<tr>
<td>$6^3D$</td>
<td>$-2.013 898 227 424 286(4)$</td>
<td>$6^3D$</td>
<td>$-2.013 901 415 453 793(7)$</td>
</tr>
<tr>
<td>$7^3D$</td>
<td>$-2.010 210 028 457 978(11)$</td>
<td>$7^3D$</td>
<td>$-2.010 212 105 955 595(3)$</td>
</tr>
<tr>
<td>$8^3D$</td>
<td>$-2.007 816 512 563 811(6)$</td>
<td>$8^3D$</td>
<td>$-2.007 817 934 711 076(3)$</td>
</tr>
<tr>
<td>$9^3D$</td>
<td>$-2.006 175 671 437 641(7)$</td>
<td>$9^3D$</td>
<td>$-2.006 176 684 884 697(3)$</td>
</tr>
<tr>
<td>$10^3D$</td>
<td>$-2.005 002 071 654 250(6)$</td>
<td>$10^3D$</td>
<td>$-2.005 002 818 080 233(10)$</td>
</tr>
<tr>
<td>$4^3F$</td>
<td>$-2.031 255 144 381 749 6(17)$</td>
<td>$4^3F$</td>
<td>$-2.031 255 168 403 245 6(7)$</td>
</tr>
<tr>
<td>$5^3F$</td>
<td>$-2.020 020 937 158 742 7(5)$</td>
<td>$5^3F$</td>
<td>$-2.020 027 957 377 369(4)$</td>
</tr>
<tr>
<td>$6^3F$</td>
<td>$-2.013 890 683 815 549 7(4)$</td>
<td>$6^3F$</td>
<td>$-2.013 890 698 348 532(0)$</td>
</tr>
<tr>
<td>$7^3F$</td>
<td>$-2.010 205 248 074 011 7(3)$</td>
<td>$7^3F$</td>
<td>$-2.010 205 258 374 864(1)$</td>
</tr>
<tr>
<td>$8^3F$</td>
<td>$-2.007 813 297 115 014 6(6)$</td>
<td>$8^3F$</td>
<td>$-2.007 813 304 535 090 84(1)$</td>
</tr>
<tr>
<td>$9^3F$</td>
<td>$-2.006 173 406 897 324 6(8)$</td>
<td>$9^3F$</td>
<td>$-2.006 173 412 365 043(7)$</td>
</tr>
<tr>
<td>$10^3F$</td>
<td>$-2.005 000 417 564 669 0(17)$</td>
<td>$10^3F$</td>
<td>$-2.005 000 421 686 604(11)$</td>
</tr>
<tr>
<td>$11^3F$</td>
<td>$-2.004 132 547 315 130 38(38)$</td>
<td>$11^3F$</td>
<td>$-2.004 132 550 488 506 7(35)$</td>
</tr>
</tbody>
</table>

*Numbers in parentheses indicate the uncertainties in the final one or two figures quoted.*
potential is no longer of use due large core penetration effects. The coefficient $a_3$ even becomes negative for the $n^3P$ states. Despite this, the Ritz estimate of $a_4$ remains remarkably accurate when applied separately to the $^1P$ and $^3P$ states, giving relatively small Ritz defects. However, the Ritz relation (22) becomes strongly violated if a singlet-triplet average of the variational eigenvalues is first formed, and the Ritz defect rises dramatically. The numerical values are $a_4, _{\text{variation}} = 3587.86(4) \times 10^{-6}$ a.u. and $a_4, _{\text{Ritz}} = 1183.60 \times 10^{-6}$ a.u., for a Ritz defect of 2404.26(4) $\times 10^{-6}$ a.u. The Ritz expansion will clearly be much more successful if the singlet-triplet average is not performed.

Table IV compares the results using the $T_n$, $T_n''$, and $T_n'''$ as eigenvalue functions for the quantum defect expansion. The quantity $\chi$ measures the goodness of fit (in the $\chi^2$ sense [30]), normalized to unity for an adequate fit. The uncertainties in the fitting coefficients were estimated both from the statistics of the least-squares fit [30] and by actually varying the input data up and down by the amount of their uncertainties. Both methods gave about the same results. Although some of the $\chi$ values appear to be rather large, the fits are still extremely good by usual standards. For example, the largest deviation from the data listed in Table II for the $T_n$ fit to the $n^3P$ states is only 242 kHz at $n = 5$. The largest deviations become much less for the higher $L$ states. The $T_n''$ and $T_n'''$ fits are true quantum-defect expansions in that $n^* \equiv n - \delta_0$ is determined by solving iteratively the equation

$$n^* = n - \delta_0 - \delta_1/n^* - \delta_2/n^{*2} - \cdots$$

for a given set of the $\delta_i$ coefficients, and then iterating the least-squares fit for the $\delta_i$ to convergence. The iteration automatically takes into account the expansion of the denominators discussed in Secs. II and III. Approximating the denominators by, for example, $n^* = n - \delta_0$ causes a remarkable deterioration of the accuracy of the fits by several orders of magnitude in the case of $T_n''$, and leads to unphysically large values for the higher-order coefficients. The $T_n''$ fits are only slightly affected. For the $D$ and $F$ states, the $\delta_{10}$ term in $T_n''$ is not included in the fits because it does not produce a further improvement in the quality as measured by $\chi$, and the values of $\delta_{10}$ are statistically consistent with zero.

At first sight, the $T_n'''$ fits appear to be clearly superior (as measured by $\chi$), especially for the $P$ states. However, close examination reveals subtle problems related to the Ritz defect. First, the values of $\delta_0$ are not consistent with the values obtained from $T_n$ and $T_n''$. The differences are much larger than the apparent uncertainties. The reason is that accuracy estimates for the coefficients in a least-squares fit are not reliable if the functional form is not correct. The high even powers of $1/n^*$ in $T_n''$ give it an important functional flexibility in fitting the low-$n$ data, but the absence of the $\delta_1/n^*$ term becomes relatively more serious in fitting the higher-$n$ data where the value of $\delta_0$ is most strongly determined. Second, $T_n'''$ seriously overestimates the accuracy of extrapolated term energies. For intermediate $n$ ($n < 100$), the error in $T_n''$ is dominated by the term $\delta_1/n^*$, while the error in $T_n'''$ decreases more rapidly as $\delta_2/n^{*2}$. However, this more rapid decrease is not real—it is a consequence of not including the Ritz defect in $T_n'''$. Notice that the values of $\delta_1$ in $T_n''$ are approximately equal to the Ritz defects listed in Table III.

The $T_n'''$ form has long been used to provide a compact and accurate representation of atomic term values. However, the $T_n$ form rapidly improves in accuracy with increasing $L$. The above results suggest that if the accuracy of the data is sufficient to reveal a statistically significant
Ritz defect, then the $T'_n$ form should be used, especially for purposes of extrapolating the term energies to higher $n$. This point should always be checked before adopting the Ritz expansion as a suitable functional form. The remarkable thing is that the Ritz expansion works so well. It is only the very high precision of the variational eigenvalues that makes the Ritz defect significant.

### B. Alkali-metal quantum defects and polarizabilities

As a final application of the results of this paper, we examine the effect of the second-order dipole polarization correction $E^{(2)}$ on polarizabilities extracted from the quantum defects of the alkali metals. Since $E^{(2)}$ grows in proportion to $a_1^2$, it becomes much more important for these atoms. Keeping only terms up to $(r^{-6})$ and neglecting relativistic corrections, core penetration effects, etc., $\delta_0$ can be written in the form

\[
\delta_0(D \text{ states}) = \frac{a_1}{105} + \frac{a_2 - 6\beta_1}{324} + \frac{10 127 a_1^2}{12\left[(3)(5)(7)\right]^3},
\]

\[
\delta_0(F \text{ states}) = \frac{a_1}{630} + \frac{a_2 - 6\beta_1}{17820} + \frac{941 69 a_1^2}{352\left[(5)(7)(9)\right]^3},
\]

These equations are based on the earlier ones of Swainson and related to the present work through the $T'_n$ form.
\[
\delta_0 (G \text{ states}) = \frac{\alpha_1}{2310} + \frac{\alpha_2 - 6\beta_1}{231660} + \frac{7461549\alpha_1^2}{52000[(7)(9)(11)]^2}. \tag{61}
\]

Similar equations for higher \( L \) can be immediately written down using Eqs. (17) and (41). Assuming that \( \delta_0 \) and the small correction term \( \alpha_2 - 6\beta_1 \) are known, the above are quadratic equations that can be solved for \( \alpha_1 \). A summary of experimental values of \( \delta_0 \) in the alkali metals has been given by Lorenzen and Niemax [9]. More accurate values for the cesium \( S, P, D, F \) and \( G \) states have been obtained by Weber and Sansonetti [31], who also point out significant systematic errors in the quantum defect fits obtained in Ref. [9]. Values for \( \alpha_1 \) and \( \alpha_2 \) calculated in the relativistic random phase approximation (RRPA) have been tabulated by Johnson, Kolb, and Huang [28]. For \( \text{Li}^+ \), the nonadiabatic correction factor \( \beta_1 \) is 0.03526\( \alpha_0^2/\varepsilon^2 \) [32]. For \( \text{Na}^+ \) and \( \text{K}^+ \), values of \( \beta_1 \) have been estimated by Eissa and Opik [33] from oscillator strength sum rules, using the method of Dalgarno and Kingston [34]. In this method, the input data are \( S(-1) = (\sum r_i^2)/3 \) and \( S(-2) = \alpha_1/4 \), where \( S(k) \) is the usual energy-weighted oscillator strength sum and the expectation value is with respect to the closed-shell ground state. These are used to calculate the constants \( a \) and \( b \) in the extrapolation formula [34]

\[
S(k) = \left[ \frac{\Delta E + a}{2.5-k} + \frac{b}{(2.5-k)^2} \right] S(0), \tag{62}
\]

where \( S(0) \) is the number of electrons and \( \Delta E \) is the smallest excitation energy (in Rydbergs) which contributes to the oscillator strength sums. Then \( S(-3) = \beta_1/4 \).

The required input values of \( S(-1) \) for the alkali-metal ions were estimated by scaling the isoelectronic "experimental" inert gas values derived by Dalgarno and Kingston [34] (their Table 3) according to

\[
S_{\text{alkali-metal ion}} = S_{\text{inert gas}} \left( \frac{\chi_m(\text{alkali-metal ion})}{\chi_m(\text{inert gas})} \right), \tag{63}
\]

where \( \chi_m \), the diamagnetic susceptibility, is proportional to the closely related quantity \( \langle \sum r_i^2 \rangle \). Using the Hartree-Fock values for \( \chi_m \) [35] and the polarizabilities in Table V, this procedure, together with Eq. (62), yields essentially the exact \( \beta_1 \) for \( \text{Li}^+ \) (\( \pm 0.2\% \)), and it reproduces the estimates of Eissa and Opik [33] for \( \text{Na}^+ \) and \( \text{K}^+ \) to within 5\% (when their older polarizabilities are used). The scaling factor in Eq. (63) ranges from 0.3776 for \( \text{Li}^+ \) to 0.8386 for \( \text{Cs}^+ \). The final results for \( \beta_1 \) are listed in Table V. The accuracy is adequate for purposes of the illustration to follow.

To illustrate the influence of the quadratic term in \( \alpha_1 \) in Eqs. (59)–(61), Table V shows the values of \( \delta_0 \) calculated with and without this term included [called \( \delta_0 \) (quad.) and \( \delta_0 \) (linear), respectively]. The effect is particularly large for the \( D \) states, where core penetration and exchange effects are also large [36]. In the case of Cs, the quadratic term increases \( \delta_0 \) by about 70\% to 0.442, but this is still much less than the observed value of 2.4663... due to core penetration. The \( G \) states of Cs are particularly interesting because core penetration and exchange effects become small. Estimates of these effects have been calculated by Sansonetti, Andrew, and Verges [36]. Extracting the coefficient of the leading \( 1/n^2 \) term from

**TABLE V.** Comparison of quantum defects for the alkali metals calculated from Eqs. (59)–(61) with the experimental values. \( \delta_0 \) (quad.) includes the term quadratic in the polarizability while \( \delta_0 \) (linear) does not. All quantities are in atomic units.

<table>
<thead>
<tr>
<th>States</th>
<th>( \alpha_1^a )</th>
<th>( \alpha_2^a )</th>
<th>( \beta_1 )</th>
<th>( \delta_0 ) (linear)</th>
<th>( \delta_0 ) (quad.)</th>
<th>( \delta_0 ) (expt.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \text{Li}^2D )</td>
<td>0.1894(^c)</td>
<td>0.1127</td>
<td>0.03526(^d)</td>
<td>0.001497</td>
<td>0.001523</td>
<td>0.002129</td>
</tr>
<tr>
<td>( \text{Na}^2D )</td>
<td>0.9457</td>
<td>1.521</td>
<td>0.206(^e)</td>
<td>0.00989</td>
<td>0.01054</td>
<td>0.015543</td>
</tr>
<tr>
<td>( \text{Na}^2F )</td>
<td>0.9457</td>
<td>1.521</td>
<td>0.206(^e)</td>
<td>0.001517</td>
<td>0.001525</td>
<td>0.001453</td>
</tr>
<tr>
<td>( \text{K}^2D )</td>
<td>5.457</td>
<td>16.27</td>
<td>2.39(^e)</td>
<td>0.0579</td>
<td>0.0796</td>
<td>0.2770</td>
</tr>
<tr>
<td>( \text{K}^2F )</td>
<td>5.457</td>
<td>16.27</td>
<td>2.39(^e)</td>
<td>0.0087</td>
<td>0.009025</td>
<td>0.01010</td>
</tr>
<tr>
<td>( \text{Rb}^2D )</td>
<td>9.076</td>
<td>35.41</td>
<td>4.40(^e)</td>
<td>0.1142</td>
<td>0.1743</td>
<td>1.34716</td>
</tr>
<tr>
<td>( \text{Rb}^2F )</td>
<td>9.076</td>
<td>35.41</td>
<td>4.40(^e)</td>
<td>0.01491</td>
<td>0.01562</td>
<td>0.01631</td>
</tr>
<tr>
<td>( \text{Cs}^2D )</td>
<td>15.81</td>
<td>86.4</td>
<td>8.50(^f)</td>
<td>0.2598</td>
<td>0.4421</td>
<td>2.466315f</td>
</tr>
<tr>
<td>( \text{Cs}^2F )</td>
<td>15.81</td>
<td>86.4</td>
<td>8.50(^f)</td>
<td>0.02708</td>
<td>0.02922</td>
<td>0.033414f</td>
</tr>
<tr>
<td>( \text{Cs}^2G )</td>
<td>15.81</td>
<td>86.4</td>
<td>8.50(^f)</td>
<td>0.006997</td>
<td>0.007105</td>
<td>0.007039f</td>
</tr>
</tbody>
</table>

\(^a\)RRPA results from Johnson, Kolb, and Huang (Ref. [28]).
\(^b\)Lorenzen and Niemax (Ref. [9]), except as noted.
\(^c\)The nonrelativistic variational value for \( \alpha_1 \) is 0.1924\( \alpha_0^2 \) (Ref. [32]).
\(^d\)Drake (Ref. [32]).
\(^e\)Oscillator strength sum rule extrapolation (see text).
\(^f\)Weber and Sansonetti (Ref. [31]).
their results, the contribution to δ0 is 0.000 026 to give a corrected quantum defect of 0.007 013 for the G states. Using the corrected value in Eq. (61) and solving for αt yields αt = 15.61a0. This is somewhat smaller than the RRPA value of 15.81a0 but within the range of probable accuracy. It is also smaller than the value 15.770(3)a0 obtained by Weber and Sansonetti [31], who did not include the quadratic term in their fit to the measured frequencies. The quadratic term accounts for most of the difference. For the F states, the corresponding penetration and exchange correction [36] to δ0 is 0.007 88. This is the same order of magnitude as the difference between the calculated and measured values shown in Table V. However, Eqs. (59)–(61) are based on just the leading two terms in the asymptotic potential, together with the quadratic term. For the F states, and possibly also for the G states, the leading two terms may not be sufficient at this level of accuracy.

V. DISCUSSION

The central result of this paper is the general expression (41) for the second-order dipole polarization correction E(2). The coefficients of the even powers of 1/n have been shown to be in accord with what one would expect from quantum-defect theory, using the Ritz expansion for the quantum defect. In fact Hartree's proof that the Ritz expansion is an exact analytic result for local central potentials provides a powerful method for finding at least the even terms in a 1/n expansion for the second-order energy resulting from any short-range perturbation. For example, the same techniques could be applied to the 1/r3 term in the asymptotic potential.

The exact effective potential experienced by the Rydberg electron is known to be nonlocal and energy dependent because of short-range and exchange effects. In this case, one would expect to see deviations from the Ritz expansion as an exact functional form for the n dependence of Rydberg energies. A 1/n expansion fit to high precision variational eigenvalues in fact shows a well-defined "Ritz defect," which represents the degree to which the Ritz expansion for the quantum defect is not valid. In this case, the odd as well as the even powers of 1/n should be included in the quantum-defect expansion. The lowest-order Ritz defect then reappears as the coefficient of 1/n. A particularly important point is that unless this term is known to vanish, it dominates the error of extrapolations of Rydberg energies up to moderately high values of n. Omitting it gives unrealistically small error estimates, even though a fit to the data may appear to be quite adequate. It is also important to bear in mind that the quality of the fit improves dramatically if a full iterative definition of the quantum defect is used. This implicitly takes into account the quantum-defect expansion in the denominators of Eq. (2).

The final conclusion of this paper is that the second-order polarization energy has a significant effect on polarizabilities extracted from experimental quantum defects. The results of Table V show that the accuracy of polarizabilities obtained in this way rapidly improves with increasing L. The G state result for Cs is particularly accurate, and provides an important test of calculations. However, a complete reanalysis of the experimental data with the quadratic term included would be desirable.

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