E1 amplitudes, lifetimes, and polarizabilities of the low-lying levels of atomic ytterbium

S. G. Porsev, Yu. G. Rakhlina, and M. G. Kozlov
Petersburg Nuclear Physics Institute, Gatchina, Leningrad district, 188350, Russia
(June 20, 2021)

Abstract

The results of ab initio calculation of E1 amplitudes, lifetimes, and polarizabilities for several low-lying levels of ytterbium are reported. The effective Hamiltonian for the valence electrons $H_{\text{eff}}$ has been constructed in the frame of CI+MBPT method and solutions of many electron equation $H_{\text{eff}}\Phi_n = E_n\Phi_n$ are found.

PACS. 31.15.Ar, 32.10.Dk, 32.10.Fn
I. INTRODUCTION

In this paper we report results of an *ab initio* calculation of E1 amplitudes, lifetimes, and polarizabilities for several low-lying levels of ytterbium. In Ref. [1] we calculated the energies and hyperfine structure (hfs) constants of low-lying levels of ytterbium. In that calculations the accuracy of atomic wave functions was tested at the short distances by comparison of the calculated hfs constants with the experimental ones. The latter are usually known to a very good accuracy, providing a good test of the quality of the wave function near the nucleus.

E1 amplitudes, in contrast, are determined by the wave function behavior at large distances. Usual experimental accuracy for the oscillator strengths and scalar polarizabilities is on the level of few percent. This is close or even less than the accuracy of precise atomic calculations (see, e. g., calculations for Ba [2] and Cs [3]). Tensor polarizabilities can be measured with the accuracy of 1% or better [4,5]. Thus, it is possible to test an atomic wave function at large distances at 1% level. Note that 1% accuracy is crucial for calculations of parity nonconservation effects in atoms, because it allows to test predictions of the Standard model at small momentum transfer [3,6]. So far such precision has been achieved only for one-electron atoms Cs and Fr [7–9]. In this work we deal with a much more complicated Yb atom.

We consider ytterbium as a two electron atom with the core [1s2,...,4f14]. Valence-valence correlations are taken into account by the configuration interaction (CI) method, while core-valence and core-core correlations are treated within the second order of the many-body perturbation theory (MBPT). The latter is used to construct an effective Hamiltonian for the CI problem in the valence space. The details of the method can be found in the papers [10,11]. Application of this method to calculations of hfs constants has been discussed in [1,2,12]. In Ref. [2] the method has been extended to the calculations of polarizabilities. Here we apply this technique for calculations of lifetimes, E1 amplitudes, and polarizabilities of ytterbium.

II. GENERAL FORMALISM

Let us write out several formulae that will be used in the following. The expression for oscillator strength for $a, J \rightarrow a', J'$ transition has the form [13] (atomic units $m = \hbar = e = 1$ are used throughout the paper):

$$f(a, J, a', J') = \frac{2}{3} \frac{\omega_{aJ, a'J'}^2}{(2J + 1)} \langle a, J || D || a', J' \rangle^2,$$

(1)

where $\omega_{aJ, a'J'} = E_{aJ} - E_{a'J'}$, $D$ is the dipole moment operator, and reduced matrix elements (MEs) are defined as follows:

$$\langle a', J', M' | D_q | a, J, M \rangle = (-1)^{J' - M'} \left( \begin{array}{c} J' \\ -M' \end{array} \begin{array}{c} J \\ q \end{array} M \right) \langle a', J' || D || a, J \rangle.$$

(2)

The lifetime $\tau$ of a level is the inverse of the total transition rate. The probability for $a, J \rightarrow a', J'$ transition is given by:
where \( c \) is the speed of light.

Static polarizability of the sublevel \(|a, J, M\rangle\) in a DC electric field \( \mathbf{E} = \mathcal{E}\mathbf{\hat{z}} \) is defined as:

\[
\Delta E_{a,J,M} = -\frac{1}{2} \alpha_{a,J,M} \mathcal{E}^2
\]

where \( \Delta E_{a,J,M} \) is the energy shift and \( \alpha_0 \) and \( \alpha_2 \) define the scalar and tensor polarizabilities, correspondingly. Being a second order property, \( \alpha_{a,J,M} \) can be expressed as a sum over unperturbed intermediate states:

\[
\alpha_{a,J,M} = -2 \sum_n \frac{|\langle a, J, M | D_z | n, J_n, M \rangle|^2}{E_a - E_n},
\]

where \( E_n \) is an unperturbed energy of a level \( n \), and the sum runs over all states of opposite parity. The formalism of the reduced MEs allows to write explicit expressions for the scalar and tensor parts of the polarizability:

\[
\alpha_{0,a,J} = -2 \frac{1}{3(2J+1)} \sum_n \frac{|\langle a, J, M | D | n, J_n \rangle|^2}{E_a - E_n},
\]

\[
\alpha_{2,a,J} = \left( \frac{40J(2J-1)}{3(2J+3)(2J+1)(J+1)} \right)^{1/2} \times \sum_n (-1)^{J+J_n+1} \begin{pmatrix} J & 1 & J_n \\ 1 & J & 2 \end{pmatrix} \frac{|\langle a, J | D | n, J_n \rangle|^2}{E_a - E_n}.
\]

In order to use Eqs. \( (3) \)–\( (7) \) in calculations one needs to know a complete set of eigenstates of the unperturbed Hamiltonian. It becomes practically impossible when dimension of a CI space exceeds few thousand determinants. It is known, that it is much more convenient to solve inhomogeneous equation instead of the direct summation over the intermediate states \[14,15\]. Indeed, let us consider the solution of the following equation:

\[
(E_a - H)|X_{a,M'}\rangle = D_q|a, J, M\rangle,
\]

where \( q = 0, \pm 1 \) and \( M' = M + q \). Obviously, the right hand side in Eq. \( (3) \) can be expressed in terms of the function \( X_{a,M} \) (note that \( D_0 \equiv D_z \)):

\[
\alpha_{a,J,M} = -2\langle a, J, M | D_0 | X_{a,M} \rangle.
\]

If we want to rewrite Eqs. \( (4) \) and \( (7) \) in terms of the function \( X_{a,M'} \), we need to decompose the latter in terms that correspond to particular angular momenta \( J_i \). Generally speaking, there can be three such terms with \( J_i = J, J \pm 1 \):

\[
X_{a,M'} = X_{a,J-1,M'} + X_{a,J,M'} + X_{a,J+1,M'}.
\]
Now, with the help of the functions $X_{a,j',m'}$ Eqs. (6) and (7) are reduced to:

$$\alpha_{0,a,J} = (-1)^{q+1} \frac{2}{3(2J+1)} \times \sum_{J'} \left( \begin{array}{ccc} J' & J & q \\ -M' & M & \end{array} \right)^{-2} \langle a, J, M | D_{-q} | X_{a,j',m'} \rangle,$$

$$\alpha_{2,a,J} = (-1)^{q+1} \frac{40J(2J-1)}{3(2J+3)(2J+1)(J+1)} \left( \begin{array}{ccc} J & 1 & J' \\ M & J & \end{array} \right)^{-2} \times \sum_{J'} (-1)^{J+J'} \left( \begin{array}{ccc} J & 1 & J' \\ M & J & \end{array} \right)^{-2} \langle a, J, M | D_{-q} | X_{a,j',m'} \rangle,$$

where sums run over $J' = J, J \pm 1$. Note, that these equations are valid only if all $3j$-symbols on the right hand side do not turn to zero. One has to take it into account when choosing for what spherical component $q$ to solve Eq. (8).

If we know the solution of Eq. (8) and its decomposition (11), then expressions (11) and (12) allow us to find both scalar and tensor polarizabilities of the state $|a, J\rangle$. Moreover, the same functions $X_{a,j',m'}$ can be also used to find other second order atomic properties, such as amplitudes of the Stark-induced $E1$ transitions or parity nonconserving $E1$ transitions between the states of the same nominal parity (see, for example, Ref. [16]).

III. CALCULATION DETAILS AND RESULTS

A. Orbital basis set and CI space.

The calculation procedure is quite similar to that of Ref. [2]. For this reason we give here only a brief description of its features. This calculation is done in the $V^N$ approximation, that means that core orbitals are obtained from the Dirac-Hartree-Fock (DHF) equations for a neutral atom (we use the DHF computer code [17]). The basis set for the valence electrons includes 6s, 6p, 5d, 7s, 7p, 6d DHF orbitals and 8s–15s, 8p–15p, 7d–14d, 5f–10f, and 5g–7g virtual orbitals. The latter were formed in two steps. On the first step we construct orbitals with the help of a recurrent procedure, which is similar to that suggested in Ref. [18] and described in Refs. [11,16]. After that we diagonalize the $V^N$ DHF operator to obtain the final set of orbitals.

For this orbital basis set the complete CI is made for both even-parity and odd-parity levels. Two-electron wave functions are the linear combinations of the Slater determinants with a given $J_z$. It means that no symmetrization with respect to angular momentum $J$ is made.

B. Effective operators.

Within the CI+MBPT method the wave function of the valence electrons is found from the eigenvalue equation:
\[ H_{\text{eff}} |a, J, M\rangle = E_a |a, J, M\rangle. \]  
(13)

Eq. (8) is rewritten as equation for valence electrons only:

\[ (E_a - H_{\text{eff}}) |X_{a,Mr}\rangle = D_{\text{eff},q} |a, J, M\rangle, \]  
(14)

with the effective operators, which are found by means of the MBPT. The effective Hamiltonian for two valence electrons is formed within the second order MBPT [10]. We used RPA for the effective dipole moment operator (see, for example, Ref. [19]). We have checked that MBPT corrections to \( D_{\text{eff}} \), which are not included in RPA, are small if RPA equations are solved with 6s electrons excluded from the self-consistency procedure. That means that RPA equations have the same form as in the \( V^N-2 \) approximation. The more detailed description of the effective operator formalism is given in [12].

C. Transition amplitudes and lifetimes.

We first solve eigenvalue Eq. (13) with the effective Hamiltonian for low-lying even- and odd-parity states. Strictly speaking, the effective Hamiltonian can be safely used only for the energy levels below the core excitation threshold. For Yb this threshold lies at 23189 cm\(^{-1}\) above the ground state [20]. However, it was shown in [1] that theoretical spectrum is quite good up to \( \sim 40000 \) cm\(^{-1}\). Correspondingly, we can work (with some caution) with the states lying slightly above the core excitation threshold. In our approach we fail to reproduce the states with unfilled \( f \) shell and correspondingly to account properly for the interaction with such states. For this reason we restrict ourselves to the consideration of the states lying sufficiently far from those with unfilled \( f \) shell. We consider \( E1 \) transitions between four low-lying odd-parity states \( (^3P_{0,1,2} \ (6s6p) \) and \( ^1P_1 \ (6s6p) \) and seven even-parity states \( (^1S_0 \ (6s^2), \ ^3D_{1,2,3} \ (5d6s), \ ^1D_2 \ (5d6s), \ ^3S_1 \ (6s7s), \) and \( ^1S_0 \ (6s7s) ) \). The state \( ^1P_1 \ (6s6p) \) requires special attention. The nearest \( f^{13} 5d \) \( 6s^2 \) state lies only 3800 cm\(^{-1}\) above the latter and their interaction is not negligible. We estimated that configuration \( f^{13} 5d \) \( 6s^2 \) contributes on the level of several percent to the wave function of \( ^1P_1 \ (6s6p) \) state. We do not take into account this configuration mixture. This reduces the accuracy of the calculated \( ^1P_1 \ (6s6p) \rightarrow ^1L_J \ E1 \) amplitudes.

When eigenfunctions for the valence electrons are found, we can calculate transition amplitudes and lifetimes. The results of calculations are presented in Table I. The magnitudes of the \( E1 \) amplitudes vary in a wide range. These variations correspond in part to the approximate selection rules \( \Delta S = 0 \) and \( \Delta J = \Delta L \), which are easily traced through Table I. For large amplitudes we estimate the accuracy of our calculation to be 3-5%. For the reason discussed above the amplitudes \( \langle ^1L_J |D| ^1P_1 \ (6s6p) \rangle \) do not follow this rule. The accuracy for these amplitudes, as well as for small amplitudes (\( \leq 0.5 \) a.u.), is about 15-20%.

Where available, we compare our results with those of other theoretical [21][22] and experimental [23][28] groups. For the convenience of comparison we recalculated the oscillator strengths and transition probabilities to the reduced MEs. Calculations in Ref. [21] were performed in the \( L-S \) coupling scheme. The simplest semiempirical method [23] was used then to evaluate the radial parts. In Ref. [22] the multiconfiguration Dirac-Fock method was used. The valence-core electronic correlations were included semiempirically. Comparing our results with the results of other theoretical works, one can see that it was important to account for the valence-core correlations.
Now, using Eq. (3) we can find the transition probabilities and the lifetimes of the levels (see Table III). In these calculations we used experimental transition frequencies. Therefore, the accuracy of these numbers depends only on the accuracy of the dominant transition amplitudes. As a result, the largest error (40\%) takes place for the states \( ^1S_0(6s7s) \) and \( ^1D_2(5d6s) \) where the transition to the state \( ^1P_0(6s6p) \) is dominant. For other states we estimate theoretical accuracy for the lifetimes as 10\% or better.

### D. Polarizabilities.

In order to find the polarizabilities we substitute eigenfunctions in the right hand side of Eq. (14) and solve corresponding inhomogeneous equation. After that Eqs. (11) and (12) give us \( \alpha_0 \) and \( \alpha_2 \). Results of these calculations are presented in Table III. It is seen that, unlike barium (see Ref. [2]), \( \alpha_2 \) has typically the same order of magnitude as \( \alpha_0 \). For this reason the theoretical accuracy for \( \alpha_2 \), as a rule, is similar to that for \( \alpha_0 \). In contrast, experimental data for \( \alpha_2 \) are usually much more precise and complete.

There are several sources of errors in the calculations of polarizabilities. Some of them are the same as for hfs calculations, and are connected with the inaccuracy in the wave functions and the effective operators (note, that RPA corrections to the dipole operator are much smaller than for hfs operators). The additional source of errors is the inaccuracy in eigenvalues. Finally, solving Eq. (14) we do not account for configurational 4-3 a.u.

The final accuracy of calculations is very different for different levels. For instance, the 95\% of the polarizability of the ground state \( ^1S_0(6s^2) \) is due to the ME \( \langle ^1S_0 | D | ^1P_0(6s6p) \rangle \). Supposing that this ME is calculated within the accuracy of 20\%, the latter for \( \alpha_0(1S_0) \) will be about 40\% (corresponding transition frequency is reproduced almost ideally \( \alpha_0(1S_0) = 266 \text{ a.u.} \) obtained in Ref. [30] where the Hartree-Fock method was used.

For the \( D_J(5d6s) \) states the situation is more complicated. There are large cancellations between contributions of \( P_J(6s6p) \) states and higher-lying states. For this reason their polarizabilities are small and the role of different small contributions is enhanced. Thus, analysis of the accuracy becomes difficult; only for the tensor polarizability of \( ^3D_3(5d6s) \) state we can estimate the accuracy to be 20\%. All other values of \( \alpha_0 \) and \( \alpha_2 \) for \( D_J(5d6s) \) states presented in Table III are rather estimates by the order of magnitude.

The scalar polarizabilities of the levels \( ^3S_1(6s7s) \) and \( ^1S_0(6s7s) \) are basically determined by the MEs \( \langle S_J(6s7s) | D | P_J(6s6p) \rangle \). Because of the closeness of \( f^{13}5d^26s \) states we failed to obtain the reliable wave functions for \( P_J(6s6p) \) states. Correspondingly, the values for \( \alpha_0(3S_1(6s7s)) \) and \( \alpha_0(1S_0(6s7s)) \) are also only the estimates.

Now let us go over to the odd-parity states. The accuracy of \( \alpha_0 \) and \( \alpha_2 \) for \( ^3P_0(6s6p) \)-triplet is 6-10\%. The main contribution here comes from the \( ^3D_J(5d6s) \) multiplet and there are no cancellations because all important levels of opposite parity lie above and contribute
with the same sign. The accuracy for $\alpha_0$ of $^1P^o_1(6s6p)$ state is about 40% and for $\alpha_2$ even worse (50%). This is due to the large contribution of the intermediate state $^1S_0(6s7s)$ to these polarizabilities (see above).

In Ref. [31] the Stark shift of the $^1S_0(6s^2) \rightarrow ^3P^o_1(6s6p)$ transition in ytterbium was measured. The Stark shift rate was found to be $K = -61.924 (0.193)$ a.u. In terms of polarizabilities this magnitude can be written as:

$$K = -\frac{1}{2} \left\{ \alpha_0 (^3P^o_1) - 2\alpha_2 (^3P^o_1) - \alpha_0 (^1S_0) \right\}.$$ 

Using the numbers from Table III, we find that $K = -55 (9)$ a.u., in good agreement with the experimental result [31].

In Ref. [23], the Stark shifts for $^1S_0(6s^2) \rightarrow ^3D_{1,2}(5d6s)$ transitions were observed. These shifts depend on the differences in scalar polarizabilities ($\alpha_0 (^1S_0) - \alpha_0 (^3D_{1,2})$):

$$\alpha_0 (^1S_0) - \alpha_0 (^3D_1) = \begin{cases} 71 & \text{theory}, \\ 86 (3) & \text{experiment}, \end{cases}$$

$$\alpha_0 (^1S_0) - \alpha_0 (^3D_2) = \begin{cases} 82 & \text{theory}, \\ 80 (4) & \text{experiment}, \end{cases}$$

where theoretical values are taken from Table III.

The method used here allows us to calculate not only static polarizabilities, but, also the Stark-induced amplitudes for different transitions. For instance, the magnitude of the vector transition polarizability $|\beta|$ for $^1S_0(6s^2) \rightarrow ^3D_1(5d6s)$ transition was calculated to be $122 (12)$ a.u., in good agreement with our previous calculation $138 (30)$ a.u. [32] and experimental result $113 (15)$ a.u. [23].

E. Conclusion.

Application of the effective operator technique to the Yb atom is hampered by the existence of the shallow 4f shell. Nevertheless, it is possible to make reliable calculations of different atomic properties including transition frequencies, hyperfine constants, E1 amplitudes, lifetimes and polarizabilities for many low-lying energy levels. It is of a particular importance, that with some caution calculations can be done even for levels above the core excitation threshold, which for Yb lies at 23189 cm$^{-1}$.

I. ACKNOWLEDGMENTS

This work was supported in part by Russian Foundation for Basic Research, Grant No. 98-02-17663. One of us (SP) is grateful to the St. Petersburg government for financial support, Grant No. M98-2.4P-522.
REFERENCES

[1] S. G. Porsev, Yu. G. Rakhlina, and M. G. Kozlov, J. Phys. B 32 (1999) (in press).
[2] M. G. Kozlov and S. G. Porsev, Eur. Phys. J. D 5, 59 (1999).
[3] V. A. Dzuba, V. V. Flambaum, and O. P. Sushkov, Phys. Rev. A 56, R4357 (1997).
[4] A. Kreutztrager and G. von Oppen, Z. Phys. 265, 421 (1973).
[5] A. Fukumi, I. Endo, T. Horiguchi, Y. Ishida, T. Kondo, T. Kuwamoto, H. Matsuzaki, T. Nakamura, and T. Takahashi, Z. Phys. D 42, 243 (1997).
[6] W. Marciano and J. L. Rosner, Phys. Rev. Lett. 65, 2963 (1990).
[7] V. A. Dzuba, V. V. Flambaum, and O. P. Sushkov, Phys. Lett. A 141, 147 (1989).
[8] S. A. Blundell, W. R. Johnson, and J. Sapirstein, Phys. Rev. A 43, 3407 (1991).
[9] V. A. Dzuba, V. V. Flambaum, and O. P. Sushkov, Phys. Rev. A 51, 3454 (1995).
[10] V. A. Dzuba, V. V. Flambaum, and M. G. Kozlov, Pis'ma Zh. Eksp. Teor. Fiz. 63, 844 (1996) [JETP Lett. 63, 882 (1996)]; Phys. Rev. A 54, 3948 (1996).
[11] M. G. Kozlov and S. G. Porsev, Zh. Eksp. Teor. Fiz. 111, 838 (1997) [JETP 84, 461 (1997)].
[12] V. A. Dzuba, M. G. Kozlov, S. G. Porsev, and V. V. Flambaum, Zh. Eksp. Teor. Fiz. 114, 1636 (1998) [JETP 87, 885 (1998)].
[13] I. I. Sobelman, Atomic Spectra and Radiative Transitions (Springer-Verlag, Berlin, 1992).
[14] R. M. Sternheimer, Phys. Rev. 80, 102 (1950); 84, 244 (1951); 86, 316 (1952).
[15] A. Dalgarno and J. T. Lewis, Proc. Roy. Soc. 233, 70 (1955).
[16] M. G. Kozlov, S. G. Porsev, and V. V. Flambaum, J. Phys. B 29, 689 (1996).
[17] V. F. Brattsev, G. B. Deineka, and I. I. Tupitsyn, Izv. Akad. Nauk SSSR 41, 2655 (1977) [Bull. Acad. Sci. USSR, Phys. Ser. 41, 173 (1977)].
[18] P. Bogdanovich and G. Zukauskas, Sov. Phys. Collection 23, 13 (1983).
[19] A. M. Mårtensson-Pendrill, in Methods in Computational Chemistry, Vol. 5: Atomic and Molecular Properties, edited by S. Wilson (Plenum Press, New-York, 1992).
[20] W. C. Martin, R. Zalubas, and L. Hagen, Atomic Energy Levels – The Rare Earth Elements, Natl. Bur. Stand. Ref. Data Ser., Natl. Bur. Stand (U.S.) Circ. No. 60 (U.S. GPO, Washington, DC, 1978).
[21] M. D. Kunisz, Acta Phys. Polon. A 62, 285 (1982).
[22] J. Migdalek and W. E. Baylis, J. Phys B 24, L99 (1991).
[23] C. J. Bowers, D. Budker, E. D. Commins, G. Gwinner, S. J. Freedman, J. E. Stalnaker, and D. DeMille, Lawrence Berkeley National Laboratory Report No. 42454 (1998) (unpublished).
[24] M. Baumann and G. Wandel, Phys. Lett. 22, 283 (1966).
[25] N. P. Penkin, K. B. Blagoev and V. A. Komarovskii, Atomic Physics VI (Riga, 1978).
[26] T. Andersen, O. Poulsen, P. S. Ramanujam, and A. Petrakiev-Petkov, Solar Physics 44, 257 (1975).
[27] C. J. Bowers, D. Budker, E. D. Commins, D. DeMille, S. J. Freedman, A.-T. Nguyen, and S.-Q. Shang, Phys. Rev. A 53, 3103 (1996).
[28] Y. S. Bai and T. W. Mossberg, Phys. Rev. A 35, 619 (1987).
[29] D. R. Bates and A. Damgaard, Phil. Trans. Roy. Soc. A 242, 101 (1949).
[30] K. M. S. Saxena and S. Fraga, J. Chem. Phys. 57, 1800 (1972).
[31] J. Li and W. A. van Wijngaarden J. Phys B 28, 2559 (1995).
[32] S. G. Porsev, Yu. G. Rakhlina, and M. G. Kozlov, Pis’ma Zh. Eksp. Teor. Fiz. 61, 449 (1995) [JETP. Lett. 61, 459 (1995)].

[33] W. Lange, J. Luther, and A. Steudel, in Proceedings 2nd Conference of European Group for Atomic Spectroscopy (Hanover, 1970), p.31.

[34] M. Baumann, M. Braun, A. Gaiser, and H. Liening, J. Phys. B 18, L601 (1985).

[35] K. B. Blagoev and V. A. Komarovskii, At. Data Nucl. Data Tables 56, 1 (1994).

[36] R. H. Rinkleff, Z. Phys A 296, 101 (1980).

[37] P. Kulina and R. H. Rinkleff, Z. Phys A 304, 371 (1982).
TABLES

TABLE I. Reduced MEs $|\langle J_{r}||L_{J}^r||L'_{J}^r \rangle|$ (a.u.). Calculations were made in the $L$-gauge. Other theoretical and experimental results are given where available. The uncertainties are indicated in the parentheses.

| State Config. | This work | Other data |
|---------------|-----------|------------|
| $^3P^o_0(6s6p)$ | 0.54 (8) | 4.40 (80) |
| $^1S^o_0(6s^2)$ | — | 0.44$^a$ |
| | | 0.549 (4)$^c$ |
| | | 0.553 (13)$^d$ |
| $^3D_1(5d6s)$ | 2.26 (10) | 2.2 (1)$^g$ |
| | | 0.24$^a$ |
| $^3D_2(5d6s)$ | 4.03 (16) | 2.39 (10) |
| | | 0.32 (6) |
| $^3D_3(5d6s)$ | — | 6.12 (30) |
| $^1D_2(5d6s)$ | — | 0.54 (10) |
| $^3S_1(6s7s)$ | 3.53 (15) | 5.05 (20) |
| | | 0.73 (15) |
| $^1S_0(6s7s)$ | 0.22 (4) | — |
| | | 0.22 (2)$^h$ |

Theory: $^a$Ref. [22]; $^b$Ref. [21];
Experiment: $^c$Ref. [23], $^d$Ref. [24], $^e$Ref. [25], $^f$Ref. [26], $^g$Ref. [27], $^h$Ref. [28].

TABLE II. Lifetimes (nsec) of the low-lying levels for Yb calculated with the reduced MEs from Table I and experimental transition frequencies.

| State Config. | This work | Other data |
|---------------|-----------|------------|
| $^3P^o_1(6s6p)$ | 875 (250) | 760–875$^e$ |
| | | 1294$^f$ |
| $^3P^o_2(6s6p)$ | 15.0 (1.5) sec | 14.5 sec$^f$ |
| $^1P^o_1(6s6p)$ | 5 (2) | 5.1–6.4$^e$ |
| | | 4.78$^f$ |

Experiment: $^a$Ref. [27], $^b$Ref. [33], $^c$Ref. [34], $^d$Ref. [28], $^e$see Ref. [35] and references therein;
Theory: $^f$Ref. [22].
TABLE III. Scalar and tensor polarizabilities (a.u.) of low-lying levels of Yb. Theoretical accuracy is indicated where analysis was possible, otherwise the numbers should be considered as estimates.

| level | config. | Theory | Experiment |
|-------|---------|--------|------------|
| $^1S_0$ | $6s^2$ | $\alpha_0$ | $\alpha_2$ | $\alpha_2$ |
| $^3D_1$ | $5d6s$ | 47 | 22 | 28 (4) $^a$ |
| $^3D_2$ | $5d6s$ | 36 | 17 | 28 (8) $^a$ |
| $^3D_3$ | $5d6s$ | $-9$ | 118 (24) | |
| $^1D_2$ | $5d6s$ | 4 | 150 | |
| $^1S_0$ | $6s7s$ | 2072 | |
| $^3S_1$ | $6s7s$ | 2030 | 0.8 | |
| $^3P_0^o$ | $6s6p$ | 252 (25) | | |
| $^3P_1^o$ | $6s6p$ | 278 (15) | 24.3 (1.5) | 24.06 (1.37) $^b$ |
| | | | | 24.26 (0.84) $^c$ |
| | | | | 23.33 (0.52) $^d$ |
| $^3P_2^o$ | $6s6p$ | 383 (32) | $-76$ (6) | |
| $^1P_1^o$ | $6s6p$ | 501 (200) | $-118$ (60) | $-57.4$ (5.6) $^b$ |

$^a$Ref. [23], $^b$Ref. [30], $^c$Ref. [37], $^d$Ref. [31].