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Theoretical studies of the long lifetimes of the $6d\,^2D_{3/2,5/2}$ states in Fr: Implications for parity-nonconservation measurements

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Lifetimes of the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states in Fr are determined from calculations of the radiative transition amplitudes of the allowed electric dipole ($E1$) and the forbidden electric quadrupole ($E2$) and magnetic dipole ($M1$) channels which were performed using the second-order many-body perturbation theory and the coupled-cluster method at different levels of approximations in the relativistic framework. The values obtained for these two quantities are 540(10) and 1704(32) ns, respectively. These relatively long lifetimes and the large electric dipole parity-non-conserving amplitudes of $7s\,^2S_{1/2} \to 6d\,^2D_{3/2,5/2}$ transitions strongly favor Fr as a leading candidate for the measurement of parity nonconservation arising from the neutral-current weak interaction and the nuclear anapole moment. In another important application, these $6D$ states in Fr can be used efficiently for resonance ionization spectroscopic techniques to carry out precise measurements of the properties of the higher excited states due to the long lifetimes of these states.

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I. INTRODUCTION

Francium is considered to be a promising candidate for the measurements of the electric dipole moment (EDM) arising from the violations of parity and time-reversal symmetries [1–4], parity-nonconservation (PNC) effects due to the neutral-current weak interaction [2,5], and the nuclear anapole moment [6–8] as it is the heaviest alkali-metal atom. All the ongoing Fr PNC experiments involve $S-S$ transitions [2,5,7]. However, relativistic many-body calculations show that the PNC amplitudes in the $7s\,^2S_{1/2} \to 6d\,^2D_{3/2,5/2}$ transitions in Fr are about three times larger than that of the $7s\,^2S_{1/2} \to 8s\,^2S_{1/2}$ transition [8–10]. The PNC studies of the $S-D$ transitions of singly ionized Ba, Ra, and Yb have been the subject of theoretical investigations [8,11–13], and the principles of their measurements have been discussed [14,15]. It has also been highlighted that the PNC measurements for the S-$D_{5/2}$ transitions in these ions would provide unambiguous signatures of the existence of the nuclear anapole moment (NAM) [16,17], which is still an open question [18–21]. Apart from exhibiting large PNC effects, another important aspect of these transitions is that the excited D states in these ions are metastable stables, and they provide long interrogation times which are very useful for carrying out high-precision measurements of the small PNC effects [14]. In this paper, we present the results of our theoretical studies of the lifetimes of the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states in Fr, which were undertaken to assess the feasibility of the measurement of PNC in this atom using the $7s\,^2S_{1/2} \to 6d\,^2D_{3/2,5/2}$ transitions.

The resonance ionization spectroscopy (RIS) technique is a very efficient multistep photon absorption process to access a higher excited state above the ionization threshold by driving through pulsed lasers [22,23]. It helps in the complete conversion of a quantum-selected excited-state population to ionization, making it possible to perform sensitive and absolute measurements of the selected population within a very large collection of atoms. When a pulsed laser is tuned to promote an electronically excited atomic state to an intermediate bound state, a second photon from the pulsed laser completes the resonance ionization process. This technique has been employed in Fr recently to measure lifetimes and hyperfine splittings of the excited $8s\,^2S_{1/2}$ and $9s\,^2S_{1/2}$ states considering the $7P$ and $8P$ states as the intermediate bound states [24,25]. However, it would be desirable in the RIS process to consider intermediate bound states that have longer lifetimes (approximately microseconds) than the P states. It has been demonstrated that the $5D$ states of a Cs atom, which belongs to the same group in the periodic table as Fr, have long lifetimes and have been efficiently used in the RIS techniques [26]. Thus, it would be reasonable to expect that the $6D$ states of Fr, with their long lifetimes, would also be suitable for RIS techniques to carry out measurements or improve the accuracies of many spectroscopic properties such as isotope shifts, hyperfine splittings, lifetimes, etc., of higher excited states of Fr.

II. THEORY

An electron from the $6d\,^2D_{3/2}$ state can decay to its low-lying $7p\,^2P_{1/2}$ and $7p\,^2P_{3/2}$ states by the electric dipole ($E1$), forbidden magnetic dipole ($M2$), and forbidden electric octupole ($E3$) channels and to the ground state by the forbidden magnetic dipole ($M1$) and electric quadrupole ($E2$) transitions. We neglect contributions due to the $M2$ and $E3$ channels as their transition probabilities are very weak. Similarly, an electron from the $6d\,^2D_{5/2}$ state can decay to its

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fine-structure partner $6d\,^2D_{3/2}$ state via both the $M1$ and $E2$ channels, to the low-lying $7p\,^2P_{3/2}$ state by the $E1$ channel, and to the ground state by the $E2$ transition. In this case too, we have omitted contributions due to the $M2$ and $E3$ channels. The transition probabilities due to the above $E1$, $E2$, and $M1$ channels for a transition, say, $|\Psi_i\rangle \to |\Psi_f\rangle$, are given by

$$A_{if}^{E1} = \frac{2.0261 \times 10^{-6}}{\lambda_{if}^3 g_i} S_{if}^{E1}, \tag{1}$$

$$A_{if}^{E2} = \frac{1.1195 \times 10^{-22}}{\lambda_{if}^5 g_i} S_{if}^{E2}, \tag{2}$$

$$A_{if}^{M1} = \frac{2.6971 \times 10^{-11}}{\lambda_{if}^3 g_i} S_{if}^{M1}, \tag{3}$$

where the quantity $S_{if}^O =| \langle \Psi_i || O || \Psi_f \rangle |^2$ is known as the line strength of the corresponding reduced matrix element $|\langle \Psi_i || O || \Psi_f \rangle|$ of the transition operator O. These quantities are later given in atomic units. In the above expressions, $g_i = 2J_i + 1$ is the degeneracy factor of state $|\Psi_i\rangle$ with angular momentum J_i, and the transition wavelength λ_{if} is given in nanometers (nms); when substituted, the transition probabilities A_{if}^O are obtained in s^{-1}. The lifetime τ of the atomic state $|\Psi_i\rangle$ is determined by taking the reciprocal of the total emission transition probabilities involving all the possible spontaneous transition channels (in seconds), i.e.,

$$\tau_i = \frac{1}{\sum_{O,f} A_{if}^O}, \tag{4}$$

where the summations over O and f correspond to all probable decay channels and all the lower states, respectively. We have attempted to obtain accurate results for the transition probabilities, hence the lifetimes of the atomic states from performing relativistic many-body calculations of the line strengths and using wavelengths that are determined from the experimental transition energies given in the National Institute of Science and Technology (NIST) database [27].

III. METHODS FOR CALCULATIONS

To investigate the role of the electron correlation effects in the evaluation of the radiative transition amplitudes, we employ the second-order many-body perturbation theory [MBPT(2)] and the coupled-cluster (CC) method in the relativistic framework. Further, we take different levels of approximation in the CC method to see the convergence of the results. We give below a brief description of these methods using Bloch's prescription [28], in which the atomic wave function of state $|\Psi_n\rangle$ is expressed as

$$|\Psi_n\rangle = \Omega_n |\Phi_n\rangle, \tag{5}$$

where Ω_n and $|\Phi_n\rangle$ are known as the wave operator and reference state, respectively. The ground and the considered excited $6D$ states of Fr have electronic configurations $[6p^6]7s\,^2S_{1/2}$ and $[6p^6]6d\,^2D_{3/2,5/2}$, respectively. To reduce the computational effort, we construct these states by creating a common reference state $|\Phi_c\rangle$ with the $[6p^6]$ configuration using the Dirac-Hartree-Fock (DHF) method. In this approach the atomic Hamiltonian H in the Dirac-Coulomb interaction approximation is divided into the DHF Hamiltonian H_0 and residual Coulomb interaction V_r. For the calculation of the exact states with a valence orbital, we define new working reference states as $|\Phi_n\rangle = a_n^\dagger |\Phi_c\rangle$. Here $a_n^\dagger$ appends an electron from the respective valence orbital denoted by an index n. As a consequence, Ω_n can now be divided as

$$\Omega_n = 1 + \chi_c + \chi_n, \tag{6}$$

where χ_c and χ_n are responsible for carrying out excitations (generating configuration state functions) from $|\Phi_c\rangle$ and $|\Phi_n\rangle$, respectively, due to V_r. In a perturbative series expansion, we have

$$\chi_c = \sum_k \chi_c^{(k)}, \quad \chi_n = \sum_k \chi_n^{(k)}. \tag{7}$$

In these expressions, the superscripts refer to the number of times V_r is considered in a particular calculation, and they represent the order of perturbation; for example, second-order MBPT, i.e., MBPT(2), has terms up to two V_r ($k = 2$).

Using the generalized Bloch equation, kth-order amplitudes for the χ_c and χ_n operators are obtained by [28]

$$\left[\chi_c^{(k)}, H_0\right] P = Q V_r \left(1 + \chi_c^{(k-1)}\right) P \tag{8}$$

and

$$\left[\chi_n^{(k)}, H_0\right] P = Q V_r \left(1 + \chi_c^{(k-1)} + \chi_n^{(k-1)}\right) P - \sum_{m=1}^{k-1} \chi_n^{(k-m)} \times P V_r \left(1 + \chi_c^{(m-1)} + \chi_n^{(m-1)}\right) P, \tag{9}$$

where the projection operators $P = |\Phi_c\rangle\langle\Phi_c|$ and $Q = 1 - P$ describe the model space and the orthogonal space of the Hamiltonian H_0, respectively. Note that here $\chi_c^{(0)} = 0$ and $\chi_n^{(0)} = 0$. Using these amplitudes, the energy of the state $|\Psi_n\rangle$ is evaluated by using an effective Hamiltonian,

$$H_n^{\mathrm{eff}} = P H \Omega_n P. \tag{10}$$

Using the CC ansatz, the above expressions can be put together to construct a wave operator to infinite order as

$$|\Psi_n\rangle = \Omega_n |\Phi_n = e^T \{1 + S_n\} |\Phi_n\rangle, \tag{11}$$

such that $\chi_c = e^T - 1$ and $\chi_n = e^T S_n - 1$. Here T and S_n are the CC excitation operators that excite electrons from the core and core along with the valence orbitals, respectively. In this work, we have accounted for only the single and double excitations, which are denoted using the subscripts 1 and 2, respectively, in the CC operators as

$$T = T_1 + T_2, \qquad S_n = S_{1n} + S_{2n}. \tag{12}$$

This is referred to as the CCSD method in the literature. When only the linear terms are retained in Eq. (11) with single and double approximation, it is referred to as the LCCSD method. The amplitudes of the T and S_n operators are determined using the expressions

$$H_N \chi_c P = Q H_N P \tag{13}$$

and

$$H_N \chi_n P = Q H_N (1 + \chi_c) P - \chi_n H_N^{\mathrm{eff}}, \tag{14}$$

where we have defined the normal-order Hamiltonian $H_N = H - PHP$ and the effective Hamiltonian to evaluate the

TABLE I. Magnitudes of the reduced matrix elements $\langle J_i||O||J_f\rangle$ of the transition operators O (in a.u.) from different many-body methods. Uncertainties from the finite-size basis set and noninclusion of the Breit interaction and due to QED effects are quoted using the MBPT(2) method. Wavelengths λ_{if} from the NIST database [27] (in nm) are given for the respective transitions.

									Uncertainties			Others
$J_i \to J_f$	λ_{if}	DHF	MBPT(2)	LCCSD	CCSD$^{(2)}$	CCSD$^{(4)}$	CCSD$^{(\infty)}$	CCSD(T)	Basis	Breit	QED	[9]
					*E*1 matrix elements							
$6d_{3/2} \to 7p_{1/2}$	2504.7	9.22	7.73	6.81	7.46	7.47	7.47	7.43	0.05	−0.01	0.001	7.174
$6d_{3/2} \to 7p_{3/2}$	4336.7	4.28	3.57	3.12	3.44	3.45	3.45	3.42	0.02	−0.01	∼0.0	3.301
$6d_{5/2} \to 7p_{3/2}$	3991.0	12.80	10.83	9.68	10.54	10.55	10.55	10.51	0.07	−0.02	0.001	10.156
					*E*2matrix elements							
$6d_{3/2} \to 7s_{1/2}$	616.5	43.10	33.74	31.39	34.02	34.06	34.06	33.78	0.20	−0.03	−0.02	33.367
$6d_{5/2} \to 7s_{1/2}$	608.7	52.74	41.69	39.31	42.24	42.27	42.27	41.96	0.24	−0.04	−0.02	41.568
$6d_{5/2} \to 6d_{3/2}$	50057.6	47.70	32.14	27.18	32.01	32.03	32.03	31.49	0.55	−0.16	0.01	
					*M*1 matrix elements							
$6d_{3/2} \to 7s_{1/2}$	616.5	∼0.0	0.001	0.002	0.001	0.001	0.001	0.001	∼0.0	∼0.0	∼0.0	
$6d_{5/2} \to 6d_{3/2}$	50057.6	1.549	1.550	1.548	1.547	1.552	1.552	1.552	∼0.0	0.001	∼0.0	

ionization potential (IP) of an electron from the valence orbital n of the respective state is given by $H_N^{\text{eff}} = PH_N(1+\chi_c+\chi_n)P$. We have also included contributions from important triple excitations perturbatively from $\chi_c^{(2)}$ and $\chi_n^{(2)}$ in the construction of H_N^{eff}, and this approach is referred to as the CCSD(T) method.

After obtaining amplitudes using the above equations, the transition matrix element of an operator O between states $|\Psi_i\rangle$ and $|\Psi_f\rangle$ is evaluated using the expression

$$\frac{\langle\Psi_f|O|\Psi_i\rangle}{\sqrt{\langle\Psi_f|\Psi_f\rangle\langle\Psi_i|\Psi_i\rangle}} = \frac{\langle\Phi_f|\Omega_f^\dagger O\Omega_i|\Phi_i\rangle}{\sqrt{\langle\Phi_f|\Omega_f^\dagger\Omega_f|\Phi_f\rangle\langle\Phi_i|\Omega_i^\dagger\Omega_i|\Phi_i\rangle}}. \tag{15}$$

This gives rise to a finite number of terms for the MBPT(2) and LCCSD methods, but it involves two nonterminating series in the numerator and denominator, which are $e^{T^\dagger}Oe^T$ and $e^{T^\dagger}e^T$, respectively, in both the CCSD and CCSD(T) methods. In order to evaluate all the significant contributions from these series, we have used Wick's generalized theorem [28] to divide these terms into the effective one-body, two-body, and three-body terms. The effective one-body terms are the dominant ones; they are computed first considering the CC terms with the approximations $e^{T^\dagger}Oe^T \simeq O + OT + T^\dagger O + \frac{1}{2}OT^2 + \frac{1}{2}T^{\dagger 2}O + T^\dagger OT$ and $e^{T^\dagger}e^T \simeq T^\dagger T + \frac{1}{2}T^\dagger T^2 + \frac{1}{2}T^{\dagger 2}T$. Then, they are stored and contracted with the T_2 and $T_2^\dagger$ operators, avoiding repetitions of the diagrams in a self-consistent procedure to account for the higher-order one-body terms from the nonterminating series. They are again stored as an intermediate step for further contraction with the S_n and $S_n^\dagger$ operators. Similarly, the effective two-body and three-body terms are computed after contracting with the above effective one-body terms with the T_2 and $T_2^\dagger$ operators, but they are computed directly contracting with the S_n and $S_n^\dagger$ operators. Obviously, these effective two-body and three-body terms also have contributions from the nonterminating series in this way. To see the convergence of the results with the series expansion, we present contributions containing k T and/or $T^\dagger$ operators from these nonterminating series, which we refer to as the CCSD$^{(k)}$ method considering terms up to $k \to \infty$ in a self-consistent procedure as described above. Our final CCSD results correspond to the CCSD$^{(\infty)}$ method. The same procedure is also adopted for the CCSD(T) method. The contribution from the normalizations of the wave functions $\mathcal{C}_{\text{norm}}$ is estimated explicitly using the expression

$$\mathcal{C}_{\text{norm}} = \left[\frac{\langle\Psi_f|O|\Psi_i\rangle}{\sqrt{\langle\Psi_f|\Psi_f\rangle\langle\Psi_i|\Psi_i\rangle}} - \langle\Psi_f|O|\Psi_i\rangle\right]. \tag{16}$$

IV. RESULTS AND DISCUSSIONS

In Table I, we give the radiative transition matrix elements for all the considered channels for the DHF, MBPT(2), LCCSD, CCSD, and CCSD(T) methods to analyze the propagation of the correlation effects through various levels of approximation in the many-body theories and the experimental values of the transition wavelengths from the NIST database [27] that we use later. We also give contributions from the CCSD method by truncating the nonlinear terms with $k = 2$, $k = 4$ and from a self-consistent ($k = \infty$) calculation. For $k = 2$, the property-evaluating expression given by Eq. (15) has the same number of terms as the LCCSD method. Therefore, differences in the results from the LCCSD and CCSD$^{(2)}$ methods imply the correlation contributions arise through the nonlinear terms in the wave-function-determining equations of the CCSD method and are found to be quite large. Often, these contributions are neglected in the calculations prohibitively large computational resources are required to evaluate them. We observe from the trends that the correlation effects for the MBPT(2) method are large and that there are strong cancellations in the LCCSD approximation and the results almost converge for $k = 4$ when the nonlinear terms are included in the CCSD method. The discrepancy between the results of the CCSD$^{(4)}$ and CCSD$^{(\infty)}$ methods is beyond the second significant digit, implying that the results have converged within the precision of interest. The valence triple excitations seem to change the results slightly. We also give uncertainties associated with these results by estimating contributions due to the finite size of our basis set and neglected contributions from the Breit interaction and corrections from

TABLE II. Contributions to the reduced matrix elements $\langle J_i||O||J_f\rangle$ from various terms of the CCSD method (in a.u.). Differences between these values from the CCSD results quoted in Table I correspond to those nonlinear terms that are not mentioned explicitly here.

$J_i \rightarrow J_f$	O	OT_1	$T_1^\dagger O$	OS_{1f}	$S_{1i}^\dagger O$	$S_{1i}^\dagger OS_{1f}$	OS_{2f}	$S_{2i}^\dagger O$	$S_{1i}^\dagger OS_{2f}$	$S_{2i}^\dagger OS_{1f}$	$\mathcal{C}_{\rm norm}$
					$E1$ matrix elements						
$6d\,^2D_{3/2} \rightarrow 7p\,^2P_{1/2}$	9.22	~ 0.0	0.019	-0.437	-0.877	0.246	-0.208	-0.248	-0.020	-0.020	-0.324
$6d\,^2D_{3/2} \rightarrow 7p\,^2P_{3/2}$	4.28	$\sim$0.0	0.002	-0.132	-0.522	0.089	-0.081	-0.097	-0.005	-0.007	-0.139
$6d\,^2D_{5/2} \rightarrow 7p\,^2P_{3/2}$	12.80	$\sim$0.0	0.006	-0.381	-1.372	0.238	-0.233	-0.292	-0.008	-0.020	-0.342
					$E2$ matrix elements						
$6d\,^2D_{3/2} \rightarrow 7s\,^2S_{1/2}$	43.10	$\sim$0.0	0.013	-6.306	-2.683	1.547	-0.139	-0.111	-0.027	-0.186	-1.618
$6d\,^2D_{5/2} \rightarrow 7s\,^2S_{1/2}$	52.74	$\sim$0.0	0.015	-7.644	-2.871	1.666	-0.125	-0.148	-0.017	0.004	-1.6
$6d\,^2D_{5/2} \rightarrow 6d\,^2d_{3/2}$	47.70	$\sim$0.0	$\sim$0.0	-8.994	-8.116	3.264	-0.079	-0.010	-0.012	-0.021	-1.947
					$M1$ matrix elements						
$6d\,^2D_{3/2} \rightarrow 7s\,^2S_{1/2}$	$\sim$0.0	$\sim$0.0	$\sim$0.0	$\sim$0.0	$\sim$0.0	0.0004	0.0005	0.0001	$\sim$0.0	$\sim$0.0	$\sim$0.0
$6d\,^2D_{5/2} \rightarrow 6d\,^2d_{3/2}$	1.549	$\sim$0.0	$\sim$0.0	0.003	-0.002	0.078	0.003	$\sim$0.0	$\sim$0.0	$\sim$0.0	-0.093

the quantum electrodynamics (QED) effects. These estimates are carried out using the MBPT(2) method, which gives the largest correlation effects. The $E1$ and $E2$ matrix elements reported in [9] using the correlation potential method to include the core-valence correlations and Brueckner orbitals for the valence electrons are also given in Table I. The CCSD results for the $E1$ matrix elements are found to be larger than these values owing to significant contributions from the nonlinear terms in the wave function.

After analyzing the trends in the correlation effects at different levels of approximation, we now focus on the contributions from the different terms of the CCSD method. We present these results in Table II along with the contributions from $\mathcal{C}_{\rm norm}$. Contributions from the corresponding radiative operator O are the DHF results, OT_1 and its complex-conjugate terms give core-valence correlations, OS_{1f} and $S_{1i}^\dagger O$ give the pair-correlation effects involving the valence orbitals, OS_{2f} and $S_{2i}^\dagger O$ give the core-polarization correlation effects involving the valence orbitals, etc. [29]. Contributions from the other nonlinear terms such as those representing the core pair-correlation effects coming from the $T_2^\dagger OT_2$ term are not given explicitly in Table II; however, their contributions can be obtained by taking the differences of the contributions given in Table II and the final CCSD results given in Table I. As can be seen from Table II, core-valence correlations are small, and the largest correlation effects come from the pair-correlation effects in the $E1$ and $E2$ matrix-element calculations. Nevertheless, the core-polarization effects are also very significant and are the dominant ones in the calculations of the $M1$ matrix elements. We also find contributions from $\mathcal{C}_{\rm norm}$ to be fairly large.

We now use the matrix elements from the CCSD(T) method and the experimental wavelengths mentioned in Table I to evaluate the transition probabilities due to different radiative decay channels from the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states. These values are given in Table III along with their branching ratios for the individual transitions. As can be seen from Table III, the branching ratios are dominated by the $E1$ transitions, and they are entirely responsible for determining the lifetimes of the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states. Using these transition probabilities, we estimate the lifetimes of the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states as $\tau_{6d3/2} = 540(10)$ ns and $\tau_{6d5/2} = 1704(32)$ ns, respectively. These values are very large compared to the other low-lying excited $7p\,^2P_{1/2,3/2}$, $8s\,^2S_{1/2}$, $7p\,^2P_{1/2,3/2}$, $7d\,^2D_{3/2,5/2}$, and $9s\,^2S_{1/2}$ states of Fr, which are measured to date of writing as $\tau_{7p1/2} = 29.45(11)$ ns, $\tau_{7p3/2} = 21.02(15)$ ns, $\tau_{8s1/2} = 53.30(44)$ ns, $\tau_{7d3/2} = 73.6(3)$ ns, $\tau_{7d5/2} = 67.7(2.9)$ ns, $\tau_{8p1/2} = 149.3(3.5)$ ns, $\tau_{8p3/2} = 83.5(1.5)$ ns, and $\tau_{9s1/2} = 107.53(90)$ ns, respectively [24]. The large PNC amplitudes in the $7s\,^2S_{1/2} \rightarrow 6d\,^2D_{3/2,5/2}$ transitions [8,9] compared to the $7s\,^2S_{1/2} \rightarrow 8s\,^2S_{1/2}$ transition in Fr [10] and the corresponding transitions in the Ra$^+$ ion [12,17] and the long lifetimes of the excited $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states make Fr a potentially attractive candidate for a PNC experiment.

TABLE III. Transition probabilities A_{if}^O due to different transition decay channels O (in s^{-1}) and their branching ratios from the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states of Fr. Uncertainties are given in the parentheses.

$J_i \rightarrow J_f$	O	A_{if}^O	Branching ratio
$6d\,^2D_{3/2} \rightarrow 7p\,^2P_{1/2}$	$E1$	1779511(33688)	0.96
$6d\,^2D_{3/2} \rightarrow 7p\,^2P_{3/2}$	$E1$	72637(1280)	0.04
$6d\,^2D_{3/2} \rightarrow 7s\,^2S_{1/2}$	$E2$	35.96(60)	$\sim$0.0
$6d\,^2D_{3/2} \rightarrow 7s\,^2S_{1/2}$	$M1$	$\sim$0.0	$\sim$0.0
$6d\,^2D_{5/2} \rightarrow 7p\,^2P_{3/2}$	$E1$	586776(11219)	$\sim$ 1.0
$6d\,^2D_{5/2} \rightarrow 7s\,^2S_{1/2}$	$E2$	39.33(19)	$\sim$0.0
$6d\,^2D_{5/2} \rightarrow 6d\,^2d_{3/2}$	$E2$	$\sim$0.0	$\sim$0.0
$6d\,^2D_{5/2} \rightarrow 6d\,^2d_{3/2}$	$M1$	$\sim$0.0	$\sim$0.0

V. CONCLUSION

We have performed relativistic many-body calculations of the lifetimes of the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states of Fr, and they are found to be large. These results favor the measurements of PNC in the $7s\,^2S_{1/2} \rightarrow 6d\,^2D_{3/2}$ and $7s\,^2S_{1/2} \rightarrow 6d\,^2D_{5/2}$ transitions of Fr, for which calculations predict large PNC effects. Also, both the $6d\,^2D_{3/2}$ and $6d\,^2D_{5/2}$ states seem to be very suitable as intermediate states in the RIS measurements of spectroscopic properties of some of the higher excited states of Fr because of their long lifetimes. To evaluate of these lifetimes, we have calculated radiative transition matrix elements using the relativistic CC method. We have also investigated the role of the electron correlation effects in the determination of these quantities systematically by approximating many-body methods at different levels and have given contributions explicitly from various CCSD terms.

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