PHYSICAL REVIEW C 73, 014612 (2006)

α decay half-lives of new superheavy elements

P. Roy Chowdhury,[1,*] C. Samanta,[1,2] and D. N. Basu[3]
[1]Saha Institute of Nuclear Physics, 1/AF Bidhan Nagar, Kolkata 700 064, India
[2]Physics Department, Virginia Commonwealth University, Richmond, Virginia 23284-2000, USA
[3]Variable Energy Cyclotron Centre, 1/AF Bidhan Nagar, Kolkata 700 064, India
(Received 17 October 2005; published 26 January 2006)

The lifetimes of α decays of the recently produced isotopes of elements 112, 114, 116, and ^{294}118 and of some decay products are calculated theoretically within the WKB approximation by use of microscopic α-nucleus interaction potentials. We obtain these nuclear potentials by folding the densities of the α and the daughter nuclei with the M3Y effective interaction, supplemented by a zero-range pseudopotential for exchange along with the density dependence. Spherical charge distributions are used for calculating the Coulomb interaction potentials. These calculations provide reasonable estimates for the observed α decay lifetimes and thus provide reliable predictions for other superheavies.

DOI: 10.1103/PhysRevC.73.014612 PACS number(s): 27.90.+b, 23.60.+e, 21.30.Fe, 25.55.Ci

I. INTRODUCTION

The main features that determine the fusion process for the production of superheavy elements (SHEs) are the fusion barrier and related beam energy and excitation energy, the ratio of surface energy versus Coulomb repulsion that determines the fusion probability and that strongly depends on the degree of asymmetry or the reactants (the product Z_1Z_2 at fixed $Z_1 + Z_2$), the impact parameter and related angular momentum, and the ratio of neutron evaporation versus fission probability of the compound nucleus. In fusion of heavy elements, the product Z_1Z_2 reaches extremely large values and the fission barrier extremely small values. In addition, the fission barrier is fragile at increasing excitation energy and angular momentum, because it is built up solely from shell effects. For these reasons the fusion of heavy elements is hampered, whereas the fusion of lighter elements is advanced through the contracting effect of surface tension. Recently isotopes of elements 112, 114, 116, and ^{294}118 have been produced in the fusion-evaporation reactions, keeping low excitation energies by irradiations of ^{233,238}U, ^{242}Pu, ^{248}Cm [1], and ^{249}Cf targets [2] with a ^{48}Ca beam at various energies. The observed decays reveal that the dominant decay mode is the α emission. The α decay energies and half-lives of 14 new α decaying nuclei have been measured. Incidentally, questions have been raised [3] about some of the SHE findings [4]. In fact, in similar sophisticated experiments at other places [5,6], the α cascades were not observed. While one awaits for further experimental verification of such an important discovery, theoretical predictions already exist for such SHEs [7] along with their α decay lifetime predictions [8].

In this work, the half-lives of new SHEs have been determined with microscopic potentials and compared with the existing theoretical and experimental results to test the extent of validity of this formalism. In view of the excellent agreement of this work with the available experimental data, half-lives of ~80 new SHEs have been predicted. In this framework, the nuclear potentials have been obtained by double folding of the α and daughter nuclei density distributions with a density-dependent effective interaction. This nuclear interaction energy for the α-nucleus interaction has therefore been obtained microscopically. A double-folding potential obtained with the M3Y [9] effective interaction supplemented by a zero-range potential-for the single-nucleon exchange is more appropriate because of its microscopic nature [10]. A potential-energy surface is inherently embedded in this description. The semirealistic explicit density dependence [11] into the M3Y effective interaction has been employed to incorporate the higher-order exchange and Pauli blocking effects. The penetrability of the prescission part of the potential barrier provides the α cluster preformation probability [12]. Theoretical calculations in terms of quantum-mechanical barrier penetrability by use of microscopically obtained nuclear potentials are provided in this work. Observed lifetimes of the 14 α decays originating from the isotopes of the synthesized new elements 112, 114, and 116 are in reasonable agreement with the theoretical estimates. Recent theoretical predictions [13] for the lifetimes of the α decay chains of SHE 115 also agree with the present calculations [14], which provided consistent estimates for the observed lifetimes [15] of the consecutive α decay chains of SHE 115.

Based on the present calculations that provide reasonable estimates for the observed α decay lifetimes of many newly synthesized elements and therefore are expected to be effective predictors of the half-lives in the region of the heaviest elements, values from years to microseconds have been calculated for various isotopes. This wide range of half-lives encourages the application of a wide variety of experimental methods in the investigations of SHEs from the investigation of chemical properties of SHEs by use of long-lived isotopes, to the atomic physics experiments on trapped ions, and to the safe identification of short-lived isotopes by recoil separation techniques.

II. THE DENSITY-DEPENDENT EFFECTIVE INTERACTION

The M3Y interaction was derived by the fitting of its matrix elements in an oscillator basis to those elements of the G matrix

*Electronic address: partha.roychowdhury@saha.ac.in

[16] obtained with the Reid-Elliott soft-core nucleon-nucleon (*NN*) interaction. The ranges of the M3Y forces were chosen to ensure a long-range tail of the one-pion exchange potential as well as a short-range repulsive part simulating the exchange of heavier mesons. The zero-range potential represents the single-nucleon exchange term while the density dependence accounts for the higher-order exchange effects and the Pauli blocking effects. The general expression for the density-dependent M3Y (DDM3Y) effective interaction supplemented by a zero-range potential for the single-nucleon exchange is given by

$$v(s,\rho,E)=t^{\mathrm{M3Y}}(s,E)g(\rho,E)=Ct^{\mathrm{M3Y}}[1-\beta(E)\rho^{2/3}], \quad (1)$$

where ρ is the nucleonic density and the M3Y effective interaction potential supplemented by a zero-range potential t^{M3Y} is given by Ref. [11]

$$t^{\mathrm{M3Y}}=7999\frac{e^{-4s}}{4s}-2134\frac{e^{-2.5s}}{2.5s}+J_{00}(E)\delta(s), \quad (2)$$

where the zero-range potential $J_{00}(E)$ representing the single-nucleon exchange is given by

$$J_{00}(E)=-276(1-0.005E/A)\ (\mathrm{MeV\ fm}^3). \quad (3)$$

This DDM3Y effective *NN* interaction supplemented by the zero-range potential is used to determine the nuclear matter equation of state. The equilibrium density of the nuclear matter is determined by minimization of the energy per nucleon. The density-dependence parameters are fixed by reproducing the saturation energy per nucleon and the saturation density of spin and isospin symmetric cold infinite nuclear matter. Although the density-dependence parameters for single folding can be determined from the nuclear matter calculations and used successfully for proton radioactivity and scattering [17], the transition to double folding is not straightforward. The parameter β can be related to the mean free path in nuclear medium; hence its value should remain the same, $\sim$1.6 fm^2, as that obtained from nuclear matter calculations [18], while the other constant C, which is basically an overall normalization constant, may change. The value of this overall normalization constant is kept equal to unity, which has been found $\sim$1 [19] from an optimum fit to a large number of α decay lifetimes. That the density dependence of the effective projectile-nucleon interaction has been found to be fairly independent of the projectile, as long as the projectile-nucleus interaction is amenable to a single-folding prescription, implies that, in a double-folding model, the density-dependent effects on the nucleon-nucleon interaction can be factorized into a target term times a projectile term [20]. The general expression for the DDM3Y realistic effective *NN* interaction to be used to obtain the oft-quoted double-folding nucleus-nucleus interaction potential is given by

$$v(s,\rho_1,\rho_2,E)=t^{\mathrm{M3Y}}(s,E)g(\rho_1,\rho_2,E), \quad (4)$$

where the density-dependence term $g(\rho_1,\rho_2,E)$ has now been factorized into a target term times a projectile term [20] as

$$g(\rho_1,\rho_2,E)=C\big[1-\beta(E)\rho_1^{2/3}\big]\big[1-\beta(E)\rho_2^{2/3}\big]. \quad (5)$$

The folding model potentials thus obtained by double folding the density distributions ρ_1 of the α and ρ_2 of the daughter nuclei with such a factorized DDM3Y–Reid-Elliott effective interaction, along with a zero-range potential representing the potential arising because of the single-nucleon exchange, have been used successfully to estimate the half-lives of the α radioactivity lifetimes of the newly synthesized elements and their isotopes.

III. THE DOUBLE-FOLDED NUCLEAR POTENTIALS AND THE HALF-LIVES OF α RADIOACTIVITY

The double-folded nuclear interaction potential between the daughter nucleus and the emitted particle is given by Ref. [16]

$$V_N(R)=\int\int\rho_1(\vec{r_1})\rho_2(\vec{r_2})v[|\vec{r_2}-\vec{r_1}+\vec{R}|]d^3r_1d^3r_2, \quad (6)$$

where ρ_1 and ρ_2 are the density distribution functions for the two composite nuclear fragments. The density distribution function in the case of an α particle has the Gaussian form

$$\rho(r)=0.4229\exp(-0.7024r^2), \quad (7)$$

whose volume integral is equal to $A_\alpha(=4)$, the mass number of the α particle. Since the experimental charge density distributions in the case of heavier nuclei can be well described by the two-parameter Fermi function [21] and since the charge that means that the proton (p) and the neutron (n) density distributions should have similar forms because of the same strengths of the *n-n* and *p-p* nuclear forces, the matter density distribution for the daughter nucleus can be described by the spherically symmetric Fermi function,

$$\rho(r)=\rho_0/\{[1+\exp[(r-c)/a]\}, \quad (8)$$

where the equivalent sharp radius r_ρ, the half-density radius c, and the diffuseness for the leptodermous Fermi density distributions are given by [20,22]

$$c=r_\rho\big(1-\pi^2a^2\big/3r_\rho^2\big),\quad r_\rho=1.13A_d^{1/3},\quad a=0.54\,\mathrm{fm}, \quad (9)$$

and the value of the central density ρ_0 is fixed by equating the volume integral of the density distribution function to the mass number A_d of the residual daughter nucleus.

The distance s between any two nucleons, one belonging to the residual daughter nucleus and the other belonging to the emitted α, is given by $s=|\vec{r_2}-\vec{r_1}+\vec{R}|$ while the interaction potential between these two nucleons $v(s)$ appearing in Eq. (6) is given by the factorized DDM3Y effective interaction described by Eqs. (4) and (5). The total interaction energy $E(R)$ between the α and the residual daughter nucleus is equal to the sum of the nuclear interaction energy, the Coulomb interaction energy, and the centrifugal barrier. Thus

$$E(R)=V_N(R)+V_C(R)+\hbar^2l(l+1)/(2\mu R^2), \quad (10)$$

where $\mu=M_eM_d/M$ is the reduced mass, M_e, M_d, and M are the masses of the emitted particle, the daughter nucleus, and the parent nucleus, respectively, all measured in the units of MeV/c^2. Assuming a spherical charge distribution for the residual daughter nucleus and the emitted nucleus as a point particle, the Coulomb interaction potential $V_C(R)$ between them is given by

$$V_C(R) = \left(\frac{Z_e Z_d e^2}{2R_c}\right)\left[3 - \left(\frac{R}{R_c}\right)^2\right] \quad \text{for} \quad R \leqslant R_c,$$
$$= \frac{Z_e Z_d e^2}{R}, \quad \text{otherwise}, \tag{11}$$

where Z_e and Z_d are the atomic numbers of the emitted cluster and the daughter nucleus, respectively. The touching radial separation R_c between the emitted cluster and the daughter nucleus is given by $R_c = c_e + c_d$, where c_e and c_d are obtained by use of Eqs. (9). The energetics allow spontaneous emission of a particle only if the released energy

$$Q = [M - (M_e + M_d)]c^2 \tag{12}$$

is a positive quantity.

The half-life of a parent nucleus decaying by means of α emission is calculated with the WKB barrier penetration probability. The assault frequency ν is obtained from the zero-point vibration energy $E_v = (1/2)h\nu$. The decay half-life T of the parent nucleus (A, Z) into a α and a daughter (A_d, Z_d) is given by

$$T = [(h \ln 2)/(2E_v)][1 + \exp(K)]. \tag{13}$$

The action integral K within the WKB approximation is given by

$$K = (2/\hbar)\int_{R_a}^{R_b} [2\mu(E(R) - E_v - Q)]^{1/2} dR, \tag{14}$$

where R_a and R_b are the two turning points of the WKB action integral determined from the equations

$$E(R_a) = Q + E_v = E(R_b) \tag{15}$$

whose solutions provide three turning points. The α particle oscillates between the first and the second turning points and tunnels through the barrier at R_a and R_b representing the second and the third turning points, respectively. Since the released energy Q enters in the action integral that goes to the exponential function in Eq. (13) and the zero-point vibration energy E_v is proportional to Q, the calculations for the lifetimes become very sensitive to the released energies involved in the decay processes.

IV. CALCULATIONS AND RESULTS

We obtained the two turning points of the action integral given by Eq. (14) by solving Eqs. (15) using the microscopic double-folding potential given by Eq. (6) along with the Coulomb potential given by Eq. (11) and the centrifugal barrier. Then we evaluated the WKB action integral between these two turning points numerically by using Eqs. (6), (10), and (11). The zero-point vibration energies used in the present calculations are the same as those described in Ref. [23] immediately after Eq. (4), and experimental Q values are used. Moreover, the shell effects are implicitly contained in the zero-point vibration energy because of its proportionality with the Q value, which is maximum when the daughter nucleus has a magic number of neutrons and protons. Values of the proportionality constants of E_v with Q are the largest for the even-even parent and the smallest for the odd-odd one. Other conditions remaining same, one may observe that, with a greater value of E_v, the lifetime is shortened, indicating a higher emission rate. Finally, we calculated the half-lives by using Eq. (13) and tabulated in Tables I and II.

The value of the normalization constant C used in the calculations is kept fixed and equal to unity. All the calculations

TABLE I. Comparison between experimental and calculated α decay half-lives for zero angular-momentum transfers, by use of spherical charge distributions for the Coulomb interaction and the DDM3Y effective interaction. The lower and upper limits of the theoretical half-lives corresponding to the upper and lower limits of the experimental Q values are also provided. Present theoretical predictions are compared with those of the GLDM [7,8] and with VSS [24] predictions.

Parent	Nuclei	Expt.	Assault frequency (this work)	Theory Ref. [25]	Expt.	DDM3Y (this work)	GLDM	VSS
Z	A	Q(MeV)	10^{20} s^{-1}	Q(MeV)	$T_{1/2}$	$T_{1/2}$	$T_{1/2}$	$T_{1/2}$
118	294	11.81 ± 0.06	5.968	12.51	$1.8^{+75}_{-1.3}$ ms	$0.66^{+0.23}_{-0.18}$ ms	0.01 ms [8]	$0.64^{+0.24}_{-0.18}$ ms
116	293	10.67 ± 0.06	4.680	11.15	53^{+62}_{-19} ms	206^{+90}_{-61} ms	18.2 ms [8]	1258^{+557}_{-384} ms
116	292	10.80 ± 0.07	5.458	11.03	18^{+16}_{-6} ms	39^{+20}_{-13} ms	6.9 ms [8]	49^{+26}_{-16} ms
116	291	10.89 ± 0.07	4.777	11.33	$6.3^{+11.6}_{-2.5}$ ms	$60.4^{+30.2}_{-20.1}$ ms	7.2 ms [8]	$336.4^{+173.1}_{-113.4}$ ms
116	290	11.00 ± 0.08	5.559	11.34	15^{+26}_{-6} ms	$13.4^{+7.7}_{-5.2}$ ms	1.3 ms [8]	$15.2^{+9.0}_{-5.6}$ ms
114	289	9.96 ± 0.06	4.369	9.08	$2.7^{+1.4}_{-0.7}$ s	$3.8^{+1.8}_{-1.2}$ s	51.5 min [8]	$26.7^{+13.1}_{-8.7}$ s
114	288	10.09 ± 0.07	5.099	9.39	$0.8^{+0.32}_{-0.18}$ s	$0.67^{+0.37}_{-0.27}$ s	63 s [8]	$0.98^{+0.56}_{-0.40}$ s
114	287	10.16 ± 0.06	4.456	9.53	$0.51^{+0.18}_{-0.10}$ s	$1.13^{+0.52}_{-0.40}$ s	2.1 min [8]	$7.24^{+3.43}_{-2.61}$ s
114	286	10.35 ± 0.06	5.230	9.61	$0.16^{+0.07}_{-0.03}$ s	$0.14^{+0.06}_{-0.04}$ s	14.5 s [8]	$0.19^{+0.08}_{-0.06}$ s
112	285	9.29 ± 0.06	4.075	8.80	34^{+17}_{-9} s	75^{+41}_{-26} s	83.5 min [8]	592^{+323}_{-207} s
112	283	9.67 ± 0.06	4.241	9.22	$4.0^{+1.3}_{-0.7}$ s	$5.9^{+2.9}_{-2.0}$ s	3.8 min [8]	$41.3^{+20.9}_{-13.8}$ s
110	279	9.84 ± 0.06	4.316	9.89	$0.18^{+0.05}_{-0.03}$ s	$0.40^{+0.18}_{-0.13}$ s	0.03s[7]	$2.92^{+1.4}_{-0.94}$ s
108	275	9.44 ± 0.07	4.141	9.58	$0.15^{+0.27}_{-0.06}$ s	$1.09^{+0.73}_{-0.40}$ s	0.05s[7]	$8.98^{+5.49}_{-3.38}$ s
106	271	8.65 ± 0.08	3.794	8.59	$2.4^{+4.3}_{-1.0}$ min	$1.0^{+0.8}_{-0.5}$ min	14.8s[7]	$8.6^{+7.3}_{-3.9}$ min

TABLE II. Comparison between different theoretically predicted α decay half-lives for zero angular-momentum transfers by use of theoretical Q values from the macroscopic-microscopic (M-M) model. Present calculations using spherical charge distributions for the Coulomb interaction and microscopic nuclear potentials from double-folding nuclear densities with DDM3Y effective interaction are compared with the VSS [24] predictions and with the Viola-Seaborg estimates used in Ref. [26].

Parent Nuclei		VSS Ref. [24]	DDM3Y (this work)	Viola-Seaborg Ref. [26]	M-M model Ref. [26]	Parent Nuclei		VSS Ref. [24]	DDM3Y (this work)	Viola-Seaborg Ref. [26]	M-M model Ref. [26]
Z	A	$\log_{10} T$ (s)	$\log_{10} T$ (s)	$\log_{10} T$ (s)	Q (MeV)	Z	A	$\log_{10} T$ (s)	$\log_{10} T$ (s)	$\log_{10} T$ (s)	Q (MeV)
104	274	9.21	8.75	9.35	6.56	104	276	12.02	11.55	12.18	6.02
104	278	14.80	14.31	15.00	5.55	104	280	17.32	16.80	17.56	5.17
104	282	17.88	17.34	18.13	5.09	104	284	21.42	20.87	21.74	4.63
104	286	23.21	22.65	23.57	4.42	104	288	24.94	24.36	25.28	4.23
104	290	14.67	14.01	14.88	5.57	104	292	17.95	17.28	18.25	5.08
106	278	7.92	7.49	8.03	7.02	106	280	10.50	10.03	10.62	6.48
106	282	12.58	12.09	12.74	6.09	106	284	12.75	12.23	12.94	6.06
106	286	15.61	15.06	15.85	5.58	106	288	17.26	16.70	17.53	5.33
106	290	18.45	17.87	18.71	5.16	106	292	10.60	9.98	10.77	6.46
106	294	12.86	12.21	13.03	6.04						
108	282	7.13	6.72	7.17	7.39	108	284	8.63	8.18	8.73	7.05
108	286	8.59	8.11	8.69	7.06	108	288	11.25	10.74	11.40	6.51
108	290	12.74	12.20	12.92	6.23	108	292	13.58	13.03	13.73	6.08
108	294	7.35	6.78	7.43	7.34	108	296	8.91	8.30	9.02	6.99
110	286	5.38	5.00	5.40	8.02	110	288	5.38	4.98	5.37	8.02
110	290	8.08	7.64	8.11	7.36	110	292	9.15	8.67	9.23	7.12
110	294	9.67	9.15	9.73	7.01	110	296	4.96	4.47	4.96	8.13
110	298	6.08	5.54	6.12	7.84						
112	288	2.44	2.14	2.35	9.06	112	290	3.07	2.75	2.98	8.87
112	292	5.57	5.20	5.56	8.17	112	294	6.03	5.63	6.02	8.05
112	296	6.27	5.83	6.26	7.99	112	298	2.77	2.34	2.70	8.96
112	300	3.65	3.19	3.59	8.70						
114	290	.02	−.17	−.16	10.08	114	292	1.52	1.28	1.38	9.57
114	294	2.84	2.55	2.73	9.15	114	296	2.91	2.59	2.77	9.13
114	298	2.98	2.63	2.84	9.11	114	300	.45	.12	.28	9.93
114	302	1.03	.67	.87	9.73						
116	284	−6.19	−6.04	−6.57	12.96	116	286	−4.92	−4.84	−5.26	12.34
116	288	−3.18	−3.18	−3.48	11.56	116	290	−2.24	−2.30	−2.51	11.17
116	292	−1.99	−2.09	−2.26	11.07	116	294	−1.15	−1.28	−1.40	10.74
116	296	−.99	−1.15	−1.25	10.68	116	298	−.99	−1.18	−1.24	10.68
116	300	−1.02	−1.23	−1.26	10.69	116	302	−2.68	−2.87	−2.96	11.35
116	304	−2.24	−2.47	−2.52	11.17						
118	288	−5.97	−5.79	−6.39	13.11	118	290	−4.64	−4.53	−5.02	12.46
118	292	−4.23	−4.15	−4.61	12.27	118	294	−4.05	−4.00	−4.42	12.19
118	296	−3.79	−3.77	−4.15	12.07	118	298	−3.54	−3.56	−3.90	11.96
118	300	−3.56	−3.61	−3.91	11.97	118	302	−3.61	−3.68	−3.98	11.99
118	304	−4.77	−4.82	−5.15	12.52						
120	292	−6.40	−6.14	−6.88	13.59	120	294	−6.07	−5.85	−6.55	13.42
120	296	−6.03	−5.84	−6.51	13.40	120	298	−5.95	−5.79	−6.43	13.36
120	300	−5.42	−5.31	−5.87	13.09	120	302	−5.38	−5.29	−5.83	13.07
120	304	−5.48	−5.41	−5.93	13.12	120	306	−6.28	−6.21	−6.76	13.53

Note: All the nuclei listed above are either spherical or have very small deformations [26].

are performed with zero angular-momentum transfer. The experimentally measured values for the released energy Q are used in the calculations. In general, the E and A appearing in Eq. (3) are the laboratory energy of the projectile in mega-electron-vlots and the projectile mass number, respectively. However, for a decay process, E/A can be shown to be equal to the energy measured in mega-electron-volts in the center of mass of the emitted particle-daughter nucleus system/(μ/m), where m is the nucleonic mass in MeV/c^2 and for the decay process the energy measured in the center of mass is equal to the released energy Q in mega-electron-volts. Since the released energies involved in the α decay processes are very small compared with the energies involved in high-energy α scattering, the zero-range potential $J_{00}(E)$ is also practically independent of energy for the α decay processes and can be taken as $-276\,\mathrm{MeV\,fm^3}$.

The results of the present calculations with DDM3Y for the lifetimes of α decays of recently produced isotopes of new elements 112, 114, 116, and $^{294}118$ and of some decay products are presented in Table I. The quantitative agreement with experimental data is reasonable. The result for $^{294}118$ is almost underestimated, possibly because the centrifugal barrier required for the spin-parity conservation could not be taken into account because of nonavailability of the spin-parities of the decay chain nuclei. The term $\hbar^2 l(l+1)/(2\mu R^2)$ in Eq. (10) represents the additional centrifugal contribution to the barrier that acts to reduce the tunneling probability if the angular momentum carried by the α particle is nonzero. The hindrance factor, which is defined as the ratio of the experimental $T_{1/2}$ to the theoretical $T_{1/2}$, is therefore larger than unity since the decay involving a change in angular momentum can be strongly hindered by the centrifugal barrier. However, as one can see in Table I, the theoretical Viola-Seaborg systematics (VSS) with Sobiczewski constants [24] largely overestimate the half-lives for as many as eight cases, showing inconsistencies, while the present calculations slightly overestimate for only three cases but still provide much better estimates than those made by the VSS. For the rest of the cases, the experimental uncertainties in the Q values associated with the α decays can almost account for the overestimations of theoretical lifetimes if the upper limits for the experimental Q values instead of the mean values are used for the calculations. Very recent theoretical predictions of the generalized liquid-drop model (GLDM) [7,8] for these decay lifetimes are listed in Table I, and the disagreements of the results with the experimentally observed half-lives are primarily due to use of theoretical Q values that do differ from the experimental ones.

The theoretical Q values calulated with 28 mass excesses from the latest mass table [25] are listed in Table I for comparison with the experimental ones. It is very obvious from the table that the results for the half-lives are quite sensitive to the uncertainties involved in the experimental Q values used in the present calculations. The theoretical Q values differ substantially from the experimental ones for higher Z, A nuclei, and they are therefore not used for the calculating the lifetimes. Although the recent theoretical mass table [25] used for calculating the theoretical Q values provides excellent estimates for normal nuclei, better mass predictions for superheavies are needed for the successful predictions of possible decay modes and their lifetimes.

In Table II we provide predictions for the α decay lifetimes for a large number of SHEs [26], though there exist many more [27], which are expected to live long enough to be detected after the synthesis in the current day experimental setup. The theoretical Q values are calculated based on the macroscopic-microscopic model [26]. The lifetime values from years to microseconds are calculated for various isotopes. It is easy to observe that the predictions for the half-lives by the present calculations are lower than those made by VSS and by the VSS of Ref. [26].

V. SUMMARY AND CONCLUSION

The half-lives for α radioactivity were analyzed with microscopic nuclear potentials we obtained by the double-folding procedure by using the DDM3Y effective interaction. This procedure of obtaining nuclear interaction potentials has a profound theoretical basis. The results of the present calculations made with DDM3Y are in good agreement with the published experimental data for the half-lives of the α decays from the isotopes of elements 112, 114, 116, and $^{294}118$ and from some decay products. As some of these experimental data await further experimental verification, these theoretical predictions are expected to provide useful guideline. Lifetime estimates from present calculations are lower than those of VSS. The released energies Q, to which the calculations are quite sensitive, when calculated from the microscopic-macroscopic model masses [25] do not provide excellent agreement with those observed for superheavies. Nevertheless, the positive decay Q values [25] support these α decay modes. Present calculations demonstrate its success of providing reasonable estimates for the lifetimes of nuclear decays by α emissions for the domain of superheavy nuclei.

[1] Yu. Ts. Oganessian *et al.*, Phys. Rev. C **70**, 064609 (2004); **71**, 029902(E) (2005).
[2] Yu. Ts. Oganessian *et al.*, JINR Communication (2002) D7-2002-287; Lawrence Livermore National Laboratory Rep. UCRL-ID-151619 (2003).
[3] P. Armbruster, Eur. Phys. J. A **7**, 23 (2000).
[4] Yu. Ts. Oganessian *et al.*, Phys. Rev. Lett. **83**, 3154 (1999); Eur. Phys. J. A **5**, 63 (1999); Nature (London) **400**, 242 (1999).
[5] W. Loveland *et al.*, Phys. Rev. C **66**, 044617 (2002).
[6] K. E. Gregorich *et al.*, Phys. Rev. C **72**, 014605 (2005).
[7] G. Royer and R. A. Gherghescu, Nucl. Phys. **A699**, 479 (2002).
[8] G. Royer, K. Zbiri, and C. Bonilla, Nucl. Phys. **A730**, 355 (2004).
[9] G. Bertsch, J. Borysowicz, H. McManus, and W. G. Love, Nucl. Phys. **A284**, 399 (1977); H. J. Gils, Nucl. Phys. **A473**, 111 (1987).
[10] C. Samanta, Y. Sakuragi, M. Ito, and M. Fujiwara, J. Phys. G **23**, 1697 (1997).
[11] A. M. Kobos, B. A. Brown, R. Lindsay, and G. R. Satchler, Nucl. Phys. **A425**, 205 (1984).
[12] D. N. Poenaru and W. Greiner, Phys. Scr. **44**, 427 (1991).
[13] H. Zhang, J. Li, W. Zuo, Z. Ma, B. Chen, and S. Im, Phys. Rev. C **71**, 054312 (2005).
[14] D. N. Basu, J. Phys. G **30**, B35 (2004).
[15] Yu. Ts. Oganessian *et al.*, Phys. Rev. C **69**, 021601(R) (2004).

[16] G. R. Satchler and W. G. Love, Phys. Rep. **55**, 183 (1979) and references therein.
[17] D. N. Basu, P. R. Chowdhury, and C. Samanta, Phys. Rev. C **72**, 051601(R) (2005); D. Gupta and D. N. Basu, Nucl. Phys. **A748**, 402 (2005).
[18] D. N. Basu, J. Phys. G **30**, B7 (2004).
[19] D. N. Basu, Phys. Lett. **B566**, 90 (2003).
[20] D. K. Srivastava, D. N. Basu, and N. K. Ganguly, Phys. Lett. **B124**, 6 (1983).
[21] K. W. Ford and J. G. Wills, Phys. Rev. **185**, 1429 (1969).
[22] D. K. Srivastava, N. K. Ganguly, and P. E. Hodgson, Phys. Lett. **B51**, 439 (1974).
[23] D. N. Poenaru, W. Greiner, M. Ivascu, D. Mazilu, and I. H. Plonski, Z. Phys. A **325**, 435 (1986).
[24] A. Sobiczewski, Z. Patyk, and S. Cwiok, Phys. Lett. **B224**, 1 (1989).
[25] W. D. Myers, W. J. Swiatecki, Lawrence Berkeley Laboratory preprint LBL-36803 (1994); Nucl. Phys. **A601**, 141 (1996).
[26] R. Smolanczuk, Phys. Rev. C **56**, 812 (1997).
[27] G. Royer, J. Phys. G **26**, 1149 (2000).