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α-nucleus potentials, α-decay half-lives, and shell closures for superheavy nuclei

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Systematic α-nucleus folding potentials are used to analyze α-decay half-lives of superheavy nuclei. Preformation factors of about several percent are found for all nuclei under study. The systematic behavior of the preformation factors and the volume integrals of the potentials allows predictions of α-decay energies and half-lives for unknown nuclei. Shell closures can be determined from measured α-decay energies using the discontinuity of the volume integral at shell closures. For the first time a double shell closure is predicted for $Z_{\text{magic}} = 132$, $N_{\text{magic}} = 194$, and $A_{\text{magic}} = 326$ from the systematics of folding potentials. The calculated α-decay half-lives remain far below 1 ns for superheavy nuclei with double shell closure and masses $A > 300$ independent of the precise knowledge of the magic proton and neutron numbers.

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The α decay of superheavy nuclei has been studied intensively in the past few years [1–10]. In many papers a simple two-body model was applied [11], and in most papers a potential was derived that was able to fit the measured α-decay half-lives of the analyzed nuclei. However, most of the studies (with the exception of [2]) did not attempt to use these potentials for the description of other experimental quantities such as, for example, α scattering cross sections or (n,α) or fusion reaction cross sections.

Therefore, an alternative approach was followed in Ref. [12], where the simple two-body model was combined with systematic α-nucleus folding potentials that are able to describe various properties, and the ratio between the calculated half-life $T^{\text{calc}}_{1/2,\alpha}$ and the experimental half-life $T^{\text{exp}}_{1/2,\alpha}$ was interpreted as preformation factor P of the α particle in the decaying nucleus. In addition to the systematic behavior of the volume integrals of the folding potentials, preformation factors of the order of a few percent were found for a large number of nuclei [12,13].

Only for very few light nuclei have some levels been found where a simple two-body model can exactly reproduce the experimental half-lives or widths (e.g., for $^{6}\text{Li} = {}^{2}\text{H} \otimes \alpha$ [14] or for $^{8}\text{Be} = \alpha \otimes \alpha$ [12,15]). Already for $^{20}\text{Ne} = {}^{16}\text{O} \otimes \alpha$ the calculated widths are somewhat larger than the experimentally observed ones [16]. Any simple two-body model with a realistic potential must provide half-lives identical or shorter than the experimental value, because the two-body model assumes a pure α cluster wave function by definition, whereas any realistic wave function is given by the sum over many different configurations. Thus, preformations of a few percent are a quite natural finding for superheavy nuclei.

The following ingredients were used in this Rapid Communication. The α-nucleus potential was calculated from a double-folding procedure with an effective interaction [12,17,18]. The nuclear densities were taken from [19] for the α particle and derived from the two-parameter Fermi distributions for ^{232}Th and ^{238}U in Ref. [19] with properly scaled radius parameter $r \sim A_T^{1/3}$. The total potential is given by the sum of the nuclear potential $V_N(r)$ and the Coulomb potential $V_C(r)$:

$$V(r) = V_N(r) + V_C(r) = \lambda V_F(r) + V_C(r). \tag{1}$$

The Coulomb potential is taken in the usual form of a homogeneously charged sphere, where the Coulomb radius R_C has been chosen identically with the rms radius of the folding potential V_F, and the folding potential V_F is scaled by a strength parameter λ, which is of the order of 1.0–1.3. This leads to volume integrals of about $J_R \approx 300\,\text{MeV fm}^3$ for all nuclei under study and is in agreement with systematic α-nucleus potentials derived from elastic scattering [20–26]. (Note that, as usual, the negative sign of J_R is omitted in this work.) Bound-state properties of $^{212}\text{Po} = {}^{208}\text{Pb} \otimes \alpha$ have been analyzed successfully using the same potential [27]. The centrifugal potential has been omitted for $L = 0$ decays. The following study is restricted to even-even nuclei because the additional centrifugal barrier may influence the α-decay half-life for decays with $L \neq 0$.

The quotations of the volume integral J_R and the potential strength parameter λ are practically equivalent for this paper. If one wants to compare this work to folding potentials with a different nucleon-nucleon interaction or even to potentials with a different parametrization (e.g., Woods-Saxon potentials), the volume integral J_R has to be used. Therefore the following discussion is restricted to volume integrals. Nevertheless, most figures provide both quantities J_R and λ.

The α-decay width Γ_α is given by the following formula [11]:

$$\Gamma_\alpha = P\,F\,\frac{\hbar^2}{4\mu}\exp\left[-2\int_{r_2}^{r_3} k(r)dr\right], \tag{2}$$

with preformation factor P, normalization factor

$$F\int_{r_1}^{r_2}\frac{dr}{k(r)} = 1 \tag{3}$$

and wave number

$$k(r) = \sqrt{\frac{2\mu}{\hbar^2}|E - V(r)|}. \tag{4}$$

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μ is the reduced mass and E is the decay energy of the α decay, which was taken from the computer files based on the mass table of Ref. [28] or from Table 1 of Ref. [1]. The r_i are the classical turning points. For $0^+ \to 0^+ s$-wave decay the inner turning point is at $r_1 = 0$. r_2 varies around 9 fm, and r_3 varies strongly depending on the energy. The decay width Γ_α is related to the half-life by the well-known relation $\Gamma_\alpha = \hbar \ln 2 / T_{1/2,\alpha}$. Following Eq. (2), the preformation factor may also be obtained as

$$P = \frac{T^{\rm calc}_{1/2,\alpha}}{T^{\rm exp}_{1/2,\alpha}}, \tag{5}$$

where Γ_α or $T^{\rm calc}_{1/2,\alpha}$ are calculated from Eq. (2) with $P = 1$. For completeness, the here predicted half-life for unknown nuclei will be defined here as $T^{\rm pre}_{1/2,\alpha} = T^{\rm calc}_{1/2,\alpha}/P$.

The potential strength parameter λ was adjusted to the energy of the α particle in the α emitter $(A+4) = A \otimes \alpha$. The number of nodes of the bound-state wave function was taken from the Wildermuth condition

$$Q = 2N + L = \sum_{i=1}^{4} (2n_i + l_i) = \sum_{i=1}^{4} q_i, \tag{6}$$

where Q is the number of oscillator quanta, N is the number of nodes, L is the relative angular momentum of the α-core wave function, and $q_i = 2n_i + l_i$ are the corresponding quantum numbers of the nucleons in the α cluster. A value of $q = 5$ was taken for $82 < Z, N \leqslant 126$ and $q = 6$ for $N > 126$, where Z and N are the proton and neutron number of the daughter nucleus. This definition of Q deviates slightly from the semiclassical Bohr-Sommerfeld quantum number G as mostly used. One finds $G \approx 22.5$ for all nuclei with $Q = 22$.

Various attempts have been made to determine the preformation factors P experimentally or theoretically [29–31]. The usage of a simple two-body wave function in connection with the Wildermuth condition [Eq. (6)] is obviously a quite simple approximation for the description of the complex many-body wave function of a superheavy nucleus that was analyzed, for example, in Refs. [32–36]. Nevertheless, the preformation factor defined as the ratio $P = T^{\rm calc}_{1/2,\alpha}/T^{\rm exp}_{1/2,\alpha}$ in Eq. (5) may be understood as an effective preformation factor. The obtained values for P do only show small variations and can thus be used for the prediction of half-lives of unknown superheavy nuclei in a consistent way. A full discussion of preformation factors is beyond the scope of the present Rapid Communication.

The resulting preformation factors P for even-even nuclei are shown in Fig. 1. An average value of $P \approx 8\%$ is found. Almost all results lie within a bar of uncertainty of a factor of 3. This uncertainty is identical to the results of Refs. [1,2]. However, the values for P are much smaller in this work (see previous discussion). A table of the results will be given in a forthcoming paper.

There are two different ways in this simple two-body model to obtain larger α-decay half-lives $T^{\rm calc}_{1/2,\alpha}$ and thus larger preformation factors P as derived from Eq. (5). First, very narrow potentials can be used. In this case the attractive nuclear potential becomes negligible in the region of the Coulomb barrier, thus effectively increasing the barrier and increasing the α-decay half-life. This idea was followed, for example, in Ref. [10], and the differences to the systematic folding potential in the present work are illustrated in Fig. 1 of Ref. [12]. A very narrow potential as used in Ref. [10] is probably unable to describe quantities other than the α-decay half-life. Second, a smaller quantum number (G or Q) can be used. This idea was followed in Ref. [1]. In this case the attractive nuclear potential is reduced at all radii, thus again effectively increasing the Coulomb barrier and the α-decay half-life. Many quantities are mainly sensitive to the tail of the wave functions at large radii outside the nuclear potential, which leads to discrete ambiguities for the volume integral J_R of α-nucleus potentials (the so-called family problem). However, it has been found in the past few years that systematic α-nucleus folding potentials have volume integrals J_R around 300 MeV fm^3 [20–26] compatible with the quantum number Q used in the present work and incompatible with the smaller G used in Ref. [1].

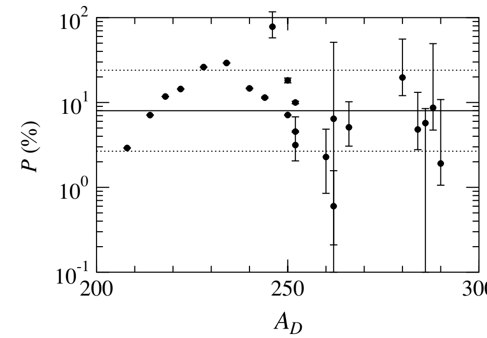

FIG. 1. The preformation factors P for several superheavy α emitters. The horizontal lines indicate an average value of $P \approx 8\%$ (full line) and typical uncertainties of a factor of 3 (dotted lines).

In principle, the application of a semiclassical model is not necessary for the calculation of α-decay half-lives or widths. From the potential in Eq. (1) one can directly calculate the wave function and the width of the decaying state. However, in practice this is difficult because of the low energies and extremely small widths of the states. For ^{212}Po $=$ ^{208}Pb $\otimes\, \alpha$ such a fully quantum-mechanical calculation is possible at the limits of numerical stability. Figure 2 shows the scattering phase shift δ_L for the $L = 0$ partial wave as a function of energy, which is given by $E = E_0 + i \times \Delta E$ with $E_0 = 8.954088523002$ MeV and $\Delta E = 2 \times 10^{-15}$ MeV. The points are the results of a phase-shift calculation; the line is a

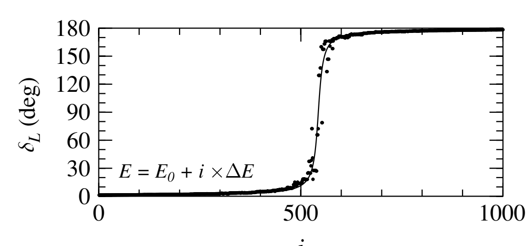

FIG. 2. Phase shift δ_L for the $L = 0$ partial wave for the system ^{212}Po $=$ ^{208}Pb $\otimes\, \alpha$. The derived width from Eq. (7) is $\Gamma = 3.52 \times 10^{-14}$ MeV. Note the extremely small step size of the calculation of $\Delta E = 2.0 \times 10^{-15}$ MeV.

fit to the points using the formula for narrow resonances

$$\delta_L(E) = \arctan \frac{\Gamma}{2(E - E_R)}, \quad (7)$$

with $E_R = 8.9541$ MeV and $\Gamma = 3.52 \times 10^{-8}$ eV, which translates to $T_{1/2,\alpha}^{\text{calc}} = 13.0$ ns. The semiclassical approximation yields $T_{1/2,\alpha}^{\text{calc}} = 8.7$ ns, which is about 30% lower than the precise quantum-mechanical value. A similar result is obtained for the decay of ^{8}Be, where one finds $\Gamma_\alpha = 6.7$ eV for the quantum-mechanical calculation and $\Gamma_\alpha = 10.5$ eV for the semiclassical approximation. As already pointed out, the preformation factors are close to unity for ^{8}Be with $P = 100\%$ (65%) for the quantum-mechanical (semiclassical) calculation and of the order of a few percent for ^{212}Po with $P = 4.3\%$ (2.9%). These results confirm the applicability of the semiclassical model within uncertainties of about 30%.

It is interesting to use the systematic folding potentials for the prediction of properties of unknown superheavy nuclei such as α-decay energies, α-decay half-lives, and shell closures above $N, Z = 126$. The basic building block is the smooth behavior of the strength parameter λ of the folding potential and the resulting volume integrals J_R (see Fig. 3 of this work and Table I of Ref. [12]).

Within one major shell, one finds variations of J_R from about $J_R \approx 335$ MeV fm^3 at the lower end of a shell to about $J_R \approx 280$ MeV fm^3 at the upper end of a shell. Between neighboring nuclei the variation in J_R is below $\Delta J_R < 5$ MeV fm^3. This allows first the determination of α-decay energies for unknown nuclei. As an example, one finds for the decay of $^{298}120 \to {}^{294}118$ a volume integral of $J_R \approx 296$ MeV fm^3, corresponding to $\lambda = 1.138$. This leads to a decay energy of $E = 12.87$ MeV. The α-decay half-life can be estimated using the given energy and an average preformation factor of $P \approx 8\%$, leading to $T_{1/2,\alpha}^{\text{pre}} \approx 8\ \mu$s. Whereas the uncertainties for the volume integral J_R and the derived α-decay energy are small, the uncertainty of the α-decay half-life is strong because of the exponential energy dependence. For a potential strength enhanced (reduced) by 2% one finds the α-decay energy $E = 10.70$ MeV ($E = 14.98$ MeV) and $T_{1/2,\alpha}^{\text{pre}} = 0.97$ s ($T_{1/2,\alpha}^{\text{pre}} = 1.6$ ns) again using $P = 8\%$. A variation of the potential strength of 1% corresponds to a variation of the α-decay energy of about 1 MeV, which is comparable to the uncertainties of mass formulas [37]. As usual, the reliability of such an extrapolation decreases for nuclei with masses far above the heaviest known nuclei. However, the uncertainties for closed-shell nuclei remain small because of the well-defined volume integral J_R for such nuclei, which can be studied at the shell closures at $N = 82$, $Z = 82$, and $N = 126$.

Shell closures can be seen as discontinuities in the volume integrals (see Figs. 3 and 4). Whereas the variation between neighboring nuclei remains below $\Delta J_R < 5$ MeV fm^3, at shell closures one finds a strong increase of J_R that is directly related to the increase of the quantum number Q. Because shell closures are not known a priori for superheavy nuclei, Fig. 4 analyzes the volume integrals around the shell closure at $N = 82$ for Xe, Ba, Ce, and Nd isotopes. Below $N = 82$, the wave functions are characterized by $Q = 16$ (full black symbols), and the volume integrals are slightly above $J_R = 280$ MeV fm^3. Above $N = 82$ one finds volume integrals around 310 MeV fm^3 for wave functions with $Q = 18$ (open symbols; see also Ref. [12]). The small gray symbols are calculated above the shell closure at $N = 82$ without an increase of the quantum number Q. Here one finds low volume integrals significantly below $J_R = 280$ MeV fm^3 that differ by more than $\Delta J_R = 10$ MeV fm^3 from the neighboring values. The behavior of the potential strength parameter λ is similar to that of J_R (see Fig. 4).

A similar behavior is found at the shell closures at $Z = 82$ and $N = 126$ around ^{208}Pb. Compared to neighboring values, the volume integrals J_R are reduced by more than $\Delta J_R = 10$ MeV fm^3 if one neglects the increase of the quantum number Q, and the absolute values of the volume integrals drop below $J_R = 280$ MeV fm^3. Consequently, changes in J_R by about 10 MeV fm^3 or values of J_R significantly below 280 MeV fm^3 are clear indications for the crossing of a major shell. Because the determination of the volume integral J_R requires only the knowledge of the α-decay energy, the measurement of one single quantity may be sufficient for the determination of a double closure: As soon as J_R drops below 280 MeV fm^3, magic neutron or proton numbers have been crossed!

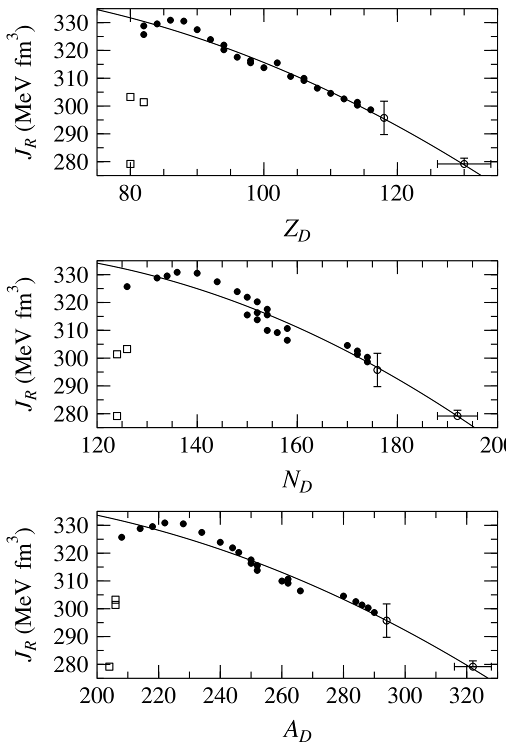

FIG. 3. Volume integrals J_R for superheavy nuclei as a function of Z_D (upper), N_D (middle), and A_D (lower). Within a major shell one finds a smooth decrease of J_R with a minimum value around $J_R \approx 280$ MeV fm^3. From the fits to the data points one can directly see that J_R drops below 280 MeV fm^3 at $Z \approx 130$, $N \approx 192$, and $A \approx 322$, leading to the magic numbers $Z_{\text{magic}} = 132$, $N_{\text{magic}} = 194$, and $A_{\text{magic}} = 326$. The open circles are extrapolations for ^{294}118 and the lower limit of $J_R > 279$ MeV fm^3 close to the next closed shells (see text). The open squares are nuclei with $Q = 18$ ($^{208}\text{Pb} = {}^{204}\text{Hg} \otimes \alpha$) and $Q = 20$ ($^{210}\text{Pb} = {}^{206}\text{Hg} \otimes \alpha$, $^{210}\text{Po} = {}^{206}\text{Pb} \otimes \alpha$).

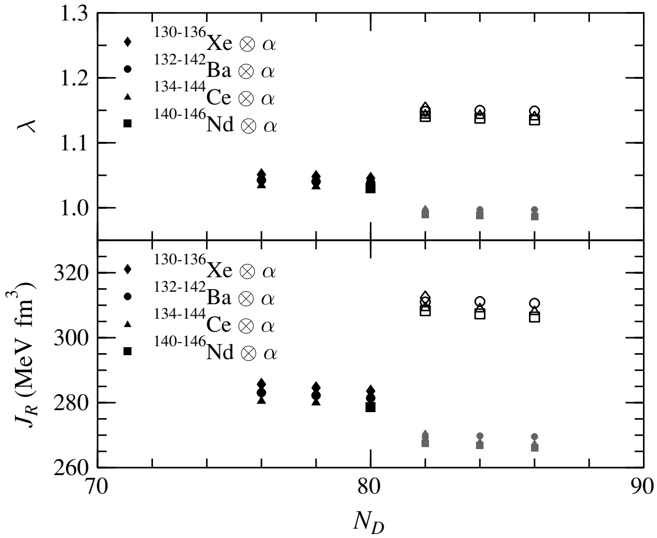

FIG. 4. Potential strength parameter λ (upper) and volume integrals J_R (lower) around the shell closure $N = 82$ for $^{130-136}\mathrm{Xe} \otimes \alpha$ (diamonds), $^{132-142}\mathrm{Ba} \otimes \alpha$ (circles), $^{134-144}\mathrm{Ce} \otimes \alpha$ (triangles), and $^{140-146}\mathrm{Nd} \otimes \alpha$ (squares) isotopes. One can clearly see the discontinuity at the shell closure $N = 82$ (see text).

This systematic behavior of α-nucleus potentials allows further a prediction of magic numbers in a yet unknown mass region above $A > 300$. The smooth energy dependence of the volume integrals J_R in Fig. 3 is fitted using a second-order polynomial for all nuclei with $Q = 22$ (full lines in Fig. 3). J_R drops below $280\,\mathrm{MeV\,fm^3}$ at $Z = 130$, $N = 192$, and $A = 322$, which means that the nucleus $^{326}132 = {}^{322}130 \otimes \alpha$ is the heaviest nucleus that can be described using a potential with $J_R > 280\,\mathrm{MeV\,fm^3}$ and $Q = 22$. Increasing Z or N leads to J_R below its lower limit, and thus the magic numbers $Z_{\mathrm{magic}} = 132 \pm 4$, $N_{\mathrm{magic}} = 194 \pm 4$, and $A_{\mathrm{magic}} = 326 \pm 6$ can be derived from Fig. 3.

The α-decay half-life of the doublymagic nucleus with $Z_{\mathrm{magic}} = 132$, $N_{\mathrm{magic}} = 194$, and $A_{\mathrm{magic}} = 326$ can be calculated using the volume integral $J_R = 279.2\,\mathrm{MeV\,fm^3}$ (taken from $^{208}\mathrm{Pb} = {}^{204}\mathrm{Hg} \otimes \alpha$). One finds the energy $E = 18.26\,\mathrm{MeV}$ and the corresponding half-life $T_{1/2,\alpha}^{\mathrm{calc}} = 1.16 \times 10^{-12}$ s with $P = 1$. Again by using $P = 8\%$, a realistic prediction of the half-life is $T_{1/2,\alpha}^{\mathrm{pre}} = 1.5 \times 10^{-11}$ s. Including the uncertainty of P, the half-life remains below 10^{-10} s. The uncertainty of the volume integral J_R at closed shells is smaller than 1%. Increasing J_R by 1% reduces the α-decay energy by about 1 MeV and increases the α-decay half-life by about a factor of 20. In any case, the half-life remains below 1 ns.

The lower limit of J_R has also been applied for the calculation of the half-life of $^{310}126 \rightarrow {}^{306}124$ with the widely discussed shell closures at $Z = 126$ and $N = 184$ (e.g. [38, 39]); but also other magic numbers have been discussed (e.g. [40]). Here one obtains the α-decay energy $E = 18.82$ MeV and the corresponding α-decay half-life $T_{1/2,\alpha}^{\mathrm{calc}} = 2.1 \times 10^{-14}$ s ($P = 1$). The realistic prediction using $P = 8\%$ is $T_{1/2,\alpha}^{\mathrm{pre}} = 2.6 \times 10^{-13}$ s, and including all uncertainties the half-life remains far below 10^{-11} s. These calculations indicate clearly that one cannot expect that any superheavy nucleus above $A > 300$ with magic proton and neutron numbers (whatever these numbers are) has a half-life significantly above 1 ns.

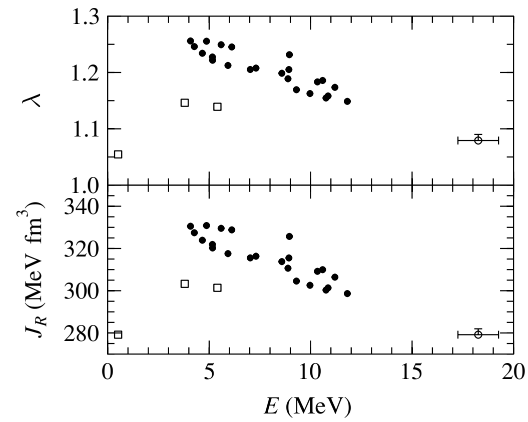

FIG. 5. Potential strength parameter λ (upper) and volume integrals J_R (lower) vs decay energy E for superheavy nuclei. Known nuclei are shown as full circles. The extrapolated doublymagic nucleus with $Z_{\mathrm{magic}} = 132$, $N_{\mathrm{magic}} = 194$, and $A_{\mathrm{magic}} = 326$ is shown as an open circle. The open squares are nuclei with $Q = 18$ and $Q = 20$ (see also Fig. 3).

Because of the significant variation of the decay energy E from about 4 MeV to about 12 MeV for known superheavy nuclei and up to about 20 MeV for the predicted but yet unknown doublymagic superheavy nuclei one may expect a correlation between the potential strength parameter λ or the volume integral J_R and the decay energy E. This relation is analyzed in Fig. 5 for superheavy nuclei and in Fig. 6 for nuclei around $N = 82$.

Figure 5 seems to indicate that larger decay energies E are correlated to smaller volume integrals J_R. However, the

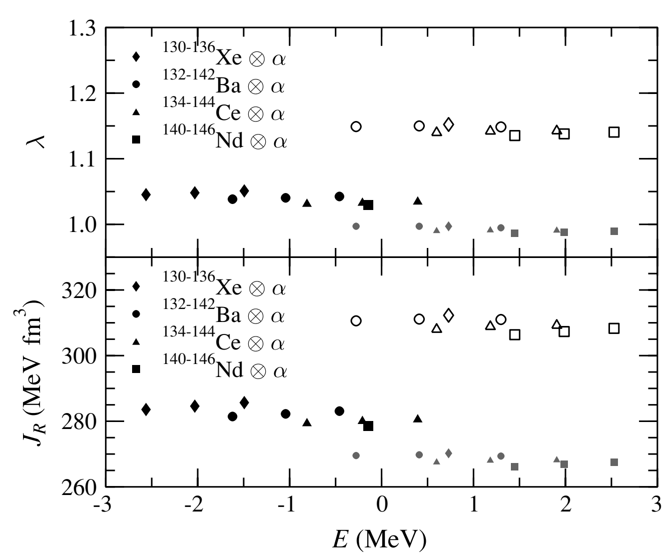

FIG. 6. Potential strength parameter λ (upper) and volume integrals J_R (lower) vs decay energy E around the shell closure $N = 82$ for $^{130-136}\mathrm{Xe} \otimes \alpha$ (diamonds), $^{132-142}\mathrm{Ba} \otimes \alpha$ (circles), $^{134-144}\mathrm{Ce} \otimes \alpha$ (triangles), and $^{140-146}\mathrm{Nd} \otimes \alpha$ (squares) isotopes (see text; for description of symbols see Fig. 4).

underlying reason for this energy dependence of J_R is the smooth variation of J_R within a major shell (see previous discussion). At very small energies one finds again a small volume integral of $J_R \approx 280\,\text{MeV}\,\text{fm}^3$ for $^{208}\text{Pb} = {}^{204}\text{Hg} \otimes \alpha$. As can also be seen from Fig. 6, the volume integrals do not depend on the energy E: Above $N = 82$ one finds $J_R \approx 310\,\text{MeV}\,\text{fm}^3$ for bound ($E < 0$) and unbound ($E > 0$) nuclei, and below $N = 82$ one finds $J_R \approx 280\,\text{MeV}\,\text{fm}^3$, again for bound and unbound nuclei.

In conclusion, systematic folding potentials can be used for the calculation of α-decay half-lives of superheavy nuclei. Additionally, the systematic behavior of the volume integrals allows predictions of α-decay energies and half-lives of yet unknown nuclei. The magic numbers $Z_{\text{magic}} = 132$, $N_{\text{magic}} = 194$, and $A_{\text{magic}} = 326$ have been derived from the discontinuities of the volume integrals at shell closures. There is strong evidence that α-decay half-lives remain far below 1 ns even for doublymagic superheavy nuclei above $A > 300$.

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