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Present view of stability of heavy and superheavy nuclei

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Contents

Abstract. Recent theoretical studies of the stability of heavy and superheavy nuclei are shortly reviewed. Even-even nuclei with proton number $Z = 82 - 120$ and neutron number $N = 126 - 190$ are considered. The important role of the shell structure in nuclear stability was illustrated. Much attention is given to deformed superheavy nuclei, which are expected to be on the way to the long-discussed, hypothetical spherical superheavy nuclei.

1. Introduction

An intensive activity, both experimental and theoretical, is evident in the field of synthesis and also in studying the properties of heaviest nuclei. A review of somewhat earlier experimental results, as well as findings of more recent experiments, may be found, for instance, in Refs [1 – 11].

The objective of the present paper is to give a short review of recent theoretical studies on heaviest nuclei. A survey of earlier results may be found in Refs [12, 13]. The studies presented here are based on the macroscopic-microscopic description of nuclear properties. A discussion of the results obtained in a fully microscopic (Hartree – Fock – Bogolubov) approach has been presented in Ref. [14].

The theoretical studies described in this review are closely connected with the experimental research on the heaviest nuclei. They aim at describing the existing experimental results and also at a prediction of the properties of nuclei not yet observed. It will be seen that they mainly concentrate on solving the problem of stability of these nuclei.

2. Essential role of shell effects

It is known that nuclei, much like the atoms, possess shell structure. Effects of this structure are important for all nuclei. Their role for heaviest nuclei is, however, all the more essential, as many of them would not simply exist without these effects.

The objective of this section is to illustrate the important contribution of shell effects to the stability (half-lives) of heaviest nuclei. The illustration is based on the results of Ref. [15] wherein an extensive quantitative analysis of shell effects in these nuclei has been performed. Even-even transthorium nuclei have been examined in that paper.

Figure 1, taken from Ref. [15], shows the experimental, $T_\alpha^{\rm expt}$, and smooth (macroscopic), T_α, α-decay half-lives (given in seconds), both on the logarithmic scale. The smooth half-life T_α was calculated by a model (of the liquid drop type) of a nucleus, which did not allow for any shell structure.

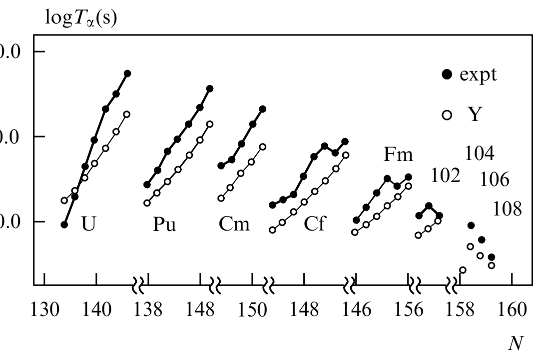

Figure 1. Logarithm of experimental (expt) and smooth (Y) α-decay half-lives (in seconds) [15].

Thus, the difference between two half-lives is indicative of the shell effect in the α-decay of a nucleus. One can see that for all heavy nuclei, excepting the two lightest uranium isotopes, the shell effect elongates the half-life. The latter enhances by 2 – 5 orders of magnitude for most of the nuclei considered.

Even larger shell effects were found in the spontaneous-fission half-lives $T_{\rm sf}$. This is shown in Fig. 2, where logarithm of $T_{\rm sf}$, both experimental and calculated in a macroscopic model without considering any shell effects, is given. It is seen that the shell effect delays the fission process in all the nuclei

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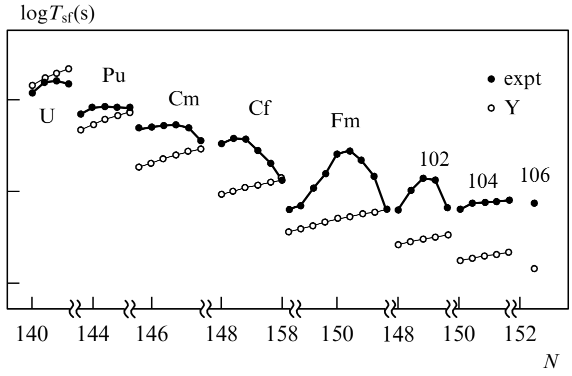

Figure 2. Logarithm of experimental (expt) and macroscopic (Y) spontaneous-fission half-life $T_{\rm sf}$, given in seconds [15].

considered, except only for a few lightest ones (isotopes of uranium). The delay increases from several orders (Pu isotopes) to about 15 orders of magnitude for the nucleus $^{260}106$†, which has the largest Z among the even-even nuclei with measured $T_{\rm sf}$. For such a heavy nucleus like $^{260}106$, with $T_{\rm sf}$ of the order of a few milliseconds, this elongation makes up practically the whole half-life of these nuclei. In other words, they would not exist without shell effects, as already mentioned above.

The mechanism by means of which practically the whole half-life of a very heavy nucleus is made up by shell effects is illustrated in Fig. 3. The figure displays the spontaneous-fission barrier of the $^{264}108$ nucleus, i.e. the dependence of the ground-state energy of this nucleus on its quadrupole-deformation parameter β_2. For each β_2, the energy is minimized with respect to the hexadecapole-deformation parameter β_4. The total fission barrier (Y + SHELL), having regard to shell effects, is shown by solid line and its smooth part (obtained by the Yukawa-plus-exponential model (Y) [16]), by dashed line. The smooth barrier calculated with another macroscopic model (liquid drop, LD [17]) is also shown (dotted line) for comparison. One can see that a

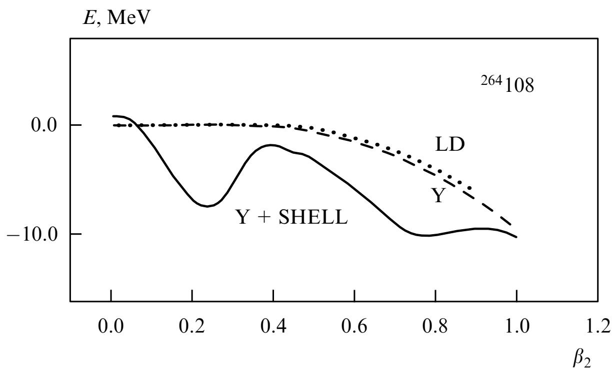

Figure 3. Total fission barrier (Y + SHELL) and its smooth part, obtained by the Yukawa-plus-exponential (Y) and by the liquid-drop (LD) models, for the $^{264}108$ nucleus [15].

significant height (about 6 MeV) of the fission barrier is obtained only after inclusion of shell effects. Without them, no fission barrier (Y and LD) appears. It should be added here that shell effects are also important after the fission barrier penetration, down to the scission point, as has been shown in Refs [18–21].

Figures 1–3 illustrate a very important part of shell effects in the properties of heaviest nuclei, particularly in their stability. Simultaneously, they point to a strong dependence of shell effects on the proton Z and neutron N numbers. This gives an implication for the theory that each nucleus should be treated individually (i.e. without any averaging over a number of nuclei) under a theoretical analysis. The strong dependence of shell effects also on the deformation of a nucleus, illustrated in Fig. 3, requires a careful treatment of this deformation in the analysis. In other words, the consideration of the properties of a heavy nucleus should be performed in a sufficiently large, multidimensional deformation space [22–24].

3. Theoretical methods

As already stated in the Introduction, extensive studies of stability of heavy and superheavy nuclei are based on the macroscopic-microscopic approach. Although relatively simple, this approach allows one to describe a number of nuclear properties, in particular the nuclear mass [25], quite well. The macroscopic-microscopic approximation has underlain the works reviewed in this paper. However, pure microscopic approaches consisting in the self-consistent Hartree–Fock–Bogolubov calculations with the use of effective two-body forces, were also taken in the literature (e.g., [26, 14]).

In the macroscopic-microscopic calculations reviewed in this article, mass of a nuclide is defined as a sum of the macroscopic part, described by the Yukawa-plus-exponential model [16], and the microscopic part, which allows for the shell correction. The latter is obtained by the Strutinsky procedure [27] and is based on the Woods–Saxon single-particle potential [28].

The α-decay half-life T_α is calculated by the phenomenological formula of Viola and Seaborg [29] with four adjustable parameters refitted to account for new data [30].

Finally, the spontaneous-fission half-life $T_{\rm sf}$ is calculated on a basis of the dynamical approach [31–33]. It consists in the search for a one-dimensional fission trajectory in a multidimensional deformation space, which minimizes the action integral corresponding to the penetration of a nucleus through the fission barrier. The inertia tensor appearing in the integral and describing the inertia of a nucleus with respect to its deformation is calculated in the cranking approach (e.g., Refs [34–36]).

A more detailed description of these methods may be found, for instance, in Refs [23, 12, 37].

4. Main results

4.1 Shell correction to a mass

As described in Section 2, shell correction refers to the main factor influencing the stability of heaviest nuclei. This phenomenon has been discussed in a number of papers, e.g., Refs [38, 39, 15, 13]. In particular, shell correction to the ground-state mass of a nucleus forms a first notion about the stability of this nucleus.

† Based on the IUPAC recommendations of 1994, the nomenclature of some transfermium elements looks as follows: No($Z = 102$), Db(104), Rf (106), and Ha(108). (*Scientific editor's note*)

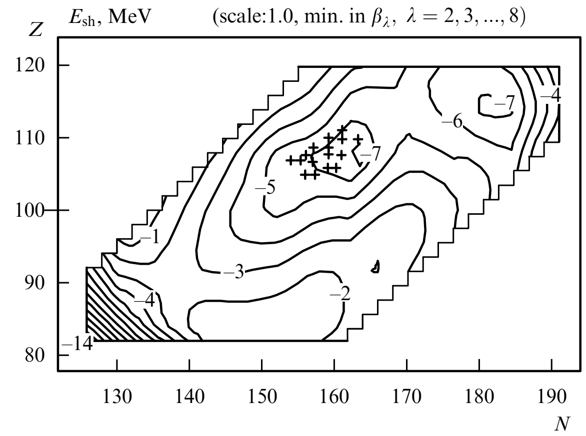

Figure 4. Contour map of the shell correction to energy E_{sh}. Crosses denote the heaviest nuclides synthesized up to date [40].

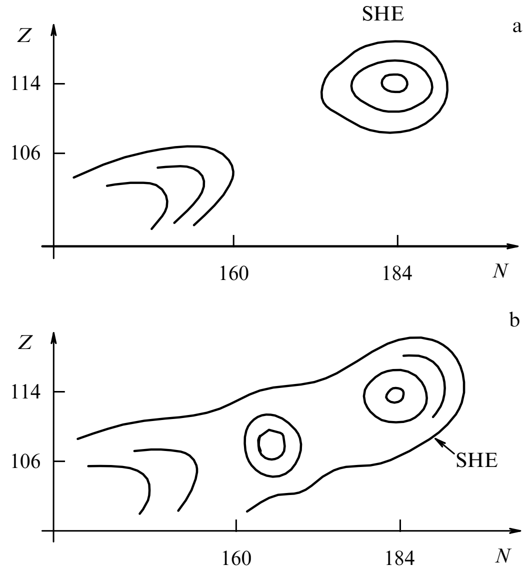

Figure 5. Regions of relatively long-lived nuclei: as believed earlier (a), and expected presently (b) [33].

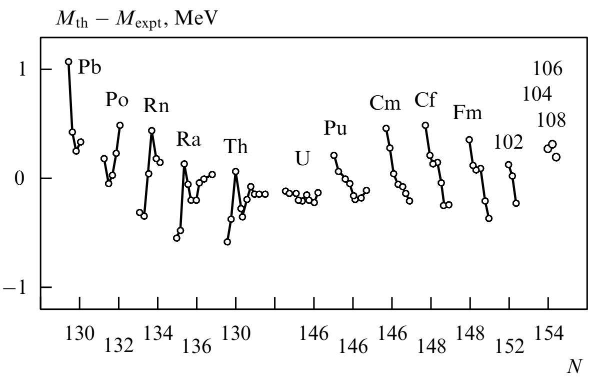

Figure 6. Discrepancies between calculated (th) and experimental (expt) masses [47].

Figure 4, taken from Ref. [40], shows the shell correction to the mass E_{sh} calculated for a large region of heaviest nuclei. One can see that E_{sh} has three minima in the considered region of nuclei. The first one, being the deepest ($E_{sh} = -14.3$ MeV), is obtained for the doubly magic spherical nucleus ^{208}Pb. The second one ($E_{sh} = -7.2$ MeV) appears at the nucleus $^{270}108_{162}$, which was predicted [41, 23] to be a doubly magic deformed nucleus. The third minimum, with the same depth ($E_{sh} = -7.2$ MeV) as that of the second minimum, is obtained for the nucleus $^{296}114_{182}$, which is close to the nucleus $^{298}114_{184}$ predicted [42, 43] to be a doubly magic spherical nucleus, the next one after the last experimentally known ^{208}Pb nucleus. Besides these three minima, there appears a rather wide plateau around the ^{252}Fm nucleus, which, although having a smaller (in absolute value) shell correction ($E_{sh} = -5.2$ MeV) than the $^{270}108$ nucleus, may also be considered as a doubly magic deformed nucleus [41, 23]. Crosses in the figure denote the heaviest nuclides synthesized up to now. The heaviest isotopes of the element $Z = 106$ have been obtained in Ref. [6], those of $Z = 108$ in Ref. [11], that of $Z = 109$ in Refs [44, 9], those of $Z = 110$ in Refs [8, 45] and, finally, that of $Z = 111$ in Ref. [9]. The recently observed $^{277}112$ nuclide [46] is not yet marked in the figure.

One can see in Fig. 4 that some of the already synthesized nuclei feature profit by 6 – 7 MeV in their mass from the shell correction. Without this profit they could not exist, as discussed in Section 2.

The appearance of the region of deformed superheavy nuclei around the predicted doubly magic $^{270}108$ nucleus (^{270}Ha) constitutes the main change in our notion about stability of heaviest nuclei in recent years. Before, it was believed for a long time that spherical superheavy nuclei, predicted to be situated around the doubly magic $^{298}114$ nucleus, would constitute an island separated from the usual peninsula of relatively long-lived nuclei by an 'ocean' of full instability. After the appearance of deformed superheavy nuclei, however, the peninsula is expected to be extended, to include also the spherical superheavy nuclei. This is illustrated qualitatively in Fig. 5, taken from Ref. [33].

4.2 Mass

It is interesting to see how well are the experimental masses reproduced by the theoretical ones, calculated with the shell correction given in Fig. 4. This is illustrated in Fig. 6, taken from Ref. [47], which shows the discrepancy between the calculated and experimental masses. One can realize that for the most of the nuclei considered this discrepancy falls within the limits ± 0.25 MeV, i.e. it is not significant. The largest discrepancy is obtained for the doubly magic ^{208}Pb nucleus. The theoretical binding energy is too small for this nucleus by about 1 MeV. We can also see that the isotopic dependence of the theoretical mass is not correct, excepting only the isotopes of uranium, and it varies from one element to another.

4.3 Half-lives of deformed superheavy nuclei

Figure 7, taken from Ref. [37], shows the α-decay and spontaneous-fission half-lives T_α and T_{sf}, respectively, calculated for deformed superheavy nuclei situated around the $^{270}108$ nucleus. One can clearly see the effect of the deformed $N = 162$ shell. A weaker effect of the $N = 152$ shell is also observed, especially for lighter elements. These effects make the systematics of the half-lives quite complicated.

A comparison between the calculated T_{sf} and T_α shows that, for $Z = 104$, T_{sf} is smaller than T_α at all N. For $Z = 106$, T_{sf} is comparable with T_α for a large number of isotopes ($N = 154 - 164$). For higher Z, it is even larger than T_α and for a larger number of isotopes. This seems to be the effect of

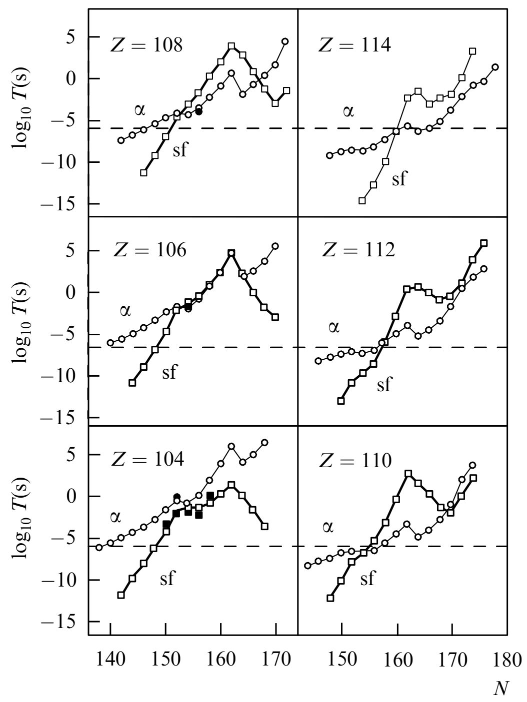

Figure 7. Logarithm of calculated spontaneous-fission (sf) and α-decay (α) half-lives (given in seconds) as functions of the neutron number N for the elements 104 – 114. Experimental values are given as full symbols. The horizontal dashed line indicates approximately the lowest half-life (1 μs) of a nucleus, which can be detected in a present-day setup, after its synthesis [37].

shells, mainly of that at $N = 162$, to which $T_{\rm sf}$ is more sensitive than T_α. Only for the lightest isotopes, $T_{\rm sf}$ is shorter than T_α for all elements investigated.

4.4 Alpha-decay half-lives for deformed and spherical superheavy nuclei

As α-decay is the main decay mode for many nuclei analyzed in Fig. 7, especially those with largest Z, it is interesting to extend the calculation of T_α to even heavier nuclides, to cover also the region of spherical superheavy nuclei. The results of such extension are shown in Fig. 8, taken from Ref. [40]. A rather large region of nuclei with $Z = 100 - 120$ and $N = 146 - 190$ was considered.

One can clearly see the effects of the neutron shells at $N = 152$, 162 and 184. The effect of the spherical shell at $N = 184$ (especially for $Z = 110$) is the strongest, the effect of the deformed shell at $N = 162$ (especially for $Z = 108$) is not much weaker. The effect of the deformed shell at $N = 152$ is the weakest. In addition, the effects of the proton shells are clearly seen. The effect of the spherical shell at $Z = 114$ (especially for isotopes with $N \approx 184$) is about the same as that of the deformed shell at $Z = 108$ (especially for isotopes with $N \approx 162$). The effect of the deformed shell at $Z = 100$ is the weakest.

It might be well to point out in Fig. 8 that due to large shell effects of the doubly magic deformed $^{270}108$ nucleus, its T_α (about 6 s) is not so much shorter than T_α (about 700 s) of the doubly magic spherical $^{298}114$ nucleus, although the latter is much more rich in neutrons. It is also interesting to note that the dependence of $\log T_\alpha$ on Z for the neutron deformed shell at $N = 162$ is much different from that for the neutron spherical shell at $N = 184$. It is less uniform, less smooth at $N = 162$.

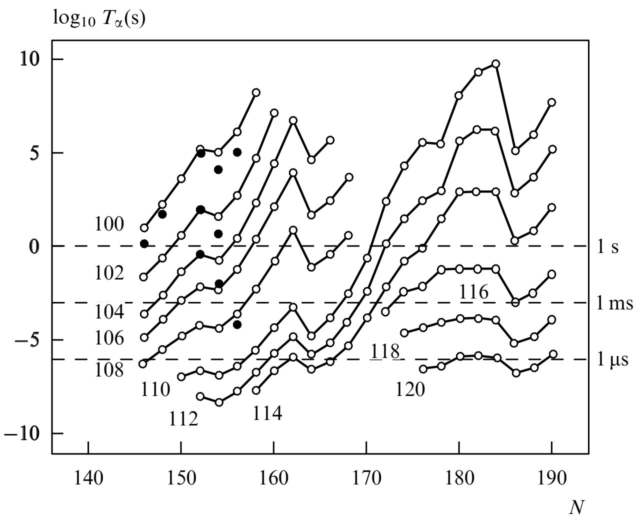

Figure 8. Logarithm of the calculated α-decay half-life T_α (given in seconds) as a function of the neutron number N for nuclei with the proton number $Z = 100 - 120$. Experimental values are also shown (by full circles) [40].

The experimental values of T_α known for 10 nuclei among those considered in Fig. 8, are reproduced by the calculations within a factor of 3, on the average. The largest discrepancy was obtained for the $^{256}102$ nucleus. The calculated value is about 8 times larger than the experimental one, for this nucleus.

Comparison with experimental values also shows that the calculated T_α underestimates the effects of the shells at $N = 152$ and at $Z = 100$.

To recognize the relation between the calculated α-decay, T_α, and the spontaneous-fission, $T_{\rm sf}$, half-lives, we show them in Fig. 9 [47] for isotopes of the 114 element. This element is planned to be synthesized in a near future, both in Darmstadt

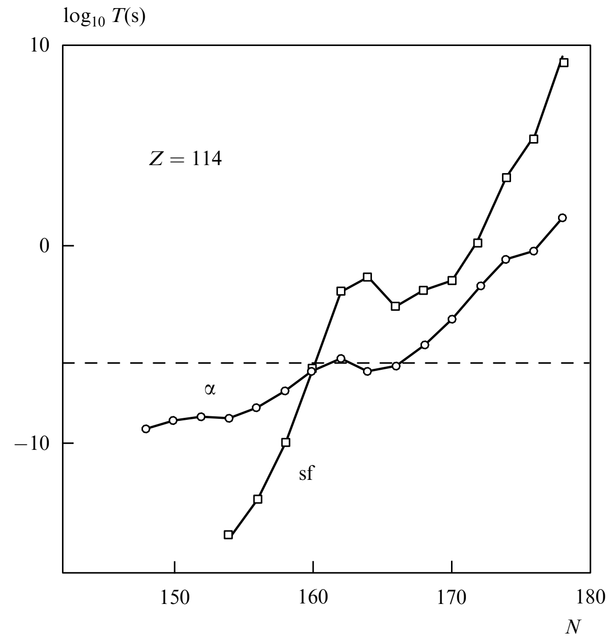

Figure 9. Logarithm of both the α-decay (α) and spontaneous-fission (sf) half-lives (given in seconds) calculated for isotopes of the 114 element [47].

[48] and in Dubna [49]. The figure shows that T_{sf} is larger than T_α for a rather large number of nuclei considered. One can really see that starting from the neutron number $N = 162$, we have: $T_{sf} > T_\alpha$. For the heaviest isotope shown ($N = 178$), T_{sf} is larger than T_α by about 8 orders of magnitude. Additionally, except for a low local maxima of T_α at $N = 162$ and of T_{sf} at $N = 164$, both the half-lives increase with increasing N. The total half-life (equal to T_α) for the heaviest isotope ($N = 178$) amounts to 24 s.

5. Conclusions

In conclusion of this short review of recent theoretical studies on stability of heavy and superheavy nuclei we can say the following:

(1) Shell effects are very important for the stability of heaviest nuclei. According to theoretical analysis, all nuclei with atomic number Z larger than about 105 – 106 exist or are expected to exist only due to these effects.

(2) Shell effects in deformed superheavy nuclei are large. They are comparable with the same effects in spherical superheavy nuclei.

(3) In particular, a large region of deformed superheavy nuclei, situated around the predicted doubly magic deformed $^{270}108$ nucleus, is expected to exist. A number of nuclei in this region have already been observed. Existence of this region changes our previous view of the stability of heaviest nuclei. In particular, spherical superheavy nuclei situated around the hypothetical doubly magic spherical $^{298}114$ nucleus is not expected any more to form an island in the 'ocean' of full instability, but rather to belong to the extended usual peninsula of relatively long-lived nuclides. Thus, one expects presently that all the nuclei on the way to spherical superheavy nuclides can be observed, if synthesized at a laboratory.

(4) Many nuclei in the superheavy region are expected to decay mainly by α-emission. This is important for the experimental studies of these nuclei, as it makes their identification easier and more certain. The experimental observations done up to now support this expectation.

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