Superheavy magic structures in the relativistic Hartree-Fock-Bogoliubov approach

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Abstract

We have explored the occurrence of the spherical shell closures for superheavy nuclei in the framework of the relativistic Hartree-Fock-Bogoliubov (RHFB) theory. Shell effects are characterized in terms of two-nucleon gaps $\delta_{2n(p)}$. Although the results depend slightly on the effective Lagrangians used, the general set of magic numbers beyond $^{208}$Pb are predicted to be $Z = 120, 138$ for protons and $N = 172, 184, 228$ and $258$ for neutrons, respectively. Specifically the RHFB calculations favor the nuclide $^{304}$120 as the next spherical doubly magic one beyond $^{208}$Pb. Shell effects are sensitive to various terms of the mean-field, such as the spin-orbit coupling, the scalar and effective masses. In addition, the breaking of relativistic pseudo-spin symmetry is shown to influence the level structure.

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Preprint submitted to Elsevier December 2, 2013

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For a fairly long period, it remains a challenging issue in nuclear physics to explore the existence limit of very heavy nuclei, i.e., the superheavy elements (SHE) with $Z \geq 104$ and the so-called stability island of superheavy nuclei (SHN). Far from being simply large clusters of nucleons, these fascinating species owe their very existence to subtle contributions to the nuclear binding energy. Experimentally, the discoveries of new elements up to $Z = 118$ have been reported in Refs. \textsuperscript{2,3}. The increasing survival probabilities with increasing proton number of SHE from $Z = 114$ to 118 seem to indicate enhanced shell effects with increasing $Z$ and therefore a possible proton magic shell may emerge beyond $Z \geq 120$ \textsuperscript{4}.

On the other hand, theoretical studies have provided a large amount of valuable information for the exploration of SHN. These studies can be separated into different categories: Microscopic - Macroscopic (Mic-Mac) models \textsuperscript{5,6}, non-relativistic mean field \textsuperscript{7} (S), covariant mean field \textsuperscript{7} (S) \textsuperscript{8,9} approaches. The extrapolation towards the superheavy region challenges the predictivity of nuclear models. The Mic-Mac methods, although generally successful in describing the nuclear binding, require preconceived knowledge of the expected densities and single-particle (s.p.) potentials \textsuperscript{5}, which fades away when stepping into new regions where stronger polarization effects and more complicated functional forms of the densities may occur \textsuperscript{8,9,10}. The stability of nuclei is mostly driven by shell effects and therefore, self-consistent mean field methods are probably the best conceptual tool to explore the superheavy region, although the Mic-Mac models still give a better quantitative description of heavy nuclides.

We are searching for doubly closed-shell systems and we assume spherical symmetry. Then, the shells are essentially determined by the spin-orbit (SO) splittings, and by the effective masses. Another effect which affects the shell structure is related to the pseudo-spin symmetry (PSS) and its breaking \textsuperscript{12,13}. In the non-relativistic self-consistent mean field theory \textsuperscript{14,15}, the SO splittings depend directly on an extra SO parameter in the energy density functional. In the superfluid covariant density functional (CDF) theory, like the relativistic Hartree-Bogoliubov (RHB) \textsuperscript{16,17} or the relativistic Hartree-Fock-Bogoliubov (RHFB) \textsuperscript{18} approaches, the SO splitting depends directly on the Lorentz scalar and vector mean fields without additional term. The SO splitting is not adjusted and can be considered as a prediction of relativistic Lagrangians, even in ordinary nuclei. This might be an advantage for exploring unknown regions. Furthermore, in the more complete RHFB version of the CDF theory the SO splittings can be affected by meson-nucleon couplings like Lorentz $\rho$-tensor couplings \textsuperscript{13} not present in the simple RHB. This is one of the main motivations for undertaking the present study in the framework of the RHFB approach.

In this work we investigate the superheavy nuclides covering $Z = 110 - 140$. In the pairing channel, the finite-range Gogny force DIS \textsuperscript{19} renormalized by a strength factor $f$ is adopted as the effective pairing interaction. The strength factor $f$ is introduced to compensate level-density differences among various mean field approaches. It was indeed shown that pairing related quantities, such as odd-even mass differences and moments of inertia, are systematically overestimated in the RHFB calculations of heavy nuclei with the original Gogny pairing force \textsuperscript{20}. The strength factor $f = 0.9$ is therefore adjusted to reproduce the odd-even mass differences of odd Pb isotopes. Concerning the relativistic Hartree-Fock (RHF) mean field, the adopted effective Lagrangians are PKA1 \textsuperscript{11,13} and the PKOi series (i=2, 3) \textsuperscript{21,22}. To compare with approaches neglecting the Fock term (RHB), we also use PKDD \textsuperscript{23} and DDM-E2 \textsuperscript{24} Lagrangians. The integro-differential RHFB equations are solved by using a Dirac Woods-Saxon basis \textsuperscript{25} with
a radial cutoff $R = 28$ fm. The numbers of positive and negative energy states in the basis expansion for each s.p. angular momentum $(l, j)$ are chosen to be 44 and 12, respectively.

Let us first discuss extrapolations to SHE of mean field models which are well constrained on medium and heavy nuclei. For instance, due to the high level density in SHE, small variations in the s.p. level spacings due to different SO splitting predictions of various models, can have a large effect on magicity. The SO force is therefore a crucial ingredient of nuclear structure models, especially when it comes to extrapolations to SHN. Fig. 1 shows the relative differences between calculated and experimental SO splittings for a selection of levels having well controlled spectroscopic factors. The relative differences are typically ~ 20% when both partners are particle or hole states, but they become larger otherwise. This is not surprising since polarization and correlation effects tend to shift unoccupied and occupied s.p. states into opposite directions [27, 28]. If one compares the results of Fig. 1 with those from non-relativistic mean field models such as Skyrme-Hartree-Fock (SHF) [8] it appears that the latter give systematically larger deviations. Fig. 1 provides therefore a good motivation for predictions of SHE based on relativistic Lagrangians.

SHE predictions have been carried out using relativistic mean field (RMF) models [17] or RHB models [10]. In such Hartree-type approaches, the contribution of the Fock term is disregarded, at variance with RHF, leading to a renormalisation of the coupling constants. It is an approximation which forbids the inclusion of the $\pi$ and the $\rho$-tensor mesons. While RMF models are as predictive as RHF ones for medium and heavy nuclei, it is preferable to base extrapolations to SHE on calculations including correctly the contribution of the Fock term. It is also a motivation of the present study.

Magicity in SHN might not be as well-marked as in the ordinary nuclei [9]. To identify the magic shells, we will employ the so-called two-nucleon gaps, $\delta_{2p}$ (proton) and $\delta_{2n}$ (neutron), i.e., the difference of two-nucleon separation energies of neighboring isotopes or isotones, which provides an efficient evaluation of the shell effects [7, 10].

$$\delta_{2p}(N, Z) = S_{2p}(N, Z) - S_{2p}(N, Z + 2), \quad (1a)$$

$$\delta_{2n}(N, Z) = S_{2n}(N, Z) - S_{2n}(N + 2, Z). \quad (1b)$$

Fig. 1: Relative differences between the theoretical SO splittings $\delta_{\text{calc}}$ and the experimental ones $\delta_{\text{exp}}$ [25] in the (semi)-doubly magic nuclei indicated on the horizontal-axis. Particle and hole SO partners are shown on the left while particle-hole ones are on the right. See the text for details.

The peak values of the two-nucleon gaps are essentially determined by the sudden jump of the two-nucleon separation energies, which can be taken as a clear evidence of the magic shell occurrence.

Fig. 2 presents the two-proton (left panels) and two-neutron (right panels) gaps for the $Z = 110 - 140$ even-even isotopes calculated with the selected effective Lagrangians. We have adopted the presentation of Ref. [7] so that the similarities and differences in the predictions of the earlier study can be more easily seen. The red-solid lines stand for the two-proton drip lines defined as the change in sign of the two-proton separation energy. Nuclei that are stable with respect to $\beta$-decay or fission are represented with filled green stars or filled blue circles, respectively. For a given $A$ (resp. $Z$), the $\beta$-stability (resp. fission-stability) line is located at the maximum of the binding energy per nucleon, and corresponds as well to the minimum of the $Q$-value for $\beta$-decay (resp. fission) [31]. The dashed blue line represents the $\beta$-stability line given by the empirical formula $Z = A/(1.98 + 0.0155A^{2/3})$ [29]. Experimental data taken from the NUBASE2012 evaluation of nuclear properties [30], including the extrapolated SHN, are located below $Z = 118$ and are shown in Fig. 2 with empty red squares. It is observed from Fig. 2 that these nuclei coincide largely with the nuclei which are stable with respect to fission (filled blue circles), as predicted by our models, especially by PKA1. The role of deformation, although not included in our calculations, may play a significant role in stabilizing SHN [32, 33]. On the neutron-rich side, the large Coulomb barrier existing in SHN pushes further down the two-proton drip line. The effect is expected to change by a few units the position of the drip line.

In Fig. 2 the squares are filled in proportion of the gap, which varies from 1 to 5 MeV, as shown in the grey-scale index. Structures with large gaps between 3 and 5 MeV appear clearly in Fig. 2. From the comparison of the different models shown in Fig. 2, it is clear that PKA1 is the Lagrangian which predicts the larger gaps for $Z = 120, 126, 138$ and $N = 184, 258$. These numbers are thus the predicted magic numbers in neutron-rich SHN based on the PKA1-RHFB model. The other effective Lagrangians also present a fairly remarkable proton shell at $Z = 120$. In addition, $Z = 132$ for PKDD-RHB and $Z = 138$ for both RHFB (PKA1 and PKOi) and RHB (PKDD and DD-ME2) approaches are found to be possible proton magic numbers, consistent with the predictions in Ref. [10]. Concerning the neutron shells, besides $N = 184$ and 258, PKA1 also presents a fairly distinct shell structure at $N = 172$, which is also present in the predictions of the other Lagrangians. Fairly distinct shell effects at $N = 184$ and 258 are also found with the other parametrizations, except with PKO2. Remarkable shell effects are found at $N = 228$, although less pronounced compared to those at $N = 184$ and 258 predicted by PKA1. Furthermore, a neutron shell is predicted at $N = 164$ with PKO2, PKDD and DD-ME2 models, and another is predicted at $N = 198$ with RHB models (PKDD and DD-ME2).

We have checked that the neutron and proton pairing gaps are also quenched for the same proton and neutron magic numbers as those obtained in Fig. 2 for each considered Lagrangian. Combined with the two-nucleon gaps, it is found that the proton shell $Z = 120$ is predicted by PKA1 as well as by the other Lagrangians. It is also predicted by some SHF models such
as SLy6, SkI1, SkI3 and SkI4 [7], but it must be stressed that the SHF models can give different predictions for \(Z = 114\) and \(Z = 126\), see for instance Ref. [7]. \(Z = 120\) can however be considered as a fairly good candidate for proton magic number. In Ref. [7] SHF forces such as SkM* or SkP predict \(Z = 126\) as a magic number for neutron-poor isotopes. \(Z = 126\) is also predicted as a magic number by PKA1 model, but not by the other Lagrangians considered in this work which predict a weak SO splitting for high-\(j\) states.

On the other hand, the situation for the neutrons is more complex. Although \(N = 172\) and 228 magic numbers seem to be generally predicted by the selected effective Lagrangians, the corresponding shell effects are rather weak. Except for PKO2, \(N = 184\) and 258 are also generally predicted as candidates for neutron magic numbers. Let us notice that a large number of SHF models considered in Ref. [7] as well as Gogny forces [9] have also a large gap for these neutron numbers. Specifically, PKA1 can provide a better description of the nuclear shell structure than the others [13] and a better agreement on the fission stability of observed SHN (see Fig. 2), and it leads to pronounced shell effects. In fact, as indicated by SHF investigations [34] \(N = 184\) is also favored evidently to be a spherical neutron magic number and the \(N = 184\) isotones are expected to have spherical shapes. By comparing the predictions between the various models discussed here, we conclude that \(^{304}_{120}^{184}\) is a most probable doubly magic system in the SHN region, and \(^{292}_{120}^{172}\) might be another candidate with less stability.

Nevertheless, from Fig. 2 one can find distinct deviations among the models in predicting the magic numbers. \(Z = 120\) can be considered as a reliable prediction of proton magic number and \(Z = 138\) could be another candidate with more model dependence. The neutron shells \(N = 172, 184, 228\) and 258 are common to several models. Other shells, e.g., \(N = 198\), appear essentially model dependent. Among the present results,

![Figure 2: Contour plots (in MeV) for the two-proton gaps \(\delta_{2p}\) (left panels) and the two-neutron gaps \(\delta_{2n}\) (right panel) as functions of \(N\) and \(Z\). The two-nucleon gaps are obtained with PKA1, PKO2 and PKO3 parametrizations for RHFB, and PKDD and DD-ME2 parametrizations for RHB. The red-solid lines represent the two-proton drip lines. Nuclei stable with respect to \(\beta\) gaps are obtained with PKA1, PKO2 and PKO3 parametrizations for RHFB, and PKDD and DD-ME2 parametrizations for RHB. The red empty squares indicate the experimental SHN from NUSEBASE2012 [30]. See text for more details.](image)

| Force | \(\rho_0\) | \(E_p/A\) | \(K\) | \(J\) | \(M^*_S\) | \(M^*_E\) |
|-------|----------|-----------|-------|------|-----------|-----------|
| PKA1  | 0.160    | -15.83    | 229.96| 36.02| 0.547     | 0.681     |
| PKO1  | 0.152    | -16.00    | 250.24| 34.37| 0.590     | 0.746     |
| PKO2  | 0.151    | -16.03    | 249.60| 32.49| 0.603     | 0.764     |
| PKO3  | 0.153    | -16.04    | 262.47| 32.98| 0.586     | 0.742     |
| PKDD  | 0.150    | -16.27    | 262.19| 36.79| 0.571     | 0.651     |
| DD-ME2| 0.152    | -16.11    | 250.30| 32.27| 0.572     | 0.652     |
one may notice that RHB calculations (PKDD and DD-ME2) predict more shell closures than RHF, and PKO2-RHF predicts fewest candidates. To interpret such distinct deviations, Table 1 shows the bulk properties of symmetric nuclear matter determined by the present sets of Lagrangians. In general the occurrence of superheavy magic shells is closely related with both the scalar mass \( M^s \) and effective mass \( M^e_{NR} \) \(^{[21]}\), which essentially determine the strength of SO couplings and level densities, respectively. Among the present models, the effective Lagrangian PKO2 predicts the largest values of both masses, leading to relatively weak SO couplings and high level density on the average. As a result there remains little space in the spectra for the occurrence of magic shells. On the other hand, the RHB models (PKDD and DD-ME2) predict more magic shells due to the relatively small masses. In fact, as seen from Fig. 2, PKO2 also presents weaker shell effects than the others. For PKA1 the situation is different. Although it has a larger effective mass \( M^e_{NR} \) than PKDD or DD-ME2, PKA1 gives a smaller scalar mass \( M^s \) and shows stronger shell effects than the others. These may partially explain why PKA1 does not suffer from the common drawback of the CDF calculations — the so-called artificial shell closures induced by low \( M^s \) and \( M^e_{NR} \) — and it better preserves the PSS \(^{[13,35]}\).

Similarly to the situation in the stable region \(^{[13]}\), the model deviations originating from the relativistic PSS can also be found in the s.p. spectra of SHN. Taking the doubly magic SHN \(^{304}120\) as an example, Fig. 3 shows the proton (left panel) and neutron (right panel) canonical s.p. spectra provided by selected models. It is found that PKA1 provides the most evident magicity at \( Z = 120 \) and \( N = 184 \), respectively, although these shell closures are much weaker than in ordinary nuclei. For the neutron shell \( N = 184 \), it is essentially determined by the degeneracy of two pseudo-spin partners \([2h_{11/2},1j_{13/2}]\) and \([4s_{1/2},3d_{5/2}]\), respectively above and below the shell. For the latter, the PSS is well preserved in all the calculations while for the former with high angular momentum the symmetry is only weakly restored by PKA1 while seriously violated by the others, leading to the occurrence of the shell closure \( N = 198 \).

In fact, a similar emergence of shell closure can also be found in mid-heavy and heavy regions of the nuclear chart. For instance, the proton shell \( Z = 82 \) in \(^{208}\)Pb can also be interpreted as the result of the degeneracies of two pseudo-spin partners \([2f_{7/2},1h_{9/2}]\) and \([3s_{1/2},2d_{5/2}]\). In the CDF calculations (except PKA1) there is a clear gap between \( 2f_{7/2} \) and \( 1h_{9/2} \), i.e., the artificial shell closure \( Z = 92 \) \(^{[13,35]}\), which somewhat compresses the magic shell \( Z = 82 \). A similar mechanism can also be found in the formation of the sub-shell 64 due to the degeneracy of the pseudo-spin partners \([3s_{1/2},2d_{5/2}]\) and \([2d_{5/2},1g_{7/2}]\) \(^{[13,35]}\). In the CDF calculations (except PKA1) the sub-shell 64 is compressed by the violation of PSS on the partners \([2d_{5/2},1g_{7/2}]\), which induces the so-called artificial shell closure \(^{[13,35]}\).

Coming back to SHN, a similar phenomenon is found for \( N = 184 \) for which the gap is related to the splitting between the pseudo-spin partners \([2h_{11/2},1j_{13/2}]\) and \([4s_{1/2},3d_{5/2}]\). The restoration of PSS in PKA1 increases the gap while the breaking of the PSS among the pairs \([2h_{11/2},1j_{13/2}]\) decreases it (see the right panel of Fig. 3). From this point of view, the conservation or the breaking of PSS play a delicate role for the occurrence of shell closures in nuclei — ordinary as well as SHN. This phenomenon could be another origin for the deviations between the models. For instance it is found in general that PKA1 presents a strong SO splitting for high-\( j \) states and leads to a better PSS, compared with the other effective Lagrangians. This goes along with the general finding that SO effects help to restore the PSS \(^{[38,39]}\). It is however difficult to predict the occurrence of the PSS, which can be broken or restored in the same nucleus, for different s.p. levels.

At variance with \( N = 184 \), which is stabilized by the restoration of the PSS, the proton shell closure \( Z = 120 \) emerges from the PSS violation. As shown in the left panel of Fig. 3, the proton shell closure \( Z = 120 \) is determined directly by the large splittings of two pseudo-spin partner states, \([3p_{3/2},2f_{5/2}]\), whereas the spin doublet \( 3p \) above the shell is almost degenerate. The shell gap at \( Z = 120 \) can therefore be interpreted as a manifestation of the PSS violation and a weak SO splitting. Below the shell \( Z = 120 \), the protons filling in the high-\( j \) states will be driven towards the surface of the nucleus due to the strong centrifugal potential and large repulsive Coulomb field in SHN. Both effects lead to an interior depression of the proton distributions and consequently the interior region of the mean potential is not flat any more \(^{[9]}\). As a result the SO splitting is reduced, particularly for the low-\( j \) states \( 3p \) and \( 2f \) which have more overlap with the interior depression. Consequently the splitting between neighbouring pseudo-spin partners (i.e., \( 3p_{3/2} \) and \( 2f_{5/2} \)) is somewhat enlarged \(^{[39]}\). In Ref. \(^{[40]}\) it is also pointed out that the distinct central depressions on the densities lead to the spherical shell gaps at \( Z = 120 \) and \( N = 172 \) as a direct consequence of PSS breaking, whereas a flatter density profile favors the shell occurrence at \( N = 184 \) as well as the proton one at \( Z = 126 \). This can happen not only for SHN, and the emergence of new shell closure at \( Z = 120 \) and \( N = 16 \) and \( N = 32 \) \(^{[41,42]}\) can also be related with the violation of PSS in light exotic nuclei.

In summary, we have explored the occurrence of spherical shell closures for superheavy nuclei (SHN) and the physics therein using the relativistic Hartree-Fock-Bogoliubov (RHB) theory with density-dependent meson-nucleon couplings, and compared the predictions with those of some relativistic Hartree-Bogoliubov (RHB) models. The shell effects are quantified in terms of two-nucleon gaps \( \delta_{2nij} \). To our knowledge, this is the first attempt to perform such extensive calculations within the RHB scheme. The results indicate that the nuclide \(^{304}120\) could be the next spherically doubly magic nuclide beyond \(^{208}\)Pb. It is also found that the shell effects in SHN are sensitive to the values of both scalar mass and effective mass, which essentially determine the spin-orbit effects and level density, respectively. In addition, the breaking or restoration of relativistic pseudo-spin symmetry (PSS) is found to modify the level structure, and therefore to contribute to the emergence or disappearance of shell closure. Experimental measurement of \( Q_{zz} \), for at least one isotope of \( Z = 120 \) nucleus would help us to set a proper constraint in determining the shell effects of SHN and to test further the reliability of the models as well. One also has to admit that...
for a more extensive exploration one needs to take into account the deformation effects.

We would like to thank Haozhao Liang for enlightening discussions on PSS and its breaking. One of us (J. Li) also thank Haifei Zhang for his help in the initial stage of this study. This work is partly supported by the National Natural Science Foundation of China under Grant Nos. 11075066 and 11375076, and the Program for New Century Excellent Talents in University under Grant No. NCET-10-0466.

References

[1] M. Block, et al., Nature 463 (2010) 785.
[2] Y. T. Oganessian, et al., Phys. Rev. C 74 (2006) 044602.
[3] Y. T. Oganessian, et al., Phys. Rev. Lett 104 (2010) 142502.
[4] G. G. Adamian, N. V. Antonenko, V. V. Sargsyan, Phys. Rev. C 79 (2009) 054608.
[5] P. Möller, J. R. Nix, W. D. Myers, W. J. Swiatecki, At. Data Nucl. Data Tables 59 (1995) 185.
[6] A. Baran, Z. Lojewski, K. Sieja, M. Kowal, Phys. Rev. C 72 (2005) 044310.
[7] K. Rutz, M. Bender, T. Bürvenich, T. Schilling, P.-G. Reinhard, J. Maruhn, W. Greiner, Phys. Rev. C 56 (1997) 238.
[8] M. Bender, K. Rutz, P.-G. Reinhard, J. A. Maruhn, W. Greiner, Phys. Rev. C 60 (1999) 034304.
[9] J. Dechargé, J.-F. Berger, M. Girod, K. Dietrich, Nucl. Phys. A 716 (2003) 55.
[10] W. Zhang, J. Meng, S. Q. Zhang, H. Toki, Nucl. Phys. A 753 (2005) 106.
[11] J. Dechargé, J.-F. Berger, K. Dietrich, M. S. Weiss, Phys. Lett. B 451 (1999) 275.
[12] J. N. Ginocchio, Phys. Rep. 414 (2005) 165.
[13] W. H. Long, H. Sagawa, N. Van Giai, J. Meng, Phys. Rev. C 76 (2007) 034314.
[14] M. Bender, P.-H. Heenen, P.-G. Reinhard, Rev. Mod. Phys. 75 (2003) 121.
[15] J. R. Stone, P.-G. Reinhard, Prog. Part. Nucl. Phys 58 (2007) 587.
[16] D. Vretenar, A. V. Afanasjev, G. A. Lalazissis, P. Ring, Phys. Rep. 409 (2005) 101.
[17] J. Meng, H. Toki, S. G. Zhou, S. Q. Zhang, W. H. Long, L. S. Geng, Prog. Part. Nucl. Phys 57 (2006) 470.
[18] W. H. Long, P. Ring, N. Van Giai, J. Meng, Phys. Rev. C 81 (2010) 024308.
[19] J. F. Berger, M. Girod, D. Gogny, Nucl. Phys. A 428 (1984) 23.
[20] L. J. Wang, B. Y. Sun, J. M. Dong, W. H. Long, Phys. Rev. C 87 (2013) 054331.
[21] W. H. Long, N. Van Giai, J. Meng, Phys. Lett. B 640 (2006) 150.
[22] W. Long, H. Sagawa, J. Meng, N. Van Giai, Europhys. Lett 82 (2008) 12001.
[23] W. Long, J. Meng, N. Van Giai, S.-G. Zhou, Phys. Rev. C 69 (2004) 034319.
[24] G. A. Lalazissis, T. Nikšić, D. Vretenar, P. Ring, Phys. Rev. C 71 (2005) 24312.
[25] S.-G. Zhou, J. Meng, P. Ring, Phys. Rev. C 68 (2003) 034323.
[26] NuDat Databases, National Nuclear Data Center, http://www.nndc.bnl.gov/nudat2/.
[27] K. Rutz, M. Bender, P.-G. Reinhard, J. Maruhn, W. Greiner, Nucl. Phys. A 634 (1998) 67.
[28] E. Litvinova, P. Ring, Phys. Rev. C 73 (4) (2006) 044328.
[29] P. Mariner, E. Sheldon, Physics of Nuclei and Particles, Vol. 1, Academic Press, New York, 1971.
[30] G. Audi, F. G. Kondev, M. Wang, B. Pfeiffer, X. Sun, J. Blachot, M. MacCormick, Chinese Phys. C 36 (2012) 1157.
[31] C.-L. Wu, M. Guidry, D. H. Feng, Phys. Lett. B 387 (1996) 449.
[32] G. A. Lalazissis, M. M. Sharma, P. Ring, Y. K. Gambhir, Nucl. Phys. A 608 (1996) 202.
[33] S. K. Patra, C.-L. Wu, C. R. Praharaj, R. K. Gupta, Nucl. Phys. A 651 (1999) 117.
[34] A. T. Kruppa, M. Bender, W. Nazarewicz, P.-G. Reinhard, T. Vertse, S. Čwiok, Phys. Rev. C 61 (2000) 034313.
[35] L.-S. Geng, J. Meng, H. Toki, W.-H. Long, G. Shen, Chin. Phys. Lett 23 (2006) 1139.
[36] W. Long, T. Nakatsukasa, H. Sagawa, J. Meng, H. Nakada, Y. Zhang, Phys. Lett. B 680 (2009) 428.
[37] W. H. Long, P. Ring, J. Meng, N. Van Giai, C. A. Bertulani, Phys. Rev. C 81 (2010) 031302.
[38] H. Liang, S. Shen, P. Zhao, J. Meng, Phys. Rev. C 87 (2013) 014334.
[39] S. Shen, H. Liang, P. Zhao, S. Zhang, J. Meng, Phys. Rev. C 88 (2013) 024311.
[40] A. V. Afanasjev, S. Frauendorf, Phys. Rev. C 71 (2005) 024308.
[41] A. Ozawa, T. Kobayashi, T. Suzuki, K. Yoshida, I. Tanihata, Phys. Rev.

Figure 3: Proton (left panel) and neutron (right panel) canonical s.p. spectra of superheavy nuclei $^{304}_{120}$. The results are extracted from the RHF calculations with PKO1 series and PKA1, and compared to the RHB ones with PKDD and DD-ME2. In all cases the pairing force is derived from the finite range Gogny force D1S with the strength factor $f = 0.9$. See the text for details.
Lett. 84 (2000) 5493.

[42] R. Kanungo, I. Tanihata, A. Ozawa, Phys. Lett. B 528 (2002) 58.