Momentum distribution in heavy deformed nuclei: role of effective mass

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Abstract

The impact of nuclear deformation on the momentum distributions (MD) of occupied proton states in \textsuperscript{238}U is studied with a phenomenological Woods-Saxon (WS) shell model and the self-consistent Skyrme-Hartree-Fock (SHF) scheme. Four Skyrme parameterizations (SkT6, SkM*, SLy6, SkI3) with different effective masses are used. The calculations reveal significant deformation effects in the low-momentum domain of $K^\pi = 1/2^\pm$ states, mainly of those lying near the Fermi surface. For other states, the deformation effect on MD is rather small and may be neglected. The most remarkable result is that the very different Skyrme parameterizations and the WS potential give about identical MD. This means that the value of effective mass, being crucial for the description of the spectra, is not important for the spatial shape of the wave functions and thus for the MD. In general, it seems that, for the description of MD at $0 \leq k \leq 300$ MeV/c, one may use any single-particle scheme (phenomenological or self-consistent) fitted properly to the global ground state properties.
I. INTRODUCTION

The momentum distribution (MD) of nucleons in nuclei is a basic observable carrying important information on the single-particle aspects of nuclear structure (see [1] and references therein). In spite of intense studies, there remain several open points which deserve closer inspection as, e.g., the impact of nuclear deformation on the MD. The deformation mixes different spherical components in single-particle wave functions [2] and thus can affect the MD. The question is how strong the deformation impact is and what pattern it produces.

In light and rare-earth nuclei, the deformation effects in MD were earlier studied with the Nilsson model and the Skyrme-Hartree-Fock (SHF) approach with SIII forces [3]. Therein, it was also discussed how the MD signal can be extracted from ($e,e'p$) reaction. In the meantime, both experiment and theory have made substantial progress. Modern ($e,e'p$) experiments start to deal with heavy (actinide) deformed nuclei (see, e.g., [4]). At the theoretical side, the self-consistent SHF approach came into the focus of MD studies [1]. Thus a state-of-art theoretical analysis has to cover heavy nuclei and include the comparison between a variety of phenomenological and self-consistent MD descriptions.

In a recent letter [5], we have presented first results on MD in $^{238}$U by using the WS model and SHF with the SkM* force. The letter focused on bound $K^\pi = 1/2^+$ states with strong $l = 0$ contributions at zero momentum. It was shown that a finite deformation can enhance the number of such states. These states can be discriminated in knock-out reactions. Their observation in heavy deformed nuclei can give a valuable information on the underlying mean field, which, in turn, may be very helpful for clarifying some hot problems in nuclear structure (see, e.g., [6]). The brief study in [5] calls for a deeper analysis. First of all, the role of the effective mass in SHF models has to be clarified. The effective mass is known to have a dramatic effect on single-particle spectra [7, 8] but its influence on MD of individual states is still unclear. It is thus the aim of the this paper to investigate in detail the influence of the effective mass on the MD. We will also give a detailed analysis of general deformation effects in momentum distributions.

As test case, we will again consider the axially deformed nucleus $^{238}$U. It is a typical actinide nucleus with well known spectroscopic characteristics. Actinides are most suitable for our aims. Indeed, the heavier is the nucleus, the denser is its single-particle spectrum, and thus the better are conditions for the deformation mixing. So, actinides promise the
most strong deformation effects.

Two essentially different single-particle models will be used: i) the phenomenological Woods-Saxon (WS) potential [9] and ii) the self-consistent density-dependent Skyrme-Hartree-Fock (SHF) potential [10] with Skyrme parameterizations SkT6 [11], SkM* [12], SLy6 [13] and SkI3 [14]. These parameterizations represent different kinds of Skyrme forces and, what is important, cover a wide interval of the effective masses, from \( m^*/m = 1.00 \) in SkT6 to \( m^*/m = 0.58 \) in SkI3 (for an extensive review of self-consistent nuclear models see [15]). The smaller is the effective mass, the more stretched is the single-particle spectrum (see, e.g. discussions in Refs. [7, 8]) and so the weaker deformation mixing of spherical configurations is expected.

II. FORMAL FRAMEWORK

The calculations have been performed with the phenomenological WS and self-consistent SHF potentials.

In case of the WS potential, the deformed shape was described by the Cassini ovaloids and the potential-energy surface was calculated as a function of the elongation \( \epsilon \) and hexadecapole deformation \( \alpha_4 \) [9]. The equilibrium (ground state) deformation was found by minimizing the total energy. We used the standard set of the WS parameters [16] slightly modified for actinides [17]. The single-particle wave functions were expanded in the Nilsson basis involving 21 shells.

The SHF calculations were performed with a code using coordinate-space representation with cylindrical coordinates [18]. Four different Skyrme parameterizations were used: SkT6 [11], SkM* [12], SLy6 [13], and SkI3 [14]. Although these four forces are fitted with different bias, they all provide a good overall description of nuclear bulk properties and are equally suitable for heavy deformed nuclei. The essential feature for our study is that these four forces cover different values of the effective mass \( m^*/m \) (see Table I).

The bare G matrix theory results in \( m^*/m = 0.7 \) [21]. The same value is obtained from empirical data for the levels far beyond the Fermi energy, \( |E - E_F| \geq 20 \) MeV, in [7]. The effective masses \( m^*/m < 1 \) are known to stretch the single-particle spectra [7, 8], making them dilute as compared to the experimental data. After taking into account the correlation effects, the spectra should be more compressed and come closer to the experimental level
TABLE I: Effective masses ($m^*/m$) for Skyrme forces, quadrupole moments ($Q_2$), Fermi energies ($E_F$) and energies of the lowest ($E_0$) proton single-particle levels in $^{238}$U. Experimental estimations for the quadrupole moment in $^{238}$U lie in the interval $Q_2 = 11.1 - 11.3$ b.

| Potential | $m^*/m$ | $Q_2$[b] | $E_F$[MeV] | $E_0$[MeV] |
|-----------|---------|----------|------------|------------|
| WS        | -       | 11.66    | -6.63      | -33.69     |
| SkT6      | 1.00    | 11.10    | -6.48      | -32.75     |
| SkM*      | 0.79    | 11.11    | -6.17      | -39.80     |
| SLy6      | 0.69    | 11.06    | -7.25      | -43.12     |
| SkI3      | 0.58    | 10.89    | -7.19      | -48.53     |

It is to be noted that WS and other phenomenological potentials (Nilsson, etc) employ a "trivial" kinetic energy, i.e. $m^*/m = 1$. Thus the Skyrme forces with $m^*/m \approx 1$ should in general give spectra close to the phenomenological ones.

The pairing was treated in the BCS approximation. The SHF forces used a zero-range two-body pairing force with strengths adjusted for each parameterization separately, for details see [15]. Details of the BCS procedure in the WS case are given in [9].

Basic ground state properties for the different models are shown in Table I. The WS potential and all SHF parameterizations give a reasonable quadrupole moment and Fermi energy for $^{238}$U. At the same time, they yield different spectral stretching (defined as a difference, $|E_F - E_0|$, between the Fermi energy (chemical potential) and the energy of the lowest single-particle level). The stretching ranges from 26 to 41 MeV and, as expected, grows with decreasing the effective mass (see also Figs. 1 and 2 below). SkT6 with $m^*/m = 1$ gives an average spectral density close to the WS one.

Let’s now outline the calculation of MD. The density for the proton single-particle state...
\( \alpha \) is determined by the standard way

\[
\rho_\alpha (r) = |\psi_{\alpha \sigma}(r)|^2 = \sum_{\sigma=\pm 1} |R_\alpha^{(\sigma)}(r,z)|^2
\]  

(1)

where

\[
\psi_\alpha (r) = \sum_{\sigma=\pm 1} R_\alpha^{(\sigma)}(r,z)e^{im_\sigma^{(\sigma)}\phi}\chi_\sigma
\]  

(2)

is the single-particle wave function of the state \( \alpha \), written in cylindrical coordinates \((r, z, \phi)\).

The label \( \alpha = K^{\pi}[Nn\Lambda] \) is composed from the exact quantum numbers \( K^{\pi} \) (total angular momentum projection onto the axial symmetry axis and parity) and the asymptotic Nilsson quantum numbers \([Nn\Lambda]\). Further, \( \sigma \) is the spin and \( m^{(\pm)} = K \mp 1/2 \) is the orbital momentum projection.

In the momentum space \((k_r, k_z, k_\phi)\), the density for the state \( \alpha \) is defined as

\[
n_\alpha(k) = |\psi_\alpha(k)|^2 = \sum_{\sigma=\pm 1} |\tilde{R}_\alpha^{(\sigma)}(k_r, k_z)|^2
\]  

(3)

where \( \psi_{\alpha \sigma}(k) \) is the Fourier-transformed single-particle wave function. In the WS potential, the wave functions are expanded in the Nilsson basis whose Fourier-transformation is done analytically (for more details see, e.g. \[3\]). In SHF, the Fourier-transformed wave function reads as

\[
\psi_\alpha(k) = \frac{1}{(2\pi)^{3/2}} \int_{-\infty}^{\infty} dz \int_{0}^{\infty} dr \int_{0}^{2\pi} d\phi \cdot e^{ikr} \sum_{\sigma=\pm 1} R_\alpha^{(\sigma)}(r,z)e^{im_\sigma^{(\sigma)}\phi}
\]

\[
= \sum_{\sigma=\pm 1} \tilde{R}_\alpha^{(\sigma)}(k_r, k_z)e^{im_\sigma^{(\sigma)}k_\phi}e^{im_\sigma^{(\sigma)}k_r}
\]

(4)

where

\[
\tilde{R}_\alpha^{(\sigma)}(k_r, k_z) = \frac{1}{\sqrt{2\pi}} \int_{-\infty}^{\infty} dz \int_{0}^{\infty} dr R_\alpha^{(\sigma)}(r,z)j_{m_\sigma^{(\sigma)}}(kr) e^{ikr}
\]

(5)

and \( j_{m_\sigma^{(\sigma)}}(kr) \) is the Bessel function.

As is usually done in \((e,e')p\) calculations, we average over the nuclear symmetry axis direction:

\[
n_\alpha(k) = \frac{1}{2} \int_{0}^{\pi} d\theta \sin \theta n_\alpha(k_r, k_z)
\]  

(6)

where \( k_r = k \sin \theta \) and \( k_z = k \cos \theta \).

It is worth noting that in general single-particle models are not well suited to describe MD because of the important contributions from short- and long-range correlations \[1\].
However these perturbing effects take place mainly in the high-momentum domain with \( k > k_F \approx 1.3 \text{ fm}^{-1} \approx 260 \text{ MeV/c} \) while we will focus on MD at low \( k \), where the single-particle models are still appropriate.

III. RESULTS AND DISCUSSION

A. Single-particle levels

Before discussing the MD for single-particle states, it is instructive to have a look at the single-particle spectra. Figs. 1 and 2 compare WS and SHF single-particle proton and neutron spectra in \(^{238}\text{U}\). To make the analysis more transparent, the level schemes are presented in the spherical limit. This avoids the complexity caused by the deformation splitting of the levels and thus concentrates on the essential trends. The figures demonstrate the stretching of the single-particle spectra with decreasing the effective mass \( m^*/m \). While the Fermi energies remain basically at the similar position, the hole levels steadily dive deeper from SkT6 with \( m^*/m = 1.00 \) to SkI3 with \( m^*/m = 0.58 \). In agreement with Ref. [7], the main stretching effect takes place for the spectra far from the Fermi energy, first of all for the deeply bound levels. The spectra near the Fermi energy also show changes but not so strong and regular, thus mainly displaying the influence of other Skyrme terms, in particular of the spin-orbit force. As is expected, the SkT6 spectra with \( m^*/m = 1.00 \) are most similar to the WS ones. This is most obvious for the case of neutrons. The proton WS spectrum looks somewhat wider, which can be partly explained by the downshift of the Fermi level in the WS case. In any case, this difference is not important. It is more essential that the WS and SkT6 demonstrate very similar global spans between the Fermi and the lowest levels (see Table II).

The stretching of the spectra with decreasing \( m^*/m \) can be understood if one assumes that the system keeps the average momenta \( k_\alpha \) for single-particle states. As is shown below, this is indeed the case. Then the kinetic energies \( T_\alpha = k_\alpha^2/(2m^*) \) increase with decreasing \( m^* \), and so, to keep the same Fermi energy, the depth of the potential needs to be increased as well. Since the kinetic energy for deep hole states is smaller than for the valence states, the relative kinetic (smaller) and potential (larger) shifts result in pushing the lowest states deeper down and thus in stretching the spectrum.
Figure 3 shows the proton levels for the deformed ground state in the vicinity of the Fermi surface. We see again that the Skyrme spectra become in general more dilute with decreasing $m^*/m$. However, the trend is much weaker than for deeply lying states and concerns only the uppermost and lowest states in the plot. The intermediate states do not exhibit any clear relation to $m^*/m$, like in the spherical case. The BCS calculations of the quasiparticle spectra reveal that only the WS and SkM* reproduce the correct ground states assignment of the neighboring odd-proton nuclei $^{239}$Np ($5/2^+$) and $^{237}$Pa ($1/2^+$). This result, however, does not allow to judge on the accuracy of the single-particle schemes. As was mentioned above, the comparison with experiment for the odd nuclei requires, in principle, the inclusion of polarization effects [20, 22], in particular the coupling with core vibrations.

In that connection, it is worth noting that SHF is not always performing well with the spectra of particular nuclei (see, e.g. discussion in Ref. [8]). This can be explained (and excused) by the fact that most of the Skyrme parameterizations are tuned to optimize the basic ground state characteristics (binding energies, r.m.s. charge radii, densities, etc) instead of the spectra. Moreover, recent Skyrme forces aim to describe nuclei, both spherical and deformed, throughout the entire mass table (including those near drip-lines) as well as nuclear and neutron matters. Certainly, such universality has a price: Skyrme forces are generally not optimized to deliver optimal spectra (unlike phenomenological potentials which are usually specially fitted to low-energy spectra in particular nuclei).

B. Momentum distributions

In Fig. 4, the MD from SkM* calculated at the equilibrium shape and in the spherical limit are compared for a representative set of nine occupied proton states. Both deeply and slightly bound states are involved. The deep hole states include $1/2^-[330]$ and $1/2^- [301]$ (with the single-particle energies -26.5 and -17.2 MeV, respectively). The other seven states lie near the Fermi energy (see Fig. 3.) As is discussed below, the weakly bound states are expected to deliver the most pronounced deformation effects and thus we pay to them more attention. In the spherical limit, the number of maxima in the MD profile for the state $nlj$ is equal to the number of radial nodes $n$. Hence, the deformation effects can be easily spotted by looking at an increasing number of the maxima and/or an essential redistribution of the strength between the maxima. We present here the SkM* results though, as is discussed
below, other Skyrme parameterization might be used as well.

Figure 4 illustrates some general deformation effects in MD. First, the lower the $K$ quantum number, the stronger the deformation effect, see e.g. the $K = 1/2$ states $1/2^-[330]$ and $1/2^-[530]$ (and the state $1/2^+[660]$ in Fig. 5). This follows from the fact that spherical configurations with low $K$ have in general a denser spectrum than those with high $K$, which favors the mixing low-$K$ states due to the deformation. The exceptions (e.g., $1/2^-[301]$) mainly concern deeply bound states whose mixing is often suppressed due to a rather dilute spectrum. At the other side, the levels near the Fermi energy are more affected by the deformation because they reside in a region of higher spectral density.

Second, the deformation usually results in a shift of the MD strength to lower momenta. Note that the normalization condition $\int n_\alpha(k)k^2dk = 1$ carries a weight $k^2$ and so even a small modification of MD at high $k$ may cause considerable changes at low $k$. As a result, just the low-$k$ domain is most sensitive to deformation (see also the discussion on $K^\pi = 1/2^+$ states in [5]).

Figure 5 compares MD from the WS and SkM*. We see rather good agreement between both cases. The modest differences mainly take place in the low momentum regions where deformation effects are most strong. This result somewhat deviates from that in [5] where the deformation effects in the WS potential were overestimated because of the insufficiently accurate treating of the WS wave functions in the momentum space.

It worth noting that the calculations [3] for Ne and Nd also displayed rather modest differences between phenomenological (Nilsson) and Skyrme (SIII) momentum distributions and the deformation effect mainly in the low-momentum domain of the $K = 1/2$ states. At the same time, our calculations predict more cases of the noticeable deformation impact since we deal with the heavier nucleus $^{238}$U where the deformation mixing is generally stronger.

Altogether, one may conclude that just the $K = 1/2$ states in heavy nuclei, lying in the vicinity of the Fermi energy, are most promising for displaying the deformation effect in MD. In deeply bound states, even if they are influenced by the deformation, the momentum distributions should be considerably smeared by polarization as well as correlations, thus hiding, to a large extend, the deformation mixing. The only chance for deeply bound states to exhibit in experiment the deformation effects is offered by the $K^\pi = 1/2^+$ states, where the deformation induced $l = 0$ strength is strictly localized at $k = 0$ and so can in principle be distinguished from the $l \neq 0$ patterns [5].
Figure 6 compares MD for the four different Skyrme forces (SkT6, SkM*, SLy6, and SkI3). In spite of the much different effective masses, all the parameterizations give very similar MD. The deviations are about invisible at high momenta for all the states and at all momenta for the states with high $K$. Even for deep hole states $1/2^- [301]$ and $1/2^- [330]$, the MD are about the same. The minor deviations related to the deformation are spotted in $1/2^- [330]$ and $1/2^- [530]$. The only strong difference takes place at low $k$ in $1/2^+ [660]$. However, this case is very specific and reflects the considerable mixing of $1/2^+ [660]$ and $1/2^+ [400]$ states, which is well known in deformed nuclei. Just because of the $1/2^+ [400]$ admixture with its dominant $3s_{1/2}$ component, the state $1/2^+ [660]$ acquires a jump at $k = 0$. The mixing $1/2^+ [660]$ and $1/2^+ [400]$ levels is caused by their pseudo-crossing at the equilibrium deformation. The states at the crossing points are known to be extremely sensitive to the details of the single-particle scheme and in this sense the state $1/2^+ [660]$ is an exception from the general picture. So, we may conclude that the value of the effective mass, being crucial for the description of the spectra, turns out to be irrelevant for the momentum distributions.

The insensitivity of MD to the effective mass is the most remarkable result of our study. It means that single-particle wave functions are much less sensitive to the effective mass than the single-particle spectra. Moreover, the similarity of MD obtained with different potentials (Nilsson, WS, SHF) signifies that independence of MD to the effective mass is a signature of a more general robustness of MD. In principle, this feature is not surprising since MD are determined by the structure of the wave functions which in turn is specified by the orbital moment and number of nodes of the dominant components. All the relevant single-particle potentials evidently keep this structure in the spherical limit. In deformed nuclei, the different models should reproduce the nuclear quadrupole moment and then their eigenfunctions should have a similar composition of angular momentum components which yields, in turn, similar MD. We thus may conclude that any single-particle potential (phenomenological or self-consistent) which reproduces the basic ground state properties should accurately describe momentum distributions of individual states in the momentum domain $0 \leq k \leq 300$ MeV/c (for the exception of the cases of the level crossing).

General arguments given above are still not enough for treating so nontrivial result as the indifference of the SHF MD to the effective mass and we need here some additional comments. It would be natural to expect the similarity of MD for the different phenomenological potentials which deviate only by the potential term while the kinetic term remains to
be the same. But, in SHF forces with various $m^*$, both the potential and kinetic parts are different. And, if we get then very different SHF single-particle spectra, why not to expect also the different MD, at least for the deep hole states? Indeed, the SHF parameterizations are fitted so as to reproduce the ground state properties which are mainly determined by the nucleons from the valence shell. But deep hole states should not be so fixed by the fit and so might in principle deviate not only in spectra but also in MD. For example, the MD peaks might be somewhat shifted. However, MD of the deep hole states persist to keep their profiles at different $m^*$. It looks like the valence shell strictly determines $n_\alpha(k)$ and $n_\alpha(r)$ (and thus the single-particle wave functions) in other shells as well, and the nucleus preserves the velocity fields for all the nucleons, both valence and deep hole.

**IV. CONCLUSIONS**

The influence of the nuclear deformation on the momentum distributions (MD) of proton hole states in $^{238}_{\text{U}}$ was studied with the phenomenological WS potential and the self-consistent Skyrme-Hartree-Fock approach. Four Skyrme parameterizations (SkT6, SkM*, SLy6, and SkI3) with effective masses $0.58 \leq m^*/m \leq 1.00$ were used. Particular attention was paid to the role of the effective mass.

It was shown that the main deformation effects take place at the low-momentum domain of $K^{\pi} = 1/2^\pm$ states in the vicinity of the Fermi energy. Indeed, the lower angular momentum projection $K$ have the states, the larger is their average spectral density. Besides, the spectral density rises with approaching the Fermi energy. The high spectral density favors the deformation mixing. As a result, just $K^{\pi} = 1/2^\pm$ levels near the Fermi energy are mainly affected by the deformation.

The most striking result concerns the role of the effective mass. The calculations confirm that the effective mass strongly influences the single-particle spectrum. At the same time, different Skyrme forces with the effective masses varying trough $0.58 \leq m^*/m \leq 1$, give about identical MD. The remaining modest differences in MD are mainly connected with the deformation effects. Such a striking similarity of Skyrme MD leads to the surprising (at the first glance) conclusion that the effective mass does not influence the momentum distributions in a nucleus. The deviations in other features of the Skyrme forces also do not noticeably influence MD. Moreover, the Skyrme MD are quite similar to the WS ones. So,
for the description of MD (and the subsequent inputs for knock-out reactions) one can use any well fitted Skyrme or phenomenological potentials. We mainly explain such stability of the momentum distributions as a consequence of the fact that any single-particle potentials properly fitted to the basic ground state properties (including the nuclear shape) keep the same structure (principle components with their orbital moments and node numbers) of the wave functions.

It is interesting that, though mainly the valence shell is responsible for the ground state properties, the momentum distributions of the deep hole states also become fixed by the fit. Unlike the phenomenological potentials, SHF forces with different $m^*$ deviate not only in the potential term but in the kinetic energy as well. And this is a nontrivial and somewhat unexpected result that MD, unlike the spectra, turn out to be so stable even for the deep hole states.

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[1] A.N. Antonov, P.E. Hodgson, and I.Zh. Petkov, *Nucleon Correlations in Nuclei* (Springer-Verlag Berlin Heidelberg, 1993).

[2] F.A. Gareev, S.P. Ivanova, L.A. Malov, and V.G. Soloviev, Nucl. Phys. A171, 134 (1971).

[3] J.A. Caballero and E. Moya de Guerra, Nucl. Phys. A509, 117 (1990); E. Moya de Guerra, P. Sariguren, J.A. Caballero, M. Casas, and D.W.L. Sprung, Nucl. Phys. A529, 68 (1991).

[4] V.P. Likhachev, J. Mesa, J.D.T. Arruda-Neto, B.V. Carlson, A. Deppman, M.S. Hussein, V.O. Nesterenko, F. Garcia, and O. Rodriguez, Phys. Rev. C65, 044611 (2002).

[5] V.O. Nesterenko, V.P. Likhachev, P.-G. Reinhard, J. Mesa, W. Kleinig, J.D.T. Arruda-Neto, and A. Deppman, J. Phys. G: Nucl. Part. Phys. 29, L37 (2003).

[6] A.V. Afanasjev, T.L. Khoo, S. Frauendorf, G.A. Lalazissis, and I. Ahmad, Phys. Rev. C67,
024309 (2003).

[7] C. Mahaux, P.F. Bortignon, R.A. Broglia, and C.H. Dasso, Phys. Rep. 120, 1 (1985).
[8] B.A. Brown, Phys. Rev. C58, 220 (1998).
[9] V.V. Pashkevich, Nucl. Phys. A169, 275 (1971).
[10] T.H.R. Skyrme, Phil. Mag. 1, 1043 (1956); D. Vauterin and D.M. Brink, Phys. Rev. C5, 626 (1972).
[11] F. Tondeur, M. Brack, M. Farine, and J.M. Pearson, Nucl. Phys. A420, 297 (1984).
[12] J. Bartel, P. Quentin, M. Brack, C. Guet, and H.-B. Haakansson, Nucl. Phys. A386, 79 (1982).
[13] E. Chabanat, P. Bonche, P. Haensel, J. Meyer, and R. Schaeffer, Nucl. Phys. A643, 441(E) (1998).
[14] P.-G. Reinhard and H. Flocard, Nucl. Phys. A584, 467 (1995).
[15] M. Bender, P.-H. Heenen, and P.-G. Reinhard, Rev. Mod. Phys. 75, 121 (2003).
[16] V.A. Chepurnov, Sov. J. Nucl. Phys. 6, 696 (1968).
[17] F. Garcia, E. Garrote, M.-.L. Yoneama, J.D.T. Arruda-Neto, J. Mesa, F. Bringas, J.F. Dias, V.P. Likhachev, O. Rodriguez, and F. Guzmán, Eur. Phys. J. A6, 49 (1999).
[18] P.-G. Reinhard, unpublished.
[19] E.N. Shurshikov, Nucl. Data Sheets, 53, 601 (1988).
[20] V.G. Soloviev, Theory of Complex Nuclei (Pergamon Press, 1976).
[21] V. Bernard and N. Van Giai, Nucl. Phys. A348, 75 (1980).
[22] K. Rutz, M. Bender, P.-G. Reinhard, J.A. Maruhn, and W. Greiner, Nucl. Phys. A634, 67 (1998)
FIGURE CAPTIONS

**Figure 1**: WS and SHF proton single-particle spectra in $^{238}$U at zero deformation (spherical limit). The effective mass decreases from $m^*/m = 1.00$ in SkT6 to $m^*/m = 0.58$ in SkI3. The levels of the positive and negative parity are depicted by the solid and dotted lines, respectively. For the view convenience, the identical levels are connected by dashed lines. The chemical potentials are indicated by dotted lines with crosses.

**Figure 2**: The same as in Fig. 1 for the neutron spectra.

**Figure 3**: WS and SHF proton single-particle spectra near the Fermi energy in $^{238}$U. The effective mass decreases from $m^*/m = 1.00$ for SkT6 to $m^*/m = 0.58$ for SkI3. The level energies are given relative to the Fermi level, $1/2^+[400]$ in WS and $3/2^+[651]$ in Skyrme potentials.

**Figure 4**: Momentum distributions for nine occupied proton states in $^{238}$U calculated with SkM* in the spherical limit (solid line) and at the equilibrium deformations (dashed line). The spherical ancestors are indicated for every state.

**Figure 5**: SkM* (solid line) and WS (dashed line) proton momentum distributions in $^{238}$U.

**Figure 6**: Proton momentum distributions in $^{238}$U calculated with Skyrme potentials SkT6 (dashed line), SkM* (solid line), SLy6 (dotted line), and SkI3 (dashed-dotted line).
$10^9 n_\alpha(k) [(\text{MeV}/c)^{-3}]$

- $1/2^+ [301]$
- $2p_{1/2}$
- $1/2^+ [400]$
- $3s_{1/2}$
- $1/2^+ [660]$
- $1i_{13/2}$
- $1/2^- [330]$
- $1f_{7/2}$
- $3/2^+ [402]$
- $2d_{3/2}$
- $3/2^+ [651]$
- $1i_{13/2}$
- $5/2^+ [530]$
- $2f_{7/2}$
- $1/2^- [505]$
- $1h_{11/2}$

$k (\text{MeV}/c)$
$1/2^-[301]$

$1/2^-[530]$

$3/2^+[651]$

$1/2^+[660]$

$11/2^- [505]$

$10^9 n_{\alpha}(k) [(\text{MeV/c})^{-3}]$

$k (\text{MeV/c})$
\[ 1/2^- \, [301] \quad \text{SkT6} \quad \text{SkM}^* \quad \text{SLy6} \quad \text{SkI3} \]

\[ 1/2^- \, [330] \]

\[ 1/2^- \, [530] \]

\[ 3/2^+ \, [651] \]

\[ 1/2^+ \, [660] \]

\[ 11/2^- \, [505] \]