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Preface

The Lysekil Conference in 1966 on "Why and how should we investigate nuclides far off the stability line" came at a time when many laboratories had begun to take an active interest in this field. The conference contributed, to a very considerable extent, to the rapid development of the subject. The ISOLDE Collaboration (CERN) has found the time ripe for a "Lysekil II" conference, this time with the emphasis on results obtained more than on experimental techniques. In particular we found it interesting to help to clarify where the most important perspectives show up. The conference on "The properties of nuclei far from the region of beta-stability" was held at Leysin, Switzerland, 31 August - 4 September 1970. The programme was arranged through discussions with members of the following committee:

I. Bergström, Stockholm  R. Bernas, Oreay
G.N. Flerov, Dubna  W. Gentner, Heidelberg
P.G. Hansen, Aarhus/CERN  A. Kjelberg, CERN (Secretary)
R.D. Macfarlane, Texas A & M/Copenhagen  S.G. Nilsson, Lund
G. Rudstam, Studsvik (Chairman)  M. Veneroni, Oreay

The contributions to the conference, in some cases in the form of abstracts, have been collected into these Proceedings. They are un-edited and printed by a photographic reproduction technique directly from typescripts supplied by the authors. We chose this method in order to make the Proceedings available within as short a delay as possible. It is our hope that these Proceedings will prove to be a useful source of information and ideas.

We acknowledge the most valuable help of the CERN Conference Secretariat (Miss E.W.D. Steel and Miss Y. Henry) in preparing the conference, the efficient and good-humoured assistance, during and after the conference days, given by Mrs. R. Mohr, the perfect technical assistance of Mr. P. Gilbert de Vaulthibault during the conference, and the extensive efforts of the CERN Documents Reproduction Section in preparing the Proceedings.

The conference was sponsored and financially supported by the ISOLDE Collaboration.

G. Rudstam
for the Organizing Committee
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THEORETICAL PREDICTIONS FOR SUPERHEAVY NUCLEI *)

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1. INTRODUCTION

Today we are going to enlarge the scope of the Conference to include nuclei beyond our current periodic table, even though some of these nuclei are in the region of beta-stability. The purpose of my talk is to review our current expectations concerning superheavy nuclei, as well as to provide an introduction for the remaining talks on this subject. The main emphasis is on the expected stability of superheavy nuclei against their various modes of decay, namely spontaneous fission, alpha decay, and beta decay. For this purpose special attention is devoted to calculating the potential energy of a superheavy nucleus as a function of its shape. I will also discuss briefly methods for producing superheavy nuclei, and finally, their expected chemical and physical properties.

With the aid of Fig. 1, I would like to begin with a few qualitative remarks. One finds in nature some 300 nuclei, representing isotopes of elements containing from 1 to at most 92 protons. These naturally occurring nuclei are shown here by heavy points. Some 1200 additional nuclei have been made artificially in the past 30 years and are shown by lighter-weight points. Element 105 is the heaviest element produced to date.

This peninsula of known nuclei terminates because of nuclear fission. As we move along the peninsula toward heavier nuclei, the disruptive Coulomb forces grow faster than the cohesive nuclear forces. The Coulomb forces finally begin to overwhelm the nuclear forces, and spontaneous fission becomes a rapid mode of decay. In addition, the large Coulomb forces also cause heavy nuclei to decay rapidly by the emission of alpha particles.

*) This review was written under the auspices of the U. S. Atomic Energy Commission.
If decay by spontaneous fission and alpha emission becomes so likely for heavy nuclei, why do we expect to find an island of still heavier nuclei past the end of the peninsula? The answer lies in the extra stability arising from the closing of a proton or neutron shell. As we increase the particle number of nuclei along the peninsula, successive protons and neutrons go into definite single-particle orbits. When a given shell of protons or neutrons is completely filled, that nucleus has relatively lower energy, i.e. extra binding and hence increased stability. Elements corresponding to the closing of proton shells are indicated in the figure.

How does this extra binding due to shell closures influence a nucleus' stability? It turns out that to a fairly good approximation the

*) Heavy points are used for naturally occurring nuclei, medium-weight points for nuclei with half lives greater than 1 yr, and light-weight points for nuclei with half lives less than 1 yr.
height of the potential barrier against fission is increased just by the amount of the additional binding. This is empirically about 5 MeV for the closing of a single shell, or 10 MeV for the closing of both a proton and a neutron shell, such as occurs at $^{208}$Pb.

The next shell closures beyond lead are predicted to occur at $^{114}$ protons and $^{184}$ neutrons, corresponding to a mass number of 298. In the absence of single-particle effects, the fission barrier of such a super-heavy nucleus would be vanishingly small. But this closing of two shells leads to a fission barrier that is some 10 MeV high. Now, this is a substantial fission barrier; for comparison, the barriers of conventional nuclei like uranium are only about 5 MeV high! Thus such a nucleus should be relatively stable against spontaneous fission, even though nuclei between the end of the peninsula and the island are not.

The extra binding also increases the nucleus' stability with respect to the emission of alpha particles, but only for the closed-shell nucleus and nuclei lighter than it. The alpha-emission probability is actually increased for heavier nuclei decaying toward a closed shell; this is responsible for the absence of naturally occurring nuclei just heavier than lead. We therefore expect nuclei in the center and in the lower left-hand part of the island to be relatively stable with respect to the emission of alpha particles, but not nuclei in the upper right-hand part of the island.

Finally, from the trend of the peninsula you would expect it to pass almost precisely through the center of the island. This means that at least some of these nuclei should also be stable with respect to beta decay.

To summarize, there are simple, qualitative arguments for expecting nuclei in the region of $^{298}$114 to be relatively stable with respect to spontaneous fission, alpha decay, and beta decay, which are the three relevant modes of decay to be considered.

For those of you who have not seen Swiatecki's allegory $^1$, I will show it in Fig. 2. The closed proton and neutron shells are indicated by the lines; however, this figure was made some time ago, and you should now disregard the line at 196 neutrons.
I would like to conclude the introduction with a few historical comments. The possibility of superheavy nuclei has been considered for many years\(^2\)-\(^6\)), but the subject began to receive widespread interest five years ago as a result of two independent developments. The first of these was the estimate by Myers and Swiatecki\(^7\) that the fission barrier of a superheavy nucleus should be several MeV high, and the second was the result by Meldner and Röper\(^7,8\)) that the next closed proton shell after 82 is probably 114. It had always been thought before, in analogy with the case for neutrons, that 126 would be the next closed proton shell. The Coulomb force, which is becoming increasingly important for heavier nuclei, is responsible for this difference. Additional evidence that 114 is probably the next closed proton shell was provided soon by several other calculations\(^9\)-\(^13\)).

Shortly after the two initial developments, Strutinsky\(^14,15\)) developed an improved method for calculating the potential energy of a nucleus as a function of its shape, and he and his coworkers\(^13,15-20\)) used the method to calculate the fission barriers of several superheavy nuclei. Then, Nilsson and his group\(^21-28\)) applied Strutinsky's method to a generalized harmonic-oscillator single-particle potential to make the first systematic survey of the expected stability of superheavy nuclei.
More recently, several other groups have begun to consider this problem. From this point on, I will not try to describe the various contributions in historical order, but instead will use an order that best illustrates the points I am trying to make.

2. **NUCLEAR STABILITY**

2.1 Potential energy of deformation

As we have seen, one of the possible modes of decay of a superheavy nucleus is by spontaneous fission. During fission, the shape of the nucleus changes in a way that is illustrated in Fig. 3. To determine the stability of the nucleus against fission, it is necessary to calculate its potential energy as a function of deformations such as these.

Figure 4 shows the general dependence of the potential energy on deformation for a superheavy nucleus like $^{298}$Th. An expanded scale is used to show the large decrease in energy (about 310 MeV) as the fragments separate to infinity. The zero of potential energy is the energy of the individual nucleons at infinity. This is to emphasize that the ground-state energy of about -2130 MeV, which represents the total nuclear binding energy, is an important quantity that a proper calculation of the fission barrier should reproduce, as well as the energy release and the details of the barrier near the ground state.

What is the best approach for calculating this curve? Should we not start with a fundamental nucleon-nucleon interaction derived from scattering data and solve the appropriate many-body equations in some approximation, for example the Hartree-Fock approximation? Such a procedure is very basic in principle, but in practice it has difficulties even describing such fundamental properties as the total binding energy and radii of spherical nuclei. For example, Bassichis and Kerman find with a Tabakin interaction that the ground-state energy of $^{298}$Th is only -269 MeV rather than the expected value of about -2130 MeV. Similarly, for $^{208}$Pb the calculated ground-state energy is -395 MeV rather than the experimental value of -1636 MeV.

Although the results of Bassichis and Kerman are not directly applicable to the calculation of fission barriers, they have nevertheless
raised an important question: Does the next closed proton shell after 82 really occur at 114, or instead could it possibly be at 120, 124, or 126? Bassichis and Kerman in fact find for the Tabakin interaction that the next closed proton shell is at 120. This occurs because the resulting spin-orbit force is weaker than in other calculations, which pushes up the $1l_{3/2}$ and $2f_{7/2}$ levels and brings down the $2f_{5/2}$ level, thereby increasing the proton shell from 114 to $114 + 6 = 120$. However, because
this calculation reproduces other important quantities so poorly, the indication of a proton shell at 120 should not be regarded as a concrete prediction, but instead as an important question that has been raised.

Other selfconsistent microscopic approaches are possible, for example using a simple effective nucleon-nucleon interaction whose parameters are adjusted to reproduce gross nuclear properties rather than two-nucleon scattering data. Such approaches have been used for example by
Meldner\textsuperscript{10,30} and by Köhler\textsuperscript{31,32} and are very promising, but thus far have been applied only to spherical nuclei. Both these selfconsistent calculations indicate that the next proton shell is at 114.

In lieu of any selfconsistent calculation of this curve, several attempts\textsuperscript{11,33-37} have been made to calculate it non-selfconsistently but still microscopically. However, as emphasized by Bassichis and Wilets\textsuperscript{38}, there are fundamental objections to such direct non-selfconsistent microscopic approaches. In addition, these methods in practice have failed to reproduce such gross properties as the total binding energy and energy release, as well as other smooth trends. On the other hand, they are able to reproduce local fluctuations, and this feature of the non-selfconsistent microscopic model is very useful when combined with a macroscopic model to form a two-part approach.

Macroscopic approach: smooth trends

Microscopic approach: local fluctuations

Synthesis: both smooth trends and local fluctuations

\[ V_{\text{total}} = V_{\text{liquid drop}} + V_{\text{correction}} \]

Fig. 5

Two-part approach used for calculating
the nuclear potential energy of deformation

A macroscopic model, for example the liquid-drop model, is capable of reproducing the entire curve of Fig. 4 to within an accuracy of about 10 MeV. This is 10 MeV out of a total energy of some 2000 MeV, or a relative error of only half a per cent. But of course when one is interested in energies like 8 MeV barrier heights, an error of 10 MeV is catastrophic! Since a macroscopic model reproduces correctly the smooth trends and a microscopic model the local fluctuations, why not synthesize the two, as illustrated in Fig. 5? The combined approach then hopefully reproduces both the smooth trends and the local fluctuations. According to this
method, the total nuclear potential energy of deformation is given by the sum of two terms: a smoothly varying liquid-drop part and a fluctuating single-particle correction term. The latter term may be as large as 10 MeV for a spherical doubly closed-shell nucleus, but is typically a few MeV in magnitude.

Such a two-part approach was first proposed in 1963 by Swiatecki\textsuperscript{39}, although the importance of fluctuating single-particle effects on the potential energy had been recognized before\textsuperscript{6,40-42}. However, Swiatecki's procedure for calculating the correction term works only for small deformations. A more general method that applies to arbitrarily large deformations was developed by Strutinsky\textsuperscript{14-17} in 1966, and has since been widely used. Strutinsky's method seems to work well in practice, but there are difficulties justifying it from basic principles. Some aspects of the foundations of the method have been studied by Tyapin\textsuperscript{43}. Alternate ways to calculate the shell-correction term have been tried by Gursky\textsuperscript{44,45}, by Bolsterli\textsuperscript{46}, by Myers and Swiatecki\textsuperscript{47,48}, and by Tsang\textsuperscript{49}, but a more satisfactory method has not yet been found.

Lin\textsuperscript{50} has recently questioned the validity of Strutinsky's method when used with single-particle potentials that approach zero at large distances (for example Woods-Saxon potentials), but the objections are far from conclusive.

Rather than go into the details of Strutinsky's method, I would like now to present some fission barriers that have been calculated using it. After some examples, I will then return to the important question of how calculations within this framework can be improved.

Figure 6 shows fission barriers for several isotopes of element 114; these were calculated by Nilsson and Tsang et al.\textsuperscript{28} using a generalized harmonic oscillator to represent the potential that the nucleons move in. The deformation coordinate represents primarily a spheroidal deformation, but also includes a small contraction of the nucleus at its center. The difference between the solid and dashed curves is unimportant for our purposes here. We see that the fission barrier for the doubly closed-shell nucleus with mass number 298 is 9 MeV high, and that the barrier height decreases dramatically as neutrons are removed from the nucleus. Figure 7 shows a similar sequence of fission barriers calculated
by Strutinsky and his group\textsuperscript{18}, again using a generalized harmonic-oscillator potential. In this calculation the barrier for \textsuperscript{298} \text{Li} is 10 MeV high.

How can we improve the calculation of fission barriers such as these? Within the framework of the two-part method, improvements can be made in either the macroscopic part by refining the liquid-drop model or the microscopic part by refining the single-particle model. Substantial work has been done in both directions.
For discussing the smooth trends of nuclei, the natural choice is a statistical method of calculation. One such method is the Thomas-Fermi approximation, which has been studied recently by Bethe\textsuperscript{51}, by Brueckner and coworkers\textsuperscript{52-56}, by Myers and Swiatecki\textsuperscript{57}, and by Moszkowski\textsuperscript{58}. (For earlier studies of the Thomas-Fermi approximation in nuclei see Refs. 8-18 of Ref. 57.) In this way a large number of average nuclear properties are obtained by means of relatively simple statistical methods and phenomenological nucleon-nucleon forces.
Another approach to the study of the smooth trends is a systematic refinement of the liquid-drop model. For all but the lightest nuclei, the range of the nuclear force and hence the surface diffuseness are small in comparison with the nuclear radius. This allows nuclear properties to be calculated by expanding in powers of the small dimensionless ratio (range of force)/(nuclear radius), which is proportional to $A^{-1/3}$. Stopping the expansions at order $A^{2/3}$ leads to the liquid-drop model.

Recently the expansions have been carried to one higher order in $A^{-1/3}$ by Myers and Swiatecki\textsuperscript{57,58}, who obtain in this way a few simple algebraic expressions that describe the next level of corrections after the liquid-drop model, for example curvature and compressibility effects. Once the unknown coefficients have been determined, these algebraic expressions give results very similar to Thomas-Fermi calculations, because the latter are being applied to systems with thin surfaces, where the expansions used to derive the algebraic expressions are valid. One would therefore expect the algebraic expressions to ultimately replace the more complicated Thomas-Fermi calculations. Some predictions of this refined liquid-drop model—referred to as the droplet model—have been compared with experimental data by Myers\textsuperscript{60,61}, and fission saddle-point properties have been calculated within the model by Hasse\textsuperscript{62}.

In the second area for improving the calculation of fission barriers, a large number of studies have been directed toward refining the single-particle potential felt by a nucleon.

Nilsson’s group\textsuperscript{11,21-28,63-68} and Van Rij and Hess\textsuperscript{67} have concentrated on further modifications and generalizations of the harmonic-oscillator potential. However, objections have been raised that the ordinary harmonic-oscillator potential cannot be used for shapes close to or following scission. This has led to the study of a potential familiar from molecular physics and textbooks\textsuperscript{68}, namely a two-centered harmonic-oscillator potential, by Gursky\textsuperscript{69}, by Mosel and Greiner’s group\textsuperscript{35-37,70,71}, and by Wong\textsuperscript{72}, and to potentials with similar features by Johansson\textsuperscript{73} and by Dickmann and Dietrich\textsuperscript{74}. Infinite square-well potentials appropriate to very deformed shapes have been studied by Brandt and Kelson\textsuperscript{75} and by Gaudin and Sajot\textsuperscript{76}. All these potentials are useful for certain purposes, but the basic deficiency remains that they are infinite rather than zero at large distances.
Potentials that do go to zero at large distances and that have other expected properties, such as diffuse surfaces and invariant spin-orbit terms (for example generalized Woods-Saxon potentials), have recently received the attention of several groups\textsuperscript{19,20,77-87}. These more realistic potentials are just beginning to be used for large deformations, but future predictions should be based on them to a larger extent.

As a specific example of a diffuse-surface potential appropriate to large deformations, I would like to describe the approach that Bolsterli, Fiset, and I\textsuperscript{83,84} are taking, even though our results are still preliminary.

\[ V_1(\mathbf{r}) = \int f(\mathbf{r} - \mathbf{r}') \rho(\mathbf{r}') \, d^3\mathbf{r}' \]  

(Eq. 1)

Equation (1) shows how we generate our single-particle potential. The spin-independent part of the potential \( V_1(\mathbf{r}) \) is obtained by folding an effective two-nucleon interaction \( f(\mathbf{r} - \mathbf{r}') \) with a uniform sharp-surface pseudodensity \( \rho(\mathbf{r}') \). For both physical reasons and simplicity we take the effective two-nucleon interaction to be a Yukawa function with adjustable strength and range. By uniform sharp-surface pseudodensity we mean that \( \rho(\mathbf{r}') \) is a constant inside the surface of the shape and zero outside. In addition to this spin-independent part, the total potential felt by a nucleon includes an invariant spin-orbit term and a Coulomb potential for protons.

Figure 8 shows that for a spherical shape the folded Yukawa potential that we generate, given by the solid curve, is very close to a Woods-Saxon potential, given by the dashed curve. For a spherical shape either one of these potentials would be about equally satisfactory, but for deformed shapes the folded Yukawa potential generalizes in a more natural way than does the Woods-Saxon potential.

The method used for describing deformed shapes\textsuperscript{83} is shown in Fig. 9. The nuclear surface is specified in terms of smoothly joined portions of three quadratic surfaces of revolution, for example two
Fig. 8
Comparison of folded Yukawa potential with Woods-Saxon potential

Fig. 9
Nuclear shape described in terms of two spheroids connected by a hyperboloidal neck
spheroids connected by a hyperboloidal neck. The shape is changed by varying the positions of the centers and the major and minor axes of the three quadratic surfaces.

In Fig. 10 we show on the left-hand side some of the many shapes of interest in fission, and on the right-hand side the resulting spin-independent single-particle potentials obtained by our folding procedure. The equipotential contours represent energies that are 10%, 30%, 50%, 70%, and 90% of the well depth. For the bottom dumbbell-like shape we note that the potential is somewhat deeper in the ends of the dumbbell than in the neck region. This is what is expected physically because the ends of

![Diagram of spheroids and equipotentials](image)

**Fig. 10**

Single-particle equipotentials for various nuclear shapes
the dumbbell are surrounded by more nucleons than is the neck. Also, we see that the potential is very well behaved in the entire surface region.

We obtained a preliminary set of parameters describing this potential by calculating the single-particle energies for the four doubly closed-shell nuclei $^{208}$Pb, $^{48}$Ca, $^{40}$Ca, and $^{16}$O, and adjusting the parameters to optimally reproduce the experimental single-particle energies.

These parameters were then used to calculate the single-particle levels for the superheavy nucleus $^{298}$114, as shown in Fig. 11. The Fermi

![Diagram of single-particle energies for $^{298}$114]

**Fig. 11**

Single-particle energies calculated for a folded Yukawa potential with a preliminary set of parameters
surfaces for neutrons and protons are indicated by the numbers 184 and 114, respectively. These gaps in the spectrum lead to extra binding for nuclei with approximately 184 neutrons and 114 protons. Nuclei with approximately 124 or 126 protons also have some additional binding, owing to two moderate gaps separated by only two particles in the \(3p_\frac{7}{2}\) state. Most harmonic-oscillator potentials predict a large gap also at neutron number 196, but these results indicate little additional binding for 196 neutrons.

![Graph showing the fission barrier for 298\[114\]](image)

**Fig. 12**
Fission barrier calculated for a folded Yukawa potential with a preliminary set of parameters
We have also calculated the single-particle levels for $^{298}\text{Il}$ for deformed shapes and then used the levels to compute a preliminary fission barrier for $^{298}\text{Il}$, which is shown by the solid curve in Fig. 12. The deformation coordinate $\gamma$, which corresponds to the shapes of Fig. 10, is defined in terms of the saddle-point shapes of an idealized liquid drop. This barrier has qualitatively the same appearance as the barriers of Figs. 6 and 7 calculated with harmonic-oscillator potentials, but is slightly lower. The dashed curve is the liquid-drop contribution to the barrier and shows that in the absence of single-particle effects the barrier would be vanishingly small.

The single-particle levels calculated for $^{298}\text{Il}$ were also used to compute the fission barriers of neighboring superheavy nuclei. However, we found that it is not a very good approximation to neglect the shifts in the levels arising from changes in particle numbers, particularly protons. Consequently these possibly misleading results will not be shown here. We will instead recompute the levels for the individual nuclei and use these levels for calculating the barriers. As protons are added, the increase in the Coulomb potential causes the proton Fermi level to rise, leading ultimately to proton-unstable nuclei.

As emphasized by Muzychka, Wong, and Rost, the barrier for a superheavy nucleus depends strongly upon the values of the parameters used. In most of the studies performed thus far, the single-particle parameters have been adjusted to best reproduce single-particle levels for either spherical or deformed ground-state nuclei. For the harmonic-oscillator potentials studied by Seeger and by Nilsson and Tsang et al. this procedure also accounts fairly well for the ground-state masses and deformations of nuclei, as well as some of the properties of fission barriers for actinide nuclei. The corresponding studies have not yet been made for diffuse-surface potentials, but it is known that static potentials that best reproduce experimental single-particle levels of light nuclei are not consistent with the measured radii of these nuclei. A crucial outstanding problem is the determination of the parameters that best describe in a consistent way throughout the periodic table a wide variety of phenomena such as nuclear radii, ground-state masses and deformations, and fission-barrier
properties. Then the extrapolation to the region of superheavy nuclei can be made with more confidence than at present.

2.2 Lifetimes

Spontaneous fission is a tunneling process through potential barriers such as those in Figs. 6, 7, and 12. However, the stability of a nucleus with respect to spontaneous fission is determined not only by its potential barrier, but also by the inertia, or effective mass, associated with the deformations involved. As stressed by Swiatecki\textsuperscript{39}, it should be possible to calculate inertias by means of a macroscopic approach, but this has not yet been tried. Semi-empirical estimates yield values for average inertias that are about 6 times the irrotational values\textsuperscript{22-28,100}.

Inertias have also been calculated microscopically by means of the cranking formalism by Szymański's group\textsuperscript{101} and by Strutinsky's group\textsuperscript{102}. The validity of the adiabatic approximation employed in the cranking formalism has been questioned recently by Griffin\textsuperscript{103}, but the objections are inconclusive. When a cranking formalism is used, the inertias are found to vary strongly from nucleus to nucleus and also as a function of deformation. However, these variations have not yet been included in calculating spontaneous-fission lifetimes. Instead, the inertias have been assumed to be strictly constant. Then, the lifetimes are calculated by use of the WKB approximation for the penetrability through the barrier. In practice the barriers have always been assumed to be one-dimensional at this stage, but the penetration of two-dimensional barriers is an important question that is finally receiving some attention\textsuperscript{104,105}.

The alpha-decay lifetimes for the superheavy nuclei can be estimated\textsuperscript{24-27,32} from the energy released in the process, which is determined from the calculated ground-state masses. Similarly, the ground-state masses determine which nuclei will be beta stable\textsuperscript{24-27}.

Figure 13 shows an example of such lifetime calculations by Tsang and Nilsson\textsuperscript{26} for their generalized harmonic-oscillator potential. The solid curves are contours of constant spontaneous-fission half life, plotted versus proton number and neutron number. As we move away from the doubly closed-shell nucleus with 114 protons and 184 neutrons, the calculated spontaneous-fission half lives decrease from $10^{13}$ yr for
nuclei along the highest contour to 1 nsec for nuclei along the lowest contour.

The dashed curves represent similar contours of constant alpha-decay half lives. As we increase the number of protons the alpha-decay half lives decrease from $10^{13}$ yr for nuclei along the highest contour to 1 μsec for nuclei along the lowest contour.

The darkened squares indicate nuclei calculated to be beta stable. When one considers all three modes of decay, the nucleus $^{294}110$ has the longest total half life, which is $10^8$ yr.

The calculations shown here are for nuclei with even numbers of protons and neutrons. However, as emphasized by Meldner and Herrmann $^{108}$,
certain odd-mass nuclei will possibly be even more stable than their neighboring even nuclei, owing to the additional stability associated with an unpaired nucleon.

3. METHODS OF PRODUCTION

By this point we should all be convinced that some superheavy nuclei are going to be relatively stable if they can be formed in their ground states. Methods by which we can possibly produce superheavy nuclei is the second part of my talk. However, I won't have time to discuss any of the attempts that are being made to produce them or to find them in nature; this will be the subject of three talks later this morning\textsuperscript{107-109}. No conclusive evidence for their existence has yet resulted from these attempts.

![Graph](https://example.com/graph.png)

**Fig. 14**

Multiple-neutron-capture processes for producing superheavy nuclei

There are two general methods by which superheavy nuclei may conceivably be formed: multiple neutron capture and heavy-ion reactions.
Figure 14 illustrates the former method, for which there are two variations. In the so-called \( r \) process\(^{10-113} \), a given nucleus increases its mass by capturing one or more neutrons, then increases its proton number by emitting a beta particle, and so on, as indicated by the zigzagged line. Many of the nuclei occurring in nature were made in supernovas by this process. In a nuclear explosion\(^{112,114} \) everything occurs so quickly that a nucleus does not have time to beta decay and then capture more neutrons. Instead, the nucleus must initially capture sufficient neutrons so that subsequent beta decays will take it to the region of superheavy nuclei. Unfortunately, both these variations are fraught with the same basic limitation: heavy neutron-rich nuclei terminate the process by undergoing fission. Although this conclusion is not definitely established (see for example the opposite conclusions reached in Refs. 115 and 116), I am personally convinced that superheavy nuclei cannot be made by either the \( r \) process in nature or artificially by nuclear explosions.

This, then, leaves heavy-ion reactions as the primary hope. But for such reactions to be successful there are two important difficulties that must be overcome; these are illustrated in Fig. 15. The first difficulty is associated with the relatively small distortion required to initiate fission in a superheavy nucleus\(^{99} \). The fission barrier for a conventional nucleus like plutonium is relatively wide. This nucleus does not undergo fission until it is deformed past the second peak, corresponding to a very elongated saddle-point shape. Compare this with the situation for a superheavy nucleus, whose fission barrier is much thinner. For a superheavy nucleus to undergo fission it must be deformed only past the first peak, corresponding to a saddle-point shape of very small eccentricity. To produce a superheavy nucleus in a heavy-ion reaction, the system must somehow rearrange itself from an initial configuration approximating two spheres in contact to a final shape less deformed than the saddle-point shape. Then the system must get rid of its excitation energy while its shape is nearly spherical.

The second difficulty is the effect of the excitation energy on the potential barrier itself. The lower solid curve is the barrier for a superheavy nucleus in its ground state. Its height is large because
protons and neutrons are occupying completely closed shells. But at high
excitation energies, many particles are elevated to higher states. This
reduces the effects of the closed shells, which in turn reduces the bar-
rier height. Both these difficulties are serious, but hopefully we can
design our experiments to minimize the problems.

There are several ways to classify the different types of heavy-ion
reactions for producing superheavy nuclei. I have decided to classify
them according to the reaction mechanism involved, for which there are
three possibilities.

Figure 16 illustrates the first of these possibilities, namely
compound-nucleus formation. A target and projectile are brought together
to form an excited compound nucleus. In most cases this nucleus will
undergo fission, but with some probability it will de-excite by the

*) The two fission barriers in this figure were calculated by Nilsson
and Tsang et al. \cite{25}.\)
emission of neutrons. Provided that the nucleus is approximately spheri-
cal when this occurs, the superheavy nucleus can be caught inside its
barrier peak, and then decay by gamma emission to its ground state.

In trying to reach the island by this method, there is a choice of
using either a medium-weight target and projectile, or a heavy target and
light projectile. With either choice, it is impossible to reach the
center of the island, so that the final neutron-proton ratio is poor for
this type of reaction mechanism. However, there are many combinations of
targets and projectiles that can be used to reach nuclei in the vicinity
of $^{208}_{124}$. These nuclei will have short alpha- and beta-decay lifetimes,
but it may nevertheless be possible to detect and study them. Alternative-
ly, provided the spontaneous-fission lifetimes are sufficiently long, nu-
clei near $^{208}_{124}$ could possibly be produced first, and then successively
alpha and beta decay to more stable nuclei near the center of the island.
Examples of such overshoot reactions have been discussed by Muzychka$^{30}$,
by Tsang and Nilsson$^{28}$, and by Seaborg$^{117}$. 
It turns out that the excitation energy of the compound nucleus can be less if it is formed from a medium-weight target and projectile rather than from a heavy target and light projectile. From this point of view we would therefore prefer a medium-weight target and projectile, where fewer neutrons need be emitted to de-excite the compound nucleus. The primary disadvantage of this choice is that larger distortions are involved. By this we mean that an initial configuration of two approximately equal spheres is farther away from the final single sphere than is an initial configuration consisting of one large and one small sphere.

Figure 17 illustrates the second type of reaction mechanism, namely a direct transfer reaction. The target and projectile do not completely fuse, but instead a light particle is released during the original interaction. This light particle can carry away some energy that would otherwise remain as undesirable excitation energy. In one of the two examples shown here, release of the fairly heavy particle $^{40}$Ca produces a nucleus near the center of the island. In the other example, release
of the light particle $^9$Be produces a nucleus with 124 protons. This nucleus could possibly be studied itself, or alternatively it could possibly decay by alpha and beta emission to a more stable nucleus near the center of the island.

![Diagram of nuclei overlaps]

Heavy target and projectile
Good neutron-proton ratio
Small distortion?

Fig. 18
Pinch-off direct transfer reaction for producing superheavy nuclei

A slightly different type of direct transfer reaction is illustrated in Fig. 18. Such a pinch-off reaction has been suggested by Swiatecki\textsuperscript{99} from considerations of experiments on liquid drops performed by Thompson's group\textsuperscript{118}. When a heavy target and projectile are brought together, the
two inner portions of each can possibly fuse to form a central superheavy nucleus, with two lighter particles released. The use of a heavy target and projectile leads to a good neutron-proton ratio. In addition, such a pinch-off reaction with a particle released on each side could possibly form the superheavy nucleus with less distortion in a nearly spherical shape.

Fig. 19
Fission reaction for producing superheavy nuclei

Figure 19 illustrates the last type of reaction mechanism, namely fission, which is a method emphasized by Flerov\textsuperscript{119}. A heavy target and projectile, such as plutonium plus plutonium, fuse to form a massive excited nucleus. This nucleus then undergoes fission, the hope being that one of the fission fragments will be a superheavy nucleus. Although the mass distribution resulting from the fission of such a massive nucleus is unknown, the hatched region shows which fragments could possibly be formed. This type of reaction is seen to give an excellent neutron-proton ratio. But the fission fragments will be very elongated at birth,
and because of this large distortion, any superheavy nucleus formed as a fission fragment will itself probably undergo fission rather than de-excite to its ground state\(^\text{29}\).

4. CHEMICAL AND PHYSICAL PROPERTIES

Assume that we are successful in producing superheavy nuclei. Then what chemical and physical properties do we expect them to have? This has recently been the subject of several reviews\(^\text{117,120-122}\), and I will make only a few concluding comments here.

![Periodic table]

**Fig. 20**

Periodic table including predicted locations of new elements

Figure 20 shows a periodic table\(^\text{25}\), including elements heavier than those already discovered. Elements are arranged in a periodic table according to the order in which the electrons—rather than the nucleons—fill their orbits. For the undiscovered elements this can be learned by
doing self-consistent calculations for the electrons surrounding the nuclei, for example relativistic Hartree-Fock calculations. Such calculations have been performed by several groups at Los Alamos and elsewhere. The results indicate that the electronic orbits for elements through 120 are systematically filled in an analogous way to those in the preceding row of the periodic table; the only difference is that the next shell is being filled. Element 114 should therefore have chemical properties similar to lead, element 113 to thallium, and so on. However, the calculations of Mann\textsuperscript{123}) suggest that beginning with element 121 the order of filling the shells is slightly different from the order in the actinide elements. This difference arises from relativistic effects associated with the higher charge of superheavy elements. Therefore, the chemical properties of element 124 should be only somewhat similar to those of uranium, and so on.

| Table 1 | Chemical and physical properties |
| --- | --- | --- |
| Property | $^{208}\text{Pb}$ | $^{298}\text{114}$ |
| Electronic ground-state configuration | Xe core + $^{4f}_{14}^{14}5d_{10}^{10}6s_{2}^{2}6p_{2}^{2}$ | Rn core + $^{5f}_{14}^{14}6d_{10}^{10}7s_{2}^{2}7p_{2}^{2}$ |
| Oxidation state | +2 (+4) | +2 |
| Ionization potential (eV) | 7.4 | 8.5 |
| $K_{\alpha_{1}}$ x-ray energy (keV) | 75 | 174 |
| Density ($g/cm^3$) | 11 | 14 |
| Melting point (°C) | 327 | 67 |
| Boiling point (°C) | 1620 | 147 |

Electronic self-consistent-field calculations can also be used for predicting other chemical and physical properties, such as those shown in Table 1. Some of these values for $^{298}\text{114}$ were predicted by Mann\textsuperscript{124}).
the others by Keller et al.\textsuperscript{125}). They are to be compared with experimental values for \textsuperscript{208}Pb. As we have already discussed, the electronic ground-state configuration of element \textsuperscript{114} is analogous to that of lead; the electrons are simply filling the next shell. Whereas lead is primarily divalent but sometimes tetravalent, element \textsuperscript{114} is expected to be purely divalent. The ionization potential for element \textsuperscript{114} should be only slightly higher than that for lead, but the x-ray energies will be substantially higher. Element \textsuperscript{114} should be somewhat more dense than lead. Finally, the melting and boiling points for element \textsuperscript{114} are pre-

\textbf{Fig. 21}

Energy release in fission for binary and ternary divisions
dicted to be substantially lower than those for lead, but these predictions arise from a complicated extrapolation\textsuperscript{125} and consequently are more uncertain than the other results.

The final properties I would like to consider are those associated with the fission of a superheavy nucleus. In Fig. 21 we show by the heavy solid curve the energy released when a nucleus divides into two fragments, as a function of nuclei throughout the periodic table, or alternatively the fissility parameter $x^{88,128}$). For binary division, the energy

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig22.png}
\caption{Calculated and experimental fission-fragment kinetic energies}
\end{figure}
release is seen to increase from about 200 MeV for uranium to over 300 MeV for the nucleus 114. The larger fraction of this energy is the translational kinetic energy of the centers of mass of the fission fragments themselves. The remaining portion, given by the difference between these two curves, is excitation energy of the fragments\(^{88}\), which is dissipated primarily through the emission of neutrons.

As indicated by the dashed curve, more energy is released for nuclei heavier than rare-earth nuclei if the division is into three fragments rather than into two\(^{128}\). This of course does not mean that the nucleus 114 will predominantly divide into three fragments, but the frequency of ternary divisions should be higher for superheavy nuclei than for conventional nuclei.

The plot of the kinetic energy for binary division vs. the fissility parameter \(x\) is re-shown in Fig. 22, this time in units of the surface energy of the spherical nucleus. We also include some experimental most probable fission-fragment kinetic energies from a variety of sources\(^{88}\). The data of particular interest here are by Sikkeland\(^{127}\) for the reaction \(^{40}\text{Ar}\) incident on \(^{238}\text{U}\). The resulting kinetic energy is seen to be what we would expect for the fission of the superheavy nucleus \(^{278}\text{I10}\). This is a convincing argument that the argon and uranium nuclei have fused to form a single nucleus, which, however, almost immediately undergoes fission into two fragments of comparable size.

As shown in Table 2, substantially more neutrons per fission and neutrons of higher energy are released in the fission of superheavy nuclei\(^{128}\). Whereas on the average 2.8 neutrons are released when \(^{240}\text{Pu}\) undergoes fission, we expect some 10.5 neutrons to be released in the fission of \(^{298}\text{I14}\). This expected increase is providing the basis for experimental searches for superheavy elements in nature\(^{108,129}\). The average laboratory energy of the fission neutrons is expected to increase from 2.0 MeV for \(^{240}\text{Pu}\) to 2.8 MeV for \(^{298}\text{I14}\).
Table 2
Properties of fission neutrons

<table>
<thead>
<tr>
<th>Fissioning nucleus</th>
<th>Number of neutrons per fission</th>
<th>Laboratory energy of fission neutrons (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{240}_{\text{Pu}}$</td>
<td>2.8</td>
<td>2.0</td>
</tr>
<tr>
<td>$^{298}_{114}$</td>
<td>10.5</td>
<td>2.8</td>
</tr>
</tbody>
</table>

5. SUMMARY

Let me end by summarizing the three main points I have tried to make. First, an island of nuclei in the vicinity of 114 protons and 184 neutrons is expected to be relatively stable with respect to spontaneous fission, alpha decay, and beta decay. We still do not know the precise shape and extent of the island. Nuclei with approximately 124 protons are also possibly fairly stable with respect to spontaneous fission, but would decay rapidly by alpha and beta emission. Second, the way to produce superheavy nuclei is by heavy-ion reactions. Of the alternatives available, direct transfer reactions and compound-nucleus formation, possibly in overshoot reactions to the region of $^{208}_{124}$, offer the most hope. Third, the chemical properties of the superheavy elements through element 120 should be similar to those of the preceding group of elements in the periodic table. The fission of a superheavy nucleus should release somewhat more kinetic energy and substantially more neutrons than the fission of a conventional nucleus.

*   *   *
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The work by the Lund group has presently the primary aim of predicting the properties of possible islands of superheavy elements.

With this aim in mind we have improved on our previous calculations by studying the most critical problem involved, that of the extrapolation of the nuclear potential parameters (this problem is treated in the communication submitted by C. Gustafsson). Furthermore we have now also included in our study the problem of odd-A effects in the Z=114 region (see communication by G. Öhlén), effects of non-axial types of deformations (P. Möller), and effects of the ultimate development of a two-center-type of potential (T. Johansson).

The calculations have also been extended to the Z=164 region (see communications by R. Bengtsson), already studied by other authors. In this calculation the problem of the shell structure parameter extrapolation is additionally critical. To cast some light on the extrapolation a similar "extrapolation" is performed from the rare-earth region, where shell structure is treated as known, and into the Z=40-60 region immediately below, where shell structure in one case is treated as being entirely unknown. The contribution discussing the results in the region mentioned is submitted by I. Ragnarsson and appears under another section of this conference report.
EXTRAPOLATIONS TO THE \( Z = 164 \) REGION

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**Introduction**

From the calculations of ref.\(^1\) and others, the prominence of the gap at \( Z = 164 \) is conspicuous for the set of nuclear parameters employed for \( A \approx 300 \). Furthermore, the size of the gap appears only weakly dependent on the value of \( A \). It thus appears that the gap may persist to \( A = 400 - 500 \) when it becomes relevant. An extrapolation to the \( Z = 164 \) region has already been attempted recently by J. Grumann, U. Mosel, B. Fink, and W. Greiner\(^2\). The authors cited at the time did not employ a renormalisation to the liquid-drop model as suggested by Strutinsky\(^3\). For this reason the prolate barrier relative to the spherical ground state was grossly overestimated and it was found consequently that the decay through the oblate barrier was the dominant fission decay channel. In the present calculations the Strutinsky shell correction method has been employed resulting in a prolate barrier considerably thinner than the oblate one. The fission half lives are therefore reduced and of a drastically lower order of magnitude than in the calculation of ref.\(^2\).

**Potential**

We have employed a potential\(^1\) (see C. Gustafsson\(^6\), for further details of the notation)

\[
V = V_{\text{osc}} + V_{\text{corr}}
\]
where

\[
V_{\text{osc}} = \frac{1}{2} \hbar \omega_0 (\epsilon_4 \epsilon_4) \left( 1 - \frac{2}{3} \epsilon_4 P + 2 \epsilon_4 P_4 \right)
\]

\[
V_{\text{corr}} = -2\hbar \omega_0 \left[ \frac{\omega_z}{\omega_z} \right] \left( \frac{1}{2} + \frac{1}{c} \right) \left( c + \frac{1}{c} \right) \left( \frac{1}{3} \right) \left( -\frac{1}{4} \right) \left( c - \frac{1}{c} \right) (f^+ S^- + f^- S^+) - \mu (\langle \frac{1}{4} \rangle^2 - \langle \frac{1}{4} \rangle^4)
\]

where

\[
c = \sqrt{\frac{\omega_z}{\omega_1}}
\]

and \( S^\pm \) are the usual spin step operators while \( f^\pm \) are angular momentum like operators defined e.g. in ref. 6). The spin-orbit term has been calculated from the form \( \langle \mathbf{S} \cdot (\mathbf{S} \times \mathbf{P}) \rangle \), where \( \mathbf{S} \) is calculated from \( V_{\text{osc}} \) for simplicity while the \( V_{\text{corr}} \) part of \( V \) has there be neglected.

We have assumed for \( \kappa \) and \( \mu \) the linear expansions in \( A^{1/3} \) as given in the contribution by C. Gustafson 6).

**Extrapolations**

These functions have been fitted to reproduce the empirical level orders in the two deformed regions \( 150 < A < 190 \) and \( A > 225 \). The brave linear extrapolations have been assumed to be valid to the proton region of \( Z = 164 \) and the corresponding neutron region along the line of beta stability with \( Z = 164 \), or \( N = 300 \). As the neutron-proton ratio here is very different from that of the presently known stability peninsula, this extrapolation is particularly uncertain. Thus it is to be expected that \( \kappa \) and \( \mu \) may depend on higher powers of \( I \) in addition to \( A \).

**Results on half-lives**

Using the methods outlined in ref. 4) we have employed the potential above to calculate single-particle energies as a function of the potential
distortion parameters \( \epsilon \) and \( \epsilon_4 \). From these a shell correction energy is obtained via the Strutinsky shell correction method\(^1\). A pairing energy is calculated using

\[
G(P) = \frac{1}{\Lambda} \left( 19.2 + \frac{7.4 (N-Z)}{\Lambda} \right)
\]

and a number of levels above and below the Fermi surface equal to \( \sqrt{15N} \) for neutrons and \( \sqrt{15Z} \) for protons. This prescription assures that the average trend of the energy gap parameter \( \Delta \), equal to \( \frac{12}{\sqrt{\Lambda}} \) MeV all over the region of observable elements, is retained into the superheavy region. Furthermore it is assumed that the pairing matrix element grows proportionally to the surface area \( S \), Coulomb and surface energies are determined as in ref.\(^1\). One problem is that one might now suspect that the little known surface symmetry term may play a more significant rôle as very large values of \( I \) are obtained in this case. This fact serves additionally to make all predictions uncertain.

Results

Fission half lives are calculated as in ref.\(^1\) from the WKB expression for the barrier penetration, where we have used the semiempirical value\(^4\)

of \( BA^{-5/3} = 0.054 \) MeV/h\(^2\).

Alpha decay rates are obtained from the semiempirical formula

\[
10 \log t_\alpha^{1/2} = 1.61 \left( Z E^{-1/2}_\alpha - Z^{2/3} \right) - 28.9 \text{ years}
\]

The single-particle level schemes for protons and neutrons are given in fig. 1. Two-dimensional fission barriers are exhibited in fig. 2 for a few nuclei. These figures are constructed in such a way that for each \( \epsilon \) we have found the \( \epsilon_4 \)-value that leads to minimum energy.
We note that elements adjacent to $^{468}\text{164}$ have a well developed spherical minimum with several MeV high barriers. The prolate barrier is then much thinner than the oblate barrier. The problem raised by the calculations of ref. 2) of an oblate path to fission is therefore no longer very relevant. However, as oblate and prolate shapes for each distortion have relatively similar energies (for moderately small $|\epsilon|$) one may ask for the extension of the energy surface into the rest of the gamma plane, where possibly the barrier may be even more transparent. No gamma-asymmetric calculations are presently available, however. The problem of the apparent oblate minima of fig. 2 in all likelihood disappears as they will probably be found to be maxima with respect to the gamma-degree of freedom, judging from the computational experience in the actinide region (Solcviev 5)).

Tabulation of half-lives

In spite of the overwhelming uncertainty of the extrapolations involved we have exhibited the resulting alpha and fission half-lives in the form of fig. 3. An inspection of this figure reveals that alpha decay is largely the limiting factor for the proton rich side of the stability line while fission determines the total half lives on the neutron rich side.

Although $Z = 164$ from these extrapolations appears associated with a large shell gap, the apparent absence of a corresponding neutron gap in reasonable vicinity of the stability line gives rather small half lives of any element involved. Thus $^{468}\text{164}$ appears to be associated with the largest half-life involved, or of the order of a few tens of milliseconds.
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   A131 (1969) 1

5. D.A. Arseniev, L.A. Malov, V.V. Pashkevich and V.G. Soloviev,
   Dubna preprint E4 - 3703 (1968)

6. C. Gustafson, report at this conference
Fig. 1. Extrapolated spherical single-particle level scheme relevant for $Z=164$, $N=300$. Note that no gap corresponding to $Z=164$ is available on the neutron side.
Fig. 2. Cut through the potential-energy surface for a few "island" elements. Note that the liquid-drop model shows a maximum for spherical shape. Due to shell structure effects these elements still exhibit a spherical minimum surrounded by a barrier on the prolate and oblate side, respectively. For $\epsilon < -0.4$ the calculations are limited solely to the $\epsilon$ degree of freedom.
Fig. 3. Island of semi-stability connected with $Z=164$. Shaded squares mark beta stable elements. Where fission half-lives are the decisive factor, corresponding isochrones are drawn as solid lines. For the proton rich side instead the alpha decay process dominates, whose half-lives are marked by dashed isochrone lines.
THE SHELL MODEL PARAMETERS OF THE Z = 114 REGION

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Introduction

In a previous publication\(^1\) we have studied the single-particle levels and total energies associated with a modified oscillator potential

\[
V = V_{\text{osc}} + V_{\text{corr}}
\]

where

\[
V_{\text{osc}} = \frac{1}{2} M \omega^2 \left( \epsilon, \epsilon_4 \right) \cdot \rho^2 \left( 1 - \frac{2}{3} \epsilon P_2 \left( \cos \theta_t \right) + 2 \epsilon_4 P_4 \left( \cos \theta_t \right) \right)
\]

\[
V_{\text{corr}} = -2 k \omega^2 \left[ \vec{t}_t \cdot \vec{s} + \mu \left( \vec{t}_t^2 - \langle \vec{t}_t^2 \rangle_N \right) \right]
\]

For \( \omega^2 \), we have assumed

\[
k \omega^2 \left( N \right) = 41A^{-1/3} \left( 1 + \frac{1}{3}\frac{N-Z}{A} \right) \quad \text{(MeV)}
\]

which defines an \( A \) and \( Z \) dependence of the general size parameter \( \omega_0 \) such that

1) the average nuclear volume increase proportionally to \( A \)
2) the r.m.s. neutron and proton radii are roughly equal along the nuclear stability line.

We shall assume that these relations are adequate also in the region \( Z \approx 114, A \approx 300 \).
Decisive for the prediction of the neutron and proton gaps supposedly connected with \( Z = 114 \) and \( N = 184 \) are the \( Z \) and \( A \) dependence of the parameters \( \kappa \) (spin-orbit splitting) and \( \mu (I_t^2 \) correction).

These parameters have been determined in the previous publication\(^1\) from a fit of single-particle levels in the deformed rare-earth and actinide regions to the empirical level order.

In doing so we have replaced the first term by the natural spin-orbit generalization \( \hat{s} \cdot (\hat{W} \hat{X} \hat{p}) \). For the case that for \( V \) in \( \hat{W} V \) we content ourselves with only the dominant part of \( V_{osc} \) above, neglecting the \( P_4 \)-part as well as other parts of the potential, we obtain for the spin-orbit term the following expression

\[
V_{s.o.} = - 2\kappa h \omega_0 \left[ \frac{\omega_1}{\omega_0} \right] \times \left\{ \frac{1}{Z} \left( c + \frac{1}{c} \right) I_t \cdot \hat{s} + \left[ \frac{1}{c^2} - \frac{1}{2} \left( c + \frac{1}{c} \right) \right] l_3 s_3 - \frac{1}{4} \left( c - \frac{1}{c} \right) (f^+ s^- + f^- s^+) \right\}
\]

where

\[
c = \sqrt{\frac{\omega_1}{\omega_0}}
\]

and \( s, f \) are the usual spin step operators while \( f, f^\dagger \) are operators defined\(^2\) as

\[
f^\dagger = \frac{1}{\sqrt{P_4}} (\xi \circ \xi + \xi \circ \xi) - i (\eta \circ \xi + \xi \circ \eta) \text{ where } \xi \text{ etc. are defined as}
\]

\[
\xi = \sqrt{\frac{M_0}{\hbar}} x
\]

and where \( \circ \xi \) is short-hand for \( \frac{\partial}{\partial \xi} \).

The operators \( f^\dagger \) exhibit matrix elements only between shell \( N \) and \( N + 2 \) while all matrix elements vanish within the \( N \)-shell.
This modification requires only a very minor change in \( \kappa \), as is apparent from the values of Table 1 relative to those of ref\(^1\). The latter values are given in parenthesis in Table 1.

We have also attempted to replace the \( \frac{1}{t} \) term by a term proportional to \((\mathcal{V} \times \mathbf{p})^2\). We found, however, that it was then no longer possible to fit the empirical level order within the model employing any constant \( \kappa \) and \( \mu \)-value. The fact that Greiner et al\(^3\) have been successful in this replacement is apparently connected with the fact that the change in distortion is accompanied by a radical change in potential shape due to the two-center-oscillator formulation. The little wedge in the bottom of the well has thus a favourable effect on the single-particle level order.

In Lund, Thomas Johansson\(^4\) has obtained a satisfactory single-particle level scheme with a \((\mathcal{V} \times \mathbf{p})^2\) term and an added term \( q e^{-\beta z} \), which prescriptions bears a strong resemblance to the two-center model of ref\(^3\). In our present work we have been content with keeping the previous \( \frac{1}{t} \) term and replacing only the spin-orbit term.

**Determination of parameters**

Our first and most time-consuming task has been that of fitting \( \mu \) and \( \kappa \) such that the empirically observed level order has been reproduced for deformed regions \( A \simeq 165, A \simeq 184, A \simeq 242, A \simeq 250 \).

In each of these regions we thus obtain an optimal \( \mu_p, \kappa_p, \mu_N, \kappa_N \). These parameter values are plotted as circles in figs. 1-4. We then draw a straight line through these points as function of \( A^{1/3} \) (which corresponds to the radius parameter and should be the most relevant variable in the extrapolation problem. See remarks by A. Bohr and B.R. Mottelson, *Nuclear Structure*, Vol. 1 (Benjamin, New York, 1969, p. 218). In this way we obtain
\[ \mu_p = 0.638 + 0.003 \ A^{1/3} \]
\[ \mu_N = 0.919 - 0.095 \ A^{1/3} \]
\[ \kappa_p = 0.0606 - 0.0030 \ A^{1/3} \]
\[ \kappa_N = 0.0617 + 0.0020 \ A^{1/3} \]

One might argue that these parameters should probably also be considered functions of the isospin or of the corresponding parameter used in the liquid-drop model: \( \frac{N-Z}{A} \). As long as we stay on the stability line, we contend at present that this is a satisfactory parametrisation. From the calculations by Ragnarsson\(^5\) in the regions \( 90 < A < 150 \), using a similar linear expansion of \( \mu \) and \( \kappa \) from the rare earth region, it appears that an extrapolation in \( A^{1/3} \) over a range \( \Delta A^{1/3} \) as large as that involved in going from \( A \approx 225 \) to \( A \approx 300 \) should be fairly safe. On the other hand, from the experience in the light region, the next island of stability connected with \( A \approx 460 \) seems beyond the region of credibility of such an expansion.

One may also note that the single-particle spectrum that we obtain for \( A \approx 300 \) (using the optimal expansion) is in amazingly good agreement with that recently obtained by Bolsterli, Fiset and Nix\(^6\), which latter is based on a potential of seemingly different radial shape and with its parameters determined from a fit of single-particle level order in the double-closed shell regions as \( ^{208}\text{Pb} \) and \( ^{48}\text{Ca} \) and \( ^{40}\text{Ca} \) rather than from the deformed regions as in our case. In figures 5 and 6 we compare the level-schemes of Bolsterli et al. with those of our calculation for protons and neutrons respectively of the nucleus \( ^{298}114 \). We notice the good agreement especially in the proton case.
To obtain some kind of estimate of the problems of extrapolation involved we have decided arbitrarily on a set of extrapolated $\kappa$ and $\mu$ values deviating from the straight line of extrapolation by 1.5 times the largest deviation encountered in the rare earth and actinide regions. For each of the two $\kappa$ and $\mu$-values for neutrons as well as protons we have obtained the sets that give the largest $Z = 114$, $N = 184$ gaps and the sets that give the smallest ones. The normally extrapolated case is denoted "EXTR" and the two extreme-cases are denoted "FAV" and "UNFAV" in figs. 7 and 8, where the levels of ref. 1) are given for comparison. Actually the "unfavourable" cases is in some way unduly "unfavourable" as the role of the closed-shell of $N = 184$ is entirely replaced by $N = 196$ as is apparent from fig. 8. (It is interesting to note that one reasonable shell gap as that at $Z = 114$ is enough to ensure a spherical shape for rather different neutron level schemes.) The corresponding resulting fission barriers are exhibited in figs. 9 - 12 corresponding to $^{296,298}_{114}$, $^{294,296}_{112}$, $^{292,294}_{110}$, and $^{290,292}_{108}$, respectively. For comparison the barriers obtained by Nilsson, Tsang et al. 1) are given as dotted curves.

**Details of calculations**

The single-particle energies as obtained from the potential above are summed corresponding to optimal filling for each set of distortions. From the sum is then subtracted the corresponding "average"-value as defined in the Strutinsky shell correction method 7). The difference is by Strutinsky denoted the shell correction energy. To this latter is added the pairing energy as obtained by the BCS method and the liquid drop energy as given in the recent Myers and Swiatecki formulation 9). In this way a total energy surface is obtained as described e.g. in refs. 7, 1) and 8).
In calculating the pairing energies we have used a surface and isospin dependent pairing matrix element as in ref.\textsuperscript{1}) and a cut-off in the BCS procedure utilising $\sqrt{15N}$ and $\sqrt{15Z}$ levels respectively for neutrons and protons above and below the Fermi surface. The intention of this cut-off is to retain the over-all $A$ dependence of $\Delta$ (the odd-even mass difference parameter) empirically encountered, i.e. proportional to $A^{-1/2}$.

The minimum of the potential energy surface gives the ground state mass. Once we have the masses, we can calculate the alpha half-lives according to Taagepera and Nurmia\textsuperscript{10}):

$$10 \log \left(t_\alpha^0\right)_{1/2} = 1.51(ZE^\alpha_a^{-1/2} - Z^{2/3}) - 28.9 \text{ years}$$

where $E^\alpha_a$ is the energy of the ejected $\alpha$-particle and $Z$ the atomic number.

In order to estimate the fission half lives, we limit ourselves for the moment to the simplified problem of an equivalent one-dimensional barrier. We thus calculate the total energy as a function of $\epsilon$ and $\epsilon_4$. For fixed $\epsilon$ we minimize the energy with respect to $\epsilon_4$. This procedure then defines a curve in the $\epsilon, \epsilon_4$-plane which is assumed to be approximately the path to fission for the nucleus considered. The values of the total energy along this curve we then project on the $\epsilon$-axis, and conjecture as the fission barrier. The errors associated with this rough simplification to a one-dimensional problem are assumed to be of a similar magnitude as the errors associated with the insufficient parametrisation at large distortions, the neglect of dynamical aspects, the uncertainty in the inertial mass parameters etc. The probability for penetration of the fission barrier is approximately
\[ P = \exp \left( -2 \int_{\epsilon'}^{\epsilon''} \left( \frac{2B}{\hbar^2} (W(\epsilon) - E) \right)^{1/2} \, d\epsilon \right) \]

where \( B \) is the inertial mass parameter, \( W(\epsilon) \) the energy of the barrier as a function of \( \epsilon \) and \( E \) the zero point excitation energy assumed\(^1,8\). Usually one furthermore assumes the number of assaults of the nucleus against the fission barrier per time unit to be equal to the frequency of the beta-vibration. For the corresponding beta vibrational energy quantum we assume an average value of 1 MeV. Consequently the half-life for spontaneous fission may be obtained as

\[ \tau_{1/2} = 10^{-20.54} \cdot \rho^{-1} \text{ sec.} \]

**Calculational results**

From figs 9-12 it is apparent that the barriers as obtained in ref.\(^1\) and \(^8\) are essentially reproduced by our regular set of extrapolated values of \( \kappa \) and \( \mu \). This is the case although the spin-orbit term is modified relative to that employed in the cited reference. The fission half-lives are somewhat reduced relative to those of ref.\(^1\), see fig. 13. The fission half-lives quoted are based for simplicity on the semi-empirical value of the inertial parameter \( B \) cited in ref.\(^8\), \( BA^{-5/3} = 0.054 \, \hbar^2/\text{MeV} \). No element is by the regular extrapolation directly assigned a half-life long enough for earthly detection. Still in view of e.g. the uncertainty in the \( B \) parameter, survival is associated with some probability. For the "favourable" set of \( \kappa \) and \( \mu \)-values, on the other hand, a whole region of elements should be candidates for survival in earthly elements (especially \( Z = 114, N = 184, Z = 114, N = 186 \)) with respect to fission, alpha and beta decay, see fig. 14. For the extremely "unfavourable" variant (see fig. 15), on the other hand, of the considered
island of stability no element is expected to live as long as one year. Obviously no superheavy element is then detectable in earthly matter. On the other hand still a considerable island of fairly long-lived elements, is accessible for detection and decay studies, if reached by heavy-ion collisions.

Conclusions

One thus concludes that the shell structure and the corresponding shell correction energy are not qualitatively changed with a modest change of the potential of such a kind that the observed actinide and rare-earth single-particle level order is reproduced. Still the extrapolation with A appears a relatively difficult problem. The great sensitivity of the predicted fission half-lives in the A≈300 region to the values of the extrapolated values of \( \kappa \) and \( \mu \) is exemplified by the "favourable" and "unfavourable" cases treated in this calculation. It should be emphasised that, as pointed out above, the "unfavourable" one of the extrapolated cases is in a way unduly "unfavourable" as other neutron and even proton (Z=124) gaps then require considerations. It should be noted that in this "unfavourable" case the fission half lives are decisive relative to the alpha half lives. In this case consequently the most long-lived elements center around \( Z = 114 \) (and \( N = 182 \)) rather than \( Z = 110 \).
References


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5. I. Ragnarsson, report to this conference

6. M. Bolsterli, E.C. Fiset and J.R. Nix, to be published


Table 1. Parameters fitted to empirical level order
Values in parenthesis are those of ref. 1.

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Table 2. (parameters relevant to the $A \approx 300$ region)

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Fig. 1  The fitted values of $\mu_p$ as function of $A^{1/3}$ (circles). Also shown are the linear extrapolation of these parameters to the region. $A \approx 300$ and the adopted estimate of the uncertainty in the extrapolation.

Fig. 2  Same as fig. 1 but for $\mu_n$. 
Fig. 3  Same as fig. 1 but for $\kappa_p$. 

Fig. 4  Same as fig. 1 but for $\kappa_n$. 

Fig. 5 The single-proton level scheme for spherical shape calculated in the "folded Yukawa" well (Bolsterli et al.) and the modified oscillator potential for the nucleus $^{298}_{114}$.
Fig. 6    Same as fig. 5 but for neutrons.
The single-particle level scheme of protons for spherical shape of the modified oscillator. Notations EXTR., FAV and UNFAV refer to various combinations of extrapolated $\mu$ and $\kappa$ namely the straight linear extrapolation, the permitted deviating combinations of $\kappa$ and $\mu$ that maximize and minimize the energy gaps for spherical shape at $Z=114$ and $N=184$ (see text). For comparison the scheme of ref. \(^1\) is exhibited.
Fig. 8  Same as fig. 7 but for neutrons.
Fig. 9  The energy minimized with respect to $\varepsilon_4$ for each $\varepsilon$ as function of $\varepsilon$ for $^{290}_{108}$ and $^{292}_{108}$. The barriers for the three sets of $\mu$ and $\kappa$ and the barrier of ref. 1) are exhibited.
Fig. 10  Same as fig. 9 but for $^{292}_{110}$ and $^{294}_{110}$. 
Fig. 11    Same as fig. 9 but for $^{294}_{112}$ and $^{296}_{112}$.
Fig. 12  Same as fig. 9 but for $^{296}_{114}$ and $^{298}_{114}$.
Fig. 13 Contours of theoretical half-lives in the vicinity of Z=114 and N=184. The thick lines are contours of spontaneous fission half-lives. The broken lines are contours of alpha half-lives. Squares corresponding to beta-stable nuclei are shaded. Regularly extrapolated parameters $\mu$ and $\kappa$. 

---

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- $\beta$-stable
- $\alpha$ half life contours
- Spontaneous fission half life contours

EXTR.
Fig. 14  Same as fig. 13 but for the "favourable" set of parameters $\mu$ and $\kappa$.

Fig. 15  Same as fig. 13 but for the "unfavourable" set of parameters $\mu$ and $\kappa$. 
THE TWO-CENTRE POTENTIAL AND THE \((\mathbf{V} \mathbf{V} \times \mathbf{D})^2\) TERM

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Introduction

The deformed oscillator potential has the disadvantage that the maximum depth of the potential remains at the center also for very distorted shapes, see fig. 1. Several suggestions to avoid this disadvantage by the introduction of a so-called two-center potential have been made \(^{2,3,4,5}\). Here we do this by adding a term proportional to \(\exp(-z^2/a^2)\) to the potential. This choice makes it possible to pass smoothly from a deformed one-center potential to a two-center potential.

It would also be more satisfactory to replace \(\mathbf{t}_t \cdot \mathbf{s}\) and \(\mathbf{r}_t^2\) - terms in the deformed-oscillator-potential by a vector proportional to \(\mathbf{V} \mathbf{V} \mathbf{x} \mathbf{p}\). This replacement (in some cases limited to the \(\mathbf{t}_t \cdot \mathbf{s}\) - term) has also been investigated earlier \(^{4,5,6,7}\).

One-body potential and calculation of energy-eigenvalues.

We take as our single-particle hamiltonian:

\[
H = T + V + C \left( (\mathbf{V}_o \times \mathbf{p}) \cdot \mathbf{s} + D \right) \left[ (\mathbf{V}_o \times \mathbf{p})^2 - \langle \mathbf{s}^2 \rangle_N \right]
\]

where \(V = V(\rho, z) = \omega_o^2 \left( V_o + V_g \right) = \frac{1}{2} \omega_o^2 \left( 1 + \frac{z^2}{3} \right) \rho^2 + \rho^2 \left( 1 - \frac{2z^2}{3} \right) z^2 \)

\[
+ \omega_o^2 \cdot \mathbf{q} \cdot \exp \left[ -z^2/\left( \beta^2 \kappa/\omega_o \right) \right] \quad (\rho^2 = x^2 + y^2)
\]

Here \(C = -2\omega_o \psi\) and \(D = -\omega_o \kappa\psi\), and \(\epsilon, \psi, \text{ and } \beta\) are deformation parameters. For the effect of \(q\) and \(\beta\) on the shapes see fig. 2. The necking-in is due to the formation of two centers in the potential.
The only term in the deformed oscillator-model that is not immediately
generalizable to strongly deformed shapes is the $\langle \mathcal{L}_N^2 \rangle$ term. The latter
is meant as an approximate volume-conservation correction to the $\mathcal{L}_N^2$- term
in the spherical case 1).

An interpretation of the $\langle \mathcal{L}_N^2 \rangle$ term that could lead to a genera-
larization is the following. The "volume" of an orbital that is strongly
depressed in energy by $\mathcal{L}_N^2$ is different from the "volume" of neigh-
bouring orbitals. If we therefore compress the radial scale of this
orbital so that the "volumes" of orbitals in the same energy region become
equal the effect will be similar to that of $\langle \mathcal{L}_N^2 \rangle$, which also results
in an energy increase. In this paper, however, we follow Gustafson et al. 1)
and keep the term $\langle \mathcal{L}_N^2 \rangle = \frac{1}{2} N(N + 3)$ independent of distortion although
its meaning for deformed shapes is not quite clear.

As a basis for calculating energy-eigenvalues we have chosen the
eigenfunctions $|Nn_z\Delta\Omega>$ of a deformed, cylindrically symmetric harmonic
oscillator. The matrix elements in this basis of our hamiltonian will
be published elsewhere. In the diagonalization of each matrix of given
$\Omega$ and parity $\Pi$, coupling matrix elements between all $N$-shells are retained.

Volume conservation and definition of the shape

It is convenient to express energies in the unit $\hbar\omega_0$, which is a
function of $A,N$ and the deformation, that is $\omega_0 = \frac{\varphi}{\omega_0} (A,N) \cdot f(\text{def})$. To
calculate $\omega_0$ we employ the simple Thomas-Fermi model according to which

$$\rho(\vec{r}) = \frac{(2M)^{3/2}}{3\pi^{2/3}} \cdot (\lambda - V)^{3/2}$$

where $\lambda$ is the classical fermi-level for $N$ particles in the potential
$V$. 8) By solving the eq.
\[ N = \int \rho \, d^3r \]
\[ \lambda - V \geq 0 \]

we get \( \lambda \). For a spherical nucleus \( \omega_0 (A, N) \) is determined, for neutrons and protons separately, by the condition

\[ \sqrt{\langle r^2 \rangle} = \frac{5}{3} \frac{r_0}{A^{1/3}} \]

where

\[ \langle r^2 \rangle = \int r^2 \rho \, d^3r. \]

The dependence of \( \omega_0 \) on deformation is given by the condition that the volume inside the surface

\[ V(\rho, z) = \lambda \]

is a constant. This surface also defines the shape of the nucleus and for comparison with experiment we have analyzed it in \( \beta_2, \beta_4 \) and so on.

**Calculated energy-levels for actinides at equilibrium**

We now want to examine how the different changes in the potential affect the level order at equilibrium. Figures 3 and 4 show the level schemes for protons around \( Z = 100 \) and neutrons around \( N = 152 \) for different potentials. Note that \( \kappa \) and \( \mu \) are unchanged within each series of figures. In a) the simple modified h.o. potential is used with only \( P_2 \) distortions included and in b) we employ the same potential as under a) but with \( \mathcal{V} \tilde{\nu} \tilde{x} \tilde{P} \) replacing \( \tilde{I}_t \). It is seen that the effect of this replacement is quite dramatic, and that the higher \( \Omega \)'s are most affected. The explanation is that the high \( \Omega \) states have large \( n_\perp \)-values (\( n_\perp \geq \Lambda \)) and therefore small \( n_z \), and accordingly are concentrated around \( z = 0 \), where also \( \mathcal{V}_0 \) is largest for prolate distortions, as the equi-potential
surfaces are packed together there. For a similar reason the high $\Omega$ states
are most affected, in a direction to spread the different $\Omega$-components of
a j-subshell, by the gaussian $V_g$ which lifts the potential around $z = 0$.
The effect of this is exhibited in case (c) that shows a level order which
reproduces the empirical level order about as well as case (a).

The level scheme reported by Greiner et al in ref. 4) as a function
of d, the distance between centers in the two-oscillator model, is relatively
similar to case (c).

It will be interesting in the future to compare the results of
ref. 4) with the more physical two-oscillator model of ref. 5) having
a smooth joining of the two oscillators. Such results are, however, at
present not available to us.

Conclusion

We have shown that two changes of the modified oscillator model,
that might be important for strongly deformed nuclei, give shapes
appropriate for fission and also give realistic level-schemes at equi-
librium. However, it appears desirable in the future also to include
$VV_g$ into the generalized $\bar{\Omega}$ - term, and to obtain a better understanding
of the $\langle x^2 \rangle_N$ - term or a generalisation thereof.

The author is very indebted to Professor S.G. Nilsson for suggesting
this investigation a couple of years ago and for enlightening discussions
during the work.

I am grateful to Professor B. Mottelson for helpful advice particularly
on the formulation of the volume conservation condition.
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Fig. 1. Nuclear shapes and corresponding (static) potentials in two cases:

a) deformed h.o. with $P_2$-distortions and the gaussian $q \exp \left[-z^2/(\beta^2 M_0)\right]$;

b) deformed h.o. with $P_2^2$, $P_4$, and $P_6$-distortions;
Fig. 2. Nuclear shapes in the plane of the deformation parameters $\alpha$ and $\beta$ ; $\epsilon = 0.4$.
Fig. 3a Single-proton levels around $Z=100$;
\[ \tilde{l} \rightarrow \tilde{l}_t, \; q=0 \]
Fig. 3b Single-proton levels around Z=100;

\[ \tilde{k} = \frac{\nu_{\text{osc}}}{M_0 \omega_0^2} \vec{p} , \quad q = 0 \]
Fig. 3c  Single-proton levels around $Z=100$;

$$\bar{p} + \frac{\nu V_{osc}}{M_0^2 x_P} , \ q \neq 0$$
Fig. 4a Single-neutron levels around $N=152$ ;

$\tilde{i} + \tilde{i}_t$, $q=0$
Fig. 4b  Single-neutron levels around $N=152$ ;

$\tilde{\mathbf{I}} + \frac{\mathbf{v}_\text{osc}}{M_0^2} x^p$, $q = 0$
Fig. 4c  Single-neutron levels around $N=152$;  
\[
\tilde{l} + \frac{v_{osc}}{M_0^2} x^p, \quad q \neq 0
\]
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EFFECTS OF THE $P_3$ AND $P_5$ DISTORTIONS ON THE FISSION BARRIERS
OF THE SUPERHEAVY ELEMENTS

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Experimentally fissioning actinide nuclei are found to decay into two
fragments of unequal mass at low excitation energies [1]. In [2] we have
studied the fission process theoretically by calculating a static energy
surface (cf. fig. 1) of the nucleus for various reflection symmetric
distortions, the distortions being characterized by two parameters: $\epsilon$,
corresponding to an elongation of the nucleus, and $\epsilon_5$, positive values
of which give a neck-in of the nucleus. The black line $\epsilon_5^{24}$ in fig. 1
defines roughly, the "path to fission" from the first saddle and on. To
study asymmetric fission we also calculate the potential energy for
asymmetric distortions. We use the following single particle potential:

$$V = \frac{1}{2} \hbar \omega_o^2 (1 - \frac{2}{3} \epsilon \epsilon_2^P + 2 \epsilon_4 \epsilon_4^P_4 + 2 \epsilon_4 \epsilon_4^P_1 + 2 \epsilon_3 \epsilon_3^P_3 + 2 \epsilon_5 \epsilon_5^P_5) - V_{corr}$$

$$V_{corr} = \kappa \hbar \omega_o \left( \frac{2}{3} \hbar \omega_o \epsilon \epsilon_2^P + \mu (\epsilon_2^2 - \left \langle \epsilon_2^2 \right \rangle_{\text{shell}}) \right),$$

where $\epsilon_1$ depends on distortion to assure that the center of mass remains fixed.

For details of the calculations see for the symmetric case [2] and
the asymmetric case [3]. In [3] we made contour plots of the potential
energy with one axis corresponding to the symmetric distortions defined
by the axis $\epsilon_2^{24}$ in fig. 1 and an asymmetric axis graded in $\epsilon_5$. To
each value of $\epsilon_5$ we chose a value of $\epsilon_3$ that minimizes the energy at
$\epsilon = 0.85$ and $\epsilon_4 = 0.12$. We then found that $^{210}_{84}$Po was stable towards
asymmetric distortions whereas the second peak for $^{236}_{92}$U was lowered
by more than 2 MeV and for $^{252}$Fm by approximately 0.3 MeV. In this paper we investigate the effect of asymmetric distortions on the fission barriers for $^{236}$U and the superheavy nuclei, in the same way as we in [3] investigated the actinide barriers, with two modifications:

1) We extend the calculations to distortions including the ground state of the nuclei.

2) When we make contour plots of the potential energy we grade the symmetric distortion axis in $\epsilon_3$ as before. The corresponding value of $\epsilon_5$ may be read from figs. 2 and 3. Contrary to the procedure of ref [3], the asymmetric axis is now graded in $\epsilon_3$. The corresponding value of $\epsilon_5$ is defined in fig. 4. The function $k(\epsilon)$ is determined by keeping $\epsilon_4$, $\epsilon_5$, and $\epsilon_3$ fixed and determining the value of $\epsilon_5$ that minimizes the energy. Call this value $\epsilon_5^{\text{min}}$. Then $k(\epsilon) = -\epsilon_5^{\text{min}}/\epsilon_3$. We have found that $k$ is only a function of $\epsilon$ and does not depend significantly on $\epsilon_4$, $\epsilon_3$ or the particular nucleus considered in the actinide or superheavy region. In [3] we assumed $k(\epsilon) = +0.5$ for all $\epsilon$.

In fig. 5 we exhibit a contour plot of the static potential energy for $^{236}$U. The line $\epsilon_3 = 0.00$ corresponds to the approximate "path to fission" in $P_2 - P_4$ space illustrated in fig. 3. We see that the second saddle point occurs at a large asymmetric distortion. What in fig. 1 was a symmetric saddle point is here a mountain peak some two MeV high forcing the nucleus out into the asymmetric part of the diagram.

In fig. 6 we have plotted the fission barrier corresponding to the plot in fig. 5. For each $\epsilon$-value we determine the minimum energy in the diagram of fig. 5 and plot it as a function of $\epsilon$. The dotted line corresponds
to the barrier we obtain when we restrict ourselves to the line $\epsilon = 0$ in fig. 5.

In fig. 7 we see a topographical map of the potential energy for the superheavy nucleus $^{294}_{110}$. We see that the saddle is stable against asymmetric distortions. There is a slight instability beyond the saddle at $\epsilon \approx 0.5$. Other even-even nuclei in the region investigated ($106 \leq Z \leq 118$) exhibit a similar small instability or are very soft towards asymmetric distortions at this value of $\epsilon$.

The distance between the grid points in the $\epsilon$-direction is somewhat large, 0.20, to save computational time. This is why the plotting program when interpolating between grid points placed the ground state at $\epsilon \approx -0.05$ when it should have $\epsilon \approx 0.00$. We feel however, that the plot is accurate enough to show that the inclusion of the $P_3$ and $P_5$ terms into the nuclear potential will not decrease the height of the barrier for $^{294}_{110}$ and that the slight instability encountered beyond the saddle point will reduce the fission half life by a factor of less than 10. Fig. 8 shows the fission barrier for $^{294}_{110}$.

In the above calculations we have restricted ourselves to one combination of even multipole distortion parameters and one combination of odd multipole distortion parameters. To get better values of the distortion parameters at the second saddle point and determine more exact saddle point shapes one should vary all distortion parameters independently. With the method used in this paper it is not possible to deduce the mass of the fission fragments of $^{236}_{92}U$ and $^{252}_{100}Fm$ from the saddle point shapes. As to the latter method preliminary calculations indicate that varying all distortion parameters independently will give "smoother" saddle point shapes. A limiting factor in the calculations is that the
computation time with our program is of the order of ten hours on a UNIVAC 1108 for determining the more exact saddle point shape for one nucleus.

The author is very much indebted to S.G. Nilsson, who suggested this problem, for guidance and continued interest in this work, and to all other members at the Department of Mathematical Physics in Lund for many helpful discussions. The author also wishes to thank R. Nix and the Los Alamos Scientific Laboratory, C.F. Tsang, W.D. Myers, W. Swiatecki and S.G. Thompson and the Lawrence Radiation Laboratory for interesting discussions and the pleasant hospitality granted.

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Fig. 1. The potential-energy surface in the \((\epsilon \epsilon)\) plane for the nucleus \(^{236}\text{U}_{92}\). The axis \(\epsilon_{24}\) marked by a solid arrow indicates the definition of part of a "static path to fission".
Fig. 2. Definition of a "static path to fission for the nucleus $^{235}_{92}$U. Cf. also fig. 1.

Fig. 3. Definition of a static path to fission for the superheavy elements.
\[ \varepsilon_5 = -k(\varepsilon) \cdot \varepsilon_3 \]

Fig. 4. Definition of \( \varepsilon_5 \) as a function of \( \varepsilon \) and \( \varepsilon_3 \).
Fig. 5. The potential energy surface of $^{236}_{92}$U in terms of symmetric and asymmetric distortions.
Fig. 6. Fission barrier for $^{236}_{92}$U with and without inclusion of odd multipole distortions. The ground state is arbitrarily put at zero energy. The dashed lines are values of the extremum points taken from ref. [4].
Fig. 7. The potential energy surface of $^{294}_{110}$ in terms of symmetric and asymmetric distortions.
Fig. 8. Fission barrier for $^{294}_{110}$ with and without inclusion of odd multipole distortions. The right end of the curve is arbitrarily put at zero energy.
THE ODD-PARTICLE EFFECT ON THE FISSION BARRIER

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Introduction

Spontaneous fission half-lives of odd-A actinides are considerably longer than those of even-even ones. Wheeler 1) introduced the "specialization energy" concept to explain this empirical finding. The specialization energy is the barrier increase due to the odd particle. To the odd-particle is assigned an orbital with given parity and angular momentum component along the axis of deformation. These quantities are assumed to be conserved through the whole fission process. In this way the odd-particle strongly affects the spontaneous fission half-life, but also alpha decay rates are affected, mainly through intrinsic inhibitions associated with the odd-particle orbitals of the mother and daughter states.

The odd-A energy surface

The even-even energy surface is calculated by using the single-particle potential described in ref 2) (Nilsson-Tsang) and the Strutinsky shell correction method 3) with shell model parameters as in ref 4).

The odd-A energy surface is obtained by adding the quasi-particle energy $E_q$ to the "even-even"-energy surface calculated with an "even-even" BCS-solution, that has an average number of particles that is odd.
The quasi particle energy is given as

\[ E_v = \sqrt{(\varepsilon_v - \lambda)^2 + \Delta^2} \]

where \( \lambda \) and \( \Delta \) are obtained from the BCS calculation. The minimum energy consistent with given \( \pi \) and \( \Omega \) is obtained by taking the \( \varepsilon_v \) which is closest to \( \lambda \) with given parity and angular momentum (when \( P_3 \neq P_5 \) distortions are considered the potential is asymmetric so only \( \Omega \) is assumed to be conserved, see below). This method is satisfactory for the ground state energy. More problematic is the method in obtaining the fission barrier (at least the symmetric case) because it assumes that the odd-particle always is found in the energetically most favourable orbital at the point of the quasi-crossing of the two orbitals of the same \( \Omega \) and \( \pi \). This is probably only partially true during the fission process. An even more questionable assumption is actually made in the present calculation, namely that a change of the odd-particle orbital from a "hole" to a "particle" orbital and vice versa is made instantaneously during the change of shape whenever such a transition is energetically favourable. Within the model assumed there is actually no mechanism available for this. We would have to ascribe this to residual forces of other kinds than pairing (the only residual forces considered in the present model).

In future calculations one should as an alternative consider also the case that such a one-particle transition from one state of hole character to another one of particle character and vice versa is not permitted.

For a more realistic treatment of the fission barrier, of which these static estimates are only approximations, we will ultimately have to employ a full-fledged dynamic theory.
We have also studied the $P_3 + P_5$ degrees of freedom. These we have treated in the same way as in ref.\textsuperscript{5}), that is to say that we have taken a certain path to fission in the $P_2 + P_4$ plane that is the same for all actinides and along this path we have studied asymmetric distortions described by the most favourable linear combination of $P_3$ and $P_5$.

In this case we have made the following even more drastic barrier-reducing assumption. As very small $P_3 + P_5$ distortions are involved, the corresponding orbitals have very small admixtures of one of the parities. We have still assumed that transitions between orbitals of opposite dominant parities occur without inhibitions. It seems reasonable in forth-coming calculations to replace this assumption with the alternate assumption that the dominant parity is conserved in the change of odd-particle orbitals.

The specialisation energy obtained with the present set of assumptions should be considered a gross underestimate of and only a lower limit to the specialisation energy.

\textbf{Results}

For most actinides the specialisation energy is about 1-2 MeV and this effect is enough to explain the longer fission half-lives of the odd-$A$ nuclei. If there is a low density of states with some particular angular momentum there can be very high specialisation energy for instance $11/2^-$ (fig. 2). The asymmetric degrees of freedom are seen to reduce this effect dramatically so our predictions about a possible long-living isomeric state of Pu\textsuperscript{241} may be affected\textsuperscript{6}). From fig. 1 it can be seen that nonconservation of parity does not change the fission-barrier much (this may of course be drastically changed if the density of states are low).
In the superheavy case the calculations are additionally complicated by the fact that most of these elements are spherical and not deformed in their ground state.

The discussion above is directly tailored for the coupling scheme appropriate to the deformed case, with $\Omega$ (and $\Pi$ in the symmetric case) the only constant of motion.

As we approach sphericity from the distorted situation we must imagine a strong Coriolis coupling between orbitals of $\Omega$ different by one, which through diminishing $\epsilon$ have to connect up with the spherical region. Ultimately the spherical seniority coupling scheme is assumed to be reached for $\epsilon = 0$. In our calculations we have proceeded as follow.

With given $I$ in the ground state (assumed equal to $j$ of the last filled subshell orbital) we have subsequently assumed the fission barrier approximated by two (limiting) $\Omega$ - orbitals namely $\Omega = j = I$ and the $\Omega$ - orbital that corresponds to the lowest orbital for small prolate deformations. $j = 1/2$ subshell only one case occurs, $\Omega = 1/2$.

As a conservative estimate we have in fig. 4 taken the lowest of the corresponding fission halflives.

Furthermore we have also assumed the inertia parameter the same as in the even - even case and used the semiempirical value ($B = 0.054 A^{5/3}$ $5/3 \ h^2 \ \text{MeV}$). This should correspond to an underestimate by 10 - 30\% of the logarithmic halflives for odd - $A$ nuclei. The halflives listed should thus merely be taken as lower limits. For the $\alpha$ - halflives an expression given by Taagepera - Numia has been used.

The longest total halflives appears to be associated with the elements $^{293}_{110}$ and $^{291}_{110}$ in addition to $^{294}_{110}$ already cited in ref 2).

The author is much indebted to Prof. S.G. Nilsson, who suggested this problem and for his interest in this work, and to all other members at the Department of Mathematical Physics, Lund, especially P. Möller, whose help with computational work has been invaluable.
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Fig. 1 Fission barrier for $^{239}$Pu and $^{241}$Pu $\frac{1}{4}^+$. The solid curves correspond to a potential energy surface with the $P_2$ and $P_0$ degrees of freedom included, the angular momentum component and parity are thus assumed to be preserved. The dotted line contains the alternative (unrealistic) assumption that the parity of the odd orbital is not conserved. The dot-dashed curves correspond to the assumption of $P_3$ and $P_5$ included and no parity conservation.
Fig. 2 Fission barrier for $^{242}$Pu $11/2^-$. For notation see Fig. 1. The inclusion of $P_3 + P_5$ plus the assumption of parity-nonconservation for the odd particle is seen to erode the secondary barrier entirely. In a realistic case we expect at least partial "parity conservation".
Fig. 3 Fission barrier for 3 superheavy elements. For these nuclei only $P_2 + P_4$ distortions are considered. Furthermore it is assumed that for these nuclei, which are spherical in their ground states, $K$ is assumed to be preserved for all $\varepsilon$-values.
Fig. 4 Calculated lifetimes for some superheavy elements.
SEARCH FOR SUPERHEAVY ELEMENTS IN NATURE BY DETECTION OF
EVENTS WITH LARGE NUMBERS OF NEUTRONS

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ABSTRACT

Some preliminary results are reported on a search for the presence
of "superheavy elements" in nature. Large samples (10-50 kg) of various
substances including W and Pb ores as well as manganese nodules were
examined for the presence of spontaneous fission events in which large
numbers of neutrons are emitted. The method utilize a large Gd loaded
liquid scintillator with 60-70% efficiency for detecting neutrons. The
system was located in a tunnel of depth ~ 275 meters in order to reduce
interference from cosmic rays. In the materials tested no events above
background were detected. This limit corresponds to half lives in the
region $10^{23}$ years relative to the principle component of the samples.

INTRODUCTION

In recent years various estimates and calculations have been made
which suggest the possibility that nuclei having atomic numbers in the
region 108-114 may be sufficiently stable to exist in nature. These nuclei
are expected to decay either directly by spontaneous fission or indirectly
by other modes of decay to give products which would then undergo spontaneous
fission.

It is expected that the spontaneous fission of these "superheavy
nuclei" would be accompanied on the average by emission of about ten neu-

*Work performed under the auspices of the U. S. Atomic Energy Commission.
trons. The emission of such a large number of neutrons from each single fission event is then a very distinctive indicator of the presence of these elements because this property does not occur in any of the presently known spontaneously fissioning species in nature. Furthermore, these evaporated neutrons would have average energies in the range 1-3 MeV and thus would be able to escape the sample and enter the counting chamber. Consequently large amounts of heavy materials can be examined with great sensitivity.

The use of a large liquid scintillator is a sensitive and convenient method for detecting events in which large numbers of neutrons are emitted. The efficiency for detecting neutrons is high especially when the scintillator contains an element such as gadolinium which has a very large cross section for neutron capture (46000 barns). Finally, the liquid scintillator provides a natural means of detecting many neutrons from single events because it makes use of the wide variation and separation in thermalization times; i.e. the individual neutrons are separated according to time of thermalization and are then captured to produce the photons and light which is seen by the photomultipliers.

In the preliminary work reported here we have used a system involving a liquid scintillator to search for superheavy elements in several materials. A number of other materials will be examined before the work is completed.

2. DESCRIPTION OF THE DETECTION METHOD

The neutrons are detected by a large gadolinium loaded liquid scintillator with the capability of measuring neutrons with high efficiency. A schematic view of the detector system is shown in Fig. 1. Its main body is a tank with dimensions $62 \times 62 \times 125$ cm which holds a liquid consisting of toluene solvent, $\sim 8$ g/l gadolinium octoate, 0.1 g/l POPOP and 5 g/l p-terphenyl. Sixteen 5-inch Dumont No. 6364 photomultipliers are mounted on
glass viewing ports on two walls of the chamber. The inner walls of the
chamber are painted with white reflective coating. A tube for holding the
samples is located at the center of the chamber and has the dimensions of
11.4 cm diameter and 105 cm depth. The system is very similar to some
described already in the literature.\footnote{1}

Neutrons produced by any source placed at the center of the cham-
ber enter the liquid and are thermalized by collisions with the hydrogen
in the solution and eventually are either captured by the gadolinium or
leak out of the tank. The \((n, \gamma)\) reaction in gadolinium produces
\(\sim 9\) MeV of gamma energy, the energy usually being shared by several gamma
rays. The electrons created by the reaction of these gamma rays with the
liquid produce scintillations which are seen by a few of the photomulti-
pliers.

The distribution in time for neutron capture is broad and has a
peak at about 10 \(\mu\)sec after the neutrons are emitted. About 90\% of all the
neutrons are captured within 1-36 \(\mu\)sec after their production. In this way
individual members of a burst of energetic neutrons (\(\sim 1\) MeV) are separated
in time for convenient electronic multiplicity counting.

3. **ELECTRONICS**

A simplified block diagram of the electronics is shown in Fig. 2.
The photomultiplier pulses are summed in two banks (8 photomultipliers in
each). A "tank pulse" is obtained when coincident pulses are observed from
banks A and B and when the linear sum of all photo multiplier outputs cor-
responds in pulse height to an energy greater than 1 MeV. This signal has
the significance that the light must be seen by at least two photomultipliers
and have a sum amplitude corresponding to an energy greater than 1 MeV. Since
at least two photomultipliers must give a signal the effect of random noise
from individual tubes is minimized. A "tank pulse" triggers, after a 0.5 µsec delay, a 36 µsec gate. (The delay is to ensure that all of the prompt gamma rays from a fission event have been emitted.) During the 36 µsec gate interval all tank pulses are counted by a scaler. At the end of the 36 µsec period the digital information in the scaler is converted to analog pulse heights and transferred to a pulse height analyzer. Then the scaler is reset. The next tank pulse defines a new 36 µsec gate interval.

A spectrum of multiplicity in the range 0-15 is thus obtained in a pulse height analyzer. The system is capable of monitoring a burst of neutrons where the trigger can be activated either by the prompt γ-rays or the first neutron captured. The circuits shown by dotted lines in Fig. 2 was added in the later experiments to remove events that had cosmic rays in coincidence and will be discussed later.

4. OPERATION OF THE SYSTEM

The efficiency of the system was checked at regular intervals by placing a weak $^{252}$Cf spontaneous fission source in the center of the tank and triggering the 36 µsec gate by the fission events (using a small solid state detector to detect the fragments). The multiplicity distribution obtained in this manner for 25,000 fission events is shown in Fig. 3. Its mean value is $\bar{n} = 2.44$. This yields an efficiency of 65.5% for detection of each neutron (using 3.72 as the average number of neutrons emitted in $^{252}$Cf fission).

Long term stability and continuous monitoring of the system which is necessary for this type of an experiment was obtained in three ways:

1) The singles counting rate of tank pulses was recorded every 2 sec by a chart recorder and showed the system to be very stable. 2) The data from the multi-channel analyzer were printed out automatically at regular intervals
Fig. 1. The gadolinium loaded liquid scintillator detection system.
Fig. 2. The electronic system.
Fig. 3. The multiplicity distribution observed with a $^{252}$Cf source placed in the center of the chamber. The 36 μsec gate was triggered by a fission fragment detector. The average multiplicity was 2.44, i.e. the efficiency = $2.44/\bar{\nu} = 65.5\%$; ($\bar{\nu} = 3.72$).
and the multiplicity distributions were found to be consistent (within statistical fluctuations). 3) In recent experiments photographs were made of the tank pulses that appeared during the 36 μsec gate, whenever an event having a multiplicity of four or more occurred. In principle these events can be checked to be certain they have a normal time distribution. However, as yet, the photographic equipment has not been operating satisfactorily.

When the system operates with the neutron tank pulse as the main trigger, it is sensitive to the environmental radiation which consists of gamma rays from natural sources (e.g. U, Th, and K) and cosmic rays and the products of their reactions in matter. The gamma rays that arise from natural sources appear in general as single random pulses. Accidental coincidences between these random pulses yield a distribution which is represented by a Poisson probability function.

\[
P(N) = \frac{8.64 \times 10^4 \cdot (CT)^{N+1} \cdot \exp(-CT)}{(N! \cdot T)}
\]

where \( P(N) \) is the number of events observed per day appearing to have a multiplicity \( N \) with a single random tank pulse rate of \( C \) (per sec) and a gate length of \( T \) (sec). At the operating count rate of 600 counts/sec less than 1 event per day appearing as multiplicity four or more can be expected.

5. RESULTS

The results of the multiplicity distributions measured on samples of lead ore and tungsten in metallic form are shown in Figs. 4 and 5. In both cases events with high multiplicities were observed at a rate that could not be accounted for by the expected random occurrence of high multiplicity events. The rate of events with given multiplicity expected from probability considerations is indicated in the figures by the dark lines and was
Fig. 4. The multiplicity distribution obtained with a lead ore sample. The dark line represents the multiplicities expected on the basis of probability considerations and the dots represent the experimental results.
Fig. 5. The multiplicity distribution obtained with a sample of purified tungsten metal. The dots represent the experimental results and the dark lines the expected random behavior.
calculated from the known single counting rate. For the low multiplicities (1 to 3) the observed rate agreed with the predicted random multiplicities within the statistical uncertainty.

The exact number of neutrons that accompany spontaneous fission of a superheavy element and the shape of the neutron distribution are not yet known. Therefore, we cannot calculate exactly the probabilities associated with detecting an event with multiplicity \( N \) or more neutrons. However when the system operates with \( \sim 67\% \) efficiency and on the average 10 neutrons are emitted the probability of detecting four or more following a trigger by the first neutron is \( \geq 60\% \), depending somewhat on the shape of the assumed distribution. We have chosen then to report the results as the number of events having multiplicity four or greater. We related this to the number of atoms in the sample to obtain the half life. The half lives can be compared directly with the results of Flerov et al.\(^2\)

In Table I we present some results obtained in the various measurements. The results on the lead and the tungsten samples are rather close. In the case of tungsten it is reasonable to assume that its homologue ekatungsten would be too short lived to exist in nature. Therefore, the excess high multiplicity events present both in W and Pb can be assumed to come from the interaction of \( \mu \)-mesons with these heavy nuclei. Calculations based on the estimated flux and known cross sections confirm this assumption.

In obtaining the results described above, events following very large electronic pulses were rejected because the system was saturated by the large pulses. This caused the rejection of \( \mu \)-meson induced events with large multiplicities whenever the \( \mu \)-meson passed through the liquid scintillator. A \( \mu \)-meson loses 1.5 – 2.0 MeV/cm in the liquid, i.e., 30–250 MeV
<table>
<thead>
<tr>
<th>Sample</th>
<th>Sample weight, Kg</th>
<th>Counting time (hrs)</th>
<th>Total events of multiplicity 4+</th>
<th>Half life $\times 10^{23}$ yrs</th>
<th>Remarks</th>
</tr>
</thead>
<tbody>
<tr>
<td>tungsten metal</td>
<td>49.4</td>
<td>478</td>
<td>43 ± 8</td>
<td>1.4 ± 0.3</td>
<td>composition: pure W</td>
</tr>
<tr>
<td>lead metal</td>
<td>33.1</td>
<td>82</td>
<td>20 ± 20</td>
<td>0.31 ± 0.31</td>
<td>composition: pure Pb</td>
</tr>
<tr>
<td>lead-ore galena, PbS</td>
<td>22.0</td>
<td>158</td>
<td>13 ± 5</td>
<td>0.40 ± 0.15</td>
<td>composition: 65% lead b</td>
</tr>
<tr>
<td>manganese* nodules</td>
<td>7.5</td>
<td>247</td>
<td>13 ± 7</td>
<td>0.41 ± 0.22</td>
<td>composition: c</td>
</tr>
</tbody>
</table>

$^a$ The backgrounds due to random multiplicities have been subtracted.

$^b$ The lead ore is PbS galena from the Sullivan mine, Bunker Hill, Idaho.

$^c$ The manganese nodules (crust) came from a depth 2.74 Km at a position 136 Km west of the Mid Atlantic ridge at 45° N. More details are given by Fleischer et al. Phys. Rev. 184, 1393 (1969).

* This experiment was done with steel shielding placed around the apparatus. See text.

when it passes through the tank, whereas the energy from neutron capture gamma rays amounts to $\sim 9$ MeV.

Our lower limits which can be accounted for by the cosmic ray background amounted to $\sim 3$ counts/day/50 kg heavy sample, of multiplicities 4 or more with 60-70% neutron detection efficiency. This limit is equivalent to a half life of $\sim 10^{23}$ y relative to the principal component of the
samples. Within this limit we are unable to find any evidence for the
decay of "superheavy elements" in the samples that we have examined so
far.

These results are in agreement with the work of Price et al. 3) that set a lower limit corresponding to a half life of $3 \times 10^{23}$ years for spontaneous fissions in the lead minerals by observing fission fragment
tracks.

In Table I, the results are also given for a ~7.5 kg sample of
manganese nodules. These nodules contained 14% Mn and 20% Fe. The total
number of events of multiplicity 4 or more was $13 \pm 7$ in a run of 247 hours
length. On the basis of the total number of atoms of Mn and Fe taken
together the half life limit is $4 \times 10^{22}$ years. The measurement of man-
ganese nodules was performed after surrounding the apparatus with steel
shielding of thickness 5 cm. The effect of the shielding is discussed
below.

In the later experiments steel shielding was introduced around the
apparatus in order to reduce the effect of background gamma radiation from
the surrounding medium. As a result the single $\gamma$-ray counting rate was
reduced from ~650 counts per sec to ~250 counts per sec. This reduction
has the desirable effect of reducing the accidental rate of higher multi-
plicity events. However, we found that the addition of the shielding had
the undesirable effect of increasing the $\mu$-meson background. The $\mu$-mesons
appeared to produce events of high neutron multiplicity in the shielding.
Some of these events are produced by muons which do not pass through the
tank. These events could therefore not be eliminated from the data.

After introducing the shielding, we also attempted to reduce the
cosmic ray background by placing a liquid scintillator system above the
main tank as indicated in Fig. 1. In this case muons coming from the vertical direction pass through the scintillator and produce large identifiable pulses which were used to remove associated high multiplicity events which appear in the main tank. Also the $\mu$-mesons passing through the large tank produce a similar large identifiable pulse (30–250 MeV) that was used to remove associated high multiplicity events. Even with these systems in operation we were not able to reduce the background below the level which existed before the shielding was installed.

It now appears that the sensitivity of our apparatus can only be increased significantly by placing it much deeper underground where the $\mu$-meson flux is proportionately decreased. Perhaps an improvement by a factor of 10 to 50 could be obtained in this way.

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SEARCH FOR SUPERHEAVY ELEMENTS ON THE ACCELERATOR ALICE

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1. INTRODUCTION

The Orsay accelerator named ALICE ¹ is able to accelerate a Krypton beam of Kr²²⁺ up to 400 MeV and Kr²³⁺ up to 440 MeV. Krypton ions Kr⁸⁺ are put at 1 MeV per nucleon by a linear accelerator, transported into the central area of a cyclotron where they are stripped of a number of electrons. The carbon stripping foils are 20 µg per cm² thick. Up to now the best yield which has been obtained corresponds to a beam of 5.10⁸ particules per second, but we expect to improve the intensity by a factor 3 in the next few weeks and by a factor 10 in six month.

One of the best way to produce a Superheavy nucleus is known to be the fusion of a projectile and a target of similar masses ², with the additional requirement that the compound nucleus should be at a low excitation energy. In order to reach the region N = 184, Z = 120-124, which seems now to be the most promising for the stability against both spontaneous fission and alpha decay, Kr projectiles have to collide targets like Th₂³² and U₂³⁸, with a kinetic energy larger than 5 MeV per nucleon (Coulomb Barrier of the order of 3.8 MeV per nucleon in the c. of m.). Because of the negative Q values of fusion reactions, the expected excitation energy will be small. For example, the excitation energy of the compound nucleus obtained by the fusion of a 425 MeV Kr ions and a Th₂³² target can be estimated around 5 MeV. Therefore the excited nucleus ¹²⁶X₁⁹⁰ should decay by α or γ emission. Since the alpha decay of the resulting nucleus at the ground state is expected to be very short, it is likely that the detected product will be lighter by one, two or three alpha particles, and this so called "overshoot" bombardment will produced ¹²⁴, ¹³⁸, ¹⁴², or ¹⁴⁴.
Some reactions of this type are presented on figure 1 and table I. The reply to the question of which nucleus might exactly be

**Compound Nuclei at Low Excitation**

\[ {^{80}}\text{Kr} + {^{230}}\text{Th} \rightarrow {^{310}}_{126} + \gamma \]

\[ {^{82}}\text{Kr} + {^{226}}\text{Ra} \rightarrow {^{308}}_{124} + \gamma \]

\[ {^{84}}\text{Kr} + {^{232}}\text{Th} \rightarrow {^{316}}_{126} \stackrel{\text{fast}}{\rightarrow} {^{308}}_{122} \]

\[ {^{86}}\text{Kr} + {^{238}}\text{U} \rightarrow {^{324}}_{128} \stackrel{\text{fast}}{\rightarrow} ... \rightarrow {^{300}}_{116} \]

\[ {^{86}}\text{Kr} + {^{208}}\text{Pb} \rightarrow {^{292}}_{118} \stackrel{\text{fast}}{\rightarrow} {^{288}}_{116} \]

**TABLE I**

be detected along the line following arrows representing the \( \alpha \) decay depends on the time necessary for the detection. With the identification method which will be indicated later in section 2, any half life larger than \( 5 \times 10^{-7} \) sec is suitable. According the most recent data from Nix\(^3\), fission barriers and \( \alpha \) decays are given on Table II for the region \( Z = 120-126 \) and \( N = 184-190 \). It is obvious that nucleides like \( ^{312}_{124} \), \( ^{310}_{124} \) and \( ^{308}_{122} \) are good candidates.
TABLE II

Fission Barriers in MeV and $\alpha$ decay half-lives (in brackets) for several nuclei in the range $Z = 120$ to $Z = 126$ (after Nix).

Another possibility for the case where $Z = 114$ could be reached is the transfer reactions of big clusters which have been discussed $^2$). The formation of a very stable cluster like $^{28}_{14}\text{Si}$ may leave a neutron rich heavy partner rather close to the $\beta$ stability line. Table III.
Particular Cluster Transfers

\[ ^{40}\text{Ar} + ^{208}\text{Pb} \rightarrow ^{20}\text{Ne} + ^{228}\text{Th} \]

Hypothetic

\[ ^{86}\text{Kr} + ^{238}\text{U} \rightarrow ^{28}\text{Si} + ^{96}\text{Th} \]

\[ ^{86}\text{Kr} + ^{244}\text{Pu} \rightarrow ^{32}\text{S} + ^{298}\text{Th} \]

TABLE III


The purpose of the first experiment which is presently under performance is to measure the mass of the recoiling nucleus. Its kinetic energy should be of the order of 100 MeV for a bombarding energy of 420 MeV and the principle of the method is to measure simultaneously the momentum with a magnetic spectrometer (mV = ZeBR), the kinetic energy with solid state detectors located in the focus of the magnet, and the time of flight of the recoiling nucleus along the path between the target and the focus of the magnet. Fig. 2 illustrates very schematically the experimental arrangement. The Krypton beam produces superheavy nuclei in the target T. An electrostatic septum (160 KV between the electrodes) deviates sufficiently the superheavy recoiling nuclei out of the main krypton beam. They enter into a double focusing magnetic spectrometer.
The measurement of the energy $\varepsilon$ and of the momentum ZeBR give a unique value of the ratio $\frac{Z^2}{m}$ when $Z$ is the effective charge of the ion and $m$ its mass. For a given ion at a given energy, there are several possibilities of effective charges $Z_i$. It is therefore expected to obtain events for several different values of BR corresponding to the various values of $Z_i$, according the relation:

\[
\frac{(BR)_i}{(BR)_{i+1}} = \frac{Z_{i+1}}{Z_i}
\]

According the formula given by Schmelzer and extrapolated for superheavy elements, the different probabilities of charges are shown as an example on Table IV. The preliminary results which we have obtained on the krypton beam itself at 360 MeV have shown that the experimental most probable charge is lower than the calculated value by 2 or 3 units.

\[
\frac{310}{126} X \text{ Energy 104 MeV - Time of flight } 0.545571 \times 10^{-6}
\]

<table>
<thead>
<tr>
<th>Ch eff.</th>
<th>27</th>
<th>28</th>
<th>29</th>
<th>30</th>
<th>31</th>
<th>32</th>
</tr>
</thead>
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<tr>
<td>$Z^2/m$</td>
<td>2.35</td>
<td>2.53</td>
<td>2.71</td>
<td>2.90</td>
<td>3.1</td>
<td>3.3</td>
</tr>
<tr>
<td>Prob.</td>
<td>0.078</td>
<td>0.1</td>
<td>0.12</td>
<td>0.13</td>
<td>0.125</td>
<td>0.11</td>
</tr>
<tr>
<td>BR</td>
<td>9 576</td>
<td>9 234</td>
<td>8 916</td>
<td>8 619</td>
<td>8 341</td>
<td>8 080</td>
</tr>
<tr>
<td>Freq.</td>
<td>16 741</td>
<td>16 145</td>
<td>15 591</td>
<td>15 073</td>
<td>14 589</td>
<td>14 136</td>
</tr>
<tr>
<td>Z/E</td>
<td>0.25962</td>
<td>0.26923</td>
<td>0.27885</td>
<td>0.28846</td>
<td>0.29808</td>
<td>0.30769</td>
</tr>
</tbody>
</table>

**TABLE IV**

- Tabulated values calculated for an hypothetic nucleus. BR is the expected magnetic rigidity in gauss-meter, Freq. the resonance frequency for the resonance probe.
In the focal plan of the magnet, 16 solid state detectors made in four silicon blocks have been put for the energy measurements. Calibrations for the pulse height versus the energy have been made with a scattered krypton beam and when the intensity will be suitable, recoil scattered nuclei of gold will be used for heavier masses. The ionisation defect for such new projectiles was up to now unknown and a great care is taken for these calibrations, mainly because any error in the energy measurement will appear in the resolution of the mass spectrum. The main difficulty is to get rid of the krypton beam itself. If krypton ions enter into the aperture of the magnet they will be deflected and will make a large background. This is one of the reason for the electrostatic septum. In the same time, the switching device opens the possibility to put a thin foil only on the trajectory of the recoiling ions and therefore to have a starting time for a time of flight measurement as shown on figure 2. The distance of 440 cm between the foil and the focus of the magnet allows a time of $5.10^{-7}$ sec of travelling along for mass 310 at 100 MeV. This is a much larger time than for any other lighter ion which might be produced with the same characteristic $\frac{Z^2}{m}$ as the very heavy fusion nucleus. As an example, fission fragments which may behave like superheavy nuclei (same energy of 100 MeV, same $\frac{Z^2}{m}$ of the order of 3) travel in $2.10^{-7}$ sec, as shown on figure 3. Therefore the time of flight measurement will give a certitude for the identification of a heavy mass.

**Number of events to be expected**

If the fusion nucleus is produced at a low excitation under the fission barrier, one might expect a cross section of 10 mb. The thickness of the thorium target cannot exceed 30 $\mu$g/cm$^2$ if the energy degradation of the recoiling nuclei has to be less than 10 MeV. Because of the charge dispersion and the energy dispersion, only about 1/10 of the events will be seen by the detectors. Therefore with $10^9$ krypton per second one hopes to be able to detect 20 nuclei per minute.

In addition, the detectors are ready for measuring the $\alpha$ decay of the ions which should penetrate into the detectors at a depth of 13 microns. A 12 MeV $\alpha$ particle should have a range in the detector of 120 microns.
3. Further experiments on chemical identification with a mass spectrometer


An other experiment is being prepared which will be possible if the cross section is larger than 1 mb and if the half life of the superheavy nuclei is larger than $10^{-4}$ sec. It should in principle identify two of the possible superheavy elements: eka Francium ($Z = 119$) and eka-Radium ($Z = 120$). The targets will be $^{231}{\text{Pa}}$ and $^{232}{\text{Th}}$ and the reactions:

$$^{231}{\text{Pa}} + ^{84}{\text{Kr}} \rightarrow ^{315}{\text{I}} \quad \text{4 \alpha} \quad ^{299}{\text{Fr}}$$

$$^{232}{\text{Th}} + ^{84}{\text{Kr}} \rightarrow ^{316}{\text{I}} \quad \text{3 \alpha} \quad ^{304}{\text{Ra}}$$

The targets are put in the ion source of a mass spectrometer.

The recoiling nuclei will be stopped into a Tantalum filament which has been specially designed in shape and thickness. A proper heating of the filament gives an ionisation yield $\frac{i^+}{i_0} = A \exp \left( \frac{W - \phi}{kT} \right)$ (where $W$ is the extraction work, $\phi$ the ionisation potential of the element stopped in the filament, $T$ the temperature) which will be characteristic of an alcaline or an alcalino-earth. Preliminary experiments have been made with cesium and baryum, and there is also some on-line work being prepared for identification of Fr and Ra produced by bombardment of Pb by carbon ions.

A very sensitive device has been built for measuring the ion current at the collection plate of the spectrometer. A synchronisation system between the bias applied for a given mass and the current measured on the collector allows to scan a large range of masses. The voltage is increased channel by channel and it is possible to perform up to 10,000 scanning cycles per second. The sensitivity is of the order of 1 atom per minute.
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Fig. 1 Schematic diagram of the production of superheavy nuclei by a fusion process followed by alpha emission.
Fig. 2  Apparatus used for mass identification. Superheavy recoils are deflected into the magnetic spectrometer BR. The time of flight is measured over 440 cm between an electron lens and the energie detecteur.
Fig. 3  Mass versus time of flight for particles of a given identification parameter $Z^2/A$. 
THE DYNAMICS OF THE RAPID NEUTRON CAPTURE PROCESS

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ABSTRACT

Previous studies of the rapid neutron capture mechanism of nucleosynthesis ("r" process) have been carried out with the assumption of constant temperature and density during the reactions. The "dynamics" of such a process thus consists essentially of the beta decay of neutron-rich isotopes in local statistical equilibrium for a given Z. We are attempting to improve the treatment of this process. We employ a nuclear reaction network containing six thousand nuclei whose neutron capture cross sections, neutron binding energies, beta decay rates, and delayed neutron branching ratios are calculated from nuclear systematics based upon a mass formula. We find beta decay rates much larger than in the work of Seeger, Fowler, and Clayton\textsuperscript{1). In a previous investigation\textsuperscript{2),} it was demonstrated that charged particle reactions in an expanding neutron-rich supernova envelope are likely to form \textsuperscript{78}Ni seed nuclei. We are following the rapid neutron capture from this seed nucleus for various expanding supernova envelope adiabats. We have found that neutron captures do not proceed rapidly until the temperature falls to about \(1.5 \times 10^{9}\) K. Virtually all of the neutrons have been captured by the time the temperature falls to \(6 \times 10^{8}\) K. Delayed neutron effects tend to produce a small odd-even abundance effect in the final material, consistent with observed abundances. The influence of cooling and expansion of the material on the final abundances is demonstrated to be substantial for our choice of temperature and density conditions. Calculations are continuing in order to determine the initial conditions giving the best fits to the general r-process abundance curve.
1. **INTRODUCTION**

Following the compilation by Suess and Urey\(^3\) of their table of abundances of the elements in the solar system, it was possible to analyze these to determine the nuclear processes which had contributed to nucleosynthesis. One of the processes thus defined was rapid neutron capture, in which neutron captures occur rapidly compared to the beta decay half-lives of nuclei near the valley of beta stability. Subsequent to the neutron capture buildup, the capture products approach the valley of beta stability by a series of beta decays. This process has been called the r-process by Burbidge, Burbidge, Fowler and Hoyle\(^4\).

This rapid neutron capture process has previously been studied by quasi-static methods only. In these calculations, a specific temperature and neutron density are assumed and a steady flow path on the neutron-rich side of the valley of beta stability is then estimated. The process is approximated by assuming that for a fixed value of Z, neutron capture has proceeded until a statistical equilibrium is established along the Z-chain with respect to neutron capture and photoneutron reactions. These nuclei then wait at their equilibrium positions until beta decay takes place, following which further neutron capture occurs until the nuclei once again come into nuclear statistical equilibrium with respect to neutrons at a new "waiting point". For such a steady flow situation, the expected abundances of the final neutron capture products are proportional to the beta decay half-lives of the nuclei at the waiting points.

A proper dynamical study of the rapid neutron capture process must be based on the conditions expected in those types of supernovae for which hydrodynamic calculations predict a neutron-rich environment in which the process can occur. In practice this means that a neutron-rich fluid will be ejected and will expand approximately along an adiabat at a rate characteristic of a hydrodynamic time scale. The expanding material will form some seed nuclei, the character of which will depend upon the neutron-to-proton ratio in the material. These seed nuclei will then capture neutrons at rates which depend on the temperature and density conditions along the expanding adiabat. We discuss various features of this dynamical process in the following sections.
2. OLDER STUDIES OF THE RAPID NEUTRON CAPTURE PROCESS

In their original treatment of the r-process, Burbidge, Burbidge, Fowler and Hoyle\textsuperscript{4}) used a conventional von Weizacker mass formula without shell corrections, although they attempted to make corrections for shell effects in their calculations. Their estimates of beta decay half-lives were based upon a consideration of only the decay of the ground state of the parent nucleus to the ground state of the daughter nucleus, or to a low lying state of the daughter nucleus.

Conventional mass formulas generally predict too small neutron binding energies far off the valley of beta stability. This results from the fact that the only form of the nuclear symmetry energy contained in these formulas is the volume symmetry energy, and only the leading parabolic term in its expansion is utilized. For the range of neutron-rich nuclei of pertinence to the r-process, however, higher order symmetry terms are very important in the determination of the neutron binding energies and beta decay energies. These terms include the surface symmetry term and the next volume symmetry term, which is of the form $1^4A$ where $I = (A - 22)/A$. Indeed, in our own current study of the r-process, we have found it necessary to go one step further: we include a higher order surface symmetry term, of the form $1^{4/3}$, and a still higher order volume symmetry term, of the form $1^{6A}$.\textsuperscript{5})

These symmetry terms give higher neutron binding energies at a given distance from the valley of beta stability; hence, for a specified temperature and neutron flux, the capture path for the r-process lies farther away from the valley of beta stability.

The assumption of beta decay only between ground states grossly underestimates the rapidity of the beta decays of very neutron-rich nuclei, where the decay energies may lie in the range 10 to 15 MeV. At these energies, beta decay from the ground state of the parent nucleus to a very large number of excited states of the daughter nucleus is then possible. A step in the right direction was taken in the subsequent study of the r-process by Seeger, Fowler and Clayton\textsuperscript{1}). These authors assumed that the daughter nuclei had a uniform spacing of excited states with a value characteristic of the first MeV of excitation, and they assumed that the beta decay could take place to one out of three of
these excited states. However, this still grossly overestimates the half-life of the beta decay process. The number of excited states per unit energy interval increases at low energies proportionately to \(\exp(\text{constant } E)\), while at higher energies it increases as \(\exp(\text{constant } E^{1/2})\). Thus it may be seen that enormously more excited states are available for the beta decay than were assumed by Seeger, Fowler and Clayton.

One cannot, however, calculate the total beta decay transition probability simply by adding up the decays to the individual daughter states assuming some average transition matrix element. The wave functions of the higher excited states are exceedingly complicated, hence the beta decay matrix elements to an individual state will become progressively smaller with increasing excitation energy. One expects the beta strength function, proportional to \((ft)\) times the level density, to approach an approximately constant value. The beta decay matrix elements for the transitions between the ground state of the parent nucleus and the low lying states of the daughter nucleus are already considerably reduced compared to what they would be if the various states involved were good single-particle states. In our own work we have therefore estimated beta decay rates by assuming that the transition matrix element to the first one hundred levels in the daughter nucleus has an average \((ft)\) value characteristic of normal allowed beta decays. We have further assumed that above the one-hundredth level there is a constant beta strength function per unit energy interval in the daughter nucleus, with the beta strength function being equal to the value attained in the daughter nucleus in the vicinity of the one-hundredth excited state.

The dramatic difference brought about by these different assumptions is evidenced by the fact that while the beta decay half-lives used by Seeger, Fowler and Clayton for the most part lay in the range \(10^{-1}\) to 10 seconds, our own estimates of beta decay half-lives tend to lie in the range \(10^{-4}\) to \(10^{-3}\) seconds. We should point out that part of this difference arises from the fact that our neutron fluxes are considerably higher than those used by Seeger, Fowler and Clayton; hence the nuclei at the "waiting points" typically have larger beta decay energies.
3. **THEORIES OF SUPERNOVAE**

The first serious numerical hydrodynamic study of a supernova explosion was carried out by Colgate and White\(^6\). The basic mechanism involved in their study was the collapse of a stellar core to nuclear densities, followed by the transport of released gravitational potential energy away from the core by neutrinos and antineutrinos. These particles were found to heat the infalling matter sufficiently to create a shock wave, leading to an expansion of the outer part of the infalling material in a strong supernova shock wave which blows the remainder of the star off into space. In such a process it might be expected that a few tenths of a solar mass would be compressed to high enough densities to create a very neutron-rich fluid which, upon expanding into space, would result in the mixing of the products of the rapid neutron capture process with the interstellar medium. This was a little embarrassing with regard to nucleosynthesis, since it appeared that considerably too much production of r-process nuclei would take place\(^2,7\).

At the present time there is grave doubt that any significant amount of mass ejection can occur by this neutrino-antineutrino transport process. Studies by Arnett\(^8\) indicated that for cores of mass \(M > 4M_\odot\), no explosion would take place. Recent numerical hydrodynamical studies of a supernova implosion by Wilson\(^9\), in which the neutrino and antineutrino transport was treated somewhat more correctly, failed to give any expulsion of material. There is thus considerable reason to doubt that the formation of any significant amount of r-process elements will occur in this type of supernova event.

Meanwhile, as a result of studies by Arnett\(^10\), it appears that a wide range of supernovae may result from thermonuclear explosions. In these cases, a star does not succeed in initiating carbon or oxygen-burning processes during the normal course of its evolution. Instead, it forms a highly degenerate core which begins to contract rapidly as a result of neutrino losses, so that when the ignition of carbon or oxygen thermonuclear reactions occurs, the thermal runaway burns the carbon and oxygen thermonuclear fuels explosively and triggers the formation of a detonation wave in the core of the star. This blows the star completely apart, dispersing all of the material into space. The thermonuclear reactions
which occur during this explosive process reproduce the abundances of nuclei in the vicinity of the iron nuclear statistical equilibrium peak very well\textsuperscript{11}). At no time during the contraction of the core does the density ever become high enough for appreciable electron capture to have occurred to provide a large abundance of free neutrons; hence it is not expected that these explosive supernovae events will make any significant contribution to the r-process.

More massive stars are now expected to implode almost completely, presumably forming black holes, with possible thermonuclear expulsion of the outer layers, but without the expulsion of any of the imploding material. Thus these also cannot be the site of r-process element formation.

Nevertheless, the identification of the pulsars with rotating neutron stars indicates that some supernovae occur in which compression of matter to nuclear densities takes place. These probably result from lower mass stars in which the contraction of the carbon core to highly degenerate densities occurs more slowly. Under these conditions, the thermonuclear ignition of the carbon reactions may fail to raise the electron degeneracy, thus preventing a detonation wave from forming. It is possible that cooling by rapid electron capture will also reimplode expanding material.

The formation of a neutron star certainly appears to be a necessary prerequisite to the formation of a fluid sufficiently neutron-rich to be a possible source of r-process nuclei. However, it is also necessary to find a mechanism by which some of the neutron-rich fluid can be ejected from the star after its compression to these enormously high densities has taken place. If the neutrino-antineutrino energy transport mechanism is ineffective, then it is not possible to envisage any large amount of neutron-rich material being ejected in the process.

An interesting possible mechanism has recently been discussed by LeBlanc and Wilson\textsuperscript{12}). They have carried out a two-dimensional hydrodynamic calculation of the collapse of a star, allowing for the presence both of rotation and magnetic fields in the interior. They found that an initially co-rotating star developed differential rotation after the collapse had taken place. Consequently, the rotational shear in the inner part of the resulting disk wrapped the magnetic field lines into a very tight spiral, creating an enormous magnetic energy close to the central axis.
of the collapsed star. The excess magnetic pressure then caused an expansion of the material along the axis, which, owing to the resulting buoyancy of the material with respect to the local neighborhood, causes the ejection of a jet of material along the axis of the configuration. LeBlanc and Wilson have estimated that this jet contains about $10^{-2}$ solar masses of material, which has been compressed to a sufficiently high density for it to be very neutron-rich. This estimate is in reasonable agreement with an estimate made by two of us\textsuperscript{13}; we find that the amount of r-process material which must be ejected, when a neutron star is formed, to account for the abundances of r-process material in the solar system is $\sim 6 \times 10^{-3}$ solar masses per explosion.

Thus the present astrophysical outlook is that r-process nucleosynthesis will result from the compression of more or less normal material to a range of relatively high densities, at which varying amounts of electron capture will take place, so that a variety of neutron-to-proton ratios will be produced in the compressed material. This material will then expand roughly along an adiabat on a hydrodynamic time scale. The physics which has to be investigated is thus the behavior of the material after such compression, as it expands along the approximate adiabat and is ejected into interstellar space.

4. **THE FORMATION OF SEED NUCLEI**

We have recently carried out a variety of investigations related to the question of the formation of seed nuclei. The first set of investigations was carried out by Truran, Arnett, Tsuı́ruta, and Cameron\textsuperscript{2).} This work was based on the model of a neutrino-antineutrino energy-transport supernova, but the details of the ejection mechanism are not particularly significant insofar as the subsequent physics is concerned: the mechanism suggested by LeBlanc and Wilson\textsuperscript{12} would serve just as well as input data for these calculations. In this study, a largely neutron material was followed as it expanded along an adiabat. It was found that, provided the electrons remained degenerate during the early stages of the expansion, neutrons would remain in excess abundance over protons; further, an equilibrium ratio between neutrons and protons was maintained by weak interactions until the temperature fell to about $2 \times 10^{10}$ K.
At this point, the model calculations suggested that a neutron-to-proton ratio of about 8 would be established. At temperatures below $2 \times 10^{10}$ K, the densities were still high enough for the protons to combine with neutrons, first to form $^4$He, and subsequently to build up into heavy elements during the expansion. At the highest densities considered, the most abundant seed nucleus was found to be $^{78}$Ni. This is a nucleus having a doubly-closed shell of 28 protons and 50 neutrons. Its dominant abundance in equilibrium under these conditions follows in part from the special stability associated with the closed shells and in part from the very large neutron excess in the medium.

The equilibrium abundances of nuclei at a temperature $T = 4 \times 10^9$ K for an expansion adiabat defined by $\rho = 7 \times 10^4 T_9^3$ ($T_9$ is the temperature in billions of degrees Kelvin), is shown in Figure 1. The three most abundant nuclei $^{78}$Ni, $^{79}$Cu and $^{80}$Zn respectively, are all characterized by a "magic number" of neutrons ($N = 50$). The total ratio of neutrons to protons was taken to be $N/Z = 8$ for this material. The free neutron number density is enormous under these conditions, $n_n = 1.88 \times 10^{30}$ cm$^{-3}$.

More recently, two of us (Delano and Cameron$^{14}$) have looked in more detail at reactions involving light nuclei in an expanding medium. We have found that there are a number of reactions involving light nuclei which can take place parallel to the triple alpha reaction in the formation of heavier nuclei; the reaction $^4$He + n $\rightarrow ^9$Be (α,n) $^{12}$C is particularly important. These reactions bridge the gap between helium and carbon more rapidly than the triple-alpha reaction, thus assuring that the heavy nuclei seeds can be built up during the expansion phase of the neutronized material.

At extremely low densities, when very little neutron excess is produced, the most abundant nucleus under conditions of nuclear statistical equilibrium can be expected to be $^{56}$Fe. However, at somewhat higher densities, the most abundant nucleus under conditions of nuclear statistical equilibrium shifts to a nickel isotope somewhat more massive than mass number 56. At the highest densities of interest, the dominant seed nucleus is $^{78}$Ni. Thus, the astrophysical context to which the rapid neutron capture calculations should be related is the expansion of a neutron-rich fluid containing a seed nucleus composed of an isotope of iron or nickel with a mass number depending upon the degree of neutron enrichment.
5. **EQUILIBRIUM FLOWS**

As we have discussed in a previous section, a "steady" flow path can be roughly defined for a particular choice of both the temperature and the free neutron number density. The assumption that nuclear statistical equilibrium is achieved and maintained along a given isotope chain allows one to write the following expression for the ratio of the abundance of isotope $(A + 1)$ to isotope $A$

\[
\frac{\ln n(Z,A+1)}{n(Z,A)} = \ln \frac{w(Z,A+1)}{w(Z,A)} - 78.45 - \frac{3}{2} \ln T_9 + \frac{11.61}{T_9} Q_n(Z,A+1)
\]

where $Q_n(Z,A+1)$ is the binding energy of the last neutron in the nucleus $(Z,A+1)$ and $w(Z,A)$ and $w(Z,A+1)$ are the appropriate nuclear partition functions. For a specified temperature and free neutron number density, $n_n$, the most abundant isotope in equilibrium can be determined from the above relation for each Z-chain.

The neutron binding energies for very neutron-rich nuclei must be determined by a mass formula extrapolation. In these calculations we have employed the predictions of the mass formula of Truran, Cameron and Hilsf. As the individual neutron and proton shell corrections determined in this mass formula exhibit a rather smooth and physically understandable variation with mass number, we believe that extrapolations into the neutron-rich regions far from the valley of beta stability may be somewhat more reliable.

The "flow paths" defined by the positions of these peak nuclei in equilibrium are illustrated in Figure 2 for a temperature $T = 2 \times 10^9$ $\text{O}_K$ and two expansion adiabats:

\[
\begin{align*}
\rho & = 5 \times 10^4 T_9^3 \\
\rho & = 5 \times 10^2 T_9^3
\end{align*}
\]

The free neutron density is taken to comprise 80 percent of the mass. The experimentally determined valley of beta stability and the neutron drip line predicted by the mass formula of Truran, Cameron and Hilsf are also indicated. For both adiabats, the flow paths follow the neutron
drip line rather closely in the light element region. For the heavier
nuclei, the flow paths do not approach the position of the drip line at
this temperature, except at closed neutron shells. The position of
the major r-process peaks at $A \sim 130$ and $A \sim 195$ are also indicated
in this figure, as is that of the lesser peak in the rare earth region.

It is important to note that the positions of the neutron closed
shells at neutron numbers $N = 82$ and 126 shown in this figure are not
consistent with the production of the dominant r-process peaks following
a sequence of beta decays. In both instances, a buildup of something
more than 5 mass units is required relative to equilibrium at $T = 2 \times 10^9$ K.
This feature of the flow paths is appreciably adjusted by variations in
the temperature, as is clear from the results shown in Figure 3. Here
the flow paths defined by the positions of the peak nuclei in equilibrium
for each $Z$-chain are shown for the adiabat $\rho = 5 \times 10^3 T^3_9$ for temperatures
$T_9 = 1, 3$ and 5. At the highest temperature, the closed shell feature
at $N = 126$ does extend beyond mass number $A = 195$ but the fast capture
peak at $A = 130$ is still not explained. Furthermore, our dynamic studies
indicate that, for the adiabats we have considered, equilibrium is rather
well maintained for each isotope chain down to temperatures $T < 2 \times 10^9$ K
in the expansion. The formation of the rapid capture peaks in the vicinity
of mass numbers $A = 130$ and 195, in fact, takes place as a result of the
"freezing" characteristics of this neutron capture process. We shall
consider this problem in detail in the next section.

The important role played by these closed shell features is further
emphasized by a consideration of the average beta decay half-lives for
each $Z$-chain. The mean chain half-life is plotted as a function of
proton number in Figure 4 for the adiabat $\rho = 5 \times 10^3 T^3_9$ at a temperature
of $2 \times 10^9$ K. The longest mean half-lives, approaching $10^{-2}$ seconds,
are those corresponding to $Z$-chains for which the most abundant isotope
contains a magic number of neutrons ($N = 82$ or 126). As these nuclei
at the neutron closed shell positions are the last to undergo beta decay,
it is clear that the flow back toward the valley of beta stability takes
place rather late in the expansion of the supernova material. This
again suggests that the formation of the r-process peaks at mass numbers
$A \sim 130$ and 195 will be sensitive to the details of the late stages of expansion.
6. GENERAL TRENDS FOR R-PROCESS DYNAMICS

In order to understand some of the dynamic features of the rapid neutron capture process under changing conditions of temperature and density, three salient properties of the network must be emphasized.

First: Along a Z-chain, the neutron binding energy of each isotope, on the average, decreases as one moves from the valley of beta stability to the neutron drip line. With neutrons maintaining an approximate equilibrium along each Z-chain, this has the following consequences: a) decreasing temperatures will tend to shift the flow path toward the drip line; b) decreasing densities will tend to shift the flow path towards the valley of beta stability; c) the depletion of neutrons tends to shift the flow path towards the valley of beta stability.

Second: Along a Z-chain, the total beta decay half-life of each isotope, on the average, decreases as one moves from the valley of beta stability to the neutron drip line. Consequently, flow paths which lie far off the valley of beta stability will conduct material up the network (beta decay to successively higher proton numbers) at a much faster rate than flow paths lying close to the valley of beta stability.

Third: Along a Z-chain, beta decay with the delayed emission of up to three neutrons increasingly tends to become the dominant mode of decay for each isotope as one moves from the valley of beta stability to the neutron drip line. This means that beta decays will not proceed along isobar lines far from the valley; instead there will be a "fall back" to lower mass numbers.

Because of the abrupt decrease of neutron binding energies at neutron closed shell positions, these general features are less applicable to a description of the dynamic features near these positions.

7. TEMPERATURE AND DENSITY PROFILES AND SOME PRELIMINARY RESULTS FOR THE R-PROCESS

The temperature relaxation for the expanding r-process gas was adopted from the supernova calculations of Arnett\(^{8,15}\). It is

\[ T_9 = 4 \times \left( \frac{0.03}{0.03 \text{ \text{time}}} \right)^{0.852} \]
For a variety of adiabats in the range,

$$4 \times 10^2 \, T_9^3 \leq \rho \leq 5 \times 10^4 \, T_9^3$$

this profile is such that the temperature falls from $T_9 = 4$ to $T_9 = 2$ in $3.8 \times 10^{-2}$ sec and from $T_9 = 2$ to $T_9 = 1$ in about a tenth of a second.

Starting with free neutrons and $^{78}\text{Ni}$ seed nuclei with number density ratios in the range

$$50 \leq \frac{n}{^{78}\text{Ni}} \leq 130$$

the r-process was studied for the above temperature and density conditions.

At the initial temperature of $T_9 = 4$, the flow (beta decays) of material up the Z-chains progresses slowly. As the temperature falls, the flow path begins to shift toward the drip line and hence enters the region of increasingly rapid beta decays. This is in evidence in Figures 5, 6, and 7, where the total beta decay rate per nucleus is plotted against temperature for different adiabats and different initial ratios of neutrons to seed nuclei. Up until the depletion of the neutrons, these plots indicate that the total beta decay rates, and hence the speed with which material flows up the Z-chains, increases roughly monotonically as the temperature decreases.

In all of the model calculations we have performed to date there has been one recurring feature which persists until the neutrons begin to be depleted. This is the formation of two strong abundance peaks, one at mass number $A = 124$, and one at $A = 190$. The reason is not hard to discover. The $A = 124$ nucleus lies on the $N = 82$ closed neutron shell position and the $A = 190$ nucleus lies on the $N = 126$ closed neutron shell position.

The equilibrium flow paths, displayed in Figures 2 and 3, show that for a broad range of temperatures and densities ($1.1 \leq T_9 \leq 5$ and $4 \times 10^3 \leq \rho \leq 4 \times 10^5$) all material is essentially constrained to flow along these closed neutron shell positions. Beta decaying up the Z-chains along the neutron positions $N = 82$ and $N = 126$, material accumulates when it reaches proton number $Z = 42$ (along $N = 82$) and $Z = 64$ (along $N = 126$). Referring to Figure 4, it will be seen that at
$T_9 = 2$ these two proton numbers correspond to the two slowest beta emitters in the occupied region of the network. Consequently $^{124}\text{Mo}$ and $^{190}\text{Gd}$ have the effect of mass traps in this network.

The dynamics of the r-process is then the following: first, neutron captures on $^{78}\text{Ni}$ seed nuclei, and subsequent beta decays, lead to the formation of the first peak at $A = 124$, and then, as more and more material piles up at $A = 124$, it slowly beta decays, leading to the buildup of the second peak at $A = 190$. By the time the second peak has formed, about 90% of the free neutrons have been captured, and virtually all the nuclei in the network are becoming trapped at $^{124}\text{Mo}$ or $^{190}\text{Gd}$. This may be seen in Figures 4, 5, and 6, where the sudden drops in the beta decay rates are quite pronounced. This was not due to a sudden shifting of the flow paths towards the valley of beta stability, and hence slower beta decay rates, but rather, it results solely from the fact that all the material in the network is quite rapidly winding up in the two slowest beta emitters, $^{124}\text{Mo}$ and $^{190}\text{Gd}$.

Since presumably the two peaks at $A = 124$ and $A = 190$ are the precursors of the r-process peaks at $A = 130$ and $A = 195$, the freeze-out process of falling temperatures and neutron depletion must be critical in determining the 5 or 6 mass unit displacement needed for agreement. The depletion of the neutrons, at temperatures in the range $1.2 \leq T_9 \leq 1.6$, will tend to move the peaks toward lower mass number and toward the valley of beta stability. If the neutrons are not depleted, temperatures falling to $0.5 \leq T_9 \leq 1.1$ will move the peaks through the closed neutron shell positions towards higher mass numbers. Consequently, in order for our calculations to reproduce the two major r-process peaks, the neutrons must not be completely depleted before the temperature falls below $T_9 = 1$.

The temperature and density profiles we have been using do not appear to be capable of accomplishing this. To illustrate, attention is called to Fig.8 in which is plotted the mass fraction in each isobar at three different times in the expansion of an r-process gas. This was for an initial neutron to $^{78}\text{Ni}$ seed nucleus ratio of 83, and relaxation was down the adiabat $\rho = 4 \times 10^3 T_9^3$ from an initial temperature of $T = 4$. At $5.7 \times 10^{-2}$ sec the temperature has fallen to $T_9 = 1.63$ and about one-third of the neutrons have been captured. The first peak at $A = 124$ is forming, and it is apparent that material is being held up at $A = 190$. 
At $7.59 \times 10^{-2}$ sec, in the second section of the figure, the temperature has fallen to $T_9 = 1.37$, and about 85% of the neutrons have been captured. Here $A = 124$ is the dominant peak but the one at $A = 190$ is rapidly building up from neutron captures on the beta decay products of $A = 124$. In the final plot, at $1.01 \times 10^{-1}$ sec, the temperature has dropped to 1.15 billion degrees after the capture of 99% of the free neutrons. The peak at $A = 190$ has moved to $A = 192$, but this is due to the beta decay of $A = 190$ along the $N = 126$ closed shell. Beta decay has reduced the peak at $A = 124$ to the point where it is an order of magnitude less in abundance than the one in the $A = 190$ region.

It is clear, that in the final depletion of the neutrons and subsequent beta decay of this neutron rich material to the valley of beta stability, there will be no net movement of the peaks towards higher mass numbers.

8. **NUCLEAR ENERGY GENERATION AND THE TEMPERATURE RELAXATION**

We have stated that our calculations might produce the r-process element abundances if the temperature falls below $T_9 = 1$ with about 5 neutrons per peak nucleus remaining. Clearly the temperature relaxation we have been using must be speeded up at the point where the neutrons are being depleted. We could arbitrarily amend our temperature relaxation in an attempt to get the desired results, but they would obviously be more believable if the physics of the situation dictated the rate at which the gas cools.

In this regard we point out that nuclear energy generation has been absent from our calculations thus far. With several MeV per neutron capture and the release of up to 15 MeV per beta decay, rough estimates indicate that these sources are capable of generating an amount of energy per unit volume which is comparable to the energy density of the gas. In short, nuclear energy generation may contribute to somewhat higher temperatures and initially slower cooling rates in the expanding gas. But, as the neutrons are depleted, the photon energy from neutron capture will be cut off and as more material becomes trapped in the slow beta emitters at $A = 124$ and $A = 190$ the rate of energy generation from beta decay will die. Consequently, it is not unreasonable
to expect a more rapid drop in the temperature at the time of the neutron depletion, if nuclear energy generation is significant.

The importance of nuclear energy generation and the possibility of having an abrupt temperature drop to below $T_g = 1$ with four to five neutrons remaining per peak nucleus is presently under investigation. To test the idea, however, and also to demonstrate the effects of delayed neutron emission when material beta decays to the valley of beta stability, we have irradiated the two peaks at $A = 124$ and $A = 190$, with about three neutrons for each peak at a temperature of $T_g = 0.519$ and a mass density of $6.97 \times 10^3$ gm/cm$^3$. We have also included two minor peaks at $A = 138$ and $A = 146$, positions which on the average had high abundances in our calculations, in order to see the effects of the freezing process on the intervening material between $A = 124$ and $A = 190$.

Figure 9 displays the mass fraction in an isobar as a function of mass number at three different times in the irradiation. The first plot at 0.0 sec shows the initial abundance peaks to be subjected to neutron capture. The second plot at $10^{-9}$ sec shows these peaks after the capture of virtually all the neutrons. The $A = 124$ mass peak has been pushed to $A = 130$, where it should be, but the $A = 190$ peak has been driven to $A = 206$, 11 mass units beyond where it should be. At this point all the nuclei are very neutron-rich and are just beginning to beta decay to the valley of beta stability. In the last plot, at $4.5 \times 10^{-3}$ sec, beta decay to the valley is about two-thirds complete and several effects are quite noticeable.

First, the peaks at $A = 130$ and $A = 206$ have fallen back to $A = 128$ and $A = 204$, respectively. This is due mostly to delayed neutron emission, but photoneutron reactions may be partially responsible. Second, we have formed the r-process peak at $A = 164$. This peak has been a recurrent feature in the beta decay, with delayed neutron emission, of material to the valley of beta stability. It does not appear to be a sensitive function of the way in which material between mass numbers $A = 124$ and $A = 150$ is distributed. As long as some material exists in this region it always seem to be formed into this broad shallow peak as the valley is approached. The third point worth mentioning is the smoothness of the final mass fraction curve and the rather uniform small odd-even differences. This is consistent with the qualitative features of the r-process abundance
curves. Our results suggest that it is due almost entirely to the smearing out effects of delayed neutron emission.

9. DISCUSSION

The present studies are obviously still preparatory to a full investigation of the dynamics of the r-process. It is clear that a final study will require an improved treatment of several features of the process:

1. It will be necessary to investigate the behavior of a variety of true adiabats at a range of expansion rates. The calculation for a given adiabat must consistently determine the initial neutron-proton ratio and the abundance distribution of seed nuclei formed in nuclear statistical equilibrium, and the changes in internal energy due to nuclear reactions must be calculated.

2. While we are reasonably satisfied with the predictions of the mass formula utilized in this work, the same is not true of the calculated beta decay rates. An improved treatment of the beta strength functions is needed.

3. The effects of fission have so far been neglected. It will be necessary to estimate both the points of termination of neutron capture by neutron-induced fission, and the effects of fission competition with beta decay. The latter point is especially important for the problem of determining whether r-process products can decay into the superheavy island of stability.

The present calculations have shown an approximate, but not detailed, agreement between calculated and observed r-process abundances. One encouraging result of the calculations is the demonstration that the final frozen r-process abundances have a reduced odd-even fluctuation as a result of multiple delayed neutron emission, in accord with observation.

ACKNOWLEDGMENTS

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REFERENCES


Fig. 1  The abundances of nuclear species in nuclear statistical equilibrium in matter with a large neutron excess, peaked at $^{78}\text{Ni}$. 

$T = 4 \times 10^9$ K
$
\rho = 4.48 \times 10^6$ gm/cm$^3$

$n_n = 1.88 \times 10^{30}$ cm$^{-3}$

log ABUNDANCE (particles/cm$^3$)

MASS NUMBER
Fig. 2 R-process paths in a nuclide chart under steady flow conditions at $T_9 = 2$, for two different densities. Note that the ordinate and abscissa scales differ by a factor two.
Fig. 3 R-process paths in a nuclide chart under steady flow conditions at several different stages of an adiabatic expansion. Note that the ordinate and abscissa scales differ by a factor two.
Fig. 4 The average beta decay lifetime of the nuclei with given proton number at $T_9 = 2$ and $4 \times 10^4 \text{ gm/cm}^2$ with 80 percent free neutrons by mass.
Fig. 5 Change in the total beta decay rate per nucleus during expansion of a high adiabat for an initial ratio of neutrons to $^{79}$Ni nuclei of 122.

Fig. 6 Change in the total beta decay rate per nucleus during expansion of a low adiabat for an initial ratio of neutrons to $^{79}$Ni nuclei of 122.

Fig. 7 Change in the total beta decay rate per nucleus during expansion of a low adiabat for an initial ratio of neutrons to $^{79}$Ni nuclei of 83.
Fig. 8 Evolution of the r-process abundances during the expansion of a low adiabat.

Fig. 9 R-process freezing of a selected "seed" nuclei at a low temperature.
NUCLEI IN NEUTRON MATTER

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Most of the matter in a neutron star consists of neutrons, together with a small percentage of protons, either uniformly mixed in or concentrated in nuclei. This is a type of nuclear matter, and the standard methods of nuclear matter theory (NM) are applicable. We have used this theory, with the soft core potential of Reid \(^1\) for the interactions between nucleons. The features of this theory are as follows:

1. The Reid soft core (RSC) potential fits the nucleon-nucleon scattering data, and the properties of the deuteron, essentially within experimental error. Other potentials \(^2\) do the same, but only a few of these give also good fits to nuclear matter and those which do so, do in general not differ greatly from RSC.

2. When all known corrections are taken into account, RSC with NM theory gives the correct binding energy of nuclear matter (16 MeV per particle), and a slightly too high density (about 0.22 nucleons per fermi cube instead of 0.17).

3. In NM theory, one first calculates the correlation (wave function) of two interacting nucleons; this "pair approximation" gives a binding energy \(^3\) of 11 MeV per par-
ticle. The remaining 5 MeV come from various corrections, the most important of which are

a) the correlations between three \(^4\) and four \(^5\) interacting nucleons

b) the added forces between three nucleons \(^6\) due to meson theory

4. The theory also gives excellent results for the density distribution and energy levels of finite nuclei \(^7\).

Nemeth and Sprung \(^8\) (NS) have calculated the properties of neutron matter, using a somewhat older calculation of nuclear matter which gave only about 9 MeV binding in the "pair approximation". In order to take into account the corrections 3a, b above, they (arbitrarily) increased the attractive potential energy in the pair approximation, so as to force the binding to be 16 MeV. They used two alternative prescriptions, viz.

a. All interactions were increased by 22%

b. The interactions in the states of isospin \(T = 1\), especially the singlet even states, were left unchanged, while the interactions in \(T = 0\) states, especially triplet even, were increased by 54%.

In a first treatment \(^9\) of the subject NS choice a was adopted because it appeared more conservative.

To describe the nuclei we employed the semi-empirical mass formula \(^1\) and the equation of \(\beta\)-equilibrium, together with a minimum condition for the total energy:
\[ E = -c_1 A + c_2 Z^2 A^{-1/3} + c_3 \left( \frac{N-Z}{A} \right)^2 + c_4 A^{2/3} \] (1)

The Fermi energy, as the energy of the highest occupied state, is given by

\[ \mu_n = \frac{\partial E}{\partial N}, \quad \mu_p = \frac{\partial E}{\partial Z} \]

for neutrons and protons respectively. Neglecting the small n-p mass-difference the equation of \( \beta \)-equilibrium gives

\[ \mu_n - \mu_p = \mu_e \] (2)

Under the conditions considered the electrons form a highly relativistic Fermi gas

\[ \mu_e = m_e c^2 \sqrt{1 - \frac{1}{N}} \] (3)

\[ x = \frac{Z}{A}, \quad \rho_N: \text{number density of nucleons in nuclei.} \]

To \( E/A \) we add the electron kinetic energy per nucleon \( \frac{3}{2} \frac{c_n}{c} \mu_n \) to get the total energy per particle. For given \( k_N (\rho_N \equiv \frac{k_N}{l s^{-1}}) \) this is a function of \( x \) only and the most stable nucleus \( \Omega_N \) is defined by \( \frac{dE_{\text{tot}}}{dx} = 0 \).

Now all quantities are known functions of \( x \):

\[ A = \frac{C A}{2C_2} x^{-2} = 12.5 x^{-2} \] (4)

\[ \mu_n = -16 + 3.88 x^{2/3} + 24(1-4x^2) \] (5)

\[ \frac{d\mu_n}{dx} < 0 \quad \text{if} \quad x > 0.04, \quad \mu_n = 0 \quad \text{for} \quad x = 0.22. \]

So at this point, corresponding to a density \( \rho_1 = 2.8 \cdot 10^{11} g/cm^3 \).
free neutrons begin to appear.

Nuclei continue to exist in equilibrium with neutron matter. The properties of neutron matter are obtained from NS $\alpha$. The condition for equilibrium is the equality of the Fermi energies of neutrons in nuclei and outside: (density of neutron gas $\rho_{Gn} = \frac{Gn}{\hbar^2}$)

$$\mu_{\text{Fermi}} = \mu_{\text{Fermi}}(\rho_{Gn})$$

Since $\mu_{\text{Fermi}} > 0$, (6) indicates that not all the neutrons in a nucleus are firmly bound, some can pass freely in and out, but the nucleus still forms a compact structure in the surrounding uniform neutron gas.

We obtain from (6) a relation between the electron density of (3) and the neutron matter density, which permits us to plot the proton chemical potential (the proton Fermi energy) both in nuclei and in neutron matter against the density of electrons (fig. 1). The two graphs cross rather sharply, defining a well-marked melting point for the nuclei, for, remembering the meaning of the chemical potential, all the protons will be in the phase which has the lower $\mu_{p}$. For $k_N < 0.235$ fm$^{-1}$ all protons will be inside nuclei, since $\mu_{pN} < \mu_{pG}$. For $k_N > 0.235$ fm$^{-1}$ all of them will be in the neutron gas. Table 1 shows how the nuclei behave up to this point. $y$ gives the fraction of nucleons bound in nuclei. The melting of nuclei is a completely sharp transition with the assumptions made. If one takes the small kinetic energy of the protons into account $^9$ then there is a small region
of density where this transition takes place (1%) if, however, you consider the effect of the Coulomb lattice formed by the nuclei the transition is again completely sudden. (Baym, Bethe, Pethick, see later cit.) The lattice interaction depresses the energy much more if all charge is concentrated in nuclei; partial assembly in nuclei is a disadvantage. Since \( \frac{1}{R} \sim \frac{\varepsilon}{R}^{2/3} \) it is a relative disadvantage to "evaporate" only a fraction of the protons.

In \(^{10}\) corrections to (3) due to a) the effect of external pressure on the nuclei immersed in neutron matter, and b) the decrease of the surface energy due to the neutron matter outside, have been estimated. The variation of all interesting quantities turned out to be small, less than 10% always.

There are, however, good arguments against the conservative choice of NS \(_a\), and in favour of \(_b\). The corrections \( \xi \) a and especially \( \xi \) b (above) depend chiefly on tensor forces which are important in triplet even states occurring in n-p interactions. In pure neutron matter, only \( T = 1 \) states can occur, hence it is reasonable to take for these the results of the NM calculations without any correction. This is NS \(_b\), and this increases the energy, hence also \( \mu_{n_G} \) and pressure of the neutron gas at given density.

Therefore G. Baym, H.A. Bethe, C. Pethick reinvestigated the problem dealt with in \(^9\)\(^{10}\) x). They preferred to use the most up to date and more accurate than NS \(_b\) NM calcul-

\[x\) The permission of Baym, Bethe, Pethick to communicate some results of their work done at Copenhagen prior to any written account is gratefully acknowledged.]
lations by Siemens 3). In the course of their calculations it turned out that nuclei persist up to rather high density, that the neutron Fermi energy increased above 12 MeV. Since the semi-empirical mass formula leads to a maximum \( \mu_{NN} \approx 8.3 \text{ MeV} \), (3) with the corrections already tried in 10), was investigated and finally replaced by a new formula. NM was calculated at all densities and values of \( x \), and then corrected for finite nuclei by adding a Coulomb term and a surface term. The advantage of this procedure is that pressure at given \( x \) automatically comes out from the calculation. This allows the necessary treatment of pressure, which can of course no longer be viewed as a small perturbation, if nuclei exist up to high densities and substantial pressure.

Thus the final formula consisted of 4 terms:

1) a bulk energy \( W(k_z, x) = -W_0 + k^2 (k_z^2 - L_0^2) + C k_z^2 y^2 \)

\( k_z \): inside density in nuclei, \( y = 1 - 2x \)

2) a Coulomb energy term \( E_{\text{coul}} \approx 0.569 \frac{k_z^2}{y} \frac{x^2}{A^{2/3}} \)

3) an energy term considering the finite extension of the nuclei, which build a Coulomb lattice

4) a surface energy term \( E_{\text{surf}} \approx \frac{W}{W_0} A^{-1/3} \)

They find that most quantities depend very sensitively on the surface energy term. This term should, however, be regarded with caution. So it turns out that with the above choice \( \mu_{NN} < \mu_p G \), always, that is "proton drip" never occurs. Nuclei disappear instead by the following mechanism:

As long as the pressure of a pure neutron gas \( P_n \) is greater
than the pressure $P_{n,p}$ of a mixture of neutrons and protons

$$P_h > P_{n,p}$$

nuclei exist. They have disappeared in regions, where

$$P_n < P_{n,p}$$

The value of $Z$ remains essentially constant at about 40-42 from neutron drip ($x \approx 0.31$) down to $x \approx 0.1$, and for smaller values of $x$, $Z$ rises fairly rapidly reaching a value of 100 or so at $x = 0.06$. By the time $x$ falls to 0.06 the nuclei fill about 1/2 the space, and the total mass density is about $2 \times 10^{14}$ gms/cc. It seems quite likely that a transition to the uniform neutron-proton solution will take place at about this density; however the result is quite sensitive to the size of the surface energy. They find no proton drip, although this would be more likely to occur if they were to increase the surface energy.

At $x = 0.06$ the value of $k_z$ is 1.26, and $k_G$ is 1.15, so the interior and exterior densities are quite similar. The electron chemical potential at $x = 0.06$ is 100 MeV, $\mu_n$ is 17.8 MeV, and the number of nuclei is $3.6 \times 10^{33}$ nuclei/cm$^3$, the pressure 0.81 Mev/fermi$^3$, and the mass density in nuclei and electrons is $1.2 \times 10^{14}$ gm/cc.
Table I

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\[ k = \frac{m_{\pi} \bar{q} c}{12 c^2} \bar{N} \]

Fig. 1
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MASSES AND DEFORMATIONS OF HEAVY AND SUPERHEAVY NUCLEI

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SUMMARY

A calculation of ground state masses and deformations has been performed. It is based on the shell correction method developed by Strutinsky involving a modification of the liquid drop mass formula and shell and pairing corrections from a single-particle model. The single-particle model employed is the modified harmonic oscillator potential developed by Nilsson and coworkers with two deformation coordinates $\epsilon$, $\epsilon_4$ representing elongation and neck-in respectively. The parameters in the model are adjusted to reproduce optimally the observed single particle level order in the rare earth and actinide regions and are not fitted directly to the nuclear masses. Shell corrections are obtained by means of the sixth-order Strutinsky prescription, and the pairing corrections are calculated by the BCS method. The liquid parameters employed are those of Myers and Swiatecki. The resulting masses and deformations show very good agreement with experimental values.

Actually the agreement in the Pb region seems excessively good in view of the fact that the spherical shell model spectra around Pb are relatively poorly reproduced.
Tables of calculated masses, shell corrections, deformations as well as particle separation energies of protons, neutrons, alpha particles, and Q-values for beta-decay are presented for heavy and super-heavy nuclei (150 < A < 330).
THE SYNTHESIS AND THE SEARCH IN NATURE FOR SUPERHEAVY ELEMENTS

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Joint Institute for Nuclear Research,
Dubna, USSR.

Abstract and manuscript not submitted.

Professor Flerov covered topics which may in part be found described in the following references:

1. G.N. Flerov and V.A. Druin
   The synthesis and properties of transfermium elements,
2. G.N. Flerov and S.A. Karamyan
   B.V. Fefilov and E.D. Vorobyev
   Neutron detector for the search for superheavy elements in nature,
   JINR, Dubna, report to be published.
STABILITY OF HEAVY NUCLEI

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Abstract

The changes in binding energy of $^{298}_{114}X$ under deformation into prolate spheroidal shapes have been calculated using the statistical theory of nuclei. The density is restricted only by the requirement that each equidense surface is a spheroid. It is found that the energy is minimized when the density has an approximately uniform surface thickness around the nucleus. To first order the change in binding energy is given well by the liquid-drop model, but at large deformation the statistical theory predicts less drop in energy and, consequently, a slightly more stable nucleus.

Realistic shell-model potentials are used to test Strutinsky's method. It is found that the results are strongly dependent on the shell-smearing parameter as well as the order of the curvature correction for smearing. No unambiguous values can be obtained for the shell correction, even with the effect of the continuum properly included. As a result, the Strutinsky theory should be regarded only as a recipe to be used in combination with Nilsson potentials.
NUCLEAR DEFORMATION ENERGY IN THE TWO-CENTER SHELL MODEL

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Abstract
As a first application of the two-center single-particle shell model developed recently, we have calculated deformation energies and fission barriers for several nuclei. The deformation parameter $\xi_0$ is defined as half the distance $2z_o$ between the two wells of our single-particle potential, divided by the radius $R_o = r_oA^{1/3}$ of the original nucleus in the spherical state: $\xi_0 = z_o/R_o$. The equipotential shapes

are axially symmetric as well as reflexion-symmetric, the latter property corresponding to a symmetric fission mode. We write the deformation energy as

$$E_{\text{def}}(\xi_0) = \Delta E_p(\xi_0) + \Delta E_n(\xi_0) + \Delta E_{\text{Coul}}(\xi_0),$$

where the proton and neutron energies are single-particle sums. The Coulomb energy is that of a homogeneously charged drop with a sharp boundary coinciding with a volume-conserved equipotential. All energies are normalized so that $\Delta E = 0$ in the spherical state $\xi_0 = 0$.

In any phenomenological shell model, the volume conservation condition essentially determines the density distribution and thus the over-all behaviour of deformation-dependent properties (like total energies) of the nucleus. So we paid special attention to it. Now the double-oscillator does not permit "simultaneous volume conservation". By this we mean that the equipotential volume cannot be
kept constant, as the nucleus deforms, for all equipotentials at the same time. Hence, three different lines of approach are considered:

(i) volume conservation for the equipotential supposed to coincide with the nuclear surface,
(ii) average volume conservation, where inner equipotentials are also accounted for,
(iii) a modification of the double-oscillator, which has the property of simultaneous volume conservation.

In all cases, the Hamiltonian includes $\mathcal{F} \cdot (\not\! V \not\! \tau \not\! p)$ and $(\not\! V \not\! \tau \not\! p)^2$ terms. The results of calculations with method (i) yield curves $E_{\text{def}}(\delta \mathcal{E})$ showing an unreasonably deep minimum. Methods (ii) and (iii), however, demonstrate that fission barriers of the right order of magnitude can be obtained from a two well shell model, the most realistic ones being those computed by method (iii). In contrast to this, single well potentials give unrealistic or even diverging results for large deformations. Strutinsky's renormalization prescription [adding a shell correction to $E_{\text{def}}(\text{liquid drop})$] is usually thought to remedy this behaviour. In that case, however, care must be taken to select realistic drop shapes, which for large deformations do not coincide with single well equipotentials. If the latter are chosen as drop shapes, unrealistic deformation energies and barriers are encountered again.

We also calculated Strutinsky shell corrections $\delta U$ for the three volume conservation methods mentioned above and compared them to the single-particle results. Since the renormalized deformation energy exhibits a second minimum for $^{236}\text{U}$, while the latter method yields none, the correctness of both approaches can be checked against experimental findings.
NEW ASPECTS OF INTERMEDIATE NUCLEI

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1. INTRODUCTION

In a broad sense all even-even nuclei are intermediate nuclei since they are between the two limits of harmonic vibration and rigid rotation. However, because of limitations of time and space, this paper will be confined to those even-even nuclei which are near a region of transition or change in nuclear shapes. Shape transitions can be grouped into two categories:

i) spherical to deformed shapes, for example N = 88 - 90
   \( Z \approx 62 \).

ii) prolate to oblate shapes, for example Z = 76 - 73
    \( N \approx 114 \).

Similar transitions are expected to occur for many other combinations of proton and neutron numbers. Studies dealing only with beta-stable nuclei would give essentially a one-dimensional view of shape transitions. Studies of beta-unstable nuclei are essential for getting the complete picture; they are bound to reveal many new regions of shape transitions and intermediate nuclei.

Theoretical techniques, employed to study intermediate nuclei, are reviewed in sect. 2. Transitions from spherical to deformed shapes are discussed in sect. 3, those from prolate to oblate in sect. 4.

*) Work supported by the U.S.A.E.C. under contract with Union Carbide Corporation.
2. THEORETICAL TECHNIQUES FOR INTERMEDIATE NUCLEI

Since these nuclei exhibit large deviations from both the harmonic vibrator model and the rigid rotator model, a perturbation expansion around the equilibrium shape is not valid. Theoretical techniques for treating the intermediate nuclei in the most complete and practical manner can be divided into two categories.

i) Boson expansion method. The microscopic Hamiltonian, as well as the nuclear wavefunction, is expanded in a basis of quadrupole phonons or bosons, defined in the lab system. This method has been stimulated greatly by the work of Beliaev and Zelevinski\textsuperscript{1)} who take into account the fact that a boson built from fermions does not obey the usual commutation rules. B. Sørensen\textsuperscript{2)} has developed this method further and has applied it to isotopes of Cd, Sn, Te, Sm, and Pb. In principle, this method is capable of including adiabatic as well as non-adiabatic effects. However, it is an expansion around the spherical shape and the question of convergence is not quite settled yet.

ii) Non-linear treatment of coupling between rotations, $\gamma$-vibrations and $\beta$-vibrations. This method has the great advantage of relating complex nuclear behaviors, which require lengthy computations, to comparatively simple pictures of the nuclear potential energy surface and mass parameter functions. For this purpose, it is convenient to work in the intrinsic or body-fixed system, where the five quadrupole variables are: $\beta$, $\gamma$ (the two shape variables) and $\theta, \varphi, \psi$ (the three Euler's angles relating the intrinsic axes to the lab axes). However, one must pay a price for this
convenience. The transformation from the lab system to the intrinsic system is not unique, hence certain symmetry conditions must be imposed on the Hamiltonian and the trial wavefunctions.

These symmetry conditions were discussed already in the 1952 paper of A. Bohr. However, a method, which is applicable to harmonic vibrators as well as rigid rotators and which satisfies all the symmetry conditions, was developed only during the past few years. In this method, called "complete numerical solution of Bohr's collective Hamiltonian," the collective Hamiltonian is written in its most general form:

\[
H_{\text{coll}} = V(\beta, \gamma) + \frac{1}{\hbar^2} \sum_{k=1}^{3} \delta_k(\beta, \gamma) a_k^2 + \frac{1}{\hbar^2} B_{\beta\beta}(\beta, \gamma) \beta^2 + B_{\beta\gamma}(\beta, \gamma) \beta \gamma + \frac{1}{\hbar^2} B_{\gamma\gamma}(\beta, \gamma) \gamma^2. \tag{1}
\]

This Hamiltonian consists of seven completely arbitrary, except for the symmetry conditions, functions of deformation variables \( \beta \) and \( \gamma \); the potential energy \( V \); the three moments of inertia \( \delta_1, \delta_2, \delta_3 \); and the three vibrational mass parameters \( B_{\beta\beta}, B_{\beta\gamma}, \) and \( B_{\gamma\gamma} \).

The collective Hamiltonian is quantized and the corresponding Schrödinger equation is written as

\[
H_{\text{coll}} \Psi_{IM}(\beta, \gamma, \Theta, \Phi, \Psi) = E_I \Psi_{IM}(\beta, \gamma, \Theta, \Phi, \Psi), \tag{2}
\]

\[
\Psi_{IM} = \sum_{K} A_{IK}(\beta, \gamma) \Phi_{MK}(\Theta, \Phi, \Psi). \tag{3}
\]

\(^*) An alternative is to work in the phonon basis and transform the final results to the \( \beta-\gamma \) plane \(^2\).
where φ is a symmetrized combination of the standard ξ-functions depending on the Euler's angles and $A_{IK}^{-(\beta,\gamma)}$ is the state-dependent probability amplitude (wave function) for the shape $(\beta,\gamma)$. In the rigid rotator model$^5$, the nucleus is assumed to have a fixed shape and K is assumed to be a good quantum number. The wavefunction $A_{IK}$ is replaced by a constant and the summation over K reduces to one term. There follow the $I(I+1)$ rule for the energy levels and the Alaga rules for the branching ratios. In the harmonic vibrator model, the Schrödinger equation (2) separates into a $\beta$-dependent part and a $\gamma$, K-dependent part. In this model, K is not a good quantum number (except for I = 0 and 3). The ground state ($0^+$) wavefunction$^3$ is a Gaussian centered around $\beta = 0$. The first excited state ($2^+$) wavefunction has a K = 0 component, $\beta \cos \gamma$ times the Gaussian; and a K = 2 component, $\beta \sin \gamma$ times the Gaussian$^4,6$. Such wavefunctions will be used for comparison in sect. 4.

In general, the Schrödinger equation (2) consists of a set of coupled (due to K mixing) or non-linear partial differential equations. These equations are solved by performing numerical integrations over a $\beta,\gamma$ mesh$^4$. This method of complete numerical solution is quite general and is independent of the method of determination of the seven functions of deformation.

Two alternative methods with similar properties have recently been developed by Dussel and Bes$^7$, and by Gneuss, Mcasel, and Greiner$^8$. By expanding the collective wavefunction in a basis of about 20-30 phonons, these authors obtain good convergence for well-deformed nuclei. However, in their present forms, these
methods are limited by the assumption that all of the six inertial functions of Eq. (1) can be expressed in terms of a single mass parameter $B$, which is independent of deformation. The six functions are written in this assumption as

$$B_k(\beta, \gamma) = \mathcal{J}_k(\beta, \gamma) \left[ 4 \beta^2 \sin^2 \left( \gamma - \frac{2}{3} \pi k \right) \right]^{-1} = B \quad (4)$$

$$B_{\beta \beta}(\beta, \gamma) = B_{\gamma \gamma}(\beta, \gamma) = B \quad (5)$$

$$B_{\beta \gamma}(\beta, \gamma) = 0 \quad (6)$$

This assumption is valid only in the limit of zero deformation. True, one can reproduce energy level schemes of various types intermediate between rotational and vibrational by considering the $\beta, \gamma$-dependence of the potential function alone. But many subtle details of the wavefunctions of intermediate nuclei depend in an essential way on the $\beta, \gamma$-dependence of the mass parameters (see refs. 9, 11) and sect. 4).

Once we adopt this point of view, it becomes almost obvious that we must use a microscopic theory to determine the seven functions of deformation. One might say that we trade the parameters of the purely "phenomenological" treatments 7, 8 of the seven functions for the parameters of the microscopic treatment such as the single particle energy levels. However, the microscopic theory provides some useful restraints and connections to a variety of experimental data (for example, the single particle levels are related to the spins of odd-$A$ nuclei) which are completely missing in a non-microscopic theory.
The most widely used microscopic theories are based on the Nilsson model\textsuperscript{10}). This model allows for the treatment of configuration mixing among a large number of single particle states, hence it is not necessary to divide the nucleus into an essentially "inert" core and a cloud of "active" nucleons. However, the consequences of using a non-rotationally invariant Hamiltonian are not clear, at least to the present author. It would be interesting to calculate the complete $\beta-\gamma$-dependence of the seven functions and the spectra of intermediate nuclei.

An alternative method in this context is to employ the pairing-plus-quadrupole model\textsuperscript{11}). Actually, the expressions for the seven functions are formally similar to those obtained by combining the Mottelson-Nilsson method, pairing theory, and the Cranking Model (the most complete treatment was first given by Bès\textsuperscript{12}). However, there are several important differences in principle as well as practice.

The pairing-plus-quadrupole model is closer to "first principles" in the sense that rotationally invariant Hamiltonians are employed and that a time-dependent, self-consistent theory (called time-dependent-Hartree-Fock) is used. The condition of volume conservation is not imposed on the system, but it turns out that the expectation value $\langle r^3 \rangle$ is independent of deformation to about the same degree of accuracy as in the Nilsson Model. The practical advantage of this method is that the calculated seven functions of deformation satisfy automatically all the symmetry conditions.

We have to pay a price for this advantage. The nucleus has to be divided, as in all "shell model" calculations of spectra of heavy
nuclei, into a core of "inactive" nucleons and a cloud of "active" nucleons. Although two complete oscillator shells for neutrons and for protons are taken into account (in the calculation for $^{196}\text{Pt}$, the core has 110 particles; the cloud has 86 particles), the "core polarization" effects are substantial. The quadrupole force strength and the inertial parameters are renormalized by about a factor of two. The quadrupole force strength $\chi$ can be chosen so as to yield the "experimental" deformation of a well-deformed nucleus, but this procedure is not satisfactory for intermediate nuclei. The alternative procedure of choosing $\chi$ to fit the energy of the first $2^+$ state is somewhat risky since this energy depends also on the inertial parameters. The question of renormalization of the inertial parameters is not completely settled yet (we shall return to this question in sect. 4). Therefore, the numerical results should be considered to be somewhat tentative. However, this method has predicted many trends of properties of intermediate nuclei many of which have already been verified experimentally. Some of the results based on the pairing-plus-quadrupole (PPQ) model are discussed in the next two sections.

3. TRANSITIONS FROM SPHERICAL TO DEFORMED SHAPES

The best known example is the $N = 88 - 90$ transition in the Sm, Gd region, where nuclei with $N = 88$ have vibrational-type spectra and nuclei with $N = 90$ have rotational-type spectra. The implication is that the equilibrium shape of $N = 88$ nuclei is spherical, but that of $N = 90$ nuclei is deformed. However, the PPQ model indicates that the change in equilibrium shapes is not so sudden. Figure 1 shows the potential energy function $V(\beta, \theta)$ for the $N = 84 - 92$ isotopes for Sm. The equilibrium shape changes from spherical to deformed in going
from $N = 86 - 88$, but the $N = 88$ nucleus is not truly deformed since the energy of zero point motion (shown in Fig. 1 by a dashed line) is larger than the energy of deformation. Evidently, the "theoretical" criterion for the spherical $\rightarrow$ deformed transition is

$$
\xi_o - V_{DS} = 0
$$

(7)

where the energy of zero-point motion is given by

$$
\xi_o = \langle 0^+ | H | 0^+ \rangle - V(\beta_{\text{min}}, \nu_{\text{min}})
$$

(8)

and the energy of deformation by

$$
V_{DS} = V(0,0) - V(\beta_{\text{min}}, \nu_{\text{min}})
$$

(9)

The usual "experimental" criteria for the spherical $\rightarrow$ deformed transition are (i) a sharp increase in the $B(E2; 0^+ \rightarrow 2^+)$ value, and (ii) a sharp decrease in the energy $E_2^+$. However, these two criteria can also be satisfied by a prolate-oblate transition (see sect. 4). Therefore, an additional criterion is needed. In a region of spherical $\rightarrow$ deformed transition, both $\beta$ and $\gamma$ bands must be low lying (or all three members of the two phonon triplet must occur in a group). In a region of prolate $\rightarrow$ oblate transition, only the $\gamma$ band would be low lying (or the $2^+$, $4^+$ members of the two phonon triplet would be practically degenerate, but the $0^+$ member would lie higher).

The observed\textsuperscript{13} energy levels of $^{152}_{62}$Sm are shown to the right of Fig. 2. In contrast to the well-deformed nuclei in the middle of the rare-earth region, the $\beta$-band is lower than the $\gamma$-band. Also,
substantial $\beta$-$\gamma$ band mixing has been reported in this nucleus. The reason can be understood from the locations of different calculated levels in the potential well shown to the left of Fig. 2. Although the ground band ($0^+$ to $6^+$) is contained completely inside the potential well, the nucleus in the $\beta$,$\gamma$-band states has much more freedom to explore prolate, spherical, and oblate shapes. Also, the $2_\beta$ and $2_\gamma$ states are quite close to each other. In fact, in its present, preliminary form, the calculation gives somewhat too much $\beta$-$\gamma$ band mixing. In order to alleviate this problem, a better method of including the core contribution (see sect. 4) to the mass parameters is being investigated.

The potential curve of Fig. 2 is typical in the sense that the potential curves of most nuclei have a maximum at $\beta = 0$ and two "minima" - one for a prolate shape and another for an oblate shape. However, this one-dimensional projection of the two-dimensional $\beta$-$\gamma$-dependent potential surface is somewhat misleading. Usually, the lower of the two "minima" is a true minimum but the upper one is a saddle shape - a minimum in the $\beta$ direction but a maximum in the $\gamma$ direction. In other words, the nucleus does not have to go through the spherical "barrier"; it can go directly from a prolate shape to an oblate shape in the $\gamma$ direction. This is evident from Fig. 3 which gives the full two-dimensional potential surface (contour plot of equipotential surfaces) for the same nucleus.

The theoretical and experimental E2 matrix elements are given in Figs. 4 and 5. The major discrepancy is in the $\beta$-$\gamma$ band transitions, as

*) The quadrupole force constant used for this potential is about 3% lower than that used for the potentials of Fig. 1. This change was made in order to make a better fit to the experimental levels. However, it does not alter the qualitative conclusions based on Fig. 1.
discussed above. The static quadrupole moments, magnetic moments, 
E2/M1 mixing ratios, monopole transition moments and Isomer shifts are 
compared in Fig. 6. The major discrepancy is in the $6(E2/M1, \beta^{-} \rightarrow 4)$. 

Although the results discussed above are of a preliminary nature 
(partly because results for $^{150}\text{Sm}$ are not complete yet), it can be 
concluded that transition in the N = 88 - 90 region is not as sudden as 
expected earlier. True, the condition (7) is satisfied in this region 
(the quantity on the left hand side of the equation changes sign), but 
$^{150}\text{Sm}$ is not a harmonic vibrator, $^{152}\text{Sm}$ is not a rigid rotator. Devia-
tions from these two limits have led to the idea of "co-existence" of 
spherical and deformed minima, the former being lower in N = 88, the 
latter in N = 90. However, the present calculation indicates that such 
a picture is incorrect.

A question about the N = 88 - 90 transition arises. Does 
this transition also occur away from the region of beta-stability 
(Z < 60, Z > 64)? Since the nuclear deformability [measured, for 
extample, by the second derivative, $(\partial^2\nu/\partial\beta^2)_{\beta=0}$] is determined by neu-
trons as well as protons, we can expect this transition to be shifted 
towards smaller N for Z > 64 and towards larger N for Z < 62.

4. TRANSITIONS FROM PROLATE TO OBLATE SHAPES

These transitions are characterized by low lying $\gamma$-vibrations and 
a change of sign of $Q_{2+}$. Compared to the spherical $\gamma$-deformed transi-
tion, discussed previously, change in the magnitude of nuclear deforma-
tion is small and the $\beta$-vibrational state (0$^{+}$) lies higher. Such 
a transition on the Os-Pt region was predicted in 1966 and has since 
been confirmed by Saladin et al.\cite{17}

The calculated potential wells for $^{182}\text{W}$, $^{190}\text{Os}$, and $^{196}\text{Pt}$ are shown 
in Fig. 7. Near the line of beta-stability, tungsten nuclei are prolate,
platinum nuclei are oblate, the osmium nuclei are in between. According to the potential wells, the prolate-oblate transition takes place near $^{190}_{\text{Os}}$. However, this transition is shifted somewhat by the kinetic energy terms so that the quadrupole moment of the first $2^+$ state changes sign at $Z = 76 - 78 (N = 114, 116)$. This is shown in Fig. 8 where the ratio $Q_{2^+} / Q_R$ is plotted against $V_{PO}$ defined as

$$V_{PO} = V \text{ (oblate 'minimum') } - V \text{ (prolate 'minimum')} \begin{cases} \text{if } \beta_{\text{min}} \neq 0 \\ \text{if } \beta_{\text{min}} = 0, \quad (10) \end{cases}$$

$$\beta_{\text{rms}} = \langle 0^+ | \beta^2 | 0^+ \rangle^{\frac{1}{2}} \approx (3Zn^2/nR)^{-1} [B(E2; 0^+ \rightarrow 2^+)]^{\frac{1}{2}}. \quad (11)$$

The "rotational" quadrupole moment $Q_R$ is proportional to $\beta_{\text{rms}}$ and involves the assumption of a rigid rotator with a prolate shape. The ratio $Q_{2^+} / Q_R$ is about 1.0 for $^{182}_{\text{W}}$ which has a substantial prolate-oblate preference, as measured by $V_{PO}$; the ratio goes through zero in a region of small $V_{PO}$ and becomes negative for the oblate, Pt nuclei. Note that the ratio $Q_{2^+} / Q_R$ is not zero when $V_{PO} = 0$. The reason is that the region of prolate-type quadrupole moments is extended somewhat by the kinetic energy terms or the $\beta, \gamma$-dependence of the six mass parameter functions.

The quantity $V_{PO}$ is a theoretical one, but it can be related to another observable quantity, the splitting of the $2^+$ or $2^+_\gamma$ state and the first $4^+$ state (Fig. 9). In a model of $\gamma$-independent potentials
and constant mass parameters, first discussed by Wilets and Jean\textsuperscript{18},
the $2_\gamma$ and $4^+$ states are degenerate. As we introduce a finite $V_{p0}$,
we expect the nucleus to become stable against $\gamma$ vibrations and hence
the $2_\gamma$ state to go up, regardless of the sign of $V_{p0}$. The change of
sign of $(E_{2_\gamma} - E_{4^+})$ in Fig. 9 is another indication of the $\beta, \gamma$-dependence
of the mass parameters.

In order to study the full $\beta-\gamma$-dependence of the six inertial
functions, one needs six contour plots in the $\beta-\gamma$ plane. However, for
the sake of brevity, only the axially symmetric part is given in
Fig. 10. For axially symmetrical shapes ($\gamma = 0^\circ$), we need only one mass
parameter for rotations (see Eq. 4 and ref. \textsuperscript{4}):

$$B_R(\beta) = B_1(\beta) = B_2(\beta) = \beta(\beta) [3\beta^2]^{-1}$$

and only two for $\gamma$ vibrations,

$$B_{\beta\beta}(\beta), B_{\gamma\gamma}(\beta) = B_3(\beta); B_{\beta\gamma}(\beta) = 0.$$

All three mass parameters show substantial deviations from the assump-
tion of a constant value (see Eqs. 4-6) which is valid only in the
spherical limit. The function $B_{\beta\beta}$ fluctuates by about a factor\textsuperscript{*}) of
2 in such a way that its "minima" occur for the same deformations
as those of the potential (shown in the lower half of Fig. 10).
The function $B_{\gamma\gamma}$ decreases almost monotonically from oblate to
spherical to prolate shapes. The function $B_R$ is approximately a
Gaussian with its peak at $\beta = 0$. Near equilibrium deformations, the

\textsuperscript{*)} These fluctuations would be even larger (about a factor of four) if
we had not added a substantial core contribution (independent of
deformation) to all three functions.
average mass parameter is smaller for prolate compared to oblate shapes, hence the vibrational amplitude \( A_{IK} \), Eq. 3) for prolate shapes is enhanced. Therefore, the prolate \( \rightarrow \) oblate transition is shifted somewhat towards \( V_{PO} < 0 \). Also, the \( 2^+ \gamma \) state is pulled below the \( 4^+ \) state. Of course, for \( V_{PO} \ll 0 \) (nuclei with a strong preference for oblate shapes), the nucleus would again become \( \gamma \)-stable, hence the \( 2^+ \gamma \) state would be pushed above \( 4^+ \), and the ratio \( Q_{2^+} / Q_R \) would approach the value -1.0.

Previously, the Pt nuclei were considered to be vibrational nuclei described by the phonon model. Therefore, the results of the present model are compared with the phonon model in Figs. 11-12. The B, C parameters of the phonon model are obtained by fitting the calculated \( E_{2^+} \) and \( B(E2; 0^+ \rightarrow 2^+ \) ). The calculated potential is quite different from the phonon model one: it has a spherical maximum, an oblate minimum and a prolate saddle point. The ground state wavefunction has an oblate peak, a spherical peak, and a prolate peak. It reflects the preference for oblate shapes (Fig. 11). This preference is enhanced in the wavefunction of the first \( 2^+ \) state (Fig. 12) since such a wave function vanishes at \( \beta = 0 \), reflecting the fact that the \( I \neq 0 \) states of all nuclei are deformed\(^1\). The wavefunction of the second \( 2^+ \) state must be orthogonal, hence the prolate peak is higher and the sign of \( Q_{2^+} \) is opposite to that of \( Q_{2^+} \).

The PPQ wavefunctions can be represented as linear combinations of only a few phonons\(^{11}\). However, the potential function (Fig. 11) and the mass parameter functions (Fig. 10) are quite different. The situation is somewhat analogous to that of the oscillator model used for the single particle states. The oscillator wavefunctions
are quite good, but there are important additional terms in the potential. The treatment of quadrupole motion is, however, much more complicated since it is a five-dimensional problem rather than a three-dimensional one.

Another indication of the somewhat complicated $\beta,\gamma$-dependence is provided by the mixing ratio $\delta(E2/M1)$. In the simple collective model, this ratio is infinitely large for the $2' \rightarrow 2$ transition (and any others for which an E2 transition is "allowed") since the gyromagnetic ratio is a constant and the M1 transition probabilities vanish - no matter how strong the rotation-vibration coupling or the phonon mixing may be. A non-zero M1 transition or non-infinite $\delta$ value must be ascribed to the $\beta,\gamma$-dependence of the $g$-value tensor: $g_k(\beta,\gamma)$ ($k = 1,2,3$). Provided the derivatives $\partial g_k/\partial \beta$, $\partial g_k/\partial \gamma$ do not change sign during a prolate $\rightarrow$ oblate transition, the sign of the mixing ratio is expected to change since the E2 operator is proportional to $\beta$ (the higher order terms are included in the PPQ calculation but are small). As can be seen from Fig. 13, some of the mixing ratios follow the change of sign from $^{192}$Os to $^{192}$Pt, but others do not. Evidently, the E2/M1 mixing ratio depends in a subtle way on the details of nuclear wavefunctions and hence is a sensitive test of nuclear models. The large discrepancy in $^{196}$Pt reflects the fact that the numerical accuracy of this calculation is not completely satisfactory.

The prolate $\rightarrow$ oblate transition discussed above occurs at $Z = 76 - 78 (N = 114, 116)$. When we go away from the line of beta-stability, a similar transition occurs in the Pt isotopes in the vicinity of $N = 108$. The evidence is based partly on the PPQ calculation $^{11,19}$ for $^{186,192-196}$Pt and partly on the experimental
information for the Pt isotopes, obtained by the ISOLDE group at CERN\textsuperscript{20}). The energy levels are given in Fig. 14. As we go towards smaller \( N \) and away from the magic shell at \( N = 126 \), the nucleus becomes somewhat more deformed and \( E_{2^+} \) is lowered. However, this is not a spherical \( \gamma \)-deformed transition since the \( 0_\beta(0'\!+) \) state lies higher than the \( 2_\gamma \) state throughout this region. Instead, the \( 2_\gamma, 4^+ \) levels cross at \( N = 108 - 110 \). Hence, these isotopes lie in a region of prolate-oblate transition. This conclusion is supported by the calculated potential wells (Fig. 15) which show a transition from a prolate minimum at \( ^{186}\text{Pt} \) to an oblate one at \( ^{192}\text{Pt} \).

Further comparisons between theory and experiment are given in Fig. 16 for the \( B(E2) \) ratios and in Fig. 17 for the electric monopole transitions. For further confirmations, it will be necessary to look for a change of sign of \( Q_{2^+} \). Discussion given above suggests that such a change of sign may be shifted slightly to the right of \( N = 108 \).

Application of the ideas discussed above suggests that there are many other regions of prolate-oblate transition. A tentative list is suggested in Fig. 18. It is hoped that this list will be made more complete in the near future. The study of nuclei away from the line of beta-stability opens up a whole new dimension and a fertile region for intermediate nuclei in various stages of transition from spherical to deformed and prolate to oblate shapes.
REFERENCES

2) B. Sørensen, Nucl. Phys. A142, 411 (1970) and earlier refs. cited there.
15) R. M. Diamond et al., pr. comm.
17) J. E. Glenn and J. X. Saladin, Phys. Rev. Letts. 20, 1298 (1968); R. J.
18) L. Wilets and M. Jean, Phys. Rev. 102, 788 (1956).
19) K. Kumar, Proc. Int. Conf. on Properties of Nuclear States, Montreal,
     1969, and to be published.
20) M. Foucher et al., Proc. Int. Conf. on Properties of Nuclear States,
     Montreal, 1969; pr. comm.
Fig. 1. Spherical to Deformed Shape Transition in the $N = 88 - 90$ Region. The potential energy $V(\beta,0)$ has been calculated by using the pairing-plus-quadrupole model. The location of the ground state is indicated by dashed lines.
Fig. 3. Contour Plot of $V(\beta, \gamma)$ of $^{152}$Sm. The potential surface has a maximum for the spherical shape, a minimum for a prolate shape and a saddle point for an oblate shape.
Comparison of Theoretical and Experimental E2 Matrix Elements.
The table gives $\langle f | \mathcal{M}(E2) | i \rangle$ in e b.
Static moments are also given.

<table>
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<tr>
<th>Transition</th>
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<th>f</th>
<th>Th.</th>
<th>Ex.</th>
<th>i</th>
<th>f</th>
<th>Th.</th>
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Fig. 4
Table (Cont'd.)

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<td>1.71</td>
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<td>6</td>
<td>2.69</td>
<td>(2.40)</td>
</tr>
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</table>

3. G. Stockstad et al., Phys. Rev. Letts. 23 (1969) 1047 and 1051, and pr. comm. Note that the values given in parentheses are not determined well by experiment, hence rotational model values are used.
Comparison of Theoretical and Experimental Quadrupole Moments, Magnetic Moments, E2/M1 Mixing Ratios, Monopole Transition Moments and Isomer Shifts for $^{158}$Sm

<table>
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<th>Quantity</th>
<th>Th.</th>
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<td>$Q_{2+}$ in e.b.</td>
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<td>-1.8 ± 0.6</td>
<td>a</td>
</tr>
<tr>
<td>$\mu_{2+}$ in n.m.</td>
<td>0.29</td>
<td>0.27 - 0.415</td>
<td>b</td>
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<tr>
<td>$\delta$ (E2/M1, $2_\gamma \rightarrow 2$)</td>
<td>36.7</td>
<td>$27^{+55}_{-11}$</td>
<td>c</td>
</tr>
<tr>
<td>$\delta$ (E2/M1, $4_\gamma \rightarrow 4$)</td>
<td>18.3</td>
<td>$12.9^{+\infty}_{-7.6}$</td>
<td>c</td>
</tr>
<tr>
<td>$\delta$ (E2/M1, $2_\beta \rightarrow 2$)</td>
<td>109.7</td>
<td>$(6^{+54}_{-18})$</td>
<td>c</td>
</tr>
<tr>
<td>$\delta$ (E2/M1, $4_\beta \rightarrow 4$)</td>
<td>59.3</td>
<td>-2.94 ± 0.85</td>
<td>c</td>
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<td>$\rho$ (E0; $2_\beta \rightarrow 2$)</td>
<td>0.11</td>
<td>0.26</td>
<td>d</td>
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<tr>
<td>$\Delta \langle r^2 \rangle / \langle r^2 \rangle$</td>
<td>$6.8 \times 10^{-4}$</td>
<td>$(5.7^{+1.2}_{-1.5}) \times 10^{-4}$</td>
<td>e</td>
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</table>

c R. G. Stockstad et al., Phys. Rev. Letts 23 (1969) 1047 and pr. comm. The Rose-Brink definition of $\delta$ (E2/M1) has been used.

Fig. 6
Fig. 7. Prolate to Oblate Shape Transition in the $Z = 74 - 78$ Region. The potential wells for other, beta-stable isotopes of $\text{W}$, $\text{Os}$, and $\text{Pt}$ are qualitatively similar.
Fig. 8. The Ratio $Q_{2+}/Q_R$ as a Function of the Calculated Prolate-Oblate Difference, $V_{P0}$. The experimental values\cite{17} are also shown for comparison. Note that the smooth curve drawn in the figure should not be extrapolated for strongly oblate nuclei with $V_{P0} \ll 0$. 
Fig. 9. The Splitting $E_{2^+} - E_{4^+}$ as a Function of the Calculated Oblate-Oblate Difference, $V_{PO}$. Note that the straight line drawn in the figure should not be extrapolated for nuclei with large $|V_{PO}|$. 
Fig. 10. The $\beta$-dependence of the Mass Parameter Functions. The mass parameter $B_R$ is related to the moment of inertia. The lower part of the figure shows the potential function and the associated oblate minimum and prolate saddle point.
Fig. 11. Comparison with the Potential and Ground State Wavefunction in the Phonon Model. Note that the PPQ model wavefunction is larger for oblate shapes compared to oblate.
Fig. 12. Comparison with the $K = 0$ Components of the First and Second $2^+$ Wavefunctions in the Phonon Model. In the PPQ model for this nucleus, the first $2^+$ wavefunction is larger for oblate shapes, but the second $2^+$ wavefunction is larger for prolate shapes.
<table>
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<th>Nucleus</th>
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<th>$E_v$ in MeV</th>
<th>$\delta$ (E2/M1)</th>
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<td>$10^{+5}_{-10}^a$</td>
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<td>$3^+ \to 2^+$</td>
<td>0.773</td>
<td>$12^{+5}_{-10}^a$</td>
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<tr>
<td>$^{188}$Os</td>
<td>$2^+ \to 2^+$</td>
<td>0.478</td>
<td>$12.3 \pm 2.8^a$</td>
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<tr>
<td></td>
<td>$3^+ \to 2^+$</td>
<td>0.635</td>
<td>$6.9 \pm 3.2^a$</td>
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<tr>
<td>$^{190}$Os</td>
<td>$2^+ \to 2^+$</td>
<td>0.371</td>
<td>$9.5 \pm 2.0^a$</td>
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<td>$3^+ \to 2^+$</td>
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<tr>
<td>$^{192}$Os</td>
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<td>0.486</td>
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<tr>
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<tr>
<td>$^{194}$Pt</td>
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<td>$3^+ \to 2^+$</td>
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$^c$Z. W. Grabowski, pr. comm.


Comparison with the experimental values of the E2/M1 mixing ratio. For sign conventions and other details see ref. e.

Fig. 13
Fig. 14. Comparison of Theoretical and Experimental Energy Levels of Even Pt Nuclei. The experimental levels are taken from R. Poucher et al. [20]
Fig. 15. Prolate to Oblate Shape Transition in the Pt Isotopes. While the Pt isotopes near the line of beta-stability ($^{192-198}$Pt) are oblate, the beta-unstable isotope, $^{196}$Pt, is prolate.
### B(E2) Branching Ratios in Even Pt Nuclei

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<th>Theo. $^{186}_{Pt}$</th>
<th>Expt. $^{188}_{Pt}$</th>
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<td>$100 \pm 13$</td>
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<td>$\approx 11$</td>
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<td>(100)</td>
</tr>
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\(^{a}\text{R. Foucher et al., ISOLDE Collaboration at CERN, pr. comm.}\)

Fig. 16
<table>
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<td>$2^+$</td>
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- $^{188}$Pt values of $X = 100 \frac{E_{f'}^2}{e_R^2} / B_{f'}$, (E2)

*Fig. 17*

R. Foucher et al., ISOLDE Collaboration at CERN, pr. comm.
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</table>

\(a\) Ref. 16 or 11.

\(b\) Refs. quoted in Ref. 11.

\(c\) Ref. 20.

\(d\) Ref. 19.

Regions of Prolate to Oblate Transition. The symbols Y, N, and P refer respectively to Yes, No, and Probable.

Fig. 18
SOME INTERESTING ASPECTS OF THE VMI MODEL

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Once upon a time there were two limiting models which described collective states of even-even nuclei and the transition probabilities between these states quite satisfactorily. These were the strongly deformed nucleus model proposed by A. Bohr in collaboration with Mottelson, and the vibrational model, suggested by Weneser and myself, which was also based on the Bohr description of the nucleus. Fig. 1 shows level schemes for the two models: the deformed nucleus at the right has a ground-state rotational band (0,2,4,6,8, even parity), with energies proportional to $J(J+1)$, the $\gamma$-vibrational band (2,3,4, etc., even parity) and the $\beta$-vibrational band (0,2,4, etc., even parity). Of course, as you know, in the Davydov-Filipov asymmetric model, the second 2+ and 0+ states are not vibrational states but also rotational states. At the left we have the level scheme of a spherical nucleus undergoing quadrupole vibrations of the surface: it has a 0+ ground state, a 2+ one phonon state, a 0+, 2+, 4+ two phonon triplet and higher phonon multiplets. The level scheme on the right is known to occur within the large proton and neutron shells, i.e., the rare earths and heavy element regions, and the level scheme on the left within the smaller shells, e.g. in Cd and Te nuclei. The study of the even-even Os and Pt level schemes indicated that a gradual transition from the deformed nucleus level scheme to the vibrational level scheme takes place in the manner indicated by dashed lines.
A number of years ago the Argentinian physicist Mallmann made a very interesting observation; namely he showed that if one plots the energy ratios $E_6/E_2$ and $E_8/E_2$ as functions of $E_4/E_2$, these ratios seem to lie on two universal curves for $E_4/E_2$ ranging from 1 to 3-1/3, the rotational limit. He tried hard to find a model giving this result but did not succeed.

In the meantime various developments took place which threw doubt on the validity of the limiting models. First: Thanks to the wide use of the Morinaga method of populating extended ground state bands, sometimes up to $J = 16$, by heavy ion reactions, a large number of data on these bands in stable and neutron deficient deformed nuclei have become available. These bands show systematic deviations from the $J(J+1)$ rule which cannot be corrected by adding one or two terms in
\[ J(J+1) \]. Second: A growing number of results, especially the appreciable values found for static quadrupole moments of the lowest 2+ states in vibrational nuclei, have invalidated the notion that these states have spherical symmetry.

I should like to show now that it is possible to describe ground state bands of all even-even nuclei of at least moderate mass, which do not contain a filled nucleon shell, by means of a simple semiclassical model. The first part of this work was published last year.\(^1\)

We started out from the beta stretching model,\(^2\) the model of a diatomic molecule, which was shown to fit ground state bands of strongly deformed nuclei satisfactorily. By means of a least squares analysis of 98 ground state bands ranging from "rotational" to "vibrational", we arrived at a different model, the "variable moment of inertia" (VMI) model, which for strongly deformed nuclei, however, is not easily distinguishable from the beta-stretching model. I shall first describe the model and its consequences and later point out the differences between it and the beta stretching model. We obtained from the level energies the expression

\[
E(\delta) = \frac{C}{2} (\delta - \delta_o)^2 + \frac{J(J+1)}{2\delta} \tag{1}
\]

Here a potential term is added to the usual rotational term given at the right. \( C \) and \( \delta_o \) are parameters characteristic for each nucleus. Where \( \delta_o > 0 \), it is denoted as ground state moment of inertia, and \( C \) is denoted as stiffness parameter. To fix the moment of inertia \( \delta \), we require that it minimize the energy for each \( J \):

\[
\frac{\delta E}{\delta \delta} = 0, \quad \text{and} \quad C > 0 \tag{2}
\]
From the equilibrium condition follows a cubic equation for the moment of inertia:
\[ \mathcal{J}^3 - \mathcal{J}_o \mathcal{J}^2 = \frac{J(J+1)}{2C} \]  
(3)

One finds that for \( \mathcal{J}_o = 0.4E \sim (J(J+1))^{2/3} \) and hence the energy ratio \( E_4/E_2 = 2.23 \).

Fig. 2 shows the quality of the fits for various types of bands. For each nucleus the measured energy levels are given at the left and the computed energies at the right. Almost without exception the fits agree within considerably less than 1%. The third figure shows how the parameter \( \mathcal{J}_o \) varies as a function of neutron and proton number: The parameter is shown to increase smoothly with the distance from magic numbers. It has roughly the appearance of a dome, with the highest point just in the middle between magic numbers. The stiffness parameter \( C \) reaches the highest values for stable nuclei of each element. It is easy to prove that the VMI law yields Møllmann type curves for \( E_J/E_2 \) (\( J = 6,8, \ldots \)) as function of \( E_4/E_2 \).

We have also shown that the parameters \( \mathcal{J}_o \) and \( C \) can be given as "universal" functions of \( E_4 \), multiplied by \( 1/E_2 \) and \( E_2^3 \) respectively. These quantities can thus be determined if only the two lowest states of a band are known.

The fourth figure shows by means of a few representative cases how the moment of inertia varies with angular momentum. Starting from the value \( \mathcal{J}_o \), one finds a smooth increase for all nuclei. The relative increase is the larger, the smaller the value for the stiffness parameter \( C \) (given in parenthesis for each nuclide). It is seen that for well-deformed stable nuclei, e.g., for \( \text{Hf}^{180} \), \( \mathcal{J} \) is almost constant as \( J \) increases. For nuclei with the most deformed ground states,
Fig. 2
Examples of least squares analyses of ground state bands
based on VMI model
namely those occurring approximately in the center of their proton and neutron shells, one finds $\delta_0 \propto A^{5/3}$, just as the rigid moment of inertia. This relation does not hold, however, for the transition nuclei, as is immediately evident from the coincidence of the curves for Xe$^{120}$ ($Z = 54$, i.e., four protons beyond the closed shell) and Pt$^{194}$ ($Z = 78$, four proton holes). For these nuclei the moments of inertia increase steeply from their value at $J = 0$. The most dramatic relative increase of $\delta$ with $J$, however, is found in nuclei with very small values of $\delta_0$, e.g., Cd$^{110}$. The realization that nuclei with a spherical ground state and $R_4 \sim 2.23$ are extremely "soft" removes the difficulty for the simple vibrational model presented by the large static quadrupole moments found for $2^+$ states.

Excited states of many nuclei with $R_4 < 2.23$ have recently been populated, either by heavy-ion reactions or by inelastic scattering.$^{3-9}$ There appear to be regular "ground-state bands" in all even-even nuclei investigated. The good fits obtained for "vibrational" nuclei ($\delta_0$ small, but positive) invite the extension of the WMI model$^{10}$ to permit negative values of the parameter $\delta_0$.

$$r^2(r-1) = -X,$$  \hfill (4)

with

$$r = \delta/\delta_0; \ X = |J(J+1)/2\delta_0^3|.$$  

For $r \neq 0$ one can write $r = 1 - X/r^2$ and find graphic solutions for $r$ by plotting each side of the equation separately. For all $X > 0$ there is one real negative root which goes smoothly from $r \approx -\sqrt[3]{X}$ to $r \approx -\sqrt{X}$ as $X$ goes from large to small. For this region one finds $\delta_{J=0} = 0$ and therefore, according to Eq. (1), $E_0 = \frac{1}{2}C\delta_0^2$. The limiting value for $R_4$
Fig. 3

The parameter $\delta_0$ as function of $N$ and $Z$.

(Magic numbers are indicated by double lines)
Fig. 4
Moment of inertia $\mathcal{J}$ plotted vs angular momentum $J$
for some representative examples
is again 2.23 as $\delta_0 \to 0$. For large negative $\delta_0$ one obtains $R_4 \to (20/6)^{1/2} = 1.82$. We have included in Fig. 4 an example of a nucleus with $\delta_0 < 0$, $\text{Te}^{120}$, which has $R_4 = 1.99$. Despite the negative $\delta_0$, the $\delta(J)$ curve resembles that of the isobar $\text{Xe}^{120}$.

From here on, the interval $2.23 < R_4 < 3.33$ will be denoted as the deformed region, the interval $1.82 < R_4 < 2.23$ as the spherical region, and the interval $1 < R_4 < 1.82$, which contains only single and double magic nuclei, as the magic region. In the deformed region the ground-state moment of inertia $\delta_{J=0} = \delta_0$, while in the spherical region $\delta_{J=0} = 0$. This means that as $\delta_0$ changes sign, its relation to the character of the ground state changes in a manner reminiscent of a second-order phase transition. Negative values for $\delta_0$ require the introduction of a new physical concept into the model. This is the notion of "internal stress" or "rigidity": The larger the negative value of $\delta_0$, the more firmly the shell structure resists departure from spherical symmetry. In Fig. 5 the established experimental values for $R_6$ and $R_8$ have been plotted against $R_4$ ($R_J = E_J/E_2$). We have included all cases with $R_4 < 2.23$ in which $\gamma$-ray cascades have been observed, including $\text{Ti}^{50}$, $\text{Cr}^{52}$, and $\text{Zr}^{90}$, whose levels are known from studies of radioactive decay. The solid curves extend the predictions of the VMI model to the spherical region. The experimental points lie along two branches: The first or VMI branch extends leftward from the deformed region. The points on this branch are well fitted by the extended VMI model and terminate precisely at $R = 1.8$, the natural limit of the model. The nuclei on this branch in the spherical region are never more than four nucleons from single magic. The second or magic branch extends to the right from $R_4 = 1$ and consists entirely of double$^5$ and
singly magic nuclei for $R_4 < 1.82$. One is tempted to include Po$^{208}$ in this branch, even though it lies in the spherical region. The nuclei in this branch, especially the doubly magic nuclei are characterized by considerably larger $E_2$'s than their nonmagic neighbors, and by a drastic reduction in level spacing for the higher states in the band. This change of scale implies that a minimum of three parameters would be required in any scheme to fit these spectra. Such a scheme would have to take account of the qualitative inference that these nuclei undergo some kind of rearrangement or "melting" when they are excited, so that further rearrangement costs much less energy: The ground states of these nuclei are not only rigid, but also "brittle". The apparent phase change between the ground state and the higher states could be interpreted as a transition from "superconducting" to "normal" associated with the symmetry breaking effect of pair excitation.

To obtain a more detailed insight into the role of angular momentum in nuclear structure, knowledge concerning transition probabilities between members of ground-state bands is of great importance. In ref. 1 it was shown that the transition quadrupole moment

$$Q_{02} \propto \sqrt{\delta_{02}},$$

where $\delta_{02} = \frac{\delta_0 + \delta_2}{2}$. Since one may assume that $Q \ll \beta$, we can now compare the potential of the VMI model

$$V_{VMI} = \frac{C}{2}(\delta - \delta_0)^2 = \frac{C'}{2}(\beta^2 - \beta_0^2)^2$$

with the potential of the beta-stretching model

$$V_{Beta-stretching} = \frac{C''}{2}(\beta - \beta_0)^2$$

The potential of eq. (5) contains a 4th power and a second power term
Fig. 5

Extension of the VMI model into the region $J_0 < 0$
in $\beta$. For deformed nuclei the coefficient of the quadratic term is negative and the potential has a maximum at $\beta = 0$, whereas for spherical nuclei one obtains a minimum at $\beta = 0$. This contrasts with the beta stretching potential, whose derivative at $\beta = 0$ has a finite negative value. I should like to mention that at the 1967 Tokyo Conference, where a preliminary version of the VMI model was presented, Frank Stephens reported the results on ground state bands from the Berkeley group in the form of Mallmann plots, in order to compare them with predictions from various models. By far the best agreement was obtained with the next higher order correction to Inglis' cranking model which had been published by Sam Harris previously. A comparison of the VMI model with the Harris model showed that the two are mathematically equivalent. All this took place, of course, before the VMI model was extended into the spherical region.

The VMI potential, both for the deformed and spherical regions, is symmetrical with respect to $\beta = 0$, i.e. it makes the prolate and oblate shapes equally probable. In view of the good fits obtained with this model it appears that odd power terms in $\beta$, e.g. a term in $\beta^3$ which will have the effect of making either the minimum for the prolate shape or that for the oblate shape deeper, should be rather small. It appears that a treatment containing such a term, either semiclassical or quantum mechanical, might make it possible to describe the $\gamma$- and $\beta$-bands, as well as the static quadrupole moments, etc.

To summarize the results presented here I should like to illustrate with the help of Fig. 6 how the energy ratio $R_4 = E_4/E_2$, which reflects important features of nuclear structure, varies with the mass number $A$. There is, as we have seen, a strong dependence on the shell structure. Closer analysis shows that there is particle-hole symmetry
Fig. 6

The variable $E_4/E_2$, plotted vs A
in experimental $R_4$ values: Whether $n$ nucleons are added to or taken from a closed shell, the increase in $R_4$ is roughly the same. It will be interesting to see whether there is a lower limit, in terms of the mass number $A$, for the validity of the model. As is seen here, for many nuclei no $E_4/E_2$ values are as yet available.

The VMI model provides an economical description of ground-state bands for an extraordinary range of even-even nuclei. One may hope that the model will be an aid in both construction and criticism of microscopic models. To complete the experimental picture, we need a better knowledge of energies of ground-state-band members in lighter nuclei in the spherical region, and transition probabilities, especially for vibrational and closed-shell nuclei.

*   *   *
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7) Pb$^{210}$: G. Igo, private communication.


SOME NEW ASPECTS OF DEFORMED NUCLEI

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I would be happy with the title of this talk, which was suggested to me by the organizing committee, if it was not for the little word "new". The deformed nuclei have been intensively studied for many years and most of the things we discuss today have been discussed earlier and are therefore not really "new".

Nevertheless, the greatly improved experimental methods which in the later years have been applied to the study of deformed nuclei have resulted in a wealth of information and have thrown light on a number of old problems. In some cases the information has also more strongly brought into focus problems for which we do not yet have a satisfactory solution. So let this be the theme of my talk: New light on old problems.

Among the most powerful tools for the study of deformed nuclei are the single nucleon transfer reactions\(^1\). Neutron hole states e.g. are populated by the \((d,t)\) reaction and particle states by the \((d,p)\) reaction. Similarly, the proton states can be studied by e.g. \((t,α)\) and \((^3\text{He},d)\) reactions. The usefulness of these reactions mostly stem from the fact that the transfer cross section in a very simple way is related to the Nilsson wave function for the orbital involved in the transfer. For transfer leading to an even-odd nucleus
the cross section to a state $I (= j$ of the transferred nucleon) is simply

$$\frac{d\sigma}{d\omega} = 2C_{jL}^2 \sigma_j(\theta)$$

(1)

The $\sigma_j(\theta)$ is a function which can be calculated from reaction theory and $C_{jL}^2$ is the square of the expansion coefficient of the Nilsson state, $\chi_\Omega$, on the spherical states $\phi_{lj}$

$$\chi_\Omega = \sum_j C_{jL} \phi_{lj}$$

(2)

The cross section (1) for the members of a rotational band form the "finger print" of the band which is the basis for its identification (fig.1). The pattern represents a direct experimental determination of a nuclear wave function.

Extensive studies of this type has first of all led to the identification of a very large number of Nilsson states not located before (fig.2). The amazing fact is that the original more than 15 years old calculations by Nilsson are in very

Fig.1. Pictorial representation of Nilsson wave function. The widths of the arrows indicate the expansion coefficients $C_{jL}$ which are proportional to the single nucleon transfer cross section.
Fig. 2. Single neutron orbitals in the mass region 150-170.
good agreement with the experimentally observed energies (fig.3) and also with the wave functions experimentally determined from (1).

I have spent some time summarizing these simple facts because they have given the Nilsson model a very strong experimental support which, in my opinion, allows us to go further in the analysis of complex phenomena in the deformed nuclei than is possible in other nuclei.

Equally important information about the deformed nuclei has been obtained from the studies of the $\gamma$-decays of states populated in (ion,xn) reactions. Several of the leading experts in this field are present at this conference and I shall therefore just mention that they, among other things,
have supplied us with a large number of accurately measured energies and transition probabilities for states extending to very high angular momenta\(^2\).

Finally, to end this survey of experimental methods and results, inelastic scattering processes have located most of the low lying (<2 MeV) collective states which are connected with the ground state by strong transitions of types E2, E3 and E4. As an example, fig.3 shows the location of collective 3- states in a number of deformed nuclei.

One of the conclusions which can be drawn from the analysis of data mentioned above is that the Coriolis coupling is of significance for the understanding of a large number of phenomena observed in deformed nuclei. This has become especially clear as states of higher and higher angular momenta have been observed.

The classical Coriolis force \(-2m\hat{\omega} \times \hat{v}\) acts on a particle with mass m moving with velocity \(\hat{v}\) in a coordinate system with angular velocity \(\hat{\omega}\). It represents the reaction of the particle against being dragged around by the rotation. Its quantum mechanical counterpart is the Coriolis interaction \(H_C\) which is of the general form \((\hbar^2/2\sigma) I \cdot j\) or more exactly

\[
H_C = \frac{\hbar^2}{2\sigma} (I_+ j_+ + I_- j_-) \tag{3}
\]

where \(I\) refers to the total angular momentum and \(j\) to the instantaneous particle angular momentum. The interaction \(H_C\) connects states differing one unit in \(K\)-quantum number and can be expressed as
\[
\langle \text{IK} | H_{\text{C}} | \text{IK} + 1 \rangle = A_{K} \sqrt{(I-K)(I+K+1)}
\] (4)

where \( A_{K} \) with the inclusion of pairing is given by

\[
A_{K} = \frac{\hbar^2}{2\gamma} \sum_{j} C_{j,k} C_{j,k+1} \sqrt{(j-k)(j+k+1)} \left( U_{K,k+1} + V_{K,k+1} \right)
\] (5)

when the coefficients \( C_{j,k} \) are the same as in (2).

With the data now available it has been possible to put the Coriolis coupling formalism to a rather detailed test. This is especially so for the coupling of the neutron states 1/2+|660⟩, 3/2+|651⟩, 5/2+|642⟩, 7/2+|633⟩, 9/2+|624⟩, 11/2+|615⟩ and 13/2+|606⟩ which originate in the \( i_{13/2} \) shell model state (high \( j \)) and which consequently are coupled with matrix elements which often exceed the unperturbed energy.

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**Fig. 4.** Strong Coriolis coupling completely destroys the rotational band structures and results in strongly mixed wave functions.
Fig. 4 illustrates the general effects on two strongly coupled rotational bands. If more bands are coupled, the general effect is to concentrate most of the transfer intensity (constructive interference of the \(C_{jl}\)'s) in one or a few low lying bands which more resemble the spherical shell model state (see e.g. \(^{181}\)W in fig. 5).

It would be very interesting to study the corresponding coupled states in an even-even nucleus. Among others, a low lying \(K=1^+\) band with strong E2 transitions to the ground band could be expected in some cases, but has not been identified so far.

The analysis of the Coriolis coupled states is now at a stage where more definite statements can be made about the

![Diagram](image)

Fig. 5. Population of 13/2+ states by the \(^{3}\)He,\(\alpha\) reaction shows effects of strong Coriolis coupling, ref. 4).
Fig. 6. Coriolis coupling calculations indicate that the matrix elements are subject to an energy dependent attenuation.

coupling strengths. In a few cases (see fig. 6) it has been found that the theoretical matrix element (5) fully accounts for the observations. In most other cases certain adjustments, mostly reductions, of the matrix elements are necessary. Thus the extensive analysis made in Stockholm\textsuperscript{3)} indicates reductions to about 75\% of the Nilsson value. An analysis\textsuperscript{5)} of the $\text{W}$ nuclei suggests an energy dependent attenuation factor. At present the origin of such attenuation is not clear. It does not seem to be connected with truncation of the state space for the calculations and not with dilutions of the single particle wave functions as the total transfer cross section appears to be present. The pairing factor in (5) might be suspected, but there is little evidence to
support the suspicion. The need of attenuation indicated by
the experiments casts some doubt upon the completeness of
the Coriolis formalism which otherwise works so beautifully.

The detailed analysis of the coupled states has also
revealed a number of additional effects. The data on fig. 5
indicate that the unperturbed Nilsson states must be shifted
somewhat in order to explain the observations. The direction
and the magnitude of the effect is as expected for a positive
hexadecapole deformation $\varepsilon_4$.

The Coriolis coupling also affects the collective
states markedly. Thus the collective octupole vibrations in
the even nuclei are usually fewer than the four ($K = 0, 1,
2$ and $3$) expected from the simplest models.

Calculations$^6$ have shown that again the inclusion of
Coriolis coupling in a rather detailed manner accounts for
the energies and transition probabilities for the octupole
bands. In many cases the coupling matrix elements calcu-
lated from a microscopic theory approach the limit
$\sqrt{(3-K)(3+K+1)}$ which results in strongly mixed wave functions.
Also the well known quadrupole vibrations might be Coriolis
coupled to $K = 1^+$ bands, but cases of this type have not
been identified as yet.

In all discussions of the Coriolis coupling the very
strong $I$-dependence (cf. eq. 5) should be kept in mind. At
present we have no information about the breakdown of this $I-
dependence which could have important consequences for the
understanding of the energy levels in the yrast region.
A quite different type of coupling of the single particle levels is associated with the crossing, in the Nilsson diagram, of levels with the same spin and parity, but belonging to different oscillator shells. If the interaction between states of different N is taken into account such crossings do not occur, but the levels in the Nilsson diagram will be deflected and there will be a change of asymptotic quantum numbers (fig.7). The best studied examples are the crossings of the $N = 4$ and $N = 6$ states at the beginning of the rare-earth region. These crossings are easy to investigate by neutron transfer because they give rise to a splitting of the large $\ell = 0$ and $\ell = 2$ transfer cross sections. Examples of this coupling are illustrated in fig.8 for the $3/2^+|402\rangle$ and $3/2^+|651\rangle$ bands\(^6\). By the use of the two-band mixing expressions, one finds

$$\Delta E = \sqrt{\Delta H^2 + 4V^2}$$

and for the ratio of amplitudes $\alpha(N=6)/\beta(N=4)$ we have

$$|\frac{\alpha}{\beta}| = \left| \frac{2V}{\Delta H - (\Delta H^2 + 4V^2)^{1/2}} \right|$$

where $\Delta H$ is the unperturbed energy difference and $V$ the un-
known interaction. Solutions $\Delta H$ of these equations as illustrated in fig. 9 directly show the crossing of the unperturbed levels. The interaction $V$ is $\approx 60$ keV, considerably stronger than obtained from the Nilsson wave functions $^7$).

The properties of collective $\beta$-vibrations have for many years been somewhat of a puzzle. Among the best known are the "Three Musketeers" in $^{150}$Nd, $^{152}$Sm and $^{154}$Gd. It has
been attempted to ascribe deviations of the E2 branching ratios from the simple vector coupling to couplings to the ground state and other bands, but with limited success. The question is interesting because the β-vibration is directly connected with the question of nuclear centrifugal stretching, but the lack of internal consistency in data and analysis has made much of the interpretation open to doubt.

Perhaps there are better behaved cases in other nuclei. Recently\textsuperscript{8}) an investigation of the low lying β-band in $^{174}\text{Hf}$ has shown that the interband transitions are well explained in terms of rotation-vibration interaction with a strength $<H_{βg}> = -2.55 \, I(I+1) \, \text{keV}$. This is somewhat less than expected from a stretching model. Some of the results of this analysis are shown in fig.10.

It should be emphasized that this β-band so far is the only example where rotation-vibration interaction has been able to account for the branching ratios. The interaction is, however, about four times weaker than found in the beginning of the deformed region where the nuclei are softer. The interaction mentioned above accounts for about 30% of the energy depressions in the ground state band which shows that effects other than stretching are important.

This brings me to the final subject I want to mention, that of the ground state bands in even-even nuclei.

The deviations of the energies of the rotational states from the simple $I(I+1)$ rule have for many years been a puzzle. In the usual treatments, the rotational energies are expanded
Fig. 10. E2 matrix elements connecting the $\beta$-band and the ground band in $^{174}$Hf. The uncoupled case corresponds to a horizontal line on the plot. The curve through the points is a fit with an $I(I+1)$ dependent interaction.
in a power-series in $I(I+1)$

$$E(I) = AI(I+1) + BI^2(I+1)^2 + CI^3(I+1)^3 + D \ldots$$  \hspace{1cm} (8)

where, for well deformed rare earth nuclei, the orders of magnitudes are $A \approx 10$ keV, $B \approx -10$ eV, $C \approx 10^{-2}$ eV and $D \approx 10^{-5}$ eV.

Other expansions are, however, possible. Thus Harris$^9$ has expanded according to the rotational frequency

$$E(\omega) = a\omega^2 + b\omega^4 \ldots$$  \hspace{1cm} (9)

If the expansion (9) is rapidly convergent, relations are established among the coefficients in the $I(I+1)$ expansion. With two terms one gets$^{10}$

$$\frac{C}{A} = 4 \left(\frac{B}{A}\right)^2$$

$$\frac{D}{A} = 24 \left(\frac{B}{A}\right)^3$$  \hspace{1cm} (10)

Such relations have been tested in a few cases where the experimental energies are accurate enough. The following table is due to Mottelson$^{11}$.

The relationships (10) are apparently followed to an amazing accuracy which implies a rapid convergence of the frequency expansion (9). The significance of this result is not yet clear, and it is not easy to see how it connects with the rotation-vibration interactions discussed before which after all contributes significantly to the depression of the ground state energies.

At this particular conference it should be kept in
### Analysis of Ground State Bands

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<th>( A ) keV</th>
<th>( \frac{10^3}{A} \cdot B )</th>
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<th>( -\frac{D}{A} \cdot 10^9 )</th>
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mind that many of the effects touched upon could be stronger and more well developed in nuclei far from the line of stability. The analysis of such cases should be able, once the proper techniques have been developed, to throw new light on the problems which still are so plentiful in the deformed nuclei.
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"Extrapolations" to the Z = 40 - 60 Region of Elements

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Lund, Sweden.

Introduction

As a test case for the method of shell parameter extrapolations employed in the superheavy region we have used the same prescriptions to treat the regions below the rare earth region.

These calculations have also great interest in themselves although other systematic calculations, e.g. the ones by Arseniev, Sobiscevski and Soloviev\(^1\), have been performed in the fission decay product regions, previously. The calculations described in the present paper are analogous to those of Nilsson, Tsang et al as of ref.\(^2\) in the rare-earth and actinide regions, They differ from those of ref.\(^1\) only in three respects

a) we have included the \(P_4\) degree of freedom while neglecting the gamma degree of freedom (the calculations seem to bear out that neither \(P_4\) nor the rotationally asymmetric degree of freedom is a very important degree of freedom in this particular region).

b) the Strutinsky normalisation\(^3\) is employed throughout our calculations (although for comparison some cases have been studied based on the Bős-Szymanski method).

c) in these calculations we have as a first attempt assumed \(\kappa\) and \(\mu\) to be linear functions in \(A\) and the extrapolations have been made from the values of \(\kappa\) and \(\mu\) that have been fitted to data in the rare-earth and actinide regions. In addition different alternative recipes have been
Potential employed

To calculate the single-particle energies and total energies as a function of nuclear distortion we have employed the potential of ref. \(^2\)

\[
V = V_{osc} + V_{corr}
\]

where

\[
V_{osc} = \frac{1}{2} \omega^2 \left[ (\epsilon, \epsilon) \cdot \rho^2 \left( 1 - \frac{2}{3} \epsilon P_2 (\cos \theta) + 2 \epsilon P_4 (\cos \theta) \right) \right]
\]

\[
V_{corr} = -2 \hbar^2 \omega \left[ \hat{I}_t \cdot \hat{s} + \mu \left( \hat{I}^2_t - \langle \hat{I}_t^2 \rangle_N \right) \right]
\]

For \( \omega \) we have assumed

\[
\omega = 42 \cdot A^{-1/3} \left( 1 + \frac{1}{3} \frac{N-Z}{A} \right) \text{ MeV}
\]

where the last condition assumes that the average nuclear volume is proportional to A and that the r.m.s. neutron and proton radii are roughly equal along the nuclear stability line.

These parameters have in the publication cited \(^2\) been determined from a fit of single-particle levels in the deformed rare-earth and actinide regions to the empirical level order.

The single-particle calculations

In the whole region (40<Z<62) calculations have been done with \( \kappa \) and \( \mu \) obtained by linear extrapolation in A from the rare-earth and actinide regions. For 40<Z<48 also modified values of \( \kappa \) and \( \mu \) have been used. As
seems apparent from fig. 1 where "semiempirical" $\kappa$ and $\mu$ values in the regions $A \geq 25$, $150 < A < 190$, $A > 220$ are given as circles, the linear $A$-dependence is hardly credible far below the rare earth region. We have therefore alternatively interpolated from the rare earth to the Al-region ("modified parameters") as indicated in fig. 1. Values of $\kappa$ and $\mu$ are given in table 1 and single particle level diagrams for extrapolated as well as modified parameters are shown in figs. 2 and 3.

The Bés-Szymanski and Strutinsky shell correction methods

Presently we have in these equilibrium calculations applied the shell correction method due to Strutinsky\textsuperscript{3).} Essentially it uses the liquid-drop model for normalization purposes. In the alternative Bés-Szymanski\textsuperscript{4)} method, on the other hand, the single-particle energies are simply added with pairing and Coulomb effects included. The physical finding of constant nuclear density inside of the nuclear surface region is expressed by the condition of the conservation of the volume enclosed by equipotential surfaces. For the $V_{osc}$ part of the potential this condition can be enforced for all equipotential surfaces simultaneously. As presently the other terms of the potential are not included in this condition, it is not surprising that for larger/distortions\textsuperscript{3,4)} and for moderate distortions of other multipoles\textsuperscript{5)} than $P_2$ (as e.g. $P_4$) the Bés-Szymanski method has been found insufficient unless the number of shells included are limited to a small number of shells. For nuclei heavier than $A \approx 170$ a calculation by the Bés-Szymanski method gives actually oblate minima deeper than the prolate ones contrary to experimental findings in the heavy rare earth and actinide regions as noticed by C.J. Lamm\textsuperscript{6)}.

It is easy to see that the peculiarity of the $I_t^2$ term is instrumental in bringing about the undue favouring of the oblate shape.
As in the regions here under investigation the competition between oblate and prolate shapes is very keen, we considered it important to fall back on the Strutinsky method, although in the investigation by Arseniev et al. it was reported that the difference between the results obtained alternatively with the Strutinsky and the Bézs-Szymanski methods, as far as equilibrium energies were concerned, were usually below 0.5 MeV in the mass regions $Z > 50$, $N > 82$. The difference between the region under study and the rare-earth and actinide regions probably lies in the fact that in this light region the importance of the $I_t^2$ term is relatively small as very modest $I_t^2$ values are involved.

The application of the Strutinsky normalisation is described in detail elsewhere so we shall only outline its main features. By a smearing function the single-particle level density $G(e)$, as obtained from the potential described above, is averaged with the help of essentially a Gaussian function with a range parameter that is of the order of the shell spacing. In this way a new smeared level density function $g(e)$ is obtained. We then form the difference between just filling the lowest possible energy levels or $E = \sum V e_v$ and

$$<E> = \int g(e) e \, de$$

which latter is obtained from the smeared level density function $g(e)$. The difference is a measure of the shell structure involved and is denoted the shell correction function by Strutinsky

$$\delta E_{shell} = E - <E>$$

Following the Strutinsky recipes, in place of the subtracted averaged energy $<E>$ the liquid-drop model energy, $E_{L.D.}$, is substituted. Finally a fluctuating pairing energy term
\[ \delta E_{\text{pair}} = E_{\text{PAIR}} - \langle E_{\text{PAIR}} \rangle \]
is added giving
\[ E_{\text{tot}} = \delta E_{\text{shell}} + E_{\text{L.D.}} + \delta E_{\text{pair}} \]
Minimum of \( E_{\text{tot}} \) is then determined with respect to the distortion parameters \( \epsilon \) and \( \epsilon_4 \).
The term \( \delta E_{\text{pair}} \) is evaluated by the straightforward BCS method. The pairing matrix elements \( G_n \) and \( G_p \) are assumed proportional to surface area. As for \( \langle E_{\text{PAIR}} \rangle \) we assume this term to be given by its value at equilibrium distortion. Contrary to other authors we do not assume this term deformation dependent.

**Pairing calculation. Comparisons of \( \Delta \) and total mass with data**

In the treatment of the pairing energy we have followed the recipes developed in ref.\(^2\). We have used a number of neutron and proton levels above and below the Fermi surface (for no pairing) of \( \sqrt{10N} \) and \( \sqrt{10Z} \); because of the smaller number of particles we are forced to make this change from ref.\(^2\). Subsequently we must also redetermine the pairing strength parameter, for which we have assumed an isospin dependence

\[ G_n \cdot A = g_n^0 + g_n^1 \frac{N-Z}{A} \]
\[ G_p \cdot A = g_p^0 + g_p^1 \frac{N-Z}{A} \]

Note that we have relaxed the conditions \( g_n^1 = -g_p^1 \) assumed in ref.\(^1\) expected to hold for \( N \approx Z \) and the limit that the Coulomb energy being negligible. For lighter nuclei we should expect this relation to be approached.
The $G$ values are determined by fitting the theoretical $\Delta$ to the experimental odd-even mass difference. This comparison is exhibited in figs. 7a and b. The values of $g^0$ and $g^1$ which we have used are also given there. It must be pointed out that when the distances between the energy levels near the Fermi surface is not small compared to $\Delta$, the comparison above is somewhat incorrect. Therefore the fit of $\Delta_p$ must be made essentially when $Z$ is not too near to 50 and the fit of $\Delta_n$ when $N$ is not too near to 50 and 82. We then find that for the $G$-values corresponding to the linearly extrapolated $\kappa$ and $\mu$ the isospin dependence is too great but that the fit on the average is rather good. However, just a small change of $G$ will change $\Delta$ considerably and in the investigations by Arseniev et al $^1$ it is reported that such a great change of $G$ as 10% will, as a rule, change the $\epsilon$-values of the minima by less than 0.01. Therefore, if $G$ is modified to better fit $\Delta$ to the even-odd mass differences, the equilibrium distortions will be negligibly affected. For the modified $\kappa$ and $\mu$ values the $G$ values have been re-determined. We also find that the fit is better, especially for the protons.

In fig. 8 the difference between the experimental and theoretical nuclear masses are plotted. The differences are of the same order as for the heavier nuclei, in fact they are better for the modified $\kappa$ and $\mu$ values than for the linearly extrapolated ones. It can be noticed that for the magic neutron number 82, the theoretical mass is too small. For greater $N$ values this discrepancy between theoretical and experimental masses disappears. This appears to make it probable that the theoretical $N=82$ gap in the energy levels is too great.
Results of calculations

In figs. 4a, b and c some typical total energy surfaces are shown in the \((\varepsilon, \varepsilon')\)-plane. In figs. 5a and b the \(\varepsilon\) and \(\varepsilon'\)-coordinates of the minima are given. These values as well as the depth of the minima are tabulated in table 2. One may notice that for the linearly extrapolated parameters the prolate minima is nearly always the deepest one while the modified \(\kappa\) and \(\mu\)-values give rise to a dominance of oblate distortions, at least for the nuclei which can be expected to be permanently deformed. We estimate roughly that the deformation is permanent when the energy of the deepest minimum is at least about 0.5 MeV smaller than the energy at spherical shape. The depth of the minima can be studied in figs. 6a, b, c and d.

Comparison with data

A large number of new data in the deformed region have recently become available for neutron rich even-even isotopes of \(^{40}\text{Zr}\), \(^{42}\text{Mo}\), \(^{44}\text{Ru}\) and \(^{46}\text{Pd}\). Usually only the \(2+0\) and \(4+2\) rotational transitional energies are determined empirically. The \(E_{4+}/E_{2+}\) ratio is in excess of 3 only in two of the nuclei studied. Some of these spectra still indicate stable distortions. From the gamma half-life one can also deduce the \(B(E2)\) value or \(Q_0\). We list in table 2 the experimental \(E_{2+}\), \(E_{4+}/E_{2+}\), and \(|Q_0|\) values as given by Cheifetz, Jared, Thompson, and Wilhelmy\(^7\). Corresponding values of deformation energies and \(Q_0\) for the theoretical prolate and oblate cases are given in the same table for the two alternative sets of \(\kappa\) and \(\mu\).

First it is found that the correlation between empirically low excitation energies of the \(2^+\) states and large theoretical deformation
energies (deep minima) is rather good. The transition point between spherical and deformed equilibrium distortions can be decided first when the vibrational energies are added. As a semiempirical rule it appears that we must require for stable distortions that the energy of the deepest minimum be at least 1 - 1.5 MeV smaller than the energy of spherical shape. (Note that in drawing figs 5a and b we used 0.5 MeV as the threshold.) With this rule we conclude for the modified $\kappa$ and $\mu$ values that the isotopes from $^{100}$Zr, $^{104}$Mo ($^{102}$Mo), and $^{108}$Ru ($^{106}$Ru) and on are permanently deformed. If any Pd-isotopes are permanently deformed is unclear by this requirement. For the linearly extrapolated $\kappa$ and $\mu$ parameters the deformations are greater than for the modified parameters. This is the case especially for the Zr-isotopes where the theoretical results for the extrapolated parameters appear to be clearly incorrect. This is, however, reasonable as we are there farthest away from the rare-earth region and it is more probable that the linear extrapolation is no longer satisfactory.

When it comes to a detailed comparison of the magnitudes of the theoretical and experimental $Q_0$ values, one notices that the former are usually only of the order of 60 - 80 % of the latter for the well deformed nuclei for both sets of $\kappa$ and $\mu$. This discrepancy between theory and experiment for nuclei exhibiting rotational spectra is much in excess of the discrepancy encountered in the rare-earth and actinide regions of nuclei. Conceivably this can be remedied by a modification of the unknown single-particle scheme (although this hardly seems probable). As the distortion estimates obtained by the B.-S. and the Strutinsky methods of calculation give results that here are in rather poor agreement with each other, the B.-S. method giving the better results, one might be tempted to question the local applicability of the parameters of the liquid-drop model entering through the Strutinsky method. Particularly sensitive is the surface energy term. Here the sharing of strength between
the isospin independent surface energy and surface symmetry energy terms is fixed in advance for the Mayer-Swiatecki\textsuperscript{8}) liquid-drop parameter set here employed. In the M.S. parameter choice the total strength is determined by a fit to the masses along the stability line of the mass valley. The heavy region is in this fit given particular weight. The fission decay products are far off the stability line and a deficiency in the surface symmetry term might then exhibit itself and conceivably explain some of the discrepancy in the $Q_o$-values.

Therefore we have increased this term by a factor 3 and calculated the new ($\epsilon,\epsilon_4$)-values of the minima. We find (see fig. 9) that in most cases the $\epsilon$-values of the minima are changed not more than 0.01 - 0.03, which is not enough to get the theoretical $Q_o$ as large as the experimental ones.

As indicated above we have also performed a calculation without the use of the Strutinsky normalisation. Also for this case we have calculated the ($\epsilon,\epsilon_4$)-values of the minima. On the prolate side the change is not significant while on the oblate side the $\epsilon$-values of the minima might be changed as much as 0.10 (see fig. 9 and 10). If we do not use the Strutinsky method the tendency for oblate distortions is increased and we find for both of the regions, $Z < 50$ and $Z > 50$, that the oblate minima in most cases are the deepest ones. We also find that in some cases the $\epsilon_4$-distortions are very much affected by the Strutinsky normalisation. (This parallels the results of P. Möller\textsuperscript{5}) in the rare-earth region.)
Conclusions

It appears that a new deformed region is well established experimentally with the study of the neutron rich light fission products. This region of distortion is well brought out by the theoretical calculations. The magnitude of the distortions are less well reproduced and it is speculated that some of the discrepancy may be due to the Strutinsky method of calculation or rather the liquid drop parameters entering into the theoretical calculations through the employment of the Strutinsky shell correction method.

Acknowledgements

For guidance and most generous help the author wishes to express his deep gratitude to Professor S.G. Nilsson, who suggested this investigation. I also wish to thank all members at the Department of Mathematical Physics in Lund, especially P. Møller for generous access to computer programs. Comments on the work by Professor B. Mottelson are gratefully acknowledged.
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   Z. Szymanski, Nucl. Phys. 28 (1961) 63


6. C.J. Lamm, private communication


Table captions

Table 1  Values of $\kappa$ and $\mu$ employed in the single-particle calculation corresponding to different regions of mass.

Table 2a  The coordinates and corresponding intrinsic quadrupole moments of the oblate ($\varepsilon_0^-, (\varepsilon_4)_0^-, Q_0^-$) and prolate minimum ($\varepsilon_0^+, (\varepsilon_4)_0^+, Q_0^+$) for $Z=54, 56, 58, 60$ and $62$. $E^+$ is the depth of the prolate minimum and $\Delta E$ is the difference in depth between the oblate and prolate minimum. If $\Delta E > 0$ the prolate minimum is the deepest one. Linearly extrapolated $\kappa$ and $\mu$ values have been used.

Table 2b  Same as table 2a for $Z=40, 42, 44, 46$ and $48$. The results are given for linearly extrapolated as well as modified $\kappa$ and $\mu$ values.

Table 3  Experimental results taken from ref$^7$) compared to theory. $E_{\text{max}}$ is the deepest minimum. The quadrupole moments corresponding to both minima are given, for the deepest minimum no parenthesis is used.
Table 1

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- Fitted to the experimental energy levels
- Extrapolated from $A = 165$ and $A = 242$
- Interpolated from $A = 165$ to $A = 25$
<p>| 116_{Xe} | -0.16 | -0.01 | -224 | 0.20 | -0.01 | 332 | 0.8 | 0.4 |
| 118_{Xe} | -0.18 | -0.01 | -246 | 0.21 | 0.00 | 349 | 0.9 | 0.3 |
| 120_{Xe} | -0.19 | 0.00 | -260 | 0.21 | 0.01 | 358 | 1.1 | 0.2 |
| 122_{Xe} | -0.18 | 0.00 | -260 | 0.20 | 0.01 | 335 | 1.0 | 0.0 |
| 124_{Xe} | -0.18 | 0.01 | -253 | 0.17 | 0.01 | 293 | 0.8 | -0.1 |
| 126_{Xe} | -0.15 | 0.01 | -218 | 0.14 | 0.01 | 234 | 0.7 | -0.1 |
| 144_{Xe} | -0.14 | -0.03 | -226 | 0.16 | -0.06 | 322 | 2.0 | 1.0 |
| 120_{Ba} | -0.20 | -0.01 | -293 | 0.25 | 0.00 | 450 | 2.1 | 0.7 |
| 122_{Ba} | -0.20 | -0.01 | -298 | 0.25 | 0.01 | 456 | 2.2 | 0.6 |
| 124_{Ba} | -0.20 | 0.00 | -294 | 0.24 | 0.02 | 431 | 2.1 | 0.3 |
| 126_{Ba} | -0.20 | 0.01 | -295 | 0.20 | 0.01 | 360 | 0.1 | 0.0 |
| 128_{Ba} | -0.19 | 0.01 | -284 | 0.19 | 0.02 | 337 | 1.3 | 0.0 |
| 130_{Ba} | -0.15 | 0.02 | -226 | 0.24 | 0.01 | 253 | 0.8 | 0.0 |
| 144_{Ba} | -0.13 | -0.03 | -217 | 0.11 | -0.04 | 220 | 0.8 | 0.3 |
| 124_{Ce} | -0.23 | 0.00 | -346 | 0.28 | 0.01 | 537 | 3.2 | 1.1 |
| 126_{Ce} | -0.22 | 0.00 | -340 | 0.27 | 0.02 | 520 | 3.0 | 0.8 |
| 128_{Ce} | -0.22 | 0.01 | -331 | 0.24 | 0.02 | 466 | 2.5 | 0.4 |
| 130_{Ce} | -0.20 | 0.01 | -310 | 0.20 | 0.02 | 378 | 2.0 | 0.2 |
| 132_{Ce} | -0.18 | 0.02 | -279 | 0.18 | 0.02 | 339 | 1.2 | 0.1 |
| 134_{Ce} | -0.12 | 0.02 | -195 | 0.10 | 0.01 | 188 | 0.5 | 0.0 |
| 146_{Ce} | -0.14 | -0.03 | -256 | 0.14 | -0.04 | 303 | 1.0 | 0.4 |
| 148_{Ce} | -0.19 | -0.04 | -330 | 0.20 | -0.06 | 456 | 3.3 | 1.5 |
| 132_{Nd} | -0.22 | 0.02 | -346 | 0.23 | 0.02 | 454 | 2.5 | 0.5 |
| 134_{Nd} | -0.20 | 0.02 | -323 | 0.20 | 0.02 | 397 | 1.7 | 0.3 |
| 136_{Nd} | -0.14 | 0.02 | -244 | 0.14 | 0.02 | 276 | 0.7 | 0.0 |
| 148_{Nd} | -0.15 | -0.03 | -276 | 0.17 | -0.04 | 382 | 1.2 | 0.5 |
| 150_{Nd} | -0.19 | -0.03 | -357 | 0.20 | -0.05 | 475 | 3.6 | 1.8 |
| 138_{Sm} | -0.16 | 0.02 | -285 | 0.18 | 0.02 | 365 | 0.9 | 0.1 |
| 140_{Sm} | -0.10 | 0.02 | -181 | 0.06 | 0.01 | 162 | 0.1 | -0.1 |
| 150_{Sm} | -0.15 | -0.03 | -281 | 0.18 | -0.04 | 425 | 1.3 | 0.6 |
| 152_{Sm} | -0.20 | -0.03 | -378 | 0.21 | -0.04 | 515 | 3.6 | 1.8 |
| 154_{Sm} | -0.20 | -0.03 | -305 | 0.23 | -0.04 | 566 | 6.0 | 3.2 |
| 158_{Sm} | -0.21 | -0.02 | -408 | 0.24 | -0.04 | 611 | 7.9 | 4.4 |
| 158_{Sm} | -0.22 | -0.02 | -424 | 0.26 | -0.02 | 646 | 9.6 | 5.4 |</p>
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<tr>
<td>$^{108}$Ru</td>
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The variation of $\kappa$ and $\mu$ with the mass value $A$. For $A = 165$ and 242 the values have been chosen to fit the experimental energy levels as well as possible. For $A = 25$ the parameters are also considered to be known. To find the $\kappa$ and $\mu$ values for $A = 90-150$ we have used two recipes, first a linear extrapolation from the heavier regions and secondly an interpolation between the rare-earth and Al regions ("modified parameters") as shown in the figure.
Fig. 2a  Single-proton levels for $A \simeq 140$ and linearly extrapolated $\kappa$ and $\mu$ values. The levels are assigned asymptotic quantum numbers. Solid lines mark even parity while dashed lines represent odd parity. ($\epsilon_4 = 0$ is assumed in this figure.)
Fig. 2b  Single-neutron levels for A > 140 based on linearly extrapolated 
κ and μ parameters, (c₄ = 0)
Fig. 3a  Single-proton levels for the modified parameters, $A \approx 110, \ (\varepsilon_4 = 0)$
Fig. 4a  The total energy surface in the $(\varepsilon, \varepsilon_4)$-plane for $^{108}_{44}$Ru.  

The $\kappa$ and $\mu$ values are linearly extrapolated from rare-earth and actinide regions. The lines represent steps of 1 MeV. The prolate minimum is 0.4 MeV deeper than the oblate one.
Fig. 4b  Same as fig. 4a except that the $\kappa$ and $\mu$ values are interpolated between the Al and rare-earth regions and the pairing strength $G$ is correspondingly changed. The oblate minimum is found to be 0.3 MeV deeper than the prolate one.
Fig. 4c  Same as fig. 4a for the double-magic nuclieons $^{132}_{50}Sn$. 

TOT. ENERGY, SCALE 1.0 MeV, Z=50 A=132
Fig. 5a  The minima of the total energy plotted in the $(\epsilon, \epsilon_n)$-plane for $Z=52, 54, 56, 58, 60,$ and 62. For each nucleus the lowest minimum is marked by a circle. In the case that the energy of a minimum differs from the spherical-shape energy by less than 0.5 MeV it is marked by a cross, if it differs by more than 0.5 MeV it is marked by a point. The $\kappa$ and $\mu$ values used are linearly extrapolated ones.
Fig. 5b  Same as fig. 5a for Z=40, 42, 44, 46, and 48. In the upper part of the figure the linearly extrapolated $\kappa$ and $\nu$ values are used, in the lower the modified values.
Fig. 6a  Contours in the (N,Z)-plane for the depths of the prolate and oblate minima. Linearly extrapolated values of $\kappa$ and $\mu$ are used.
Fig. 6b  Contours in the (N,Z)-plane for the depths of the oblate minima.
For the dotted lines $\kappa$ and $\mu$ are linearly extrapolated while for
the solid lines modified parameters have been used.
Fig. 6c: Contours for the difference in depth between the oblate and prolate minima. $k$ and $\mu$ are linearly extrapolated. If $AE < 0$ the prolate minimum is the deepest one.
Modif. param. 
$\Delta E = E^-/\text{min} - E^+/\text{min}$

Fig. 6d  Same as fig. 6c for the modified parameters.
Fig. 7a  Experimental odd-even mass differences and calculated $\Delta$ values.
The $G$ values as shown in the figure are those used together with the linearly extrapolated $x$ and $\mu$ in the calculations.
Fig. 7b  Same as fig. 7a. The crosses represent the $\Delta$ values obtained for the $G$ values used together with the linearly extrapolated $\kappa$ and $\mu$, the points correspond to the $G$ values used together with the modified $\kappa$ and $\mu$. 
Fig. 8  The difference between the experimental and theoretical masses. The crosses is the fit achieved when $\kappa$ and $\mu$ are linearly extrapolated while modified $\kappa$ and $\mu$ parameters are associated with the points.
Fig. 9 The minima of the total energy plotted in the $(\varepsilon, \epsilon)$-plane.

Linearly extrapolated $\kappa$ and $\mu$ values have been used. The points indicate the same minima as fig. 5b, the crosses are obtained when the Bö-Szymanski method is used and the circles when the Myers-Swiatacki isospin dependent surface energy term is increased by a factor 3 and the Strutinsky method is used.
Fig. 10  The experimental intrinsic quadrupole moment, indicated by points\textsuperscript{7)}, compared to theoretical results. In the upper part of the figure the linearly extrapolated $\kappa$ and $\nu$ values are used, in the lower the modified values. The squares are the quadrupole moments corresponding to the prolate minima, the triangles are associated with the oblate minima. If a square or triangle is open, the Strutinsky method has been used, if it is filled, the Bős-Szymanski method has been used.
EXCITED STATE IN EVEN-EVEN FISSION PRODUCTS

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ABSTRACT

Experimental results are presented on the ground state bands of 31 even-even nuclei produced in the primary fission of $^{252}$Cf. Systematics of the energies and life times of the transitions are given. From these data there is evidence indicating that a region of deformation in the $A \approx 100$ region may exist.

1. INTRODUCTION

In this talk we shall present information concerning the energy levels of very neutron rich even-even isotopes of elements with $38 \leq Z \leq 62$. Fission fragments from spontaneous fission of $^{252}$Cf provided experimental access to this region. The data which in a few of the cases can be correlated with already known levels, extend the knowledge about the systematic behaviour of collective excitations to neutron rich nuclei far from the beta stability line. The data include the lowest $2^+$ states (and in many of the cases some of the higher $4^+, 6^+$ and $8^+$ levels) of all the isotopes that have calculated independent fission yield of more than 1% (with the exception of $^{136}$Te). This is shown in Figs. 1 and 2 where we present a modified chart of nuclides indicating the isotopes observed in this experiment. The systematics of the energy levels of

* Work performed under the auspices of the U.S. Atomic Energy Commission.
† On leave of absence from the Weizmann Institute of Science, Rehovot, Israel.
isotopes that are removed from closed shells are well fitted using the phenomenological model of Mariscotti et al. There is evidence from level spacings and B(E2) values that light fission fragments such as $^{102}$Zr and $^{106}$Mo have rotational like behaviour. These results support recent theoretical studies of Ragnarsson and Nilsson and Arseniev et al. which have predicted a new region of stable deformation which includes these nuclei. Another feature of the results is the evidence that the well known abrupt discontinuity of the ratio $E^+_4/E^+_2$ for isotopes with 88 to 90 neutrons reaches its maximum effect in Nd, Sm and Gd, and becomes much smoother in the Ce and Ba nuclei.

2. **THE EXPERIMENTAL TECHNIQUE**

Prompt K x-rays and/or γ-rays in coincidence with pairs of fission fragments were measured using the detector arrangement indicated in Fig. 3. Three separate experiments using different photon detectors were performed: 1) recording γ-rays with a 1 cm$^3$ Ge(Li) detector (resolution 1 keV at 122 keV) in position $γ_2$; 2) recording γ-rays and/or x-rays in coincidence using a 6 cm$^3$ Ge(Li) detector in position $γ_1$ and a 2 cm$^2$ Si(Li) detector in position $γ_2$; 3) recording γ-ray, γ-ray coincidences with a 35 cm$^3$ Ge(Li) coaxial detector in position $γ_2$ and a 6 cm$^3$ Ge(Li) detector in position $γ_1$. In all the experiments a nominally $10^5$ fission per minute source of $^{252}$Cf was electrodeposited onto the surface of fragment detector F1. Thus Doppler shifting and broadening problems were minimized for transitions from the fragments stopped in that detector. This technique, which simplified the spectra, applies to half life times longer than the stopping time of the fragments ($\sim 10^{-12}$ sec). Life time determinations in the time region 0.1 - 2.0 nsec were obtained from the
ratio of the non-Doppler shifted gamma ray intensity observed when the fragment stopped in the plated detector F1 relative to the intensity observed when the fragment stopped in the second detector F2, which was separated from the plated detector by 8 mm. The various detector systems were digitally gain stabilized using external gamma ray sources as indicated in Fig. 3.

In all the experiments the analog pulse heights were digitized and stored event by event in a PDP-9 computer. The on-line computer was programmed to monitor the resolution of the detectors and to transfer the experimental data onto magnetic tape in a compressed format. A total $2 \times 10^8$ multiparameter events were recorded and later processed on a CDC 6600 computer.

The masses of the fragments were calculated from the measured energies using the Schmitt calibration method and the known neutron corrections. Gamma-ray spectra associated with fragment masses in 2 amu wide mass intervals were obtained by sorting the three parameter data.

In Fig. 4 examples of high resolution $\gamma$-ray spectra associated with 2 amu mass intervals are shown. Each of these $\gamma$-ray spectra was then analyzed to give quantitative energies and intensities of individual transitions. This was accomplished using the on-line photopeak analysis code developed by Routti and Prusin. The widths of the mass distributions associated with single gamma transitions ranged from 0.4 to 6.5 amu (FWHM) and the mean values of the masses for these distributions were determined with standard statistical errors of less than 0.2 amu for the strong transitions; however, the absolute determination of the masses are uncertain by $\pm 1$ amu. This can be attributed mainly to the use of the average neutron corrections. The average value for a given mass is presumably too small for isotopes that are on the neutron
deficient side of the most probable isotope for a given element and conversely too large for isotopes on the neutron excess side. In some cases we have observed odd mass isotopes of even Z elements with experimental masses between those of the adjacent even isotopes although the latter were separated by less than 2 amu. There may also be a small smooth systematic error due to the uncertainties in the fragment pulse height to energy calibration procedure. The x-ray, gamma-ray coincidence data were used to obtain information on additional transitions associated with single isotopes.

3. RESULTS

The results of this investigation are summarized in Table I. For each isotope in the table we present the transition energies that were observed, the half life of the $2^+$ level and the yield per fission of the $2^+$ level (corrected for internal conversion). The considerations and criteria that have been used in assigning these levels have already been explained in two previous papers $^{6,7}$. Some of the levels have been assigned by other workers and are repeated here because the transitions have been observed in the prompt gamma spectra. The cases in which our results have been correlated with other works are explained in the comments under the table.

Some of the lowest $2^+ \rightarrow 0^+$ transitions in isotopes on the neutron deficient side of the Zp line have been observed in (t,p) reactions and in post beta decay of fission products. Our results are in good agreement with values obtained by these other methods. This agreement confirms the techniques we have utilized for making level assignments.

Most of the transitions that we have assigned as $2^+ \rightarrow 0^+$ have also
been observed (though not previously assigned) in studies of gamma rays
emitted following beta decay of unseparated prompt fission products.\textsuperscript{10)}
These $2^+ \rightarrow 0^+$ transitions occurring after beta decay are observed with
intensities which are substantial fractions of the cumulative mass chain
yields. These measurements help substantiate that we are in fact observing
ground state band transitions in the de-excitation of the prompt products.

4. **DISCUSSION**

4.1 **Systematics of Energy Levels**

The systematics of the energy levels of the ground state band in
the cases where $E4^+/E2^+ > 2.3$ are well fitted using the phenomenological
variable moment of inertia model of Mariscott \textit{et al.}\textsuperscript{1}). The results of
$E6^+/E2^+$ and $E8^+/E2^+$ vs. $E4^+/E2^+$ for the ground state bands are shown
in Fig. 5 which includes our current data and all the previously known
experimental data that were summarized in reference 1. Although our
data contains lighter nuclei than those of reference 1, no large devia-
tions from the known systematic behaviour are observed.

The energies of the $2^+$ states obey smooth systematics. The trend
shows a decrease in the $2^+$ level energies with \textit{increased displacement}
from the $N=50$ and $N=82$ shells. In the light fragments ($Z \leq 48$) there
is a decrease of $2^+$ level energies with displacement from the $Z=50$
shell. On the other hand, in the heavy fragments ($Z \geq 50$) region isotopes
with neutron number between 82 and 86 have an increasing $2^+$ level
energies with increasing displacement from the $Z=50$ shell.

4.2 **Rotational-like Light Fission Fragments**

In Fig. 6 we present the systematic behaviour of the lowest $2^+$
levels and the $E4^+/E2^+$ ratio in the Zr-Pd region. One can see
clearly that the smooth decrease in level energies and corresponding smooth increase in the E4/E2 ratio in Pd becomes more rapid in Ru and Mo isotopes and then turns out to be an extreme jump between 98Zr - 100Zr. A similar behaviour is seen in Fig. 7 that shows the B(E2)_{exp}/B(E2)_{single particle} ratio for this region. The B(E2) values were obtained from the measured energies and life time values following the formalism of Stelson and Grodzins. The $\beta_2$ values in that formalism are $\approx 0.35$ for $^{110}$Ru, 0.45 for $^{106}$Mo and $^{106}$Mo, and 0.6 for $^{102}$Zr. In principle, our reported life time values are upper limits because a hold up can occur in a transition before the $2^+ \rightarrow 0^+$ stage. This implies that the experimental B(E2) and $\beta_2$ values can only be higher.

The central question from these studies is whether the theoretical predictions for deformation can be verified. It is not possible to determine the existence of static deformations from observed energy level spacings or from measurements of B(E2; 2 + 0). However, studies of such systematics are indicative of nuclear softness and therefore it is of interest to compare these properties in this new region with the corresponding values for the rare earth and actinide regions which are the two major areas of known permanent deformation. There are several different indicators of deformation and it is informative to compare each. Figure 8 is a composite plot containing five indicators associated with deformation $[\beta_2, B(E2)/B(E2)_{SP}, E_4/E_2^+, \beta_2/\beta_{2SP}, (79.51/E_2^+) \times (158/A)^{5/3}]$ plotted as a function of mass. The last indicator represents to a first approximation a mass independent comparison of the energies of the first $2^+$ states using arbitrarily the deformed $^{158}$Cd nucleus as a reference. The nuclei presented in the plot include the current, light fission product region, and a representative sampling of isotopes in the rare earths (150 - 180) and in the actinides (224 - 244).
In this light fission-product region, of the isotopes studied, $^{102}$Zr appears as the most favorable candidate for deformation. Its value for $\beta_2$ (0.604) and for the mass independent energy parameter (1.08) are larger than any of the corresponding values found in the rare earth and actinide nuclei. Also its values for B(E2)/B(E2)$_{SP}$ (234.) and $\beta_2$/B$_{2SP}$ (15.2) are larger than for any of the rare earths though smaller than some of the actinides. The only parameter for which it has a lower value than obtained in the other regions is the $E_{4+}/E_{2+}$ ratio where the $^{102}$Zr value of 3.15 is somewhat smaller than the limiting value for a perfect rotor (3.33) which is closely approached in both the rare earth and actinides. The other new isotopes for which we present information have smaller values for these deformation indicators than $^{102}$Zr but even they have, in several instances, values comparable or larger than those typically found in the rare earth and actinide region and in all cases are larger than the values found for spherical nuclei near closed shells.

For the isotopes with higher masses the decrease in the deformation indicators is believed to be due to the approach of the Z = 50 closed shell, and for the lighter isotopes the effect of the N = 50 shell should be important. The theoretical calculations of Arseniev et al. imply that the regions of strongest deformation should be in the heavier isotopes of strontium (98 - 102) and of krypton (96 - 102) which are not produced in significant yield in the $^{252}$Cf fission process. The change in the energy of the lowest $2^+$ level between $^{98}$Zr and $^{100}$Zr is very extreme and larger than the well noted discontinuity between $^{150}$Sm and $^{152}$Sm. The 1.233 MeV first excited $2^+$ level in $^{98}$Zr is supported by recent $^{96}$Zr($t$,p)$^{98}$Zr reaction studies and also by our gamma ray spectra for which no line was found that could be attributed to a lower $2^+$ state.
For this $A \approx 100$ region the theoretically predicted trend of increasing $\beta_2$ deformation with displacement from the $Z = 50$ shell was indeed found by the experiment, however the magnitude of the experimental $\beta_2$ value are much larger than the predicted values. This poses questions regarding the parameters of the Nilsson levels in this region and regarding dependence of the surface energy coefficient on the neutron excess. The very sharp changes in deformation parameters in the Zr isotopes are also not predicted by the calculations.

Although experimental information concerning rotational bands in neutron rich Ru nuclei has been reported previously by Johansson$^{18}$ and by Zicha et al.$^{19}$ we are unable to reproduce their results. We have not been able to find any of the $\gamma$-rays reported by them in coincidence with transitions we have assigned to the ground state bands of $^{108}$Ru and $^{110}$Ru.

4.3 88-90 Neutron Discontinuity

The data described here include $Z = 56$ and $Z = 58$ nuclei having 88 and 90 neutrons. The data show that the 88-90 neutron discontinuity is smearing out as the proton number decreases below $Z = 60$. This is seen both for the energies of the $2^+$ level and the $E4^+/E2^+$ ratio. A plot of the $E4/E2$ ratio of nuclei with $56 \leq Z \leq 70$ is presented in Fig. 9. This figure clearly shows that the maximum effect of the 88-90 neutron discontinuity occurs in the region $60 \leq Z \leq 66$. The nuclei with $N = 92$, $Z = 58$, and $N = 92$, $Z = 68$ have $E4/E2$ ratios and $B(E2)$ values which indicate that they are rotational as $^{152}_{62}$Sm, which is known to have permanent quadrupole deformation, even though for $Z = 58$ and $Z = 68$ the 88-90 neutron effect is relatively rather smooth. A similar effect has been shown to occur in the 76-80 proton number region$^{20}$ where a sharp discontinuity occurs for $106 \leq N \leq 112$ and a smooth behaviour was observed outside this region. We can summarize then that the transition from a vibrational spectrum to a rotational one
can be either abrupt or smooth depending probably on a delicate balance
between proton and neutron pairing correlations. Calculations by Nilsson
\text{et al.}^{21)} indicate that deformation is expected to occur abruptly between
86 and 88 neutrons for the nuclei discussed here. These calculations are
similar to the ones described in the light fragment region. They are based
on the Nilsson model combined with the Strutinsky normalization procedure
and reproduce the general trend of decreased deformation for nuclei with
88 neutrons on both sides of $Z = 62$.

We are grateful to the following persons for their help in this work:
Elizabeth Quigg wrote the necessary programs for the PDP-9 computer. Thomas
Strong handled the processing of our data using the CDC 6600 computer.
Robert Latimer and James Harris electrodeposited the $^{252}$Cf sources on our
fission detectors. Very useful discussions with John Rasmussen, Chin
Fu Tsang, Frank Stephens, and Rand Watson are acknowledged.
REFERENCES


8. R. Foucher, Orsay, France, private communication.


Table I. Experimental results for ground state bands.

|          | Transition Energy in keV | $t_{1/2}(2 + 0)$ | Yield $2^+$
|----------|--------------------------|------------------|----------------
|          | 2$^+ + 0^+$              | 4$^+ + 2^+$      | 6$^+ + 4^+$    | 8$^+ + 6^+$ |
| $^{94}_{\text{Sr}}$ | 837.4                    |                  |                |              | 0.51          |
| $^{95}_{\text{Sr}}$ | 815.5                    |                  |                |              | 0.37          |
| $^{98}_{\text{Zr}}$ | 1223                     |                  |                |              | $\sim 0.3$    |
| $^{100}_{\text{Zr}}$ | 212.7                    | 352.1            | 497.9          |              | 0.52          | 1.3           |
| $^{102}_{\text{Zr}}$ | 151.9                    | 326.6            | 486            | (587)        | 0.86          | 1.3           |
| $^{102}_{\text{Mo}}$ | 296                      |                  |                |              | 0.06          |
| $^{104}_{\text{Mo}}$ | 192.3                    | 368.7            | 520.0          |              | 0.45          | 3.37          |
| $^{106}_{\text{Mo}}$ | 171.7                    | 350.8            | (511.8)        |              | 0.75          | 3.37          |
| $^{106}_{\text{Ru}}$ | 269                      |                  |                |              | 0.16          |
| $^{108}_{\text{Ru}}$ | 242.3                    | 423              |                |              | 0.22          | 1.44          |
| $^{110}_{\text{Ru}}$ | 240.8                    | 423.1            | 576.1          | (707.7)      | 0.23          | 3.49          |
| $^{112}_{\text{Ru}}$ | 236.8                    | 408.9            |                |              | 0.20          | 0.97          |
| $^{112}_{\text{Pd}}$ | 348.8                    | 535.8            |                |              | 0.77          |
| $^{114}_{\text{Pd}}$ | 332.9                    | 520.7            | 649.3          |              | 1.48          |
| $^{116}_{\text{Pd}}$ | 340.6                    | 538.0            |                |              | 0.87          |
| $^{118}_{\text{Cd}}$ | 488.0                    | 677.3            | (805)          |              | 0.32          |
| $^{132}_{\text{Te}}$ | 974                      | 990              |                |              | $\sim 0.2$   |
| $^{134}_{\text{Te}, d, e}$ | 1278                    | 297              | 115            |              | 1.5           |
| $^{138}_{\text{Xe}}$ | 589.5                    | 482              |                |              | 2.3           |
| $^{140}_{\text{Xe}}$ | 376.8                    | 457.9            |                |              | 1.5           |
| $^{140}_{\text{Ba}}$ | 602.2                    | 632              |                |              | 0.52          |
| $^{142}_{\text{Ba}}$ | 359.7                    | 475.7            | 632            |              | 2.90          |

(continued)
Table I. Continued

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\(^a\) The 2^+ \rightarrow 0^+ transitions in $^{94}_{\text{Sr}}$ and $^{96}_{\text{Sr}}$ have been assigned by R. Foucher et al. using an on-line mass separator.

\(^b\) The first 2^+ in $^{96}_{\text{Zr}}$ was found by Blair et al. using the $^{96}_{\text{Zr}}(t,p)^{98}_{\text{Zr}}$ reaction.

\(^c\) The first 2^+ in $^{106}_{\text{Ru}}$ was reported by Casten and was produced by the $^{104}_{\text{Ru}}(t,p)^{106}_{\text{Ru}}$ reaction.

\(^d\) The 2^+ \rightarrow 0^+ and 4^+ \rightarrow 2^+ transitions in $^{132}_{\text{Te}}$ and the 2^+ \rightarrow 0^+ transition in $^{134}_{\text{Te}}$ were observed by Bergström et al. following beta decay of fission products analyzed by an on-line mass separator.

\(^e\) The decay of $^{134}_{\text{Te}}$ has also been reported by John et al. who measured the 6^+ \rightarrow 4^+, 4^+ \rightarrow 2^+, and 2^+ \rightarrow 0^+ transition in their study on delayed gamma-rays in $^{252}_{\text{Cf}}$ fission.

\(^f\) The 2^+ \rightarrow 0^+ transitions in $^{140}_{\text{Ba}}$ and $^{142}_{\text{Ba}}$ have been observed by Alväger et al. following beta decay of mass separated fragments.

\(^g\) The 2^+ \rightarrow 0^+ transition in $^{144}_{\text{Ce}}$ has been observed by Wilhelmy et al.

\(^h\) The rotational band in $^{156}_{\text{Sm}}$ has been observed in the $^{154}_{\text{Sm}}(t,p)^{156}_{\text{Sm}}$ reaction by Bjerrgaard et al.
Fig. 1. A region of the chart of nuclides presenting the isotopes produced as light fission products. The solid line approximates the line of β stability and the heavy squares in this region are the beta stable isotopes. The dotted line represents the $Z_p$ values for the fission of $^{252}\text{Cf}$ and the heavy squares with diagonal lines in the lower right hand corner are the even-even products observed in the current experiments. The heavy squares with an additional diagonal line in the upper corner are isotopes for which our assignments are tentative.
Fig. 3. Schematic representation of detector system. Detectors F1 (with electrodeposited $^{252}$Cf) and F2 measured energies of fragments. Detectors $\gamma_1$ and $\gamma_2$ measured energies of $\gamma$-rays and/or x-rays. External sources for stabilization of the photon detectors were $^{243}$Cm ($\alpha - \gamma$ coincidence), $^{60}$Co ($\gamma - \gamma$ coincidence) and $^{241}$Am ($\alpha - \gamma$ coincidence).
Fig. 4. Gamma ray spectra recorded using a 1 cm$^3$ high resolution Ge(Li) detector. Results shown are for fission fragments that were stopped in the plated detector having masses 103-105 and 105-107.
Fig. 5. Plot of the energy ratios $E_{6+}/E_{2+}$ and $E_{6+}/E_{2+}$ versus $E_{4+}/E_{2+}$. The data presented as dots were taken from ref. 1 and those presented as open squares are the current experimental results.
Fig. 6. Energies of the lowest $2^+$ levels and $E_{4^+}/E_{2^+}$ energy ratios in the Zr - Pd region.
Fig. 7. $\frac{B(E2)_{\text{exp}}}{B(E2)_{\text{SP}}}$ in the Zr – Pd region. The data presented for the most neutron rich isotopes are from the current experimental results and the others are from ref. 17.
Fig. 8. A composite plot presenting five indicators of deformation plotted as a function of mass. The mass intervals used contain only the current light fission product experimental region and a representative sampling from the two major known regions of deformation. The values of $\beta_2$, $\beta_{2SP}$, $B(E2)$, and $B(E2)_{SP}$ were extracted from relationships presented in ref. 17; the indicator $(79.51/E_{2+}) \times (158/A)^{5/3}$, gives a relative comparison between the energies of the first $2^+$ states on a basis which removes the inherent mass-dependence from the moment of inertia. The open circles represent current results obtained using experimental energies and life times. The open squares represent current results obtained using experimental energies and calculated life times (ref. 1). The closed circles represent literature values (refs. 1 and 17).
Fig. 9. Systematic behavior of the ratio \( \frac{E_4}{E_2} \) as a function of proton number in the \( N = 86 - 92 \) region. Data presented as \( \Delta \) are from the current experimental results and the other data are from refs. 22 and 23.
LIFETIME MEASUREMENTS ON NEUTRON DEFICIENT EVEN-EVEN
ISOTOPES OF XE AND BA

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Physikalisches Institut der Universität Heidelberg,
Heidelberg, Germany

1. INTRODUCTION

In this paper we should like to report on some
lifetime measurements on neutron deficient even-even
isotopes of Xe and Ba, which give additional information
on one of the so called "new regions of deformation
far off the stability line". These regions were first
indicated by Sheline1) and the light Xe and Ba isotopes
are lying in the region with 50 < N,Z < 82. It was shown
by several groups2,3) that the neutron deficient even-
even nuclei in this region show nice quasi-rotational
ground state bands. Fig. 1 shows the level schemes of
the even-even Ba and Xe isotopes. Although the ground
state bands of the lightest isotopes look very rotational,
they are far from following the I(I+1) rule. In order
to get some information about transition probabilities
and deformation parameters we have tried to measure
the lifetime of the first 2+ states in 120,122 Xe and
126,128 Ba. Since the lifetimes were expected to be in

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the pico second range we have applied the recoil distance Doppler shift technique. Suitable reactions to produce neutron deficient nuclei in the desired region are $^{(\alpha,xn)}$ and $^{(nI,xn)}$. The latter is especially convenient for Doppler shift measurements, since the large momentum transfer of the heavy ion reaction gives large recoil velocities and thus large Doppler shifts. In addition, the evaporation of several neutrons out of the compound nucleus with relatively low energies results in a tightly collimated beam of residual nuclei recoiling along the beam axis.

2. EXPERIMENTAL PROCEDURE

The Xe and Ba nuclei were produced via the reactions $^{108,110}$Pd($^{16}$O,4n)$^{120,122}$Xe and $^{114,116}$Cd($^{16}$O,4n)$^{126,128}$Ba. The $^{16}$O was obtained from the Heidelberg Emperor tandem with an energy up to 80 MeV. Targets were prepared by evaporating approximately 500 $\mu$g/cm$^2$ onto 1 mg/cm$^2$ gold backings. The target was mounted in such a way that the beam had to pass first the gold backing and then hit the target. The recoiling nuclei were stopped in a gold covered plunger plate, which could be moved with a precision micrometer towards the target. The probability of reaching the stopper without decay in flight is simply given by the decay law $e^{-t/\tau}$, where the flight time $t$ is determined by the stopper distance $D$ and the recoil velocity $v$: $t = D/v$. This probability can be measured with a Ge(Li) detector at 0° with respect to the beam axis, since the deexciting gammarays split into two peaks, one produced by the decay of nuclei already stopped in the plunger and one by the decay in flight, which gives a Doppler shifted line. The above mentioned probability is then given through the ratio of the unshifted to the total gamma ray intensity. Fig. 2 shows a number of spectra for the reaction $^{110}$Pd($^{16}$O,4n)$^{122}$Xe taken at different plunger distances. One can clearly
see how the relative peak intensities vary with the plunger distance. A plot of the unshifted fraction versus the plunger distance D gives exponential decay curves, from which one can extract the lifetime, if the recoil velocity is also known.

3. DISCUSSION

The results of the 2⁺ lifetimes are summarized in table 1. The deduced transition strengths are in between the values of pure vibrational and rotational nuclei. The deformation parameter ε is in rather good agreement with the theoretical results of the Dubna group⁴), who calculated an oblate equilibrium shape for the nuclei in this region. Unfortunately the transition probabilities are not suited to get the sign of the deformation.

For the future it seems promising to measure the lifetimes of several members of the ground state band with high precision in order to compare relative B(E2) values with the predictions of various models⁵).

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| Nucleus | Transition | Energy [keV] | $\alpha$ ($\alpha'_K + 1.3\alpha'_L$) [%] | $\tau$ [ps] | $B(E2)/B(E2)_{sp}$ [W.U.] | $|\varepsilon_{\text{exp}}|$ | $|\varepsilon_{\text{the}}|$ |
|---------|------------|-------------|--------------------------------|-----------|-------------------------|-----------------|-----------------|
| $^{120}_{\text{Xe}}$ | $2^+ \rightarrow 0^+$ | 322 | 0.034 | 1.25 | 125$\pm$15 | 51 $\pm$ 6 | 0.20 | -0.25 |
| $^{122}_{\text{Xe}}$ | $2^+ \rightarrow 0^+$ | 331 | 0.031 | 1.12 | 77.7$\pm$8.5 | 77 $\pm$ 9 | 0.23 | -0.23 |
| $^{126}_{\text{Ba}}$ | $2^+ \rightarrow 0^+$ | 256 | 0.077 | 0.97 | 265$\pm$25 | 68 $\pm$ 7 | 0.23 | -0.25 |
| $^{128}_{\text{Ba}}$ | $2^+ \rightarrow 0^+$ | 279 | 0.060 | 1.19 | 140$\pm$30 $<190$ c) | 83 $\pm$ 17 | 0.26 | -0.22 |

a) D.A.Arseniev et al., Nucl.Phys., A126 (1969) 15
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Fig. 1. Levels of neutron deficient even-even isotopes of Xe and Ba. Data are collected from ref. 2, 3, 6, 7, 8.
Fig. 2. Gamma spectra taken with a Ge(Li) detector at different plunger distances.
ON- AND OFF-BEAM GAMMA SPECTROSCOPY ON NEUTRON DEFICIENT ISOTOPES OF Ge, Se, Br AND Kr

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1. INTRODUCTION

Until now not very much is known about the neutron deficient isotopes of Ge, Se, Br and Kr. We have tried to get some information about the level structure in this region using heavy ion compound reactions and methods of gamma spectroscopy. It is known from other regions of the nuclide chart that heavy ion reactions populate very selectively high spin states in the residual nuclei and one expects fairly clear gamma decay spectra especially for the even-even nuclei. In the rare earth region one observes mainly \((\text{HI}, x\text{n})\) reactions even for the very neutron deficient side. In the Kr region the Coulomb barrier for the emission of charged particles is less effective and complex reaction channels of the type \((\text{HI}, x\text{yn} z\text{yp})\) can be open. \(x, y\) and \(z\) can have values of 0, 1, 2 and more depending on the beam energy and the projectile-target combination. In order to get not too complicated gamma spectra we have used beam energies close to the Coulomb barrier of the compound reactions.

2. EXPERIMENTAL PROCEDURE AND RESULTS

The neutron deficient isotopes in the Kr region were produced bombarding targets of enriched nickel isotopes \(^{58,60,61,62,64}\text{Ni}\) with 42–57 MeV \(^{16}\text{O}\) and 36 MeV \(^{12}\text{C}\). The heavy ion beams were supplied by the MP and EN tandems at the Max-Planck-Institut für Kernphysik, Heidelberg. Gamma spectra were measured with two Ge(Li) detectors at \(0^\circ\) and \(90^\circ\) with respect to the
beam axis. The targets were self-supporting and about 500 g/cm² thick. Recoiling nuclei were stopped in a gold catcher placed 1 mm behind the target. Due to the Doppler shift and this target-catcher distance the spectra taken by the two Ge(Li) detectors allowed to distinguish between prompt (shifted in the 0° spectrum) and delayed gamma rays with lifetimes longer than approximately 1 nsec. In addition an off-beam spectrum of the gold catcher foil was measured in order to get the longer lived activities.

The assignment of the gamma lines to the various residual nuclei were made on the basis of their excitation functions, cross bombardments and comparison of prompt and activity spectra. The level schemes were constructed with the help of level energy systematics, gamma ray anisotropies and relativ intensities. The main reaction channels observed at beam energies close to the Coulomb barrier were (16O,2n) on 62,64Ni, (16O,pn) on 60Ni and (16O,2p) on 58Ni. In fig. 1 one can see the relative yield of gamma transitions from the reactions 60Ni(16O,2n)74Kr, 60Ni(16O,pn)74Er and 60Ni(16O,2p)74Se as a function of the beam energy.

The (16O,pn) reaction is already one order of magnitude larger than the (16O,2n) reaction. The increasing competition of the proton emission for the light Ni targets makes it quite difficult to reach any N=Z nucleus, e.g. 72Kr. Reactions involving emission of alpha particles were seen to be always of considerable strength. At higher bombarding energies the gamma spectra became very complex, since many reaction channels of about equal strength are open.

The proposed level schemes of several even-even isotopes of Kr, Se and Ge are shown in fig. 2. It is interesting to see that the light Kr isotopes have very low lying first 2⁺ states and show quasi-rotational
ground state bands. The level schemes of the Se isotopes show more vibrational structure. The new isotope $^{72}$Br was found through the reaction $^{58}$Ni($^{16}$O,pn)$^{72}$Br. It decays with a half life of approximately 1.7 min to the $2_1^+$ and $2_2^+$ states in $^{72}$Se. Recently, some ($\alpha,n\gamma$) studies were performed to obtain the levels in $^{78,80,82,84}$Kr (ref.2) and $^{72,74,76,78}$Se (ref.3). The results agree quite well with ours for $^{78}$Kr and $^{72,74}$Se.

3. DISCUSSION

The nuclei studied in the present work are of specific interest in view of the recently proposed pseudo LS coupling and pseudo SU$_3$ coupling schemes$^{5,6}$. In this theoretical framework one expects similarity between s-d shell nuclei ($^{16}$O plus nucleons) and these nuclei in the p-f shell ($^{80}$Zr minus nucleons). The conjugate of $^{74}$Kr is then $^{22}$Ne and that of $^{72}$Se is $^{24}$Ne. The ratio of the energies of the $4^+$ to the first $2^+$ state is 2.57 for $^{74}$Kr and 2.63 for $^{22}$Ne. It is 1.90 for $^{72}$Se and 1.99 for $^{24}$Ne. However, this correspondence becomes worse, when $T_z$ becomes higher.

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Fig. 1. Relative gamma ray yield as a function of the incident energy.
Fig. 2. Proposed level schemes of $^{74,76,78}$Kr, $^{72,74}$Se and $^{68}$Ge (thin lines) together with the level schemes of the known even-even Kr$^{2,3}$, Se$^{1,3,4}$ and Ge$^{4}$ isotopes (thick lines).
PROMPT GAMMA TRANSITIONS IN PRIMARY FRAGMENTS
FROM NEUTRON-INDUCED FISSION

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ABSTRACT

Prompt gamma transitions in moving fission fragments of specific masses have been investigated in thermal neutron-induced fission of $^{235}\text{U}$. The masses of coincident fission fragments were deduced from their correlated kinetic energies as determined by two Si surface-barrier detectors. Gamma spectra were measured by means of a Ge(Li) detector surrounded by an 8" $\phi$ x 9" NaI(Tl) crystal in Compton suppression mode. The high velocity of the fission fragments was used to restrict with geometric means the effective time resolution to about $10^{-9}$ sec after fission. The sign of the observed Doppler shift allowed the assignment of gamma-ray lines to single members of fragment pairs. The main objective of these experiments is a study of the properties of individual neutron-rich nuclei far off the stability line.

The gross energy tendency of 54 identified gamma-ray lines as a function of mass is in agreement with the assumption that the deexcitation is predominantly collective. In the magic regions $A = 88 - 92$ and $A = 130 - 134$ only gamma rays were observed whose energies fit well into the systematics of low lying vibrational transitions in neighbouring nuclei. In the framework of such a systematic it seems highly probable that the $2^+ \rightarrow 0^+$ transitions in the neutron-rich even-even nuclei $^{134}\text{Te}$ and $^{132}\text{Te}$ have been identified. The specific assignment of 6 other lines in these magic regions as $2^+ \rightarrow 0^+$ transitions in the possible candidates $^{88, 90, 92}\text{Kr}$, $^{94}$, $^{96}\text{Sr}$ and $^{130}\text{Sn}$ is discussed.

INTRODUCTION

There is increasing interest in studying nuclei far off the stability line which have long escaped experimental investigation because of their
very low half-life. Examining the prompt radiation emitted within a few nanoseconds after fission from specific fragments has become possible by the rapid development of semiconductor detector technology and associated electronics during the last few years. Such investigations are expected to yield e.g. information on spin states having intermediate angular momentum values between those involved in the radiative neutron capture and those excited in heavy ion induced reactions.

First studies have been performed by Bowman, Watson et al. [1,2] using the spontaneous fission of $^{252}$Cf. These studies have recently been extended to four-parameter measurements yielding extensive information on even-even nuclei [3] in mass ranges where stable deformations are expected [4] and which arise with appreciable yield in spontaneous fission [5]. It was felt desirable [6] to extend these measurements to thermal neutron-induced fission, both in order to compare the results with those obtained from $^{252}$Cf and to cover fragment mass regions where the yield in spontaneous fission is low. It has been recently demonstrated [7,8] that the various additional experimental difficulties which are inherent in experiments of this type at a reactor [9] have been successfully overcome. The present paper focuses on the discussion of lines arising from fragment masses which are scarcely covered in spontaneous fission, i.e. near the magic numbers $N = 50$ and $Z = 50$, $N = 82$.

EXPERIMENTAL PROCEDURE

A detailed description of the experiment has been given elsewhere [9,10]. Therefore only the principle of the experimental procedure shall be sketched very briefly. Fig. 1 shows a schematic representation of the experimental set-up. A carefully collimated and filtered neutron beam from the reactor FR2 struck the weightless target material $^{235}$U. The mass ratio of the correlated fission fragments was determined by the ratio of the two kinetic energies as measured by two Si surface barrier detectors. A carefully designed gamma-ray collimator restricted the solid angle in such a way that the Ge(Li) detector saw only the target and the first one to two centimeter flight path of the fission fragments travelling towards the Si surface-barrier detectors. A NaI(Tl) anti-Compton shield reduced the Compton distribution and fast fission
neutron induced lines in the gamma-ray spectra. The sign of the observed Doppler shift allowed the assignment of a specific line to one single member of a fragment pair.

The triple pulse-height data were processed in a 256 x 256 x 2048 channel matrix via the Karlsruhe Multiple Input Data Acquisition System (MIDAS). Final post neutron emission masses A were calculated offline using experimental neutron numbers in the corrections for the prompt neutrons.

Gamma-ray spectra were obtained by sorting the triple data in fragment mass intervals the centres of which differing in general by 2 amu. Four typical examples are displayed in Fig. 2.

RESULTS AND DISCUSSION

In Table I 54 gamma rays have been assigned to individual fragments. The masses were arrived at by comparing the peak intensity in adjacent mass intervals. The width (FWHM) of the mass distributions associated with single gamma transitions ranged typically between 4 and 7 amu \( \sqrt{10} \). Only such gamma transitions have been included in the table for which both Doppler shifted members were resolved enough to be identified on the basis of a congruous mass distribution. The absolute uncertainties in the mass determinations are mainly due to systematic errors in the calibration procedure and the neutron corrections.

The most probable charges \( Z_p \) were taken from the tables in \( \sqrt{11} \) starting from the original nonintegral mass values as derived from the maxima of the mass distributions. Due to the small width of the charge distributions for a given mass (\( \sim 1.5 \) charge units FWHM) the true charge should lie within \( \pm 1 \) unit of \( Z_p \) (only transitions in fragments with fairly high yield \( \geq 0.5 \% \) are probable to be resolved in the present experiment).

The gamma-ray energies are mean values of the Doppler pairs. It was strived for including not only uncertainties in the determination of the peak positions but also systematic errors in the given error bars.

For fission fragments travelling in the direction of maximum detection efficiency, i.e. approximately towards the centre of the fragment detectors,
the flight path viewed by the gamma detector was 16 mm. Therefore the experimental intensity values represent the relative numbers of quanta emitted within about 1.1 nsec by the light fragments and within about 1.7 nsec by the heavy fragments after fission.

A large number of transitions may be identified on the basis of the close agreement of energy, charge and mass assignment with those observed in fragments from spontaneous fission of $^{252}$Cf $^{[3,5,12]}$. The interpretations are supported by the fact that also in the present data the $2^+\rightarrow0^+$ assigned transitions follow the calculated independent yields (calculated using the mass yields given in $^{[13]}$ and the charge distributions of $^{[11]}$). Such a behaviour is expected on the basis of considerations involving the removal of the initial 6 - 10 units of angular momentum associated with each fragment. The calculated independent yields minus the fractions of converted transitions $^{[14]}$ are given in column 5 of Table I as "predicted intensities" (in arbitrary units; the relative variation is the most interesting feature in this connexion). The discrepancy between experimental and predicted intensities for $^{134}$Te $^{[2^+\rightarrow0^+]}$ is due to the fact that an appreciable feeding of the 1278 keV level proceeds via a 162 nsec isomeric state $^{[12]}$. In two low yield cases the predicted values are given in parenthesis where the $2^+\rightarrow0^+$ transitions were not identified among the well resolved lines.

In Fig. 3 the averaged energies of all transitions in Table I are plotted as a function of mass. The gross energy tendency is consistent with what one would expect if the observed deexcitations of the primary fragments involve predominantly collective transitions. This is an additional support for the individual interpretations given in column 6, 7 and 8 of Table I.

For the fragments A = 88 - 92 and A = 130 - 134 only gamma quantum have been observed which fit well into the systematics of known transitions between the lowest vibrational states in neighbouring even-even nuclei. It is reasonable to try first an interpretation of the observed gamma rays in the framework of ground state band transitions in even-even nuclei since they are expected to be fed very strongly, by analogy with heavy ion induced reactions $^{[15]}$ (removal of high initial angular momentum). In odd nuclei low lying single particle levels may cause a distribution of the deexcitation intensity through more than one cascade. John $^{[16]}$ has shown that an extrapolation
of the first excited $2^+$ state energies of adjacent isotones ($N = 82$) to $Z = 52$ yields an energy of 1.28 MeV. This interpretation is in agreement with the assignment of the 1278 keV transition as observed in the present experiment. It seems very probable that this line must be identified with the $2^+ \rightarrow 0^+$ transition in $^{134}\text{Te}$. In the same mass interval an 1180 keV transition with comparable intensity is observed. If this arises from an even-even nucleus it must be assigned with the highest probability to $^{134}\text{Te}$, too.

In Table II the excitation energies of the first excited $2^+$ states $^{17,18,19}$ in the even isotopes of Sn, Te, Xe, Ba and Xe are arranged in such way that they clearly exhibit their systematic behaviour as a function of $Z$ and $N$. A smooth extrapolation of these values both along $Z$ and $N$ yields in both cases the same energy of about 960 keV (Fig. 4). This extrapolation seems to be fairly reliable since in this mass range the $2^+$ excitation energies show quite generally a smooth variation $^{19}$ which is only interrupted by the energies of the $Z = 50$ and $N = 82$ nuclei. It appears conclusive that the observed 965 keV line arises from the $2^+ \rightarrow 0^+$ transition in $^{132}\text{Te}$. Using this assignment and the systematics in Table II and Fig. 4 it is tempting to interpret the 1222 keV line as deexcitation of the lowest $2^+$ state in $^{130}\text{Sn}$. The certainty of this interpretation suffers from the fact that a smooth extrapolation to $N = 82$ is not possible because of the steep rise in energy usually associated with this neutron number; secondly in the framework of the experimental uncertainties (Table I) an assignment to $^{132}\text{Sn}$ cannot be excluded. However, this later possibility seems unlikely due to the observation $^{19}$ that the first excited $2^+$ states of $N = 82$ nuclei have always higher energies than those of adjacent even isotones.

The discussion of the interpretation of some lines associated with masses between 88 and 96 will be accomplished using the systematic summary in Table III. The energies do not show a similarly smooth behaviour as in the mass range discussed above. Clearly visible is the influence of the proton number $Z = 40$. This excludes the possibility of predicting the $2^+$ energies in neighbouring nuclei by an extrapolation. However, some general systematic trends may be used to discuss the specific assignments of the observed transitions at 706 keV, 813 keV, 834 keV, 869 keV and 956 keV. Again it is assumed, because of their high relative intensity per fragment, that they arise from even-even nuclei.
It is obvious that the energies of these lines fit fairly well to the assembly of known $2^+ \rightarrow 0^+$ transitions in this mass range. Within the experimental uncertainties the 706 keV and the 956 keV line might be assigned to $^{90}$Kr or $^{92}$Kr, respectively. Since their intensities seem to be too different for a cascade fed by a high initial spin state they might belong to two different nuclei. This argument alone is too weak due to the big uncertainties in the intensities. However, it is supported by the fact that the experimental and predicted intensities agree reasonably well for an interpretation as given in Table I. The preferable assignment of the 706 keV transition to $^{90}$Kr and the 956 keV line to $^{92}$Kr is further supported by the general behaviour of the first excited $2^+$ states which have been found systematically lower in $Z = 36$ than in $Z = 40$ nuclei. This order is also in agreement with the observation that e.g. for $Z = 40$ and $Z = 42$ the $2^+$ energies of the $N = 56$ nuclei just out of those of the adjacent isotopes.

Similar arguments hold for a preferable assignment of the 869 keV transition to $^{88}$Kr (not $^{83}$Se). A rise in the $2^+$ energies from $N = 54$ to $N = 52$ appears more reasonable than that from $Z = 36$ to $Z = 34$.

The experimental uncertainties admit for the 813 keV and the 834 keV lines an assignment to $^{94}$Sr and/or $^{96}$Sr. Their intensities coincide within the error bars. This behaviour would fit well into the picture of a ground state vibrational cascade fed by an initial high angular momentum state. On the basis of the systematic compilation in Table III the assignment to $^{96}$Sr appears preferable since both energy values seem to be too low for a $2^+$ state energy in $^{94}$Sr. On the other hand the predicted independent yields of $^{90}$Kr, $^{94}$Sr and $^{96}$Sr behave as 55 : 59 : 25 preferring an assignment to $^{94}$Sr. Therefore no preferable assignment can be given without further experimental information. Interpretations given in parenthesis (Table I) are preliminary.

It should be mentioned that all nuclei with calculated yields higher than 1.2 % appear to be observed in the present investigations.

The $(\nu,p)$ and beta-decay data in Table I have been taken from ref. $^{20}$ and $^{21}$. 
CONCLUSIONS

The experiment described has demonstrated that despite the additional experimental difficulties the multi-parameter measurements of the prompt radiations from primary fragments can be successfully extended to neutron-induced fission. This makes it possible to investigate also mass regions where the yield in spontaneous fission is low. It appears especially promising to use $^{233}$U in such investigations since the mass regions covered in this neutron-induced fission partly overlap those reached in $(t, p)$ reactions. This would allow a comparison of the data and a substantiation of the mass and charge assignments.

The data do not point to any distinct differences in the excitation mechanism of the primary fragments from spontaneous and from neutron-induced fission. However, this can be only stated for the lowest lying levels which are the preferably accessible up till now because of intensity limitations. Possible differences in the deexcitation of the same nuclei excited one time in spontaneous fission another time in neutron-induced fission are expected to show up particularly in the first steps of deexcitation. Therefore further experiments with a higher sensibility for low intensity transitions are highly desirable.

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Fig. 1. Schematic representation of the experimental set-up.
Fig. 2. Prompt gamma-ray spectra for the fragment mass ranges $A_L = 88 - 90$ and $A_H = 143 - 145$ (upper figure), and $A_L = 102 - 104$ and $A_H = 130 - 132$ (lower figure). The spectra demonstrate the dependence of gamma-ray energy on the velocity and direction of the fragment motion. Each figure represents two cases:

(a) Light fragments moving towards the gamma-ray detector, and
(b) heavy fragments moving towards the gamma-ray detector.

The letters L and H indicate some assignments to the light and heavy fragments, respectively.
Fig. 3. Averaged energies of the observed transitions (see Table I) as a function of mass $A$. The length of the horizontal bars gives the magnitude of the averaging interval.
Fig. 4. Systematic display of the first excited $2^+$ states in even-even nuclei near the magic numbers $Z = 50$, $N = 82$. The broken lines indicate levels derived from the present data.
<table>
<thead>
<tr>
<th>Fragment Mass</th>
<th>Most Probable Charge</th>
<th>Gamma-Ray Energy (MeV)</th>
<th>Relative Experimental Intensity (in units of 10^14 per element)</th>
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<td>Zp</td>
<td>Bv (keV)</td>
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<td>(566 ± 3)</td>
<td>(589, 40)2</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>(566 ± 3)</td>
<td>(589, 40)2</td>
<td></td>
</tr>
<tr>
<td>130</td>
<td>1</td>
<td>51</td>
<td>1280 ± 5</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>132</td>
<td>1</td>
<td>51</td>
<td>905 ± 5</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>134</td>
<td>1</td>
<td>51</td>
<td>1180 ± 5</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>135</td>
<td>1</td>
<td>51</td>
<td>425 ± 4</td>
<td>(589, 40)2</td>
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<td>136</td>
<td>1</td>
<td>51</td>
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<td>(589, 40)2</td>
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<td>139</td>
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<td>54</td>
<td>372 ± 3</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>482 ± 3</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>482 ± 3</td>
<td>(589, 40)2</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>723 ± 4</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>723 ± 4</td>
<td>(589, 40)2</td>
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<tr>
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<td>1</td>
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<td>283 ± 2</td>
<td>(589, 40)2</td>
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<td>(589, 40)2</td>
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<td>55</td>
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<td>55</td>
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<td>660 ± 4</td>
<td>(589, 40)2</td>
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<td>56</td>
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<td>1</td>
<td>56</td>
<td>115 ± 2</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>145</td>
<td>1</td>
<td>56</td>
<td>115 ± 2</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>156 ± 2</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>156 ± 2</td>
<td>(589, 40)2</td>
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<tr>
<td></td>
<td></td>
<td>156 ± 2</td>
<td>(589, 40)2</td>
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<tr>
<td>146</td>
<td>1</td>
<td>56</td>
<td>103 ± 2</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>146</td>
<td>1</td>
<td>56</td>
<td>103 ± 2</td>
<td>(589, 40)2</td>
</tr>
<tr>
<td>147</td>
<td>1</td>
<td>56</td>
<td>598 ± 4</td>
<td>(589, 40)2</td>
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<tr>
<td>147</td>
<td>1</td>
<td>56</td>
<td>298 ± 2</td>
<td>(589, 40)2</td>
</tr>
</tbody>
</table>

*Derived from the tables given in Ref. / 11 /.

The high energy Doppler shifted member is partly screened by another line.

The spectra permit the alternative interpretation: A = 130 ± 1, Zp = 50, Bv = 585 ± 4 keV, Ref. Ins. = 13 ± 7.
### TABLE II. Energies of the first excited $2^+$ states in even-even nuclei near the magic numbers $Z = 50$ and $N = 82$. The energies for $^{132}\text{Te}$, $^{134}\text{Te}$ and $^{130}\text{Sn}$ are based upon transitions observed in the present experiment.

<table>
<thead>
<tr>
<th>N</th>
<th>Sn</th>
<th>Te</th>
<th>Xe</th>
<th>Ba</th>
<th>Ce</th>
</tr>
</thead>
<tbody>
<tr>
<td>72</td>
<td>1140</td>
<td>603</td>
<td>386</td>
<td>279</td>
<td></td>
</tr>
<tr>
<td>74</td>
<td>1131</td>
<td>667</td>
<td>441</td>
<td>356</td>
<td></td>
</tr>
<tr>
<td>76</td>
<td>1164</td>
<td>743</td>
<td>538</td>
<td>464</td>
<td>410</td>
</tr>
<tr>
<td>78</td>
<td></td>
<td>840</td>
<td>668</td>
<td>605</td>
<td></td>
</tr>
<tr>
<td>80</td>
<td>(1222)</td>
<td>965</td>
<td>850</td>
<td>818</td>
<td>790</td>
</tr>
<tr>
<td>82</td>
<td></td>
<td>1278</td>
<td>1320</td>
<td>1426</td>
<td>1596</td>
</tr>
<tr>
<td>50</td>
<td></td>
<td>52</td>
<td>54</td>
<td>56</td>
<td>58</td>
</tr>
</tbody>
</table>

### TABLE III. Energies of the first excited $2^+$ states near the closed neutron shell $N = 50$. The values for $^{88,90,92}\text{Kr}$ are preliminary assignments on the basis of the present experiment.

<table>
<thead>
<tr>
<th>N</th>
<th>Kr</th>
<th>Sr</th>
<th>Zr</th>
<th>Mo</th>
<th>Ru</th>
</tr>
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<tbody>
<tr>
<td>50</td>
<td>(1560)</td>
<td>1836</td>
<td>2180</td>
<td>1540</td>
<td></td>
</tr>
<tr>
<td>52</td>
<td>(869)</td>
<td>840</td>
<td>934</td>
<td>871</td>
<td>833</td>
</tr>
<tr>
<td>54</td>
<td>(706)</td>
<td>920</td>
<td>778</td>
<td>660</td>
<td></td>
</tr>
<tr>
<td>56</td>
<td>(956)</td>
<td>1730</td>
<td>787</td>
<td>540</td>
<td></td>
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<tr>
<td>58</td>
<td></td>
<td></td>
<td>536</td>
<td>475</td>
<td></td>
</tr>
<tr>
<td>36</td>
<td></td>
<td>38</td>
<td>40</td>
<td>42</td>
<td>44</td>
</tr>
</tbody>
</table>
STUDY OF SOME NEUTRON-DEFICIENT NUCLIDES WITH HIGH RESOLUTION X-RAY DETECTOR AND MASS SEPARATION*)

Y. Y. Chu

Chemistry Department, Brookhaven National Laboratory
Upton, L. I., New York 11973, U.S.A.

Nuclides several neutrons removed from the line of β-stability on the neutron-deficient side were generally produced by high energy spallation and fission. The complexity of the product mixture and the decrease in half-life as the product nuclides were moving away from the stability valley made the identifications difficult, and at times, ambiguous. Many of these identifications were based on the decay relationship with their respective descendants in the presence of other radioactive nuclides of the same element. Some were characterized only by positron radiations. Very few were confirmed by isotopically pure samples and not all of them were assayed with high-resolution γ spectroscopy.

The recent advent of the on-line mass separation projects extended the study of nuclides far from stability valley considerably. However, this technique is so far limited mainly to elements which could be easily differentiated from other elements produced in the same reaction in situ. Other elements from the same family of the Periodic Table are usually not interfering in the mass region of interest. For elements with low volatility or with similar chemical and physical properties and are neighboring to each other (e.g., the rare-earths), it is

*) Work performed under the auspices of the U.S. Atomic Energy Commission.
still not possible to study nuclides far from $\beta$-stability with the present state of art of the on-line techniques.

There are at least several dozens of neutron-deficient nuclides with half-life between 1 min and 1 hr in the rare earths region alone as shown in Figure 1. Among them quite a few have questionable assignments. These nuclides can be produced not only by high energy reactions, but more favorably by simple nuclear reactions, such as ($^3$He,Xn). For products with $4 \leq X \leq 7$, they are easily accessible by using helium 3 ions with energies between 30 and 60 MeV. Nuclides produced in this manner are all neutron-deficient and the composition of the product mixture is much simpler than that of high energy spallation or fission. By using a thin window Ge(Li) x-ray detector with sufficient resolving power to resolve K x-rays of adjacent elements, it is possible to differentiate products with highest atomic number (resulting from $^3$He, Xn reactions) from products with lower atomic numbers. Therefore, direct assaying the K x-rays of the irradiated sample could be used for chemical identification for short-lived nuclides. A least-squares computer program $^{1}$ can be used on the decay of a particular K x-ray to separate the various components of that element. The study of excitation functions of these products in the energy range available helped to ascertain the mass assignments and the characterization of the short-lived nuclides. For nuclides with half-life in the order of tens of minutes or longer, mass separations were employed to provide isotopically pure samples such that these nuclides could be studied unambiguously.

Several neutron-deficient nuclides of hafnium and heavy rareearths were produced and studied with the above-mentioned method.
Hafnium-169 was produced by $^{170}$Yb $\left(^{3}$He,$4n\right)$ $^{169}$Hf reaction. The previously reported 1.5 h activity$^3$ could not be found in mass separated $^{169}$Hf sample. Instead it was identified by the Lu K x-rays and the study of the excitation function of the production of $^{169}$Hf. Its half-life was found to be $3.26 \pm 0.05$ min. Several intense $\gamma$-rays were observed and a partial decay scheme of $^{169}$Hf could be proposed as shown in Figure 2.

The decay of $^{168}$Lu has been studied by a number of investigators.$^3$-$^9$ The level scheme in $^{168}$Yb has been studied by measuring the energy spectra of the exiting particles from $(d,d')^{10}$ and $(p,2n)^{11}$ reactions. The ground state rotational band, the $\gamma$- and $\beta$-vibrational bands and even higher order band were identified from these results. There was some question,$^6$ however, concerning the previously observed 2 h isomer.$^5$ This work was undertaken to clarify this point and to further study the decay of $^{168}$Lu. $^{168}$Lu was produced by $^{169}$Tm($^3$He,$4n$)$^{168}$Lu reaction. The 2 h isomer could not be found in the mass separated $^{168}$Lu sample. The $\gamma$-rays observed in $^{168}$Lu by direct assaying the irradiated sample with a Ge(Li) detector were essentially the same as those belonging to the known short-lived isomer.$^8$,$^9$

However, these $\gamma$-rays fell into two categories with half-lives 5.15 min and 6.05 min (except the 88- and 199-keV $\gamma$-rays) instead of the well known half-life of 7 min. The relative intensities of these $\gamma$ rays varied with different means of production of $^{168}$Lu as shown in Table 1. It is possible to conclude from these observations that the 5.15 min activity should be assigned to the high spin isomer and the 6.05 min activity to the low spin isomer. The orbitals available for the odd proton and odd neutron from the Nilsson's formulation$^{12}$ and the decay of


168\textsuperscript{Lu} isomers to 168\textsuperscript{Yb} and some of the intense $\gamma$-rays are given in Figure 3.

### Table 1

Relative intensities of some $\gamma$ rays from the decay of 168\textsuperscript{Lu}

<table>
<thead>
<tr>
<th>$E_\gamma$ (keV)</th>
<th>$t_{1/2}$ (min)</th>
<th>Transition</th>
<th>Mode of production of 169\textsuperscript{Tm}</th>
<th>Decay of 168\textsuperscript{Hf}</th>
</tr>
</thead>
<tbody>
<tr>
<td>88</td>
<td>6.0 m</td>
<td>2$^+ \rightarrow 0^+$</td>
<td>1.00</td>
<td></td>
</tr>
<tr>
<td>199</td>
<td>5.6 m</td>
<td>4$^+ \rightarrow 2^+$</td>
<td>3.32</td>
<td></td>
</tr>
<tr>
<td>299</td>
<td>5.15m</td>
<td>6$^+ \rightarrow 4^+$</td>
<td>0.92</td>
<td></td>
</tr>
<tr>
<td>384</td>
<td>5.15m</td>
<td>8$^+ \rightarrow 6^+$</td>
<td>0.10</td>
<td></td>
</tr>
<tr>
<td>716</td>
<td>5.15m</td>
<td>5$^+ \rightarrow 4^+$</td>
<td>0.12</td>
<td></td>
</tr>
<tr>
<td>884</td>
<td>6.05m</td>
<td>4$^+ \rightarrow 4^+$</td>
<td>0.79</td>
<td></td>
</tr>
<tr>
<td>896</td>
<td>6.05m</td>
<td>2$^+ \rightarrow 2^+$</td>
<td>0.77</td>
<td></td>
</tr>
<tr>
<td>979</td>
<td>6.05m</td>
<td>3$^+ \rightarrow 2^+$</td>
<td>1.23</td>
<td></td>
</tr>
<tr>
<td>984</td>
<td>6.05m</td>
<td>2$^+ \rightarrow 0^+$</td>
<td>0.60</td>
<td></td>
</tr>
<tr>
<td>1015</td>
<td>5.15m</td>
<td>5$^+ \rightarrow 4^+$</td>
<td>0.48</td>
<td></td>
</tr>
</tbody>
</table>

162\textsuperscript{Yb}: Ytterbium-162 was produced by 164\textsuperscript{Er}(\textsuperscript{3}{He},5n)162\textsuperscript{Yb} reaction. Mass separation was performed to provide isotopically pure 162\textsuperscript{Yb} sample. The Tm K x-rays and two $\gamma$-rays (118- and 162-keV) were observed. The half-life of 162\textsuperscript{Yb} was found to be 18.87 ± 0.14 min.

162\textsuperscript{Tm}: Thulium-162 was produced by 165\textsuperscript{Ho}(\textsuperscript{3}{He},6n)162\textsuperscript{Tm} reaction. The previously reported 77 min isomer\textsuperscript{13} was not found from an isotopically pure 162\textsuperscript{Tm} sample. The Er K x-rays and two $\gamma$-rays (102- and 226-keV) were observed. These $\gamma$-rays correspond to the transitions between 2$^+ \rightarrow 0^+$ and 4$^+ \rightarrow 2^+$ of the ground state rotational band in 162\textsuperscript{Er}\textsuperscript{14}. The half-life of 162\textsuperscript{Tm} was found to be 22.23 ± 0.50 min.

* * *
REFERENCES


2. The results on $^{169}$Hf have been published since the submission of the abstract. Y. Y. Chu and J. Reednick, Phys. Rev. C2, 310 (1970).


\[ \frac{3}{2}^-[523] \]

\[ \frac{5}{2}^-[523] \]

\[ \frac{3}{2}^+[404] \]

\[ \frac{5}{2}^+[404] \]

\[ ^{169}\text{Hf (3.26 min)} \]

\[ \beta^+ \]

\[ \text{EC} \]

\[ 491 \text{ KeV} \]

\[ 123 \text{ KeV} \]

\[ 0 \]

\[ ^{169}\text{Lu (34 hr)} \]

Figure 2.
\[ 0^+ \xrightarrow{\text{26 min}} \stackrel{^{168}}{\text{Hf}} \xrightarrow{\varepsilon, \beta^+} \]

71\textsuperscript{st} PROTON: \[ \frac{7}{2} + [404]; \frac{1}{2} - [541]; \frac{9}{2} [514]; \frac{5}{2} + [402] \]

97\textsuperscript{th} NEUTRON: \[ \frac{5}{2} - [523]; \frac{11}{2} - [505]; \frac{5}{2} + [642] \]

\[ \xrightarrow{\varepsilon, \beta^+} \]

\[ \xrightarrow{\text{6.05 min}} \stackrel{^{168}}{\text{Lu}} \]

\[ \xrightarrow{\varepsilon, \beta^+} \]

\[ \xrightarrow{\text{5.15 min}} \]

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\[ \gamma_{970} \]

\[ \gamma_{984} \]

\[ \gamma_{1067} \]

\[ \gamma_{896} \]

\[ \gamma_{984} \]

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THE DECAY OF SOME NUCLEI FAR FROM THE STABILITY LINE
STUDIED WITH STANDARD SPECTROSCOPIC TECHNIQUES

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P. Koldewijn, J. Konijn and B.J. Meijer
Institute for Nuclear Physics Research,
Amsterdam, Holland.

1. INTRODUCTION

Some nuclides in regions of deformation and far from
the region of beta stability were studied with standard
beta- and gamma-ray spectroscopic techniques for some
cases with half-lives not shorter than a few minutes. The
activities are produced by the \( \alpha \) reactions
\( (E^{(3}\text{He}) \leq 78 \text{ MeV, } x=2,3, \ldots 8) \) and the \( p \) reactions
\( (E(p) \leq 55 \text{ MeV, } x=1,2, \ldots 5) \) on targets of natural and
enriched isotopes. The activities produced this way are
limited to a very small region of \( A \) and \( Z \) and can in most
cases be assigned unambiguously even without carrying out
a chemical separation.

In addition to gamma ray and electron spectra, coinci-
cidences were studied with Ge(Li)–Ge(Li) and Ge(Li)–Si(Li)
detectors in combination with a 4k PDP-8, used as a two-
dimensional 64000 channel analyzer\(^1\). Spectra of con-
version electrons and positons were studied with Si(Li)
detectors. Some features of the isotopes being studied are
listed in the following section.
2. DECAY PROPERTIES

$^{158\text{Tm}}$: $T_\frac{1}{2} = 4.3 \pm 0.2$ min; mode of production:

$^{162\text{Er(p,5n)}}$, measured $\gamma$-ray spectrum, $\gamma$-$\gamma$ coincidences;
$\gamma$-rays: 192.6 keV (100), 335.3 keV (28), 1150.5 keV (19).
The existence of this isotope was reported independently
by Neiman and Ward$^2$).

$^{159\text{Tm}}$: $T_\frac{1}{2} = 12 \pm 1$ min; mode of production:

$^{162\text{Er(p,4n)}}$, measured $\gamma$-ray spectrum; $\gamma$-rays 84.8 keV
(57), 144.1 keV (31), 220.2 keV (96), 272.0 keV (57),
289.6 keV (100), 348.3 keV (64). The existence of this
isotope was found independently by Gromov et al.$^3$ who
did not report $\gamma$-rays.

$^{160\text{Tm}}$: $T_\frac{1}{2} = 9.2 \pm 0.4$ min; mode of production:

$^{162\text{Er(p,5n)}}$, $^{164\text{Er(p,5n)}}$, measured $\gamma$-ray spectrum,
$\gamma$-$\gamma$ coincidences; $\gamma$-rays: 126.1 keV (100), 264.2 keV (30),
727.9 keV (35), 853.5 keV (26), 860.3 keV (26). The exist-
ence of this isotope was reported independently by Neiman
and Ward$^2$).

$^{164m\text{Tm}}$: $T_\frac{1}{2} = 5.1 \pm 0.1$ min; mode of production:

$^{166\text{Er(p,3n)}}$, measured $\gamma$-ray spectrum, $\gamma$-$\gamma$ coincidences,
delayed $\gamma$-$\gamma$ coincidences, conversion electrons; $\gamma$-rays:
91.4 keV E2 (43), 139.4 keV E2 (27), 208.1 keV E2 (152),
240.5 keV E1+M2 (78), 314.9 keV E2 (100), 410.2 keV E2
(15), 547.0 keV E1 (46), 820.7 keV (14), 897.9 keV (44),
1049.8 keV (16), 1231.2 keV (42), 1364.6 keV (43) and 17 transitions of lower intensities.

The decay scheme deduced from the present measurement for this thus far unreported isotope is shown in fig. 1.

A delayed coincidence measurement with a Ge(Li)-NaI(Tl) setup was performed. The half-life of the 1985 keV level was determined as $22.7 \pm 2.0$ ns. A spectrum of the γ-rays measured in delayed coincidence with the KX radiation is shown in fig. 2.

$^{158}$Yb: $T_1 = 4.6 \pm 0.5$ min; mode of production:

$^{162}$Er($^3$He,7n), measured γ-ray spectrum; γ-rays: 173.9 keV (100), 215.7 keV (47). Our experiment does not completely
Fig. 2

γ-rays in the $^{164m}$Tm decay measured in delayed coincidence with the KX radiation. A singles spectrum of this decay is shown in the upper part of the figure.

exclude a $^{159}$Yb assignment. The existence of this isotope was reported independently by Neiman and Ward$^2$).

$^{160}$Yb: $T_{1/2} = 4.1 \pm 0.2$ min; mode of production: $^{164}$Er($^3$He,7n), measured γ-ray spectrum; γ-rays: 78.3 keV (100), 600.0 keV (91), 631.7 keV (37). Our experiment does not completely exclude a $^{161}$Yb assignment. The existence of this isotope was reported independently by Neiman and Ward$^2$).
$^{163}$Yb: $T_1 = 11.4 \pm 0.5$ min; mode of production: $^{164}$Er($^3$He,4n), $^{166}$Er($^3$He,6n), measured γ-ray spectrum; γ-rays: 64.0 keV (70), 123.5 keV (21), 130.9 keV (18), 161.6 keV (12), 326.0 keV (18), 687.4 keV (14), 860.3 keV (100), 1746.6 keV (15), 1907.6 keV (15) and 32 transitions of lower intensities. The existence of this isotope was found independently by Gromov et al. $^3$ who did not report γ-rays.

$^{170}$Ta: $T_1 = 6.3 \pm 0.4$ min; mode of production: $^{175}$Lu($^3$He,6n), measured γ-ray spectrum; γ-rays: 100.8 keV (100), 221.2 keV (156). Not reported previously.

$^{171}$Ta and $^{171m}$Ta: mode of production: $^{175}$Lu($^3$He,7n), measured γ-ray spectra. $T_1 = 6.3 \pm 0.4$ min; γ-rays: 59.3 keV (57), 87.9 keV (57), 111.9 keV (51), 198.8 keV (100) and $T_1 = 2.0 \pm 0.5$ min; γ-rays: 365.4 keV (100). The existence of both activities was not reported previously. The γ-ray spectrum of these Ta isotopes is shown in fig. 3.

$^{172}$W: $T_1 = 6.7 \pm 0.5$ min; mode of production: $^{176}$Hf($^3$He,7n), measured γ-ray spectrum, γ-γ coincidences, γ-γ delayed coincidences, conversion electrons; γ-rays: 35.9 keV (53), 39.7 keV (11), 130.4 keV (25), 175.0 keV (18), 457.6 keV (100), 623.6 keV (25) and 25 transitions of lower intensities. The existence is deduced from these
Fig. 3

γ-rays in the decay of $^{170}$Ta, $^{171}$Ta and $^{171m}$Ta observed in a source produced by the $^{175}$Lu($^3$He,7n) reaction.

measurements of levels at 130.4 keV, 166.3 keV and 623.8 keV. A half-life of 180±10 ns was found for the 166.3 keV level. The existence of this isotope was found earlier by Arlt et al.4) who did not report γ-rays.

$^{176}$Os: $T_1 = 3.6\pm0.5$ min; mode of production: $^{180}$W($^3$He,7n), measured γ-ray spectrum; γ-rays 81.5 keV (36), 775.8 keV (98), 857.2 keV (69), 1209.2 keV (71), 1290.9 keV (100). The existence of this isotope was reported recently also by Arlt et al.5).

3. CONCLUDING REMARKS

It is possible to deduce the main decay properties of nuclei far from the stability line with half-lives in the minutes region with standard spectroscopic techniques. A
serious difficulty is, however, encountered in measurements of isotopes with high decay energies. Here a substantial part of the decay intensity can be contained in a large number of very weak transitions hidden in the Compton background of the stronger lines. In a recent study of the $^{178}$Re decay (Q=4.7 MeV) a serious discrepancy from the theoretical values was found for the $K/\beta^+$ ratio for the beta decay to the lower states in $^{178}$W; this deviation is most likely caused by the above effect.
This work is part of the research program of the Institute for Nuclear Physics Research (I.K.O.), made possible by financial support from the Foundation for Fundamental Research on Matter (F.O.M.) and the Netherlands Organization for the Advancement of Pure Research (Z.W.O.).

Note: At the Conference our attention was drawn to the extensive measurements of Rezanka et al. (these proceedings) of the decay of 25 min. $^{171}$Ta.
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ROTATIONAL BANDS IN $^{171,173,175}$Hf FROM (PARTICLE, xn) REACTIONS

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The structure of these three odd-neutron hafnium isotopes $^{171,173,175}$Hf was studied by means of the gamma-ray spectroscopy of (particle,xn) reactions on self-supporting metallic monoisotopic targets. The employed targets and

<table>
<thead>
<tr>
<th>Product</th>
<th>Reaction</th>
<th>Optimum energy in MeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{175}$Hf</td>
<td>$^{174}$Yb($\alpha$,3n)</td>
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<tr>
<td>$^{173}$Hf</td>
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</tr>
<tr>
<td>$^{170}$Yb($\alpha$,3n)</td>
<td>35</td>
<td></td>
</tr>
</tbody>
</table>

beams are summarized in Table 1. The gamma-ray spectra, their relative excitation functions, angular distributions time correlations regarding the beam bursts and, in case of $^{171}$Hf, few gamma-gamma coincidence relationships were determined. Based on these data, on the tabulated radio-

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activity information and on the unpublished results on the 171,173Ta decays obtained at Yale, we identified the levels of the $1/2^-[521]$, $5/2^-[512]$ and $7/2^+[633]$ rotational bands up to high spins (see Figs. 1-3). The very important information, communicated to us most kindly prior to publication by Drs. B. Elbek and P. Kleinheinz of the Niels Bohr Institute, was the location of the $13/2^+$ states in $^{173,175}$Hf, in both nuclei at 440 keV. Our assignments agree well with these. The rotational parameters of the former 2 bands are summarized in Table 2. The even parity bands in all these 3 nuclei show large alternating energy shift terms (see Fig. 4) the origin of which has to be found in the Coriolis interaction which is mixing the even parity states originating mainly in the $i_{13/2}$ shell model state. Using a simple model and basic nuclear constants from literature we were able to reproduce the even parity bands to no
worse than 8 keV for every observed levels. Essentially, only 3 parameters, viz. the attenuation factors for Coriolis matrix elements, had to be varied to obtain this fit (see Table 3).

Table 3
Comparison of experimental and calculated levels for even-parity bands

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<th>171\text{Hf}</th>
<th>173\text{Hf}</th>
<th>175\text{Hf}</th>
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<td>Energy(keV)</td>
<td>Energy(keV)</td>
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<tr>
<td>9/2</td>
<td>61.9</td>
<td>59.7</td>
<td>+2.2</td>
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</tbody>
</table>
Fig. 1
Fig. 2
Fig. 3

\[ ^{175}\text{Hf}_{103} \]
$\frac{E_R - E_{2+}}{\hbar^2 \Omega}$

$\frac{1}{2}^+[633]$
POSSIBLE "NUCLEAR CO-EXISTENCE" IN ODD-MASS Ir NUCLEI

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Abstract:

The decays of $^{189}$Pt $\rightarrow$ $^{189}$Ir, $^{191}$Pt $\rightarrow$ $^{191}$Ir, $^{191}$Os $\rightarrow$ $^{191}$Ir and $^{193}$Os $\rightarrow$ $^{193}$Ir have been studied using solid state detectors, magnetic spectrometers and prompt and delayed coincidences. New levels with the following spin-parity assignments and half-lives were found:

$^{193}$Ir: $299.45 \pm 0.09$ keV, $I^m = 7/2^-$, $T_{1/2} = 0.19 \pm 0.03$ ns;
$598.26 \pm 0.15$ keV, $I^m = 3/2^-$; $740.46 \pm 0.15$ keV, $I^m = 5/2^-$. 

$^{191}$Ir: $391.14 \pm 0.05$ keV, $I^m = 7/2^-$, $T_{1/2} = 0.24 \pm 0.02$ ns;
$659.06 \pm 0.08$ keV, $I^m = 3/2^-$, $T_{1/2} < 0.12$ ns.

$^{189}$Ir: $615.70 \pm 0.08$ keV, $I^m = 7/2^-$, $T_{1/2} = 0.18 \pm 0.02$ ns;
$1184.54 \pm 0.11$ keV, $I^m = 5/2^-$, $T_{1/2} < 0.08$ ns.

In all three Ir nuclei the values of $B(E2; 7/2^- \rightarrow 11/2^-)$ are close to the values of $B(E2; 2^+ \rightarrow 0^+)$ in neighbouring, essentially spherical, even nuclei. Energies and transition probabilities indicate that the $7/2^-$ states are analogous to the $7/2^-$ states found in the odd mass Au nuclei, i.e. they are associated with an essentially spherical shape of the nucleus. This is in contrast to the low-lying positive parity states, which in all three nuclei can be arranged into rotational bands, which may be assigned as the Nilsson states $3/2^+$ [402] and $1/2^+$ [400]. From the $B(E2)$ values of the intra-band transitions we obtain a deformation $\delta = 0.25$ for $^{189}$Ir and $\delta = 0.12$ in $^{191}$Ir and $^{193}$Ir.

1. Introduction

It is of interest to study properties of nuclei in mass regions where the nuclear equilibrium shape changes between spherical and deformed.
Mass-separator facilities in connection with accelerators like e.g. the ISOLDE project \(^1\) offer possibilities to study such regions more systematically than has been done earlier.

We have studied levels in \(^{189}\text{Ir}\) \(^{191}\text{Ir}\) and \(^{193}\text{Ir}\) with two principal goals in mind:

1. To measure the transition probabilities between the low-lying positive parity states in these nuclei and compare the results with predictions by models of both deformed and spherical nuclei.

2. To find out the origin of a number of strong and medium-strong E2 and M1 transitions, which in each nucleus was found not to combine with other transitions and levels.

To do this we have studied the decays of \(^{193}\text{Os} \rightarrow ^{193}\text{Ir}\), \(^{191}\text{Os} \rightarrow ^{191}\text{Ir}\), \(^{191}\text{Pt} \rightarrow ^{191}\text{Ir}\) and \(^{189}\text{Pt} \rightarrow ^{189}\text{Ir}\) using Ge (Li) and Si (Li) detectors in single and coincidence modes, magnetic and solid-state electron spectrometers and delayed coincidence techniques. Experimental details have been published or are going to be published elsewhere \(^2\text{-}^7\).

2. Transition probabilities between low-lying positive parity states

A preliminary level scheme for \(^{189}\text{Ir}\) is shown in fig. 1. As is well known \(^8\) the energies of the six low-lying positive states can be grouped into \(K = \frac{3}{2}\) and \(K = \frac{1}{2}\) bands according to the figure. From the Nilsson diagram the bands are naturally interpreted as the \(\frac{3}{2}^+ [402]\) and \(\frac{1}{2}^+ [400]\) states respectively. The low energy level structure is very similar for \(^{191}\text{Ir}\) \(^3\) and \(^{193}\text{Ir}\) \(^4\).

The values for the reduced E2 transition probabilities as obtained from measurements of delayed coincidence and L subshell ratios are shown in fig. 2. We observe the following features:

1. The transitions corresponding to the intraband transitions \(\frac{5}{2}^+ \leftrightarrow \frac{3}{2}^+\) and \(\frac{3}{2}^+ \leftrightarrow \frac{1}{2}^+\) show a sharp rise when the mass-number decreases from 191 to 189.

2. The \(B(E2)\)-values in \(^{189}\text{Ir}\) corresponding to intra-band transitions are on the average very much larger than the values for the inter-band
Fig. 1 Preliminary level scheme for $^{189}$Ir as populated in the decay of $^{189}$Pt. The widths of the arrows are proportional to the transition intensities, the white part denoting the internal conversion intensity.

transitions. From the former we deduce a quadrupole moment $Q = 8 \pm 1$ b and a deformation $\delta = 0.23 \pm 0.03$, which is large enough to definitely classify $^{189}$Ir as a deformed nucleus.

2. Also in $^{191}$Ir and $^{193}$Ir the B(E2) values corresponding to the intraband transitions are all definitely larger (about a factor of 3) than the values corresponding to inter-band transitions. Assuming also $^{191}$Ir and $^{193}$Ir to be deformed we obtain for $^{191}$Ir $Q = 4.0 \pm 0.8$b, $\delta = 0.13 \pm 0.03$ and for $^{193}$Ir $Q = 4.0 \pm 0.4$, $\delta = 0.13 \pm 0.01$.

4. For comparison fig. 2 also shows B(E2) values calculated from two different theoretical models:

a) The dashed lines show B(E2)-values for inter-band transitions obtained in a simple calculation assuming the $\frac{3}{2}^+$ band and the $\frac{5}{2}^+$ band to be Coriolis coupled 4). The coupling strength $E^{\ast} < \frac{1}{2} \left| j \right| \frac{3}{2} >$ was taken as 20 keV, which is the value predicted from the Nilsson model 9).

b) The dot-dashed lines show the predictions by the pairing-plus-quadrupole force model calculations by Reehal and Sorensen 10), in which the coupling between quasi-particles in a spherical potential and quadrupole vibrations up to two phonons are considered.
None of these calculations can be said to reproduce experiments remarkably well, however, considering also the branching ratios given to the right in fig. 2, we find the deformed approach to be superior of the spherical approach, thus supporting the suggestion of $^{191}$Ir and $^{193}$Ir as essentially deformed nuclei.

![Diagram](image)

**Fig. 2** Transition probabilities and branching ratios of E2 transitions in odd mass Ir nuclei. The "primed" levels are the members of the $K = \frac{1}{2}$ band. Experimental points are connected with solid lines. The dashed lines show for the B(E2) values the results of the Coriolis coupling calculation mentioned in the text. For the ratios the dashed lines indicate the Alaga values. The dot-dashed lines show the results of the pairing-plus-quadrupole force calculations by Reehal and Sorensen.

3. Negative parity states

It has long been known\textsuperscript{11} that $11/2^-$ isomeric states are populated in Ir nuclei in the decay of $^{193}$Os and $^{191,189}$Pt. In view of the low spin of the mother nuclei ($3/2^-$ for all) and the fairly low decay energies involved ($Q_\beta \sim 1.0 - 1.5$ MeV), the population is surprisingly strong, and it is natural to ask how it takes place. As was mentioned in the introduction there is in each Ir nucleus a group of transitions, which have no connection with the established levels. From conversion electron
measurements and coincidence studies we have found strong evidence for these transitions essentially forming a cascade feeding the $11/2^−$ isomeric state in each nucleus, thereby defining $7/2^−$, $5/2^−$ and $3/2^−$ levels\(^6\). An example of a detailed partial level scheme is shown in fig. 3 and a general survey of these levels is given in fig. 4. The relative order of the E2-transitions could be determined with delayed coincidences\(^6\). From these measurements we also deduced the reduced E2 transition probabilities given in fig. 5. In all these cases only a few very weak E1 transitions connect the negative parity states with the low-lying positive parity states.

Regarding the nature of these negative parity levels it is in view of the above discussion natural to look for an interpretation in terms of the deformed model. The isomeric state is then naturally assigned as $11/2^− [505]$, which should occur at low energy for both small and large deformations. For the $7/2^−$ states, however, no Nilsson orbital is available (the $7/2^− [523]$ state is expected at considerably higher energy and would have the $9/2^− [514]$ state between it and $11/2^− [505]$).
Fig. 4  Partial level schemes showing the negative parity states and the positive parity rotational bands in odd mass Ir nuclei. For comparison the corresponding levels in Au isotopes are shown.

Another possibility would be to explain the $7/2^-$ levels as the K-2 gamma vibrational state built on the $11/2^-$ [505] orbital. Inspecting fig. 5 we find however, that the value of $B(E2; 7/2^- \rightarrow 11/2^-)$ in all three Ir nuclei is almost a factor of ten larger than $B(E2; 2^+ \rightarrow 0^+)$ for the second $2^+$-state (gamma vibrational state) in the neighbouring even Os nuclei. Thus an interpretation in terms of a gamma vibration is not resonable. The only way to explain within the deformed model the low-lying $7/2^-$ state and the value of $B(E2; 7/2^- \rightarrow 11/2^-)$ would be to assume a strong enough coupling between the Nilsson orbitals branching off from the $h_{11/2}$ shell model state to reverse the level order in the rotational bands or produce inter-band $E2$ transitions with the same strength as intra-band transitions. However, no such strong mixing has been observed for the $11/2^-$ [505] neutron state $^{124}$). Instead we would like to suggest a "simpler" way to explain the data:

In fig. 4 we show for comparison also the level schemes of odd Au nuclei, which are generally accepted as being essentially spherical. Here we also observe a $7/2^-$ state about 200 keV above the $h_{11/2}$ isomer. The similarity between $B(E2; 7/2^- \rightarrow 11/2^-)$ and $B(E2; 2^+ \rightarrow 0^+)$ for the
Fig. 5  Values of $B(E2; \, 7/2^- \rightarrow 11/2^-)$ in odd mass Ir and Au nuclei. The experimental points are connected with heavy solid lines. For comparison are shown $B(E2; \, 2^+ \rightarrow 0^+)$ for the first and second $2^+$ state in the neighbouring even nuclei.

first $2^+$ state in the even neighbours, cf. fig. 5, strongly suggests the $7/2^-$ state in Au as a member of the multiplet of states obtained by coupling the $h_{11/2}$ particle to the $2^+$ excitation of the even core. Turning to Ir we see in figs. 4 and 5 that the situation for the $7/2^-$ states here is almost identical with Au regarding both energies and intensities. It is true that the values of $B(E2; \, 2^+ \rightarrow 0^+)$ for the even Os neighbours are definitely larger (a factor of two) than $B(E2; \, 7/2^- \rightarrow 11/2^-)$ in $^{191}$Ir and $^{189}$Ir, but since both $^{188}$Os and $^{190}$Os are rotational rather than vibrational nuclei, the latter $B(E2)$-values should be compared with more neutron-rich and thus less deformed Os nuclei.
Additional experimental data support the characterization of the negative parity states as essentially spherical:

1. The transition probabilities between the negative parity states and the positive parity states are generally strongly retarded 6).
2. Also the 3/2^- states found in ^{191}\text{Ir} and ^{193}\text{Ir} have to be explained as spherical states. Their predominant decay via an E2 transition to the 7/2^- state suggests the 3/2^- state as a member of the multiplet of states obtained by coupling two quadrupole vibration phonons to the h_{11/2}^+ particle. This is supported by the fact that we find B(E2; 3/2^- \rightarrow 7/2^-) > 0.26 e^2 b^2 (> 17 s.p.u.) in ^{191}\text{Ir}.

4. Concluding remarks

Experimental evidence has been found that ^{189}\text{Ir} and possible also ^{191}\text{Ir} and ^{193}\text{Ir} in the 11/2^- state and the collective excitations coupled to this state have a definitely smaller deformation than in the ground state and the low-lying positive parity states.

It has been pointed out by Soloviev 13) that the deformation in an excited state can be smaller than in the ground state if the Nilsson orbital corresponding to the excited particle rises rapidly. This is likely to be the explanation of the situation in the Ir nuclei since the 11/2^- [505] is expected to be a particle state and it rises unusually strongly with the deformation.
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IDENTIFICATION OF ISOMERS AMONG PRIMARY FISSION FRAGMENTS

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ABSTRACT

The isomeric x-rays and γ-rays emitted by fission products from the thermal neutron fission of $^{235}_{92}$U in the time window between 1 μs and 80 μs after the fission event were investigated at the gas-filled fission product separator of the research reactor FRJ-2. Isomeric transitions were detected at the fission product masses $A = (88, 93, 96 - 100$ and $132-136)$ with about 90% of the γ-lines being emitted in the regions $A = (96 - 100)$ and $A = (132 - 136)$. An identification of the isomeric nucleides was achieved in most cases. In the mass region (96 - 100) all identified or possible isomers have 58 or 59 neutrons. The identified isomers in the region $A = (132 - 136)$ are $^{134}$Te and $^{136}$Xe with $N = 82$. The isomerism is discussed.

1. INTRODUCTION

Fission fragments are highly excited nuclei which are de-excited by neutron emission, emission of γ-radiation, and β-decays. The prompt neutrons are emitted within $10^{-14}$ s after fission, β-decays occur with half-lives $> 10^{-2}$ s. Gamma-radiation is emitted promptly as well as delayed following β-decays. The overwhelming part of the prompt γ-rays has a mean half-life of ca. $10^{-11}$ s $^1,2)$. But there is a considerable fraction of the prompt γ-radiation which is emitted from isomeric transitions with half-lives up to 100 μs, as was first shown for fission products from $^{235}_{92}$U by Maienschein et al. $^3$), who found an isomeric component with a half-life of ≈ 100 ns. In the meantime several authors have investigated such isomeric transitions. Johansson $^4$) measured isomeric γ-rays of $^{252}$Cf-fission products in the time region (10 - 100) ns using NaJ-detectors. By measuring the kinetic energies of complementary fission products with semiconductor detectors he determined the masses of the γ-emitters.
He found that the isomers were concentrated into a few narrow mass regions. Popeko et al. \textsuperscript{5)} investigated with a Ge(Li)-detector the isomeric γ-rays in the time region (10 – 100) ns after the thermal neutron fission of \(^{235}\text{U}\). Using the same method as Johansson they determined the masses of the isomers. About 20 low energy γ-rays have been identified with half-lives < 50 ns. Walton and coworkers \textsuperscript{6)} looked for isomers among the fission products of the photofission of \(^{232}\text{Th}, ~^{235}\text{U}, ~^{239}\text{Pu}\) and neutron fission of \(^{235}\text{U}, ~^{239}\text{Pu}\) in times > 2\(\mu\)s. They detected the γ-rays with plastic detectors or NaJ-crystals. No mass assignment of the isomers has been made. The authors have identified about 10 γ-rays with energies up to 1330 keV and half-lives up to 5\(\frac{1}{2}\)\(\mu\)s. Guy \textsuperscript{9)} measured the isomeric γ-rays of \(^{252}\text{Cf}\)-fission products with a Ge(Li)-detector and the masses of the emitters with the technique indicated before. This measurement was extended from 3 to 2000 ns after fission. The authors determined the energies, half-lives, and intensities of about 150 γ-rays in this time region. The short-lived γ-rays are emitted by almost all masses that are produced with sufficient yields. The long-lived radiation, however, is confined to narrow mass regions which are similar to those in ref. \textsuperscript{4}).

Besides measurements of the γ-radiation there have been done experiments on x-rays emitted from isomers among the fission products due to internal conversion of the γ-rays. H.Hohmann \textsuperscript{7)} found a delayed contribution in the x-ray spectrum from \(^{236}\text{U}\) fission, whereas Glendenin et al. \textsuperscript{8)} investigated delayed x-rays from \(^{252}\text{Cf}\). For example Watson et al. \textsuperscript{10)} measured with a Si(Li)-diode the x-rays that were emitted by \(^{252}\text{Cf}\) within 93 ns. Kapoor et al. \textsuperscript{11)} investigated x-rays from fission products up to 1\(\mu\)s after the thermal neutron fission of \(^{235}\text{U}\) using a NaJ-crystal. Assuming that the production of isomers is comparable for the different fissioning nuclei a comparison of the results of ref. \textsuperscript{10)} and of ref. \textsuperscript{11)} and of measurements of the x-rays emitted within ca. \(10^{-9}\) s after fission, see e.g. refs. \textsuperscript{8,12-14)}, gives some information concerning the isomeric transitions of the fission products.

We are doing an experiment with the gas-filled fission product separator at the research reactor FRJ-2 of the Kernforschungsanlage Jülich to investigate both the γ-rays and the x-rays of isomers with half-lives from ca. 0.2 to 100\(\mu\)s produced by the thermal neutron fission of \(^{235}\text{U}\).
The interest in the $\gamma$-radiation of fission products as a radiation emitted by very neutron rich nuclei with initially high spins and deformations has been reinforced during the recent years by the prediction of a region of stable deformations for $Z < 50$ and $50 < N < 82$ \cite{4, 15-18}. Isomeric transitions of the light fission products might give some information concerning this problem.

2. EXPERIMENTAL SET-UP

2.1 Separation of isotopes

The gas-filled separator providing beams of fission products from the thermal neutron fission of $^{235}\text{U}$ has been described in some detail in refs. \cite{19 - 21}. The peculiarity of the separator is that fission products having equal nuclear charge $Z$ and equal nuclear mass $A$ experience the same deflection in the magnetic field of the separator, in spite of the fact that they emerge with different ionic charges and energies from the fission product source. The resolution of the separator is not sufficient to resolve single values of $Z$ and $A$. At each value of the magnetic rigidity $B\cdot\varphi$, a mixture of different nuclides is obtained. A given nuclide reaches the focus of the separator within a certain range of $B\cdot\varphi$, its distribution over $B\cdot\varphi$ being a Gaussian function. In fig. 1 an example of such a distribution is shown. The nuclide has been identified by one of the isomeric $\gamma$-rays that are the subject of our investigation.

The calibration of the separator has been done by measuring the intensity distributions of known $\gamma$-rays that follow the $\beta$-decays of long-lived fission products \cite{21}. The $B\cdot\varphi$-values at which these distributions had their maxima were attributed to the masses of the $\gamma$-emitters. The
long-lived calibration nucleides themselves are not produced in the fission process and the subsequent prompt neutron emission, they are β-decay daughters of primarily produced fission products. The deflection of the calibration nucleides is determined by the mean primary charge $Z_p(A_1)$ of the corresponding isobaric chain as the time needed for the separation is equal to the time-of-flight of the unslowed fission products through the separator which amounts to 1 μs. Consequently the calibration curve which is shown in ref. 20) may be applied, without further considerations, for the identification of long-lived secondary fission products which are deflected according to $Z_p(A)$.

The identification of the isomers among the primary fission products, we are interested in here, is achieved as described in the following. The distribution of an isomeric radiation over $B_1 \cdot \varphi$ is measured and the value $B_1 \varphi_{\text{max}}$ corresponding to the maximum of this distribution is determined. To $B_1 \varphi_{\text{max}}$ corresponds a mass $A^*$ according to the calibration curve.

A mean primary charge $Z_p^*$ may be attributed to $A^*$ for example from the $Z_p$-curve of ref. 22) or by interpolation between the $Z_p$-values of Wahl 23), both $A^*$ and $Z_p^*$ being usually non-integer. With these values the candidates which are possibly the emitters of the isomeric radiation can be calculated from eq. (1). This equation is obtained by expanding the three-dimensional plane $B_1 \varphi(Z,A)$ starting from the calibration curve in the $Z$-direction with $B_1 \varphi$ having a fixed value.

$$A-A^* = (Z_p-Z_p^*) \cdot \frac{1 - \Gamma_A}{\varphi_N} - \frac{(Z_p-Z_p^*)^2}{2} \cdot \frac{d}{dA} \frac{\Gamma_A/dA}{\varphi_N^2}$$

$$\varphi_N = Z_p^*/A^* = \text{charge density of the fissioning nucleus}$$

$$\Gamma_A = (\Delta \cdot B_1 \varphi/B_1 \varphi) / (\Delta A/A) = \text{mass dispersion of the separator as derived from the calibration curve.}$$

All nucleides the charge $Z$ and mass $A$ of which fulfill the above equation have the maximum of their $B_1 \varphi$-distribution at the same value $B_1 \varphi_{\text{max}}$ and may therefore be the isomer we are looking for. As the experimental values of $A^*$, $Z_p^*$, and $\Gamma_A$ are uncertain to some extent the mass value of a nucleide with a given $Z$ has to fulfill eq. (1) within about ± 0.7
Fig. 2: Region of light fission products. The hatched bands show the nucleides focussed simultaneously in the separator.

mass units. All candidates that are determined by equation (1) lie on certain sections through the Chart of Nuclides as is demonstrated by the hatched bands in fig. 2 for light fission products. The full lines in these bands are calculated from eq. (1) without any errors of the experimental parameters. The slope of the bands are defined by $\Gamma_A$ which is determined by $A^*$ and the gas-filling of the separator. The bands in fig. 2 are calculated for 6.5 torr He-gas which is the gas-filling used during the measurements in the light fission product group. It is seen that with helium the separator at certain masses of the fission products works as mass separator ($A \approx 88$), at other masses it is almost an isotone separator ($A \approx 100$). The region of nucleides which can be investigated with the separator is indicated in fig. 2.

If isomeric x-rays are emitted from the primary fragments, the $B\varphi$-distribution and the energy of these x-rays allow an almost unambiguous identification of the emitter. From the x-ray energy the nuclear charge $Z$ of the emitter can be determined, and the isomer is located at the crossing point of the horizontal line for $Z = \text{const.}$ in the $(\mathcal{N}, Z)$ plane, and the
section determined by eq. (1). A certain ambiguity remains if several isotopes are admitted by eq. (1).

2.2 Detection of isomeric radiation

The experimental arrangement installed at the focus of the separator is shown in fig. 3. The fission products are stopped by a Hostaphan-foil which is the end-window of the vacuum system. Immediately before being stopped they pass a fast transmission counter as proposed by Muga \[24\].

![Diagram](image-url)

Fig. 3: The experimental arrangement at the focus of the separator

The fission products produce in a thin plastic scintillator foil of NE 102 A light which is seen by a 56 AVP-photomultiplier. By a coincidence between two such detectors we reduce the \(\gamma\)-induced background by two orders of magnitude. This transmission-counter system starts a time-to-pulse height converter (TPC) the time range of which can be varied between 0.05 and 80 \(\mu\)s. The TPC is stopped by the isomeric radiation. For the measurement of isomeric \(\gamma\)-rays a Ge(Li)-diode (resolution 3.0 keV for \(^{60}\)Co \(\gamma\)-rays, size 38 cm\(^3\)) and of the isomeric x-rays from internal conversion of isomeric \(\gamma\)-rays a Si(Li)-detector (resolution 300 eV for \(^{57}\)Co, area = 1 cm\(^2\)) are used. A 4096-channel analyzer is used in the two parameter mode to store the results of the measurements i.e. the time signal of the TPC and the energies of the isomeric \(\gamma\)-rays and the energies of the isomeric x-rays, respectively. Typically the 4096 channels are divided into 32
time-channels and 128 energy-channels. The time resolution of the electronic system amounts to about 50 ns with the Si-diode and to 100 ns with the Ge-diode for low energy $\gamma$-rays. For high energy $\gamma$-rays it is considerably better. ($\gamma - x$) and ($\gamma - \gamma$)-coincidences between the detectors are possible.

![Counting rate vs Time](image)

**Fig. 4:** A typical decay curve

Some typical experimental results are shown in figs. 1 and 4 - 6. As discussed before fig. 1 is the $B\gamma$-distribution of the intensity of a certain isomeric $\gamma$-transition measured to identify the isomer. Fig 4 shows the decay curve for a certain isomeric $\gamma$-transition measured for a fixed value of $B\gamma$. Several of the $\gamma$-lines have a time-independent component which is produced by random coincidences between fission products and $\gamma$-rays with the same energy as the isomeric transition but following $\beta$-decays of fission products which have arrived earlier. This time-independent component is determined by delaying the stop-signals of the TPC in order to store in the first time-channels only random coincidences. The known time-independent component is then subtracted from the whole intensity measured for a certain $\gamma$-energy to obtain the true isomeric component. Fig. 5 shows an energy-spectrum of the Ge-diode measured for two times, i.e. just after the fission products have arrived and about
8 μs after the arrival of the fission products. γ -lines are seen that disappear completely within these 8 μs. Other lines are reduced considerably or are apparently time-independent. The resolution in this picture is bad as the number of analyzer channels used is small. Fig. 6 shows the x-ray spectra just after the arrival of the fission products and random coincidence spectra.

The intensities we measure amount typically to 10/min and 1/min for isomeric γ -transitions of light and heavy fission products, respectively. For x-rays they are about an order of magnitude smaller. The (γ -x)-coincidence rates are smaller than 6/h.

3. EXPERIMENTAL RESULTS

Our results are compiled in tables I, II, and III for light and heavy fission products, respectively. Tables I and II are divided into five parts. Part 1 gives the candidates for the isomerism calculated from eq.(1). Only the possible emitters with fission yields greater than 0.1 % are listed as we estimate the absolute intensities of our γ -lines to be higher than 10^{-3} photons/fission. Part 2 shows the experimental facts about isomeric x-rays, part 3 gives our results of the γ -ray measurements. In part 4 we have added the results of Cuy, ref. 9), for isomeric γ -transitions to compare with our γ -lines. In part 5 we listed the one or the few candidates which most probably emit the isomeric radiation cf. the following section. Table III contains the information about high energy isomeric γ -rays we found for light fission products. For these γ -rays the values of B·Q_{max} have not yet been determined and thus we cannot name the possible emitters.
Fig. 6: x-ray spectra measured in the mass regions near \( A = 99 \) and \( A = 135 \) with different fission product intensities. Full lines = spectra measured just after the arrival of the fission products, dashed lines = random coincidence spectra. Duration of the measurement = 3761 min and 7528 min for light and heavy fission products, respectively.

For the light fission products we found isomeric radiation in three different mass regions as is shown by the horizontal subdivision of table I. We want to mention that the isomeric radiation of each mass group as well as all the radiation in the heavy fission product group has the same value of \( B_0 \), a fact caused accidentally by the properties of the separator for the special gas-filling, i.e. 6.5 torr helium for light fission products and 0.5 torr nitrogen for the heavy fission product group.

We have identified some 40 isomeric \( \gamma \)-rays in the range of half-lives we investigated (\( \approx 0.3 \) to 80 ms) and with intensities > \( 10^{-3} \) photons/fission. It is obvious from tables I and II that the candidates we have determined are concentrated in narrow fission product mass regions around \( A = 88, 93, 99 \) and 135. While the candidates in the regions around \( A \approx 88, 89 \) emit together only three of the \( \gamma \)-lines in these tables the bulk of the \( \gamma \)-rays comes from the candidates near \( A = 99 \) and 135.
4. DISCUSSION OF NUCLEIDE ASSIGNMENTS

As is discussed in sect. 2.1 eq. (1) admits several nucleides as candidates for the isomeric γ-lines detected at a certain value of $B\gamma_{\text{max}}$. The uncertainties of the experimental parameters in eq. (1) which cause this ambiguity cannot be reduced essentially. The identification of the emitters of several γ-lines of tables I and II and the reduction of the candidates for the others is achieved by comparing our γ-ray results with the results of our x-ray and (γ - x) coincidence measurements, and with the results of the γ-ray measurements of other authors. The most probable emitters one gets considering all data we have are listed in part 5 of tables I and II.

4.1 Light fission products

The γ-lines with 1110 keV and 1590 keV which are emitted by fission products with masses around 88 have the same half-lives within the experimental errors. We find the same half-life for the x-rays of Br which is the only x-ray emitter in this mass region. Therefore we conclude that Br is the emitter of the two γ-rays. The only Br-isotope which fulfils eq. (1) is $^{88}\text{Br}$ and thus we assign the two γ-lines to $^{88}\text{Br}$. In ref. 9 the 111.2 keV γ-line is assigned to the mass 89. From our results $^{89}\text{Br}$ cannot be the emitter unless we increase the errors of the experimental parameters we insert into eq. (1) unreasonably.

We attribute the γ-line with $E_\gamma = 257$ keV to $^{92}\text{Rb}$ as the x-ray measurement reveals a half-life $> 10$ μs for Rb.

In the mass region near $A = 99$ twelve γ-lines must be attributed to eleven possible emitters. The situation is complicated by the fact that the line $E_\gamma = 169.8$ keV as well as the x-rays of Y show two different half-lives. The two half-lives might belong to different nucleides or might be caused by two isomeric states in one nucleide.

The half-lives of the long-lived component of the line with $E_\gamma = 169.8$ keV as well as of the 120.1 keV and 203.7 keV lines agree within the errors with the long half-life of the Y x-rays. Thus, $^{97}\text{Y}$ or $^{98}\text{Y}$ is most probably the emitter of this 8 μs-radiation, in agreement with the mass assignment of ref. 9) for the 204.0 keV line. A first (γ - x)-coincidence measurement gives coincidences between these lines and the x-rays of Y.
The assignment of the short-lived γ-transitions in this mass region is complicated by the fact that the short half-lives of the γ-rays as well as of the x-rays are very similar. Thus, the comparison between the half-lives of γ-rays and x-rays provides no unique identification of the nuclear charge of the isomer emitting a certain γ-line. The γ-lines with $E_γ = 101.0, 110.7, 130.0, 158.9,$ and $185.3$ keV are in coincidence with the x-rays of Y. They are grouped in table I according to their half-lives including the short-lived component of the $169.8$ keV-line. The half-lives of these lines agree within the errors, except the line with $E_γ = 130.0$ keV. We conclude that these lines are emitted by $^{98}Y$ as the mass assignment of ref. 9) for two of them is $A = 98$. The half-life of the $130.0$ keV-line is considerably smaller. If we assume a third isomeric state in Y; the half-life of $0.6 \mu$s of the x-rays of Y may be composed of the half-life of the $130.0$ keV-line and of the half-lives of the lines assigned to $^{98}Y$. According to ref. 9) the mass of the emitter of the $130.0$ keV-line is 99 or 100 giving with eq. (1) the candidates $^{99}Zr$ and $^{100}Nb$. This mass assignment does not agree with the result of our ($γ$-x)-coincidence measurement and the assumption of a third isomeric state in Y isotopes.

The $140.5$ keV-line is in coincidence with the x-rays of Rb. The assignment of this γ-line to $^{96}Rb$ is in agreement with the half-life of the x-rays of Rb and with the mass determination of ref. 9). The $167.2$ keV-line is also attributed to $^{96}Rb$ as its half-life and the mass assignment of ref. 9) agree with the values of the $140.5$ keV-line.

The $121.5$ keV-line is in coincidence with the x-rays of Zr. Thus, $^{99}Zr$ is the emitter in agreement with the x-ray half-life of Zr and with the mass of ref. 9). As we see the $121.5$ keV-line in coincidence with the x-rays of Zr there must be another transition in Zr, which causes the corresponding x-ray emission and which we do not see as a γ-line.

### 4.2 Heavy fission products

The relations are more transparent in the heavy fission product group. The γ-lines in table II are grouped according to the four half-lives we found. In ref. 9) the lines with $E_γ ≈ 115, 297, 1280$ keV are attributed to $^{134}Te$. It is assumed that they are emitted in a cascade with the $115$ keV and the $1280$ keV line as the highest and the lowest transition, respecti-
vely. The same cascade has been detected by Ahrens et al. 25) following the β-decay $^{134}\text{Sb} \rightarrow ^{134}\text{Te}$. Our results agree with this assignment.

The γ-ray lines with $E_\gamma \approx 197, 381$ and $1313$ keV and with half-lives of about 3 μs are assigned in ref. 9) to $^{136}\text{Xe}$. This assignment was supported by the $^{136}\text{Xe (p, p')}^{136}\text{Xe}$ measurement of Moore et al. 26) who found levels in $^{136}\text{Xe}$ in which the γ-rays fit as a cascade going to the ground-state. This attribution is in agreement with our half-life for the x-rays of Xe. Walton et al. 6) found γ-ray lines with $E_\gamma = 205, 390, 1330$ and the half-life $3.4 \pm 0.4$ μs in their measurement with a NaJ-crystal. The γ-ray lines were also seen following the β-decay of 48 s-isomer of $^{136}\text{J}$ with high spin 27-30). While Lundan et al. 28) proposed a cascade 196 keV - 382 keV - 1320 keV - 1313.2 keV going to the ground-state, Carraz et al. 29) and Monnand et al. 30) found a cascade 197.5 keV - 381.5 keV - 1313.3 keV going to the ground state.

The 391 keV γ-ray line is attributed to $^{134-136}\text{J}$ according to the half-life. Not identified is the emitter of the two remaining γ-ray lines in table II with $E_\gamma = 324$ keV and 1181 keV and a half-life of about 0.65 μs. The half-lives of the x-rays of Sn, Sb, Te agree with this γ-ray half-life. According to the mass determination of ref. 9) this radiation is emitted from fission products with $A = 134, 135$. Thus $^{134,135}\text{Sb}$ and $^{134,135}\text{Te}$ are the most probable emitters.

5. CONCLUSIONS

A comparison of the detected γ-ray lines from this measurement and from ref. 9) shows good agreement for the γ-ray energies and for the half-lives shorter than 1 μs. Intensities for the different fissioning systems have not been compared. The lines with low intensities of table III have not been listed in ref. 9).

A comparison between the isomeric x-ray yields of ref. 10) and ref. 11) shows that for times > 0.1 μs x-rays may be expected from Rb, Sr, Y, Nb, Mo, Tc and Sb, Te, J, Xe. X-rays from these elements have been found, except from Mo, Tc. The time-dependent components of the x-rays with energies corresponding to Mo and Tc in fig. 6 are caused by the $K_{\beta}$-radiation of nuclei with lower Z-values. Additionally we found a strong emission of x-rays at Zr.
As may be seen from part 5 of tables I and II the isomers we found are concentrated, except of two single isomers $^{88}\text{Br}$ and $^{93}\text{Rb}$, in two mass regions, $A = (96 - 100)$ and $A = (132 - 136)$. This fact is in agreement with the results of ref. 9) for $^{252}\text{Cf}$, where an intense long-lived contribution to the isomeric $\gamma$-radiation near these masses has been detected. Moreover, Johansson 4) observed an intense delayed $\gamma$-radiation for $A \approx 96$ and $A = (130 - 135)$.

![Chart of Nuclides](chart.png)

Fig. 7: The most probable emitters (hatched) inserted in the Chart of Nuclides. The stable nucleides are shown for orientation.

In the heavy fission product region isomers are found with $N = (82 \pm 1)$. There are $(g, g)$-nuclei among the emitters, as $^{134}\text{Te}$ and $^{136}\text{Xe}$. The $\gamma$-lines emitted as isomeric radiation have been found after $\beta$-decays 25, 27-30) and some of them promptly within a few ns after fission 31). The x-ray spectrum emitted following $\beta$-decay (dashed line in fig. 6) has a similar structure as the isomeric x-ray spectrum. The isomers from even-$Z$-nuclei only seem to be populated more intensely primarily than after $\beta$-decay. The isomeric states of $^{136}\text{Xe}$ and $^{134}\text{Te}$ are located at 1851 keV and 1692 keV above the ground state, respectively. These states fit well into the systematics of isomeric states of $N = 82$ nuclei 29). A low lying two-quasi-particle neutron state $(h_{1/2})^2$ with high spin $I = (5 - 7)$ may have enhanced life times in the nuclei of low level density near the closed shells $N = 82$ 32). The isomerism found is thus explained as spin isomerism of states which are populated both primarily and after $\beta$-decay. Isomers with odd neutron or proton numbers may be understood as long-lived three-quasi-
particle states of these nucleides which should be found at lower ener-
gies.

Fig. 7 reveals a striking observation in the mass region $A = (96-100)$. All nucleides we have identified or proposed as candidates have the neutron number $N = 58$ or 59, i.e. the bulk of the isomeric $\gamma$-rays with half-lives larger than 0.3 $\mu$s among the light fission products is emitted from nucleides with $N = (58-59)$. No even $A$, even $Z$ nucleides are found among the emitters, except $^{96}$Sr as possible x-ray emitter. The isomeric $\gamma$-lines have not been observed until now among the $\gamma$-lines following $\beta$-decays of the fission products. The dashed x-ray spectra shown in fig. 6 is mainly determined by x-rays following $\beta$-decays of short-lived fission products. The x-ray yields following $\beta$-decays show no similarities with the yield we found for isomeric x-rays pointing again to the fact that the isomeric $\gamma$-transitions with their characteristic x-ray-yields are only found among the primary fission products. From the energies and yields of the primary $\gamma$-quanta we estimate that the isomeric states in the region of $N = (58-59)$ are about (300-400) keV above the ground state.

We cannot offer a definitive explanation of the isomersim among light fission products. An accidental pile-up of spin isomers, although not very probable, cannot be excluded. We want to point out an unproved hypothesis to explain the isomerism. According to calculations of Arseniev et al. $^{17}$ $N = (58-60)$ is a transition region to a new region of deformed nuclei. Nuclei with $N \approx 59$ should have a stable oblate deformation with an energy excess of about 2 MeV compared to a spherical shape. Excited states of these nuclei may have prolate shapes as has been shown by Arseniev et al. $^{33}$ for nuclei in the mass region $Z > 50$ and $N < 82$. These states may be found at low energies, if the difference between the deformation energies of prolate and oblate shapes is small. They may be observed as low lying one-
 quasi-particle states with lifetimes enhanced due to large values of $\Delta K$. If these states are the explanation for the isomerism at $N = (58-59)$, our measurements tell that this shape isomerism is restricted to the transition region and is not found with half-lives in the $\mu$s range within the region of stable deformations for $N > 60$. It should be pointed out that a further transition region exists for $Z = (45-46)$ in the mass range $A = (110-115)$.

* The confinement of the isomerism to the transition region may be explained if a co-existence of deformed and spherical states is postulated, as has been found for Ir-nuclei $^{38}$.
In this region, too, isomers are located, as has been shown in ref. 4,9) in experiments with $^{252}$cf fission products. More extensive calculations and new measurements in the transition region may help to prove or disprove the hypothesis of the new proposed shape isomerism.

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| Table I: Isomeric x- and γ-rays measured for light fission products |
|---|---|---|---|---|---|
| nucleide assignment by separator from eq. (1) | isomeric x-rays | isomeric γ-rays |
| | element | \( t_{1/2} \) [\( \mu s \)] | rel. yield | \( E_x \) + 0.5 [keV] | \( t_{1/2} \) [\( \mu s \)] | rel. yield | \( E_\gamma \) [keV] | \( t_{1/2} \) [\( \mu s \)] | mass |
| 88-Se; 88-Br; 88,89-Kr; 89-Rb | Br | 6.8 ± 1.0 | - | 111.0 | 6.2 ± 0.6 | 45 | 111.2 | 3.0 * | 89±0 | 88-Br |
| 91-Br; 92-Kr; 93-Rb; 93,94-Sr; 95-Y; 96-Zr | Rb | >10 | - | 257.0 | 57 ± 15 | - | - | - | - | 93-Rb |
| 91-Kr; 95,96-Rb; 96,97-Sr; 97,98-Y; 99-Zr; 100-Nb; 101,102-Mo | Y | 0.6 ± 0.1 | 248 | 101.0 | 0.8±0.15 | 15 | 100.7 | 0.53 | 98±0,1 | 98-Y |
| | | 0.2 ± 0.9 | 100 | 110.7 | 0.00±0.15 | 7 | 111.0 | 0.76 | 98±0 | 98-Y |
| | | | | 158.9 | 0.75±0.2 | 9 | 158.0 | 1.50 | 97±1 | 97±1 |
| | | | | 169.8 | 0.62±0.15 | 35 | 170.5 | 1.1 | 98±0 | 98±0 |
| | | | | 185.3 | 0.85±0.1 | 8 | 186.4 | 0.65 | 98±1 | 98±1 |
| | | | | 120.1 | 7.9 ± 0.3 | 100 | - | - | - | 98±1 |
| | | | | 169.8 | 8.1 ± 0.3 | 64 | - | - | - | 98±1 |
| | | | | 203.7 | 8.1 ± 0.3 | 69 | - | - | - | 98±1 |
| | | | | 121.4 | 0.36±0.05 | 110 | - | - | - | 99±1,0 |
| | Zr | 0.32±0.08 | 550 * | 121.4 | 0.36±0.05 | 110 | 121.4 | 0.36 | 99±1,0 | 99-Zr |
| | Nb | 0.32±0.08 | 345 * | - | - | - | - | - | - | 100-Nb |

* estimated value
### Table II: Isomeric x- and γ-rays measured for heavy fission products

<table>
<thead>
<tr>
<th>Nucleide Assignment by Separator from Eq.(1)</th>
<th>Isomeric X-rays</th>
<th>Isomeric γ-rays</th>
</tr>
</thead>
<tbody>
<tr>
<td>132,133Sn</td>
<td>Sn</td>
<td>0.59±0.2</td>
</tr>
<tr>
<td>134,135Sb</td>
<td>Sb</td>
<td>0.58±0.12</td>
</tr>
<tr>
<td></td>
<td>Te</td>
<td>0.69±0.15</td>
</tr>
<tr>
<td>134-136J</td>
<td>Te</td>
<td>0.18±0.1</td>
</tr>
<tr>
<td>137Cs</td>
<td>J</td>
<td>0.97±0.15</td>
</tr>
<tr>
<td></td>
<td>Xe</td>
<td>1.8±1.2</td>
</tr>
</tbody>
</table>

### Table III: Isomeric γ-rays of the light group without detailed nucleide assignment

<table>
<thead>
<tr>
<th>Eγ [keV]</th>
<th>t₁/₂ [μs]</th>
<th>Rel. Yield</th>
<th>229</th>
<th>240.5</th>
<th>252</th>
<th>275</th>
<th>308</th>
<th>325</th>
<th>345</th>
<th>432</th>
<th>522</th>
<th>661</th>
<th>772</th>
</tr>
</thead>
<tbody>
<tr>
<td>&lt; 1</td>
<td>20 ± 5</td>
<td>3.7</td>
<td>23</td>
<td>20</td>
<td>25</td>
<td>&lt; 1</td>
<td>10</td>
<td>&lt; 1</td>
<td>2 ± 1</td>
<td>&lt; 1</td>
<td>&lt; 1</td>
<td>≈ 1</td>
<td></td>
</tr>
<tr>
<td>2.8</td>
<td>3.3±0.3</td>
<td>3.7</td>
<td>23</td>
<td>20</td>
<td>25</td>
<td>&lt; 1</td>
<td>10</td>
<td>&lt; 1</td>
<td>2 ± 1</td>
<td>&lt; 1</td>
<td>&lt; 1</td>
<td>≈ 1</td>
<td></td>
</tr>
</tbody>
</table>

*related to 100 in table I
SOME STUDIES ON NEUTRON-RICH NUCLEI

G. Herrmann, N. Kaffrell, N. Trautmann, R. Denig, W. Herzog, D. Hübscher and K.L. Kratz,
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1. INTRODUCTION

In the fall of 1966, when the Lyseki meeting was held, one of our major facilities for studying neutron-rich nuclides, the Mainz Triga reactor, went into routine operation. Apart from the thermal neutrons available at this facility, we have also been utilizing for irradiations 14-MeV neutrons delivered by a 600 kV Cockcroft-Walton accelerator through the T-D reaction. At the reactor neutron-rich nuclides are produced mainly by fission of $^{235}\text{U}$ and other fissile nuclei, at the accelerator mainly by (n,p)-reactions on enriched stable isotopes.

Both facilities offer certain advantages for work on short-lived nuclides. At the reactor bursts of thermal neutrons can be produced through power excursions which deliver to the sample more than $10^{14}$ neutrons/cm$^2$ within 50 ms. This short irradiation time leads to a very desirable enhancement of the ratio of short-lived activities to the accompanying longer-lived components. At the accelerator fluxes of $1\cdot10^{11}$ 14-MeV neutrons/cm$^2$s are now achieved which permit detailed decay-scheme studies even in the region of heavy elements where the production cross sections of neutron-rich isotopes are rather small.

Part of our program is devoted to the development of rapid chemical separations which have been summarized elsewhere$^{1-3}$. Isotopic separations have not yet been applied but we are planning to install an on-line system in the near future.
The second part of our program entails the use of such techniques to identify new nuclides and to confirm the existence of controversial nuclides. Time limitation does not permit a detailed presentation of the results achieved so far. Only an example for such exploratory experiments will be given in Sect. 2. A summary of our studies on delayed-neutron precursors has been published\textsuperscript{2).}

The third part of our program comprises decay scheme studies including $\gamma$-$\gamma$ coincidence spectroscopy with two Ge(Li)-detectors. Here, we climb from the stability valley somewhat slower than in the case of exploratory experiments. This is done to avoid experimental problems and to facilitate the interpretation of results, which is often difficult if no information on adjacent nuclei is available. We present below some results obtained for two transitional and two deformed nuclei, stressing the properties which might be typical for nuclei far off the stability line.

2. **HEAVY ISOTOPES OF BROMINE AND IODINE**

In order to demonstrate that nuclides with half-lives down to a few tenths of a second can directly be observed with our techniques, some data for halogen isotopes formed in fission of $^{235}\text{U}$ are shown in Table 1. The second column gives the half-lives found by following the decay of $\gamma$-rays in samples of bromine and iodine isolated by rapid gas chromatography\textsuperscript{2)} within 0.5 s after a neutron burst. The third column shows the half-lives deduced from the growth in the halogen fractions of already known\textsuperscript{4)} $\gamma$-rays of noble gas daughters. In the last column, results of an earlier identification of these nuclides by a milking technique are given. In this method, volatile halogen compounds were swept with a known flow-rate through a tube where their descendants were collected on a series of equidistant electrodes.

The agreement between the direct and the indirect meas-
Table 1

Half-lives of neutron-rich isotopes of bromine and iodine
[in seconds]

<table>
<thead>
<tr>
<th>Nuclide</th>
<th>(\gamma)-ray spectroscopic measurements</th>
<th>Milking techniques c)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>direct a)</td>
<td>indirect b)</td>
</tr>
<tr>
<td>Br-90</td>
<td>1.63 ± 0.14</td>
<td></td>
</tr>
<tr>
<td>Br-91</td>
<td>0.64 ± 0.08</td>
<td>0.62 ± 0.12</td>
</tr>
<tr>
<td>Br-92</td>
<td>0.25 ± 0.10</td>
<td></td>
</tr>
<tr>
<td>I-139</td>
<td>2.3 ± 0.6</td>
<td></td>
</tr>
<tr>
<td>I-140</td>
<td>0.84 ± 0.14</td>
<td>0.87 ± 0.13</td>
</tr>
<tr>
<td>I-141</td>
<td>0.45 ± 0.10</td>
<td>0.55 ± 0.25</td>
</tr>
</tbody>
</table>

a) Through decay of prominent \(\gamma\)-rays
b) Via the growth of \(\gamma\)-rays of krypton and xenon daughters

Measurements are satisfying, thereby supporting the assignment of the rapidly decaying \(\gamma\)-rays to the indicated isotopes which contain 9-14 more neutrons than their heaviest stable isotopes.

3. LEVELS OF \(^{106}\)Ru

Interesting properties of \(^{106}\)Ru were reported recently by Zicha et al.\(^5\)) who observed in the decay of the 36-s \(^{106}\)Ru a \(\gamma\)-ray cascade of 94, 219 and 344 keV energy in which the 94 keV transition was delayed with a half-life of 2.7 ns. It was suggested that the \(\gamma\)-rays connect four states of a rotational band built upon the ground state of \(^{106}\)Ru. The deformation deduced for this state was surprisingly high as indicated by a deformation parameter \(\beta = 0.55\).
Fig. 1. Low-energy $\gamma$-ray spectrum of short-lived technetium fission products

Simultaneously performed studies in our laboratory, however, give no evidence for this $\gamma$-ray cascade: in the $\gamma$-ray spectra of short-lived technetium isotopes formed in the fission of $^{238}$U by 14-MeV neutrons and displayed in Fig. 1, no peaks are observed at 94 and 219 keV (indicated by arrows). Only a 344 keV $\gamma$-ray is present as a shoulder on the strong 346 keV $\gamma$-ray of 50-s $^{103}$Tc but it may also belong to $^{103}$Tc since the shoulder does not vanish in successive spectra. The upper spectrum in Fig. 1 was recorded with a 30 cc Ge(Li)-detector, the lower one—corresponding to the energy region framed in the upper spectrum—with a 0.5 cc Ge(Li)-detector. No X-rays apart from the ruthenium X-rays can be seen in the lower spectrum, demonstrating the purity of the chemically isolated technetium samples.
Additional measurements were carried out with technetium samples produced by thermal-neutron induced fission of $^{235}\text{U}$ and $^{239}\text{Pu}$ but again no indication for a 94-219-344 keV $\gamma$-ray cascade was obtained.

The low-energy $\gamma$-rays of $^{106}\text{Tc}$ observed in our work are connected in Fig. 1 by thin lines. Other strong components stem from 50-s $^{103}\text{Tc}$ and 21-s $^{107}\text{Tc}$. The most intense $\gamma$-rays of $^{106}\text{Tc}$ occur at 270 and 522 keV, as has also been found by Wilhelmy $^6$) in a study of short-lived $\gamma$-ray emitters in spontaneous fission of $^{252}\text{Cf}$ performed with $\gamma$-X-ray coincidence techniques.

A partial decay scheme deduced from our singles and coincidence $\gamma$-ray spectra is shown in Fig. 2. The level sequence up to 991 keV energy agrees with recent results of Casten et al. $^7$ obtained with the $^{104}\text{Ru}(t,p)^{106}\text{Ru}$ reaction.
These authors have also discussed the peculiarities of the low-energy states of $^{106}$Ru which constitute a transitional situation midway between the rotational and the vibrational description.

Other interesting aspects of the level scheme, Fig. 2, are the level at 1092 keV energy which requires further investigation, and the gap in the level structure. Most of the $\beta$-transitions go to levels above the gap.

4. **LEVELS OF $^{192}$Os**

The second case to be discussed deals with the region of transitional nuclei situated between the rare earth elements and the $Z = 82$, $N = 126$ shells. When $^{192}$Os is bombarded with 14-MeV neutrons, a 5.9-s activity is observed with a relatively simple $\gamma$-ray spectrum shown in Fig. 3. This activity has to be assigned either to $^{192}$Re from the $^{192}$Os(n,p)-reaction or to an isomer of $^{192}$Os formed by the $^{192}$Os (n,n') reaction. Attempts to decide between both cases by a rapid chemical separation failed so far. In both cases, however, the $\gamma$-rays displayed in Fig. 3 belong to excited states of $^{192}$Os.

Two examples for $\gamma$-$\gamma$ coincidence spectra obtained with two Ge(Li)-detectors by summing up 400 runs are shown in Fig. 4. The lower spectrum connects — in the terminology of deformed nuclei — members of the ground state rotational band, the upper spectrum those of the band built upon the $\gamma$-vibration.

The level scheme of $^{192}$Os is given in Fig. 5. Spin and parity assignments not enclosed in brackets are taken from coulomb excitation studies$^8)$. The assignments enclosed in brackets are proposed by us taking into account some analogies to the well-investigated nucleus $^{190}$Os, especially in the assignment of $J^\pi = 10^-$ to the highest level observed.

Both the ground state rotational band shown on the left
of Fig. 5 and the $K^\pi = 2^+$ rotational band on the $\gamma$-vibration shown in the middle are populated up to spin values of $J = 8$. The $J^\pi = 4^+$ level at 1069 keV may be interpreted, like a corresponding level in $^{190}$Os, as the two-phonon $\gamma$-vibration. None of both energy sequences follow the $J(J+1)$ rule for rotational bands. In the $\gamma$-band, a strong even-odd staggering effect is immediately evident in Fig. 5: the $5^+$ and $7^+$ levels are shifted to higher energies compared to the positions of their adjacent even-spin levels. The depopulation of the $\gamma$-band shows also an interesting trend: with increasing spin, deexcitation occurs preferentially within the band, not into the ground state band.

In Fig. 6, the level energies obtained experimentally (shown at the left) are compared to theoretical predictions of the asymmetric rotator model of Davydov et al.\textsuperscript{9,10} and of the rotation-vibration model of Faessler et al.\textsuperscript{11} (at the right). Whereas theoretical values from both the models agree with the experimental data in the case of the ground
state band, the $\gamma$-band and the $4^+$ state are described better by the rotation-vibration model. It predicts the upward shift of the odd-spin levels in the $\gamma$-band whereas in the asymmetric rotator model, these levels are expected to shift downward. However, the staggering effect may not only result from interactions between $\beta$- and $\gamma$-vibrations, but also from a completely $\gamma$-unstable structure of the nucleus$^{12}$).

5. **LEVELS OF $^{164}$Dy**

In even-even deformed nuclei, four-quasiparticle and two-phonon excitations are to be expected at higher excitation energies but have not yet been clearly substantiated. We have obtained evidence for such states among the levels of $^{164}$Dy fed in $\beta$-decay of the 3.0-min $^{164}$Tb which we produce by the $^{164}$Dy(n,p)$^{164}$Tb reaction.
Fig. 5. Level scheme of $^{192}$Os

The low-energy $\gamma$-ray spectrum of $^{164}$Tb is shown in Fig. 7. A total of 109 $\gamma$-rays were observed; 98% of the $\gamma$-ray intensity could be placed in a decay scheme. Part of the scheme, relevant to the aspects to be stressed here, is given in Fig. 8.

The low-lying levels of $^{164}$Dy are well known from nuclear reaction studies$^{13}$. In the decay of $^{164}$Tb, the rotational bands built on the ground state (at the left in Fig. 8), on the $K^\pi=2^+$ $\gamma$-vibration (in the middle) and on a $K^\pi=2^-$ octupole vibration (at the right) are observed. The strongest two-quasiparticle components which should contribute to these collective states are given in the inserts of Fig. 8.

An interesting level occurs at 2206 keV energy, populated by a fast $\beta$-transition of the type 'allowed-unhindered, au'. Normally, 'au' $\beta$-transitions feed rather pure two-
Fig. 6. Comparison of the observed levels of $^{192}$Os with theoretical predictions.

Quasiparticle states whose Nilsson quantum numbers [Nn Z A] are the same as those of the initial nucleus. For the ground state of $^{164}$Tb, the structure p[411↑],n[633↑] coupling to $J^\pi = 5^+$ is to be expected according to data on adjacent nuclei. For the 2206 keV state, the depopulation by γ-ray transitions supports strongly a $J^\pi = 4^+$ assignment. However, inspection of the theoretically predicted two-quasiparticle states$^{14}$ of $^{164}$Dy shows no such state which can be populated from the ground state of $^{164}$Tb by a β-transition of the type 'au'.

Therefore, the 2206 keV state is suggested to be a rather pure four-quasiparticle state of the configuration

$$p[411↑],p[523↑],n[523↑],n[633↑],$$

i.e., it is composed of the lowest pp and nn two-quasiparticle states$^{14}$ in $^{164}$Dy. Thus the β-decay occurs between the orbitals n 5/2$^-$[523↑] → p 7/2$^-$[523↑] while the component n[633↑],p[411↑] represents the configuration of the
initial state which remains unchanged. The level energy is somewhat lower than expected for a four-quasiparticle state, but this may be explained by the spin splitting effect\textsuperscript{15}). Some admixtures of other configurations should contribute to the state because of its rapid deexcitation into the $\gamma$-band and the $K^- = 4^-$ band which is strictly forbidden for a pure four-quasiparticle state. Serious deviations from the Alaga rules are found when the reduced branching ratios of $\gamma$-transitions from the 2206 keV level into the $\gamma$-band are calculated assuming pure E2 transitions.

Another state which should be mentioned is the $J^\pi = 4^-$ level at 1588 keV energy on which a rotational band is built. Here, the branching ratios of deexciting $\gamma$-rays agree with the Alaga rules. Soloviev\textsuperscript{16}) has suggested that this level is a two-phonon state consisting of quadrupole $2^+$ plus
Fig. 8. Partial decay scheme of $^{164}$Tb

octupole $2^-$ excitations. Two- and four-quasiparticle components indicated in the insert of Fig. 8 may also contribute.

6. **LEVELS OF $^{238}$U**

The last case to be presented is one of the heaviest nuclei accessible to us: $^{238}$Pa, decaying with 2.3 min half-life and produced by the $^{238}$U(n,p)-reaction. A multistep solvent extraction procedure is applied to separate $^{238}$Pa from the strong accompanying fission product activities.

The $\gamma$-ray spectrum, Fig. 9, is rather complex. About 165 $\gamma$-rays have been identified so far. Coincidence measurements are time consuming because of low source strengths. To get meaningful coincidence data, about 100 runs had to be summed up.
Fig. 9. $\gamma$-ray spectrum of 2.3-min $^{238}$Pa

About 80% of the total $\gamma$-ray intensity is placed in a tentative decay scheme shown in Fig. 10. Further coincidence measurements are in progress to reveal the open questions. We restrict ourselves to the discussion of some low-lying states; for the sake of clarity the branchings from higher levels into these states are not shown in Fig. 10. Apart from the usual rotational bands built on the ground state (at the left) and the $K^\pi = 2^+$ $\gamma$-vibration (at the right), three low-lying octupole vibrations with $K^\pi = 0^-$, $1^-$ and $2^-$ are observed. Some of the levels belonging to these bands have also been found in coulomb excitation and scattering studies\textsuperscript{17,18}). Collective levels of this kind have also been theoretically predicted in this energy region.
Fig. 10. Partial level scheme of $^{238}_{\text{U}}$.

in two studies$^{19,20}$. Both calculations predict the $K^\pi = 0^-$ band to be the lowest one but differ in the ordering of the $K^\pi = 1^-$ and $2^-$ bands.

The lowest two-quasiparticle configuration $nn[631\uparrow]+[622\uparrow]$ is expected at about 900 keV according to the calculations of Soloviev and Siklos$^{21}$. This structure is proposed for the level at 1059.5 keV energy. The level lies very close to the $\gamma$-vibration at 1060.2 keV, and could be
conclusively detected only via coincidence with the populating 502 keV $\gamma$-ray. Contrary to the usual behaviour, the $K^\pi=2^-$ band deexcites preferentially into the 1059.5 keV level and not into the $\gamma$-vibration.

7. **ACKNOWLEDGEMENT**

We wish to thank the Bundesministerium für Bildung und Wissenschaft for financial support.

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THE $^{132}$Sn-REGION STUDIED BY MEANS OF THE OSIRIS FACILITY

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Research Institute of National Defence,
Stockholm, Sweden.

Configurations which consist of a single particle or a single hole in addition to closed shells display simple low-energy nuclear spectra. Such spectra provide especially detailed evidence of the independent-particle motion. The magic numbers are 2, 8, 20, 28, 50, 82, and 126. Nuclei representing the first four of these (the $^4$He, $^{16}$O, $^{40}$Ca, and the $^{56}$Ni regions) and the two latter (the $^{82}$Pb, $^{126}$region) have been extensively studied (although the evidence on the $^{56}$Ni-region is rather incomplete). The only other combination which is physically realizable, i.e. which is inside the limits for particle emission ($S_p = 0$ and $S_n = 0$ in the nuclear chart) is the $^{132}$Sn, $^{50}$Sn-region. This region, however, constitutes a conspicuous gap, since information is almost nonexistent, although it has been shown that it is a reasonable assumption to consider the $^{132}$Sn core to be doubly magic$^1$). In view of these facts and because of a longstanding tradition$^2$) it was considered worthwhile to devote a considerable part of the OSIRIS experiment to this subject.

Fig. 1 shows a compilation of known and estimated half-lives and $Q_\beta$-values for nuclides in the region. The estimates (numbers within parenthesis) are based on energy-level, logft-value and $Q_\beta$-value systematics. Special estimates, also using neutron binding energy systematics, were made for relevant nuclei.

Based on the compilation, a programme for the investigation was drawn up in which the eventual goal of the single particle aspect is to reach the "valence nuclei" (inside fully drawn lines in fig.1). However, the drop-off of the fission yield just passes through the "valence star" and complicates the study of particularly $^{131}$In and $^{133}$Sn. Another aggravating factor is that $^{133}$In is
estimated to decay intensely via delayed neutron emission. Thus it was decided first to study nuclei which are either of the single closed shell type or which have one closed shell plus or minus a nucleon (marked by an extra square in fig. 1). Another reason for study of some of these nuclei is the necessity of mapping the daughters of the main nuclides.

Using the obtained single particle energies, interesting calculations may be performed on the nuclei having two particles (two holes or a particle-hole) combination outside closed shells. Again the area of relatively straightforward access is limited by low yields and short half-lives. The first goals of investigation are intended to be $^{134}\text{Xe}$, $^{132}\text{Xe}$, $^{130}\text{Sb}$, and $^{132}\text{Sn}$ which can be observed.
in the decay of 11 sec. $^{134}$Sb, 1 min. $^{132}$Sn, and 0.5 sec. $^{130}$In.
In connection with the study of levels in even tellurium isotopes
at the 225 cm cyclotron in Stockholm the decay of $^{128-132}$Sb has
been investigated (cf. separate contribution).

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* * *

**EXPERIMENTAL TECHNIQUES FOR NUCLEAR STRUCTURE STUDIES**
**IN THE TIN-REGION AT OSIRIS**

S. Borg, L.-E. De Geer, G.B. Holm and B. Rydberg,
Research Institute of National Defence
and Research Institute for Physics,
Stockholm, Sweden.

It is characteristic for the OSIRIS target-ion source that a
number of neighboring elements are produced with efficiencies of
the same order of magnitude. (The only exception are essentially
the transition and rare earth elements which have not been observed
so far and the alkalisides which are enhanced.) This implies that
the problems that meet the nuclear spectroscopist are entirely dif-
ferent from those which he is accustomed to and also slightly dif-
ferent from those of the usual ISOL-system\textsuperscript{1,2}) where foremost one
element is primarily produced. Whereas one can be sure of the mass
number, the element assignment is more difficult. It is especially
hard to discern short-lived, low-yield members in a fission chain
in the presence of high-yield members.

The heart of the experimental arrangement is a PDP-9 small general
purpose computer which is used both for control of equipment and
for data acquisition and handling. The radiations of the different
elements can in general be assorted into groups each having its
characteristic half-life, with the aid of a moving tape collector.
The element assignment of the various groups can be made using a
Si(Li) X-ray detector (Cf. fig. 2). The tape collector, which is
a considerably improved and extended version of the previous mo-
dels\textsuperscript{2}), is regulated by the PDP-9. Thus the tape speed, the ir-
radiation time, the background counting time etc. are all con-
trolled by the computer. This indirect method of element assign-
ment can be avoided if an extra step of chemical separation is
included. However, when one is trying to reach as short half-lives
as possible, an extra separation step is undesirable.

\[ A=131 \]
\[ I=2 \]
\[ T=100 \]

---

Fig. 2
The intensity ratios of some photo peaks at two different tape velocities, 0.8 cm/min and 40 cm/min. (The ratio units are arbitrary, since different reactor effects and measuring times were used. The error bars are based on statistical considerations). The result shows how one can easily sort, at least the strong transitions, into different elements by applying suitable measuring conditions. The ratios correspond to the following activities: S=25 min. Te; 2=23 min. Sb; 0.4=1.1 min. Sn; 0.09=0.3 sec. In.

For data acquisition, programs have been written for multiscale, Pulseheight analysis, multispectra and coincidence measurements. Half-lives and peakenergies and -intensities can both be evaluated by using simple estimates employing the light pen and by using more elaborate least-square fits.

Another difficulty is that investigating short-lived activities in the OSIRIS system intensity is usually a problem if one wants to perform more elaborate experiments like angular correlations or magnetic spectroscopy. This makes spin and parity assignments difficult. However, working in the doubly magic $^{132}$Sn-region, this obstacle may be partly overcome by the anticipated lucidity of the region.

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* * *
PROTON STATES IN THE TIN-REGION


The valence states of the 51 st proton can be studied in the decay of $^{133}$Sn. The half-life of this nuclide has previously been given as 55 sec.\(^1\). However, we have not been able to observe any activity of this half-life at mass number 133, neither has it been possible to observe any gamma rays with half-life shorter than those from 2.1 min. $^{133}$Sb, using the moving tape method, except the 9 sec $^{133}$mI transitions (see separate contribution). The apparent nonexistence of the 55 sec. activity is quite in line with systematics since the estimated half-life of this nuclide is of the order of 5 sec (as estimated in the contribution by G.B. Holm).

On account of the expected high $Q_\beta$-value of $^{133}$Sn it was decided to look for gamma rays belonging to this decay by performing beta-gamma coincidence measurements, using beta windows of different energies. With increasing beta energy only one gamma ray, at 963 keV, remains. This energy agrees well with what can be extrapolated from other N=82 nuclei and is thus interpreted as leading from the d$^{5/2}$ excited state to the g$^{7/2}$ ground state. (Cf. fig. 3). The half-life of this decay, as determined by multiscaling of the high-energy part of the beta spectrum, is 1.7 sec.

The configuration of three protons outside closed shells is represented by $^{135}$ 53\(^+\)82 whose levels are populated in the decay of 18 sec. $^{135}$Te. In this decay essentially three gamma rays are observed having energies (and intensities): 267(8 ), 604(100 ), 871 (23 ). The two first transitions are in coincidence. The relative strength of the 267 keV transition could either be due to that it is M1, whereas the 871 keV transition is E2 (this was the interpretation given to a similar situation in $^{137}$Cs\(^2\)), or to a large similarity between the first and second excited states. Since the log ft-value does not favour the former interpretation (in contrast to in the $^{137}$Cs-case), the latter is preferred. It thus appears
that the three quasiparticle 5/2-state may occur at the low exi-
tation of about 0.8 meV both in $^{135}$I and $^{137}$Cs. In $^{135}$I, however, it happens to come close to the one quasiparticle state with the result that the two states mix strongly and repel each other. Such a mechanism would explain the irregularity in the position of the first exited state in $^{135}$I.

No clear evidence of the valence proton hole states (in $^{131}$In$_{82}$) which are populated in the decay of $^{131}$Cd has been obtained so far, most probably because of low production yield.

Another way of studying proton states is in the combination one proton and two or four neutron holes. These levels are populated in the decay of 1.1 min. $^{131}$Sn and 2 and 6 min. $^{129}$Sn. These decays have been met in passing when the main object has been other decays. However, they show, as could be expected, that the remaining neutron core is far from inert, displaying level schemes which contain a number of "extra" states.

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NEUTRON HOLE STATES IN THE TIN-REGION


The following half-lives shorter than one hour were previously known at mass number 131: 25 min. Te, 23 min. Sb and 1.3 min. Sn. Since there exists a possibility to find several new activities of considerable interest - apart from two very short-lived In also a 19/2- isomer in Sb and the h\textsubscript{11/2} isomer in Sn - this mass has been subjected to a search in a wide range of half-lives. For the already known activities this has resulted in two very long lists of previously unknown transitions in the decay of Sn and Sb plus the confirmation of a large number of Te transitions found by Wallers et al. 1)

On the short-lived side of the time spectrum a new activity of 0.3 sec. was discovered. By just using different tape speeds a transition of 2431 keV could be ascribed to this activity. However, employing the beta coincidence method, in addition to the 2431 transition a weaker, shortlived gamma ray at 330 keV was discovered. Like in the case of 133\textsuperscript{a}Sn we have no positive element identification. The shortness of the halflife and the high beta energy, makes it most plausible that we deal with the element that precedes the most short-lived one previously identified in the decay chain. We thus assign the 0.3 sec to the decay of In which populates levels in the valence nucleus 131\textsuperscript{b}Sn\textsuperscript{25/2+}.

In contrast to the proton state situation this region is more isolated, the closest parallel being 125\textsuperscript{a}Sn. From systematics, such as it is, it is indicated that the 131\textsuperscript{a}In ground state can be assigned to be g\textsubscript{9/2}. Since the only In-Sn transition that can give rise to the short half-life of 0.3 sec. is g\textsubscript{9/2} - g\textsubscript{7/2}, it appears likely that the 2431 keV gamma ray represents a transition to the expected g\textsubscript{7/2} state in 131\textsuperscript{a}Sn.
The origin of the 330 keV gamma ray is more problematic. It cannot be due to an $h_{11/2} - d_{3/2}$ M4 transition because it is observed in the beta-gamma coincidence measurements. The other state which may occur at low excitation, $s_{1/2}$, can only be directly populated from a possible $p_{1/2}$ isomer in In. However, the $p_{1/2} - s_{1/2}$ transition is expected to be slower than the $g_{9/2} - g_{7/2}$ by at least an order of magnitude. Whereas no direct multiscaling measurements on the very short-lived transitions have been performed yet, the beta-gamma coincidence measurements have been performed at different tape speeds. The results show that they have half-lives of the same order of magnitude. Furthermore, gamma-gamma coincidence measurements with the gate at 2431 keV yield negative results.

Another possibility is that the 330 keV transition represents a state in In populated in the decay of Cd. The fission yield of Cd is expected to be lower than that of In by about an order of magnitude. The relative intensities of the transitions do not contradict this assumption too much. However, again there is no straightforward interpretation of the 330 keV transition. One can assume that the Cd ground state is $f_{7/2}$ which should decay to $f_{5/2}$ in In with a half-life of about the same order of magnitude as In. However, the $f_{5/2}$ can be expected to cascade to the $p_{1/2}$ and $g_{9/2}$ states via a $p_{3/2}$ state, thus giving rise to more than one gamma ray. Refined measurements with improved statistics are being prepared for this problem.

The decays of $^{129}$In and $^{127}$In have also been studied. As could be expected, the $^{131}$In simplicity of the structure has disappeared here. For example in $^{129}$Sn, the 2431 keV level has been split into at least two levels (at 2118 and 1865 keV) and a number of "extra" states have appeared (at e.g. 1009, 770 and 729 keV).

In $^{131}$Sn an isomeric state is expected. However, so far only one half-life has been found to be associated with the decay of tin (in e.g. following Sb X-rays) and in general no beta activity which is unaccounted for has been discovered. A survey of the $h_{11/2}$
and $d_{5/2}$ decays of Sn-isotopes of lower mass number indicates that at $A = 131$ the half-lives may be approximately the same and so short that the isomeric transition is virtually nonexistent. The energy of the strong 798 keV transition agrees well the $d_{5/2} - g_{7/2}$ systematics in the Sb-isotopes. It is thus likely to be a prominent member of the $d_{3/2}$ decay. There is then a possibility that the even stronger 1226 keV transition and gamma rays associated with it (like 450 keV which is in coincidence), belong to the $h_{11/2}$ branch of the decay and actually have a slightly different half-life. There is a tendency of a difference of a few seconds in half-life, but better statistics are needed to prove this hypothesis.

Another activity that has been looked for, but not been found is a $19/2^-$ isomer in Sb. Such a state would be a parallel to the 53 min state in $^{135}$Os, which decays via an $11/2^+$ level to the $7/2^+$ ground state, or to the 9 sec isomer in $^{133}$I which presumably cascades via 15/2+ and 11/2+ states to the $7/2^+$.

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**LEVELS IN DOUBLY-EVEN TELLURIUM ISOTOPES**


The level systematics of doubly-even Te has been studied by means of $(\alpha, 2n\gamma)$ reactions for the lighter isotopes ($^{114-126}$Te). The levels of $^{128-132}$Te were populated in doublyodd Sb decay.
The $2^+$ and $4^+$ levels show a vibrational pattern for $114 \leq A \leq 128$, whilst in the heavier isotopes the shell effect is noticeable. The energy of the $6^+$ state is nearly constant (1800 keV) for $118 \leq A \leq 132$. It can be explained by the two-proton configuration $[p_1(g_{7/2})p_2(d_{5/2})]_6^+$ for these states.

Two-neutron states arising from configurations $[n_1(h_{11/2})n_2(s_{1/2})]_5^-$ and $[n_1(h_{11/2})n_2(d_{3/2})]_7^-$ are observed for $126 \leq A \leq 132$.

**Levels in even-even Te-isotopes**

![Diagram of levels in even-even Te-isotopes](image)

Fig. 4
EVIDENCE FOR A HIGH-SPIN THREE QUASIPARTICLE STATE IN $^{133}$I PRODUCED IN FISSION

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Peker has pointed out that among low-lying states of $N$ or $Z = 80$ there should exist negative parity, two quasiparticle states of the type $[h_1^{1/2}, d_3^{3/2}]_7^-$, which in favourable cases can be studied in isomeric decays. In fig. 5 we plotted the energies of the 7- states in even-even nuclei having $N = 80$. The regularity of all low-lying states is remarkable. If one assumes that the addition of a proton is not going to change the level sequence very much, one might be able to find states having spin as high as $21/2$ or $19/2$ at about 2 MeV. Only one such case has so far been reported, namely $^{135}$Cs($19/2^-(840 \text{ keV})11/2^+(781 \text{ keV})7/2^+$). In fig. 6 we have used the experimental information on $^{135}$Cs and the information contained in fig. 5 so as to give us an idea of where the high-spin negative parity states may be found in other nuclei having 80 neutrons. For $Z \geq 57$ (HI, Xn, γ)-investigations are most suitable in searching for the expected 7- states. For $Z \geq 53$ one has to rely entirely on nuclei produced in fission. In the search for new short-lived activities on the interesting mass number 133 at the OSIRIS facility a component with a 9 sec. half-life was found. However, it was soon clear that this decay was not due to a hitherto unknown β-decay (like for example $^{133}$Sn), because the energies of the two gamma rays associated with the 9 sec. half-life were 647 and 912 keV, respectively. These energies coincide exactly with the two strongest gamma rays found in the decay of $^{133m}$Te. However, in this decay the 912 keV gamma ray is considerably stronger than the 647 keV gamma ray, but in our measurements they are of equal intensity, indicating an isomeric cascade. In addition to photo peaks corresponding to the 647 and 912 keV gamma rays, there is a 9 sec. peak at 1559 keV, the intensity of which is consistent with the pile up of the two gamma rays. This fact suggests that they are in
coincidence, which is further supported by coincidence measurements. Also the iodine X-rays show a 9 sec. component. Additional support for the fact that we are dealing with an iodine isomer is the observation that the relative intensity of the 647 and 912 keV gamma rays were much stronger when using a fission-ion source which favours volatile elements.

**Fig. 5.** Systematics of the $2^+$, $4^+$, and $7^-$ states in even-even nuclei for $N = 80$.

**Fig. 6.** Expected levels in odd nuclei for $N = 80$. The level positions are guessed from fig. 5 and the known levels of $^{135}$Cs.

It seems unlikely that the 647 keV gamma ray is responsible for the isomeric transition. This would imply an $M3$ or $E4$ transition and thus very high-spin, positive parity states which are hard to explain. If, on the other hand, the isomeric transition is caused by a high-spin, negative parity state, as suggested by the introductory discussion, the experimental results would be consistent assuming that the isomerism is caused by a low energy, highly converted transition. An extensive search for a gamma ray in the
region 50-100 keV yielded a negative result, employing both singles and coincidence techniques. However, in a preliminary run with a Si(Li) electron detector it was possible to detect conversion electrons corresponding to a 73 keV transition. Because of the uncertainty in the intensity calibration it is not possible to give a definite value for the K/L-ratio. The experimental value (4), however, indicates E2 (2), M2 (4) or M3 (2). On the other hand, the half-life indicates E3 (with $P_w=10$) or M2 ($P_w=5.10^5$). Thus, the only assignment that can be reconciled with the two measurements is M2, even if the consequential retardation is unusually high. If the efficiency of the detector drops by a factor of eight from 40 keV to 70 keV — which is hardly likely — the transition is instead an E3 or M4. M4 is excluded by the half-life, whereas the retardation of the E3 transition is slightly low.

It can thus be concluded that all available evidence indicates the existence of a high-spin, negative parity state in $^{133}_{\text{I}}$. If the isomeric transition is M2 the cascade may either be $19/2^- \rightarrow 15/2^+ \rightarrow 11/2^+ \rightarrow 7/2^+$ or $17/2^- \rightarrow 13/2^+ \rightarrow 11/2^+ \rightarrow 7/2^+$. The E3 alternative gives $19/2^- \rightarrow 13/2^+ \rightarrow 11/2^+ \rightarrow 7/2^+$. For the time being there are arguments for and against all these alternatives. For instance, the high M3 retardation may be explained by the expected mismatch of the wave functions of the $19/2^-$ and $15/+^+$ states. On the other hand the second excited state cannot very well be $15/2^+$, because of the log ft-value in the decay of $11/2^-{^{133}_{\text{Te}}}$. However, the main observation, namely that an isomer having a spin of the order of $19/2$ has been observed in fission, remains true.
ON THE DECAY OF $^{110}\text{Sb}$, $^{111}\text{Sb}$, $^{112}\text{Sb}$, $^{113}\text{Sb}$ AND $^{114}\text{Sb}$

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The nuclear structure of the Sb and Sn isotopes has been the subject of numerous experimental and theoretical investigations. It is of particular interest to extend these studies to the region of the neutron deficient isotopes, primarily to see how well these isotopes remain "superconducting" as one approaches the doubly magic structure $^{100}\text{Sn}$.

Two previously unknown light antimony isotopes, $^{110}\text{Sb}$ and $^{111}\text{Sb}$, have been produced through the $(p,3n)$ and $(p,2n)$ reactions on $^{112}\text{Sn}$, respectively. The gamma spectra following their beta decay have been studied, together with those of $^{112}\text{Sb}$, $^{113}\text{Sb}$ and $^{114}\text{Sb}$, produced by $(p,xn)$ reactions on $^{112}\text{Sn}$ and/or $^{114}\text{Sn}$. Enriched metallic $^{112}\text{Sn}$ (70%) and $^{114}\text{Sn}$ (60%) targets were exposed to protons of 10 to 47 MeV from the UCLA cyclotron. The targets were shuttled between the internal beam of the cyclotron and the counting area, along a pneumatically operated rabbit system. In a single counter experiment, the gamma rays were detected by two Ge(Li) detectors of 35 cc and 3 cc, respectively, with an overall energy resolution of 3 keV for $^{60}\text{Co}$. After passing through standard pile-up rejection circuits, the pulses were digitized by 1024 channel ADC's and stored as a function of the decay time, on sequential disc files of an XSDS-925 on line computer system. The two dimensional energy versus decay time spectra thus obtained were later on brought back into core memory for energy, intensity and half-life analysis, using peak fitting routines to a gaussian shape with low energy tail.

In Fig. 1a, a 128 x 16 cut of one of these two dimensional spectra is shown. In this example, a $^{112}\text{Sn}$ target had been exposed

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to a 10 sec burst of 47 MeV protons, and the decay of its activity followed in 10 sec intervals. One clearly sees the 23 sec decay of the 1213 and 1245 keV gamma rays of $^{110}$Sn, the 53 sec decay of the 1258 keV line from $^{112}$Sn, and a longer lived component consisting of a mixture of the 1300 keV and 1298 keV transitions in $^{114}$Sn and $^{116}$Sn, respectively. The accumulated sum during the first 80 sec following the bombardment is shown in Fig. 1b. The total spectrum for the next 80 sec is shown in Fig. 1c.

The isotopes produced were clearly identified by their half-life and by the threshold of the relevant (p,xn) reaction. The half-lives of the produced Sb isotopes were measured by integrating the counts below one particular gamma peak for all time intervals of the $(E_y,t)$ decay spectrum. The results are shown in Fig. 2. Only the time-analysis of the strongest observed gamma transition has been plotted, for the clarity of the figures. Our results on $^{112}$Sb, $^{113}$Sb and $^{114}$Sb agree very well with previously reported values$^1$. The half-life of $^{113}$Sb, not shown in Fig. 2, was measured to be $6.3 \pm 0.5$ min.

Coincidence experiments were then conducted using a 10 cm diameter x 10 cm thick NaI(Tl) and a 35 cc Ge(Li) detector. The output of these two counters was digitized by 1024 channel ADC's and stored on magnetic tape, together with their associated time-amplitude converted spectrum, (TAC) used to distinguish between real and accidental coincidences. The magnetic tape was played back into the on-line computer and 2 dimensional E(NaI) X E(Ge) arrays were formed, for particular decay times and TAC spectrum selections.

Combining the results of both experiments, most of the observed gamma transitions can be fit into a decay scheme. The energy levels obtained, accurate to 1 keV, agree very well with the results of Sn(p,d)$^{2,3}$ and Cd(α,xn) reactions$^4$. These preliminary decay schemes are presented in Fig. 3 and Fig. 4, together with the relative intensity of the observed transitions which were sufficiently strong to be accurately analyzed.

Although the mass difference between the relevant Sb and Sn isotopes has not yet been measured, the measured half-lives and esti-
mate from semi-empirical mass relationships\(^5\)) indicate clearly that most of the \(\beta\) decay of the investigated Sb isotopes is allowed. This calls for the following conclusions.

Levels of \(J^\pi = 3/2^+, 5/2^+\) but not \(1/2^+\) are populated in \(^{111}\)Sn and \(^{113}\)Sn. This indicates that both \(^{111}\)Sb and \(^{113}\)Sb have \(J^\pi = 5/2^+\). This assignment is in good agreement with a 2d 5/2 proton orbital, expected for a single proton outside a \(Z = 50\) closed shell. Note that no isomeric level is observed in \(^{111}\)Sb, in agreement with the behavior of the 2d 5/2, 1g 7/2 and 3s 1/2 neutron quasi-particle energies in this region\(^6,7\)). The second excited state of the even Sn isotopes around 2200 keV has most likely \(J^\pi = 4^+\), the cross over transition to the ground state not having been observed. In this case, \(^{110}\)Sb, \(^{112}\)Sb and \(^{114}\)Sb have all \(J^\pi = 3^+\), allowed beta decay being observed to \(4^+\) and \(2^+\) levels.

In order to complete the decay scheme of the Sb isotopes, \(\beta\)-\(\gamma\), conversion electrons and beta spectrum endpoints measurements are needed. Experiments along these lines are in progress in our laboratory.

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a) $E_\gamma$ spectrum of $^{110,112,114}$Sb from 1170 to 1350 keV in 16 slices of 10 sec.

b) Sum over the first 80 sec

c) Sum over the next 80 sec

Fig. 1
Fig. 2 Half-life curves for γ transitions in $^{110,111,112,114}\text{Sb}$.
Fig. 3  Partial decay scheme of $^{111}$Sb and $^{113}$Sb.
Relative transition intensities are indicated.
Fig. 4 Partial decay scheme of $^{110}_{\text{Sb}}$, $^{112}_{\text{Sb}}$ and $^{114}_{\text{Sb}}$. Relative transition intensities are indicated.
LOW-LYING STATES IN ODD-MASS POLONIUM ISOTOPES

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Mass separated sources produced at the ISOLDE facility at CERN
have been used for studies of low-lying excited states in
$^{201,203,205,207,209}$Po. These levels were reached both through the
electron capture decay of astatine and the alpha decay of radon isotopes.
Similarities to lead isotopes are expected to show up in polonium, since
only one proton pair is added to the $Z = 82$ proton core in lead. Thus
one expects the $3p_{1/2}^2$, $2f_{5/2}$, $3p_{3/2}$ and $1i_{13/2}$
eutron quasi-particle states to occur at low excitation energies. States due to the coupling of these
states to a $2^+$ phonon should also be found.

When this investigation started, the polonium ground-state spins
were known in all cases$^1)$. Isomerism in $^{205,207}$Po had been reported by
Hargrove and Martin$^2)$. These isomers were interpreted as being due to the $1i_{13/2}$
neutron hole state decaying through a $^{13/2}_9 + (M2)^{9/2} - (E2)^{9/2} -$ cascade.
In $^{207}$Po, low-lying levels had been studied$^3$,$^4$ in alpha decay from $^{211}$Rn.
Some work concerning levels in $^{203}$Po had also been done by Stoner$^5$). His
conclusions regarding the level scheme are, however, not in agreement
with our findings.

Isomerism in $^{201,199}$Po was established by Tielsch-Cassel$^6$) and by
Brun et al.$^7$). The corresponding isomers in $^{197,195}$Po were studied by
Siivola$^8$).

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Since then, two papers by a Dubna group\textsuperscript{9,10} have appeared. They observed isomerism in $^{203}$Po and proposed a partial decay scheme of $^{205}$At. Hopke et al.\textsuperscript{11} observed only the strongest (719 keV) transition in the decay of $^{205}$At, and prematurely concluded that the state at 719 keV is fed by a beta transition with a log ft value of 6.1.

In our work, alpha-gamma coincidence measurements were carried out in order to obtain the excitation energies for levels in $^{201,203,205}$Po. Alpha spectra in coincidence with gamma-rays ($E_\gamma \geq 5$ keV) were recorded with a time resolution of $\sim 100$ nsec. As the alpha branches to the excited states are of the order $10^{-3}$ - $10^{-4}$, the contribution from chance coincidences was considerable. Therefore, two alpha spectra were recorded simultaneously in two different parts of the memory of a multichannel analyser, one spectrum corresponding to real (and chance) coincidences, the other being due to chance coincidences only.

Figure 1 shows the real and chance spectra for $^{207}$Rn $\rightarrow$ $^{203}$Po. Two alpha transitions to excited states are readily observed. Levels in $^{201}$Po and $^{205}$Po were studied in the same way. Results from these measurements are included in Fig. 7 (states marked \textbullet): excited levels at 145 keV in $^{205}$Po, at 60 keV and 130 keV in $^{203}$Po, and at 140 keV in $^{201}$Po.

Levels in $^{203,205,207,209}$Po were also studied in the isobaric decay from astatine using Ge(Li) and Si(Li) detectors and a double focusing magnetic beta spectrometer. Gamma-gamma coincidence measurements with two Ge(Li) detectors were also carried out.
Fig. 2 Simplified level scheme of $^{209}$Po.

Fig. 3 Simplified level scheme of $^{207}$Po.
Figures 2, 3, and 4 show simplified level schemes of $^{209,207,205}$Po, respectively. Several weakly populated levels have been omitted for simplicity. The details of the level schemes will not be discussed here. Only a few comments will be made.

It is interesting to observe the trend of the feeding of levels in the isobaric decay from astatine. In $^{209}$Po, a $\frac{9}{2}^+$ state at 2311 keV takes a considerable part (~80%) of the feeding. It is very reasonable to explain this state as being one with a neutron excited to the $2g_{\frac{9}{2}}$ level above the $N = 126$ gap. This state has been identified in stripping experiments and occurs at approximately 2.5 MeV in $^{205,207}$Pb. The beta decay is then understood as a $P(h_{\frac{9}{2}}) \rightarrow N(g_{\frac{1}{2}})$ transition. In $^{207}$Po the strength is spread over several positive parity states, and in $^{205}$Po no such state has been established. If the strength was still concentrated on only a few levels, these would easily be found as the higher $Q$-value (4.4 MeV for $^{205}$At as compared to 3.7 MeV for $^{207}$At) would enhance these transitions. One may thus conclude that the beta strength is spread over many states, due to the increased complexity of the excitation spectrum.
This conclusion is in agreement with (and, admittedly, inspired by) the work of Hansen et al.\textsuperscript{13} on gross beta decay. They find significant structure for $^{209,207}$At but a rather constant smear for $^{205}$At.

In $^{203}$Po we found a 641 keV M4 transition (see Fig. 5) which most probably is the isomeric transition from the $^{13/2}_+^+$ level to the $^{5/2}_-^-$ ground state. The half-life of the isomer, 1.2 min, was reported by Morek et al.\textsuperscript{10}.

Figure 6 contains our information on $^{201}$Po obtained from the $^{201}$Po decay and from the $^{205}$Rn alpha decay.

The 9 min isomer in $^{201}$Po suggested\textsuperscript{6,7} as being the $^{13/2}_+^+$ single neutron state was found to decay with a $417 \pm 1$ keV M4 transition. From a closed energy cycle including $^{201,201}$Po and $^{197,197}$Pb one can deduce an excitation energy of $423 \pm 3$ keV if data from Ref. 1 are used. Since the ground-state spin is $^{7/2}_-^-$, it is reasonable to assume the existence of a state at low energy ($< 10$ keV) with spin $^{5/2}_-^-$, although we are not able to detect any transition in this energy range with our present equipment.

Figure 7 summarizes our knowledge of quasi-particle levels and also for the $^{9/2}_-^-$ state which should have ground state ($f^{5/2}_-^-$\textsuperscript{1} + phonon character. For comparison, we include the corresponding excited states in the lead isotones. One notices a strong resemblance, in particular for the
\[ N = 117 \ 119 \ 121 \ 123 \ 125 \]
\[ \text{Po} \quad \text{Pb} \]
\[ 1 \text{ MeV} \]
\[ 13^+ / 2 \]
\[ 19^+ / 2 \ 9^- / 2 \]
\[ 13^+ / 2 \]
\[ 13^+ / 2 \ 9^- / 2 \ 3^- / 2 \]
\[ 13^+ / 2 \ 9^- / 2 \ 5^- / 2 \]
\[ \left( 1 / 2 \right) \left( 5 / 2 \right) \]
\[ \star \star \ 3^- / 2 \ 3^- / 2 \]
\[ 0 \ 3^- / 2 \left( 5 / 2 \right) \ 5^- / 2 \ 5^- / 2 \ 1^- / 2 \ 1^- / 2 \]

Fig. 7 Level systematics for \(^{201,203,205,207,209}\text{Po}\) and \(^{199,201,203,205,207}\text{Pb}\). Levels marked with an asterisk (\(\star\)) were established in \(\alpha-\gamma\) coincidence measurements.

\(i_{13/2}\) states. The relationship between lead and polonium isotones is sufficiently close to make the level order and energy spacings similar at low excitation energies. Calculations performed for the lead region, well reproduce general trends for states being essentially of one-quasi-particle character (see, for example, Ref. 6 and references quoted therein). The importance of systematic experimental information about such states as a basis for calculations is obvious. Arvieu et al. \(^{14}\) also point out the importance of more complex states for testing the validity of theory. Unfortunately, the experimental results are too meagre at present and need further investigation on this point.
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PROPERTIES OF THE LOW-LYING LEVELS IN THE EVEN Pt AND Os NUCLEI

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In this contribution we present the status1-4) of the analysis of our results from the study of low-lying level characteristics in a series of even neutron-deficient platinum and osmium nuclei (182 ≤ A ≤ 192).

This work has been particularly favoured by the technique used for the separation of these isotopes. A molten lead target is bombarded by a 600 MeV proton beam, and the Hg nuclei produced are separated by the ISOLDE on-line facility at CERN5). The nuclear spectroscopic equipment at our disposal consisted of Ge(Li) and Si(Li) detectors, beta spectrometers, fast and delayed coincidence systems, including a double-lens electron-electron magnetic spectrometer, a PDP-9 computer, etc.

The availability of nuclei far from the β stability line is very favourable in the search for systematic trends in the properties of nuclei between the pure rotational and magic regions. The study of this transitional region is already the subject of several theoretical works6-8) intended to extend the basic models and to establish their common features,

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as will be mentioned in the discussion. It should be emphasized that some of our interpretations are tentative. Some of these will be checked by angular correlation measurements. Nevertheless it seems useful to present even such preliminary data at this stage.

We have particularly searched for:

- the existence or non-existence of various bands previously proposed;9-12);

- the inversion of the relative positions of collective low-energy levels;

- the variation in the energy of levels of the same characteristics;

- major variations in the value of B(E2) branching ratios.

Some results obtained by other groups13-25) have been used or confirmed, in particular those from heavy ion reactions concerning the fundamental band.

Decay schemes have been established mainly on the basis of single gamma spectra, gamma-gamma coincidences [Ge(Li)-Ge(Li)], and conversion electron data, and by using energy sum relations. In the interpretation of internal conversion data, highly converted transitions between even parity states with the same angular momentum have been assumed to be E0 or E0+E2+M1. We have also measured some half-lives of the 2^+_1 levels (Table 1). Values for B(E2, 2^+_1 → 0^+_1) have been deduced, and, when possible, compared with those calculated by Kumar and Baranger6b). Some results were also obtained at Dubna26) for ^182Os and ^184Os (Figs. 1 and 2).

In Fig. 3 we present the results for low-lying levels of Pt nuclei. A similar study is in progress27) on levels in ^182,184,186Os. From Fig. 3 and Table 2 we can recognize the smooth variation in the energies for the fundamental band [as previously observed for some of these nuclides14-16)], and also the almost constant energy of the 2^+_2 level23,24). Some new features, which are perhaps significant, should be pointed out:

- The collective character is apparent for the 2^+_1 level from the measured B(E2, 2^+_1 → 0^+_1) in the Os and Pt nuclei (Figs. 1 and 2).

- The 0^+_2 state seems to show a very low minimum energy at A=186 and an apparent maximum for A=19413,17-19,27).
- 1033 -

- Some $2^+$ levels (and in some cases $4^+$) de-excite to the $2_1^+$ and $4_1^+$ levels of the fundamental band by E0+E2+(M1) transitions (with a strong E0 component) and preferentially to the $0_2^+$ (see Table 3). This is a characteristic of a $\beta$ band in deformed nuclei.

- The behaviour of the $3_1^+$ level and of the $4_2^+$ (when known) is similar to the behaviour of the $2_2^+$ level. This, and the fact that the $4_2^+$ and $3_1^+$ de-excite preferentially to the $2_2^+$, is characteristic of a $\gamma$ band in deformed nuclei.

- The relative branching ratio $B(E2, 2_2^+ \rightarrow 0_1^+)/B(E2, 2_2^+ \rightarrow 2_1^+)$ increases from $\sim 10^{-5}$ to 0.1 for $196 \gtrsim A \gtrsim 184$ (see Fig. 4). [The vibrational value is 0.0, the rotational value 0.7 (Alaga rule).]

- If one applies the usual terms of the vibrational model, we notice that the two- and three-phonon multiplets are less mixed for $^{186}$Pt than for other masses. (This is the mass showing the minimum energy for the $0_2^+$ level.) We also notice that the $0_2^+$ and $4_1^+$ levels are closer together than $0_2^+$ and $2_1^+$ levels. Present experimental results$^{14,18-25}$ do not allow us to determine the nature of the second $0^+$ levels. (We have some candidates for $0_1^+$ levels which are still not placed in the decay schemes.) We have observed many highly converted transitions (E0+E2+M1?) which have not been placed in the schemes. These possibly result from strong coupling between $\beta$ and $\gamma$ bands, or of two-phonon $\beta-\gamma$ states.

- Finally the $3^-$ level, which is characteristic of octupole vibration, has a rather constant energy ($\approx 1400$ keV).

Three principal conclusions may be drawn:

- Low-lying levels in Pt nuclei seem to be organized in bands, and the $\beta$ band-head is strongly mixed with the ground band.

- The energy of the levels of the proposed $\beta$ and $\gamma$ bands present minima for masses $A=186-190$.

- A marked variation of the energy and of the relative positions of the levels takes place between masses 186-190. As these two last points refer to the same masses, they are perhaps correlated.
Among the theoretical interpretations proposed for the transitional nuclei, we would like to mention three which propose explanations for some of the experimental results we have found. In the model of Kumar and Baranger\textsuperscript{6}) (microscopic calculation of the Bohr-Mottelson Hamiltonian), the inversion of the position of $2^+_2$ and $4^+_1$ levels is related to a shape modification of the nucleus (prolate-oblate) and to the existence of anharmonic terms for the kinetic energy. The minimum for the $0^+_2$ level seems to be predicted\textsuperscript{6c}) for $^{186}$Pt, and also the general trend of the low energy (positive parity) levels.

In the Greiner model\textsuperscript{7}) (another approach to the Bohr-Mottelson Hamiltonian calculation) one can also find bands interacting strongly with each other, but here the crossing $2^+_2-4^+_1$ is less significant with respect to a shape modification (prolate-oblate), and the $0^+_2$ state is a three-phonon state.

In the model described by Holzwarth\textsuperscript{8}), the minimum of the $0^+_2$ state appears and is related to the anharmonic terms of the Hamiltonian expressed in boson terms (generator coordinate method).

The experimental situation as found in Pt nuclei is also found in Os nuclei\textsuperscript{6,13}) and we think there exist similar trends in Te, Pd, Gd, Sm, and Ba nuclei\textsuperscript{27}).
REFERENCES


8) G. Holzwart (to be published).


27) French-CERN Isolde Collaboration (to be published).
### Table 1

Half-life of the first $2^+$ in Os$^{182,184}$ and Pt$^{184,186,188,190}$

<table>
<thead>
<tr>
<th>Energy level (in keV)</th>
<th>Half-life B(E2 $2_1^+ \to 0_1^+$) exp. in $c^2cm^4 \times 10^{-8}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Os$^{182}$ 119.8</td>
<td>0.95 $\pm$ 0.10 $10^{-9}$ sec</td>
</tr>
<tr>
<td>Os$^{184}$ 127</td>
<td>1.1 $\pm$ 0.05 $10^{-9}$ sec</td>
</tr>
<tr>
<td>Pt$^{184}$ 162.4</td>
<td>360 $\pm$ 12 $10^{-12}$ sec</td>
</tr>
<tr>
<td>Pt$^{186}$ 191.5</td>
<td>260 $\pm$ 10 $10^{-12}$ sec</td>
</tr>
<tr>
<td>Pt$^{188}$ 265.4</td>
<td>72 $\pm$ 13 $10^{-12}$ sec</td>
</tr>
<tr>
<td>Pt$^{190}$ 296</td>
<td>45 $\pm$ 15 $10^{-12}$ sec</td>
</tr>
<tr>
<td></td>
<td>0.666 $\pm$ 0.070 (*)</td>
</tr>
<tr>
<td></td>
<td>0.671 $\pm$ 0.040 (*)</td>
</tr>
<tr>
<td></td>
<td>0.765 $\pm$ 0.038</td>
</tr>
<tr>
<td></td>
<td>0.590 $\pm$ 0.025</td>
</tr>
<tr>
<td></td>
<td>0.516 $\pm$ 0.103</td>
</tr>
<tr>
<td></td>
<td>0.502 $\pm$ 0.150</td>
</tr>
</tbody>
</table>

Table 2

Pt nuclei experimental levels energy values (in keV)

<table>
<thead>
<tr>
<th>Level</th>
<th>$^{182}_{\text{Pt}}$</th>
<th>$^{184}_{\text{Pt}}$</th>
<th>$^{186}_{\text{Pt}}$</th>
<th>$^{188}_{\text{Pt}}$</th>
<th>$^{190}_{\text{Pt}}$</th>
<th>$^{192}_{\text{Pt}}$</th>
<th>$^{194}_{\text{Pt}}$</th>
<th>$^{196}_{\text{Pt}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2^+_1$</td>
<td>154.9</td>
<td>162.4</td>
<td>191.5</td>
<td>265.4</td>
<td>296</td>
<td>316.5</td>
<td>328.5</td>
<td>355.7</td>
</tr>
<tr>
<td>$4^+_1$</td>
<td>418.7</td>
<td>434.6</td>
<td>490.1</td>
<td>670.6</td>
<td>737.5</td>
<td>784.5</td>
<td>811.2</td>
<td>877.3</td>
</tr>
<tr>
<td>$6^+_1$</td>
<td>771</td>
<td>797.1</td>
<td>877.2</td>
<td>1185</td>
<td>1290</td>
<td>1390</td>
<td>1412</td>
<td>1510</td>
</tr>
<tr>
<td>$2^+_2$</td>
<td>-</td>
<td>648.3</td>
<td>607.3</td>
<td>605.9</td>
<td>597.8</td>
<td>612.4</td>
<td>622.1</td>
<td>683.7</td>
</tr>
<tr>
<td>$3^+_1$</td>
<td>-</td>
<td>939.6</td>
<td>956.9</td>
<td>936.1</td>
<td>917.2</td>
<td>920.9</td>
<td>922.8</td>
<td>1015</td>
</tr>
<tr>
<td>$0^+_2$</td>
<td>-</td>
<td>(492.7)</td>
<td>(471.6)</td>
<td>798.2</td>
<td>921.6</td>
<td>1195.0</td>
<td>1267</td>
<td>1135</td>
</tr>
<tr>
<td>$2^+_3$</td>
<td>-</td>
<td>844.3</td>
<td>798.4</td>
<td>1115</td>
<td>1203.4</td>
<td>1576.6</td>
<td>1623</td>
<td>-</td>
</tr>
<tr>
<td>$4^+_2$</td>
<td>-</td>
<td>1236.8</td>
<td>1222.4</td>
<td>(1625)</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>$3^-_1$</td>
<td>-</td>
<td>-</td>
<td>1407.8</td>
<td>1349.6</td>
<td>1353.7</td>
<td>1378</td>
<td>1432</td>
<td>1447</td>
</tr>
<tr>
<td>$1^-_1$</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>1775.3</td>
<td>1737.3</td>
<td>1739.4</td>
<td>1803</td>
<td>-</td>
</tr>
<tr>
<td>-</td>
<td>(a)</td>
<td>-</td>
<td>(a)</td>
<td>(a)</td>
<td>(a)</td>
<td>(b)</td>
<td>(c)</td>
<td>(d)</td>
</tr>
</tbody>
</table>

a) This work.


<table>
<thead>
<tr>
<th>Levels $i_f$</th>
<th>$\beta^{184}$</th>
<th>$\beta^{186}$</th>
<th>$\beta^{188}$</th>
<th>$\beta^{190}$</th>
<th>$\beta^{192}$</th>
<th>$\beta^{194}$</th>
<th>$\beta^{196}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2^+_1 \rightarrow 0^+_1$</td>
<td>$\geq 10$</td>
<td>$11 \pm 3$</td>
<td>$3.3 \pm 0.2$</td>
<td>$1.24 \pm 0.05$</td>
<td>$0.52 \pm 0.02$</td>
<td>$0.21 \pm 0.05$</td>
<td>$7 \pm 2 \times 10^{-6}$</td>
</tr>
<tr>
<td>$2^+_1 \rightarrow 2^+_1$</td>
<td>$100 \pm 20$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 3$</td>
<td>$100 \pm 3$</td>
<td>$100 \pm 2$</td>
<td>$100 \pm 3$</td>
<td>$100 \pm 3$</td>
</tr>
<tr>
<td>$3^+_1 \rightarrow 2^+_1$</td>
<td>$\geq 9.7$</td>
<td>$15 \pm 2$</td>
<td>$4.5 \pm 0.2$</td>
<td>$2.0 \pm 0.2$</td>
<td>$1.7 \pm 0.10$</td>
<td>$0.5$</td>
<td>$0.22$</td>
</tr>
<tr>
<td>$3^+_1 \rightarrow 2^+_2$</td>
<td>$100 \pm 30$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 3$</td>
<td>$100 \pm 3$</td>
<td>$100 \pm 2$</td>
<td>$100$</td>
<td>$100$</td>
</tr>
<tr>
<td>$2^+_2 \rightarrow 0^+_1$</td>
<td>$5.5 \pm 2$</td>
<td>$7.9 \pm 0.8$</td>
<td>$0.8 \pm 0.1$</td>
<td>$\approx 0.03$</td>
<td>$\approx 1.7$</td>
<td>$\approx 1.7$</td>
<td>$\approx 1.7$</td>
</tr>
<tr>
<td>$2^+_2 \rightarrow 0^+_2$</td>
<td>$100 \pm 30$</td>
<td>$100 \pm 10$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 10$</td>
<td>$100 \pm 10$</td>
<td>$100 \pm 10$</td>
<td>$100 \pm 10$</td>
</tr>
<tr>
<td>$2^+_2 \rightarrow 2^+_1$</td>
<td>$- \leq 6$</td>
<td>$0.6 \pm 0.1$</td>
<td>$0.6 \pm 0.1$</td>
<td>$0.3 \pm 0.1$</td>
<td>$0.3 \pm 0.1$</td>
<td>$0.3 \pm 0.1$</td>
<td>$0.3 \pm 0.1$</td>
</tr>
<tr>
<td>$2^+_2 \rightarrow 2^+_2$</td>
<td>$- -$</td>
<td>$4.2 \pm 0.3$</td>
<td>$6.1 \pm 1$</td>
<td>$6.1 \pm 1$</td>
<td>$6.1 \pm 1$</td>
<td>$6.1 \pm 1$</td>
<td>$6.1 \pm 1$</td>
</tr>
<tr>
<td>$4^+_1 \rightarrow 2^+_2$</td>
<td>$- 8 \pm 1$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
</tr>
<tr>
<td>$4^+_1 \rightarrow 2^+_1$</td>
<td>$- 2.1 \pm 0.2$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
</tr>
<tr>
<td>$4^+_1 \rightarrow 2^+_3$</td>
<td>$- 100 \pm 10$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
<td>$- -$</td>
</tr>
<tr>
<td>$0^+_2 \rightarrow 2^+_1$</td>
<td>$- -$</td>
<td>$11 \pm 1$</td>
<td>$3.9 \pm 0.4$</td>
<td>$3.9 \pm 0.4$</td>
<td>$7.9 \pm 0.8$</td>
<td>$7.9 \pm 0.8$</td>
<td>$14 \pm 2$</td>
</tr>
<tr>
<td>$0^+_2 \rightarrow 2^+_1$</td>
<td>$- -$</td>
<td>$100 \pm 10$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 5$</td>
<td>$100 \pm 5$</td>
</tr>
</tbody>
</table>

(a) This work.


Fig. 1: $B(E2; 2^+ \rightarrow 0^+)$ exp. values for Pr$^{184,186,188,190}$. 

$B(E2; 2^+ \rightarrow 0^+)$ for PT NUCLEI

- exp. points (this work)
- other exp. points
- KUMAR values
$B(E2 \, 2^+_1 \rightarrow 0^+_1)$ for OS NUCLEI

- exp. points (this work)
- other exp. points
- KUMAR values

Fig. 2: $B(E2 \, 2^+_1 \rightarrow 0^+_1)$ exp. values for Os$^{182,184}$. 
Fig. 3: Energy levels for Pb\textsuperscript{182+186} nuclei.

LEVELS IN PT NUCLEI (182\leq A \leq 196)
Fig. 4 : $\frac{B(E2 \, 2^+ \rightarrow 0^+)}{B(E2 \, 2^+ \rightarrow 2^+)}$ exp. values for Pt nuclei.
ISOMERISM IN ODD AND ODD-ODD NUCLEI WITH MASS NUMBER 185 ≤ A ≤ 191

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1. INTRODUCTION

During last years, substantial progress was accomplished in the elaboration of the theoretical methods which can be applied to the description of the transitional even nuclei. The situation is much less satisfactory when odd-A nuclei are concerned, where the experimental data must be still compared with models which were constructed for strongly deformed or spherical nuclei. Although the imperfection of this kind of comparison for transitional nuclei is well known, it is certainly interesting to see how far the "extreme" models can be extrapolated.

The first important step in the study of odd-A nuclei is the location of the single-particle levels. These yield information on the average field of the nucleus. Experimentally, when radioactivity studies are concerned, one of the most unambiguous methods of locating some single-particle levels is the study of isomerism and the determination of log ft values.

In this paper we present some results concerning the study of isomeric states in a number of odd-A and doubly odd nuclei in the Pt region. In some cases these results are correlated with decay scheme investigations. Due to the experimental method used, only isomers with half-lives longer

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***) On leave from the Institute for Nuclear Research, Swierk, Warsaw, Poland (present address).
than a few seconds or in the nanosecond to microsecond range could be
detected. The assignment of the Nilsson model characteristics is attempted
for some of the observed excited states in the nuclei investigated.

2. **ISOMERS WITH HALF-LIFE LONGER THAN ≈ 10 sec**

The on-line and off-line measurements for this range of half-lives
were done using multispectrum analysis on gamma or conversion electron
spectra [further experimental details have been published elsewhere⁴).]
In some cases a chemical separation of the mass separated samples was
performed. A summary of these data is presented in Table 1, which also
includes data noted elsewhere⁵⁻⁷). Some comments and examples are dis-
cussed below.

2.1 **Isomerism in odd-A nuclei**

Heavier isotopes of mercury (A ≥ 193) possess isomeric states which
are currently assigned to the close-lying i₁₃/₂ and p₁/₂ or p₃/₂ spherical
model orbits⁸⁻¹⁰). For lighter isotopes of mercury the many different
orbitals deduced from Nilsson diagram suggest the existence of isomerism
(cf. Fig. 3), assuming the existence of deformation.

In ¹⁸⁵Hg, three different half-lives seem to be present. The A=185
chain, however, possess a rather strong α branching⁴) and at present we
cannot completely exclude that some of the half-lives attributed to the
185 chain belong to the 181 one. It is interesting to note that isomerism
is also expected to exist in ¹⁸⁹Hg and ¹⁹¹Hg. Our data and those of
Ref. 11 indicate that a high spin, probably 11/2⁻ state in ¹⁸⁹Au and
¹⁹¹Au, is strongly fed from the corresponding decay of the Hg parent.
Therefore, the spin of the 8.7 ± 0.2 min ¹⁸⁹Hg and 50 min ¹⁹¹Hg must be
high. We have found the existence of a 7.7 ± 0.2 min isomer of presumably
low spin in ¹⁸⁹Hg, by observation of the decay of the individual gamma
lines. A similar search in ¹⁹¹Hg was unsuccesful possibly due to a
similarity in half-life to that of the high-spin state.

The isomerism in odd-A gold isotopes may probably be attributed to
the close lying 11/2⁻ and 3/2⁺ states, which can be identified with h₁₁/₂
and d₃/₂ shell-model states or 11/2⁻ (505) and 3/2⁺ (402) Nilsson states.
In this work we have investigated in more detail the decay of the 4.55 min
Table 1
Isomers with $T_{1/2} \geq 10$ sec

<table>
<thead>
<tr>
<th>Nuclide</th>
<th>Half-life</th>
<th>Main transitions observed (keV)</th>
<th>Spin</th>
<th>$\alpha$-decay dataa)</th>
<th>Comments</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{185}$Hg</td>
<td>7.7 ± 0.2 min, 8.7 ± 0.2 min</td>
<td>201,229,238,248,279 and others</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{187}$Hg</td>
<td>2.4 ± 0.2 min, 1.6 ± 0.3 min</td>
<td>103,220,233,271,112,335</td>
<td></td>
<td>2.2 ± 0.3 min</td>
<td>b</td>
</tr>
<tr>
<td>$^{185}$Hg</td>
<td>50 ± 2 sec, 26 ± 3 sec, 155 ± 20 sec</td>
<td>189,222,258,211,292,243,331</td>
<td>48.0 ± 1.5 sec</td>
<td>17 ± 5 sec</td>
<td>c</td>
</tr>
<tr>
<td>$^{185m}$Au</td>
<td>4.55 ± 0.10 min</td>
<td>166,322</td>
<td>11/2+</td>
<td></td>
<td>e</td>
</tr>
<tr>
<td>$^{187}$Au</td>
<td>6.4 ± 1.3 min</td>
<td>181</td>
<td>(1,2)</td>
<td></td>
<td>d</td>
</tr>
<tr>
<td>$^{185m}$Au</td>
<td>2 min</td>
<td>191</td>
<td>(1,2)</td>
<td></td>
<td>f</td>
</tr>
<tr>
<td>$^{185}$Au</td>
<td>4.2 ± 0.3 min, 6.8 ± 0.5 min</td>
<td>311,145</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{185}$Pt</td>
<td>33 ± 5 min, 70 ± 10 min</td>
<td>120,135,197 and others</td>
<td>153 and others</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{186m}$Ir</td>
<td>1.7 h</td>
<td>137,297,986 and others</td>
<td>(2-)</td>
<td></td>
<td>f</td>
</tr>
</tbody>
</table>

b) Relative position of isomers not known.
d) Existence of isomerism is tentative at present.
f) Ground state has higher spin and longer half-life.
isomeric state in $^{189}$Au. The results of this investigation are summarized in the next section, which deals with a nanosecond isomer discovered in the decay of $^{189m}$Au.

2.2 Isomerism in odd-odd nuclei

From our data$^4$ and from data given in the literature$^{12-14}$, it is known that the radioactive decay of odd-odd nuclei in this region very often populate high-spin (up to 6+) states in even-even nuclei. It is the case, for example, in the decay of $^{184}$Au, $^{186}$Au, $^{184}$Ir, and $^{186}$Ir, and they must therefore have rather high spins. In our investigation these odd-odd nuclei are obtained from the decay of even parents, and therefore a "spin-gap" exists between the 0+ state of mother even-nuclei and presumably the ground state of odd-odd nuclei. This "spin-gap" may favour the existence of isomerism. We have found the isomerism in the case of $^{186}$Au and confirmed$^{12}$ it in the case of $^{186}$Ir. The search for isomerism in $^{184}$Au and $^{184}$Ir was unsuccessful. Our studies and those of Zaitseva et al.$^{15}$ indicate the configurations for the ground and excited states of $^{186}$Ir, as shown in Fig. 1.

![Diagram](image)

Fig. 1

3. ISOMERS IN THE NANOSECOND RANGE

At present, only off-line measurements have been performed within this part of our program. Standard nanosecond timing methods were used. A summary of the data obtained is presented in Table 2 which also includes data reported elsewhere$^{16,17}$. Some examples are discussed below.
Table 2
Nanosecond isomers

<table>
<thead>
<tr>
<th>Nuclide</th>
<th>Metastable state half-life (nsec)</th>
<th>Main transitions observed from the decay of metastable states (keV)</th>
<th>Comment</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{185m_1}$Os</td>
<td>$3000 \pm 400$</td>
<td>97</td>
<td>a</td>
</tr>
<tr>
<td>$^{185m_2}$Os</td>
<td>$780 \pm 50$</td>
<td>157 (M1)</td>
<td></td>
</tr>
<tr>
<td>$^{185}$Ir</td>
<td>$2.1 \pm 0.2$</td>
<td>$235 \pm 5$</td>
<td></td>
</tr>
<tr>
<td>$^{187m_1}$Ir</td>
<td>$19 \pm 3$</td>
<td>655 (M1)</td>
<td>b</td>
</tr>
<tr>
<td>$^{187m_2}$Ir</td>
<td>$11.5 \pm 0.3$</td>
<td>106 (E2)</td>
<td></td>
</tr>
<tr>
<td>$^{189m}$Ir</td>
<td>$155 \pm 15$</td>
<td>247 (M1)</td>
<td></td>
</tr>
<tr>
<td>$^{191m}$Ir</td>
<td>$11.5 \pm 0.3$</td>
<td>94 (E2+M1)</td>
<td>c</td>
</tr>
<tr>
<td>$^{189m}$Pt</td>
<td>$4.17 \pm 0.10$</td>
<td>82 (E2+M1)</td>
<td>d</td>
</tr>
<tr>
<td></td>
<td>$464 \pm 25$</td>
<td>166 (E2)</td>
<td></td>
</tr>
</tbody>
</table>

- a) The existence of this isomer is tentative.
- b) The metastable state can belong to the $^{185}$Pt or $A = 181$ chain.
- c) A previous measurement of this half-life was reported in:
- d) A previous measurement of this half-life was reported in:

3.1 Decay of $^{189}$Au and isomerism in $^{189}$Pt

Detailed nuclear spectroscopy of $^{189}$Hg and $^{189}$Au was performed within our program and also elsewhere$^{6,7,11,18-20}$. The gamma-ray spectrum of the 28 min $^{189}$Au ($3/2^+$) is rather complicated, but only two transitions are assigned to the decay of the 4.55 min $^{189m}$Au ($11/2^-$), i.e. 166 keV and 322 keV. By performing $X_K$-gamma coincidence we discovered that the 166 keV transition in $^{189}$Pt is delayed ($T_{1/2} = 464$ nsec). Investigation of the decay of the delayed coincidences has shown that the isomeric
state in $^{189}_{111}$Pt is populated only from the decay of the 4.55 min $^{189m}_{91}$Au although the 166 keV transition is observed from the decay of both gold isomers. Therefore we postulate that some low-energy, still undiscovered, transition is responsible for the isomerism in $^{189}_{111}$Pt. Another result from the delayed coincidence measurements was that the isomer in $^{189}_{111}$Pt is fed not only by EC + $\beta^+$ decay, but also by the 322 keV transition. The simplified decay scheme is presented in Fig. 2. In this figure we indicate that the 166 keV transition goes to the ground state, but with currently available data it can also go to a 6.3 keV or 45.7 keV state as postulated in Refs. 11 and 18. Assuming a total decay energy of $2.9 \pm 0.5$ MeV $^{21}$, the log ft values for the isomeric state and the level depopulated by the 322 keV transition are $4.7 \pm 0.3$ and $5.4 \pm 0.3$, respectively. This indicates an allowed unhindered (au) transition in both cases. The explanation of these low log ft values comes quite naturally if one accepts that $^{189}_{111}$Pt is deformed and that the isomeric state has 9/2$^-$ (505) character, and a rotational level built on it has an energy 322 keV higher. In this case the $\beta$ transition proceeds between the 11/2$^-$ (505) state in $^{189}_{111}$Au and members of the 9/2$^-$ (505) quasi-band. The inertial parameter of this band would be $\hbar^2/2y = 29.2$ keV, as compared to the values 44.4 keV and 49.3 keV, respectively, in $^{188}_{112}$Pt and $^{190}_{98}$Pt for the quasi-ground band. This reflects the well-known fact that an odd-A nucleus has a higher moment of inertia than the neighbouring even nuclei. The ratio of the experimental ft values to this band is $5 \pm 2$ and that calculated from the square of the Clebsch-Gordan coefficients is 5.4.

The predicted Nilsson states $^{22}$) for $\epsilon = 0.15$ are shown in Fig. 3. A compression factor of 2, recommended by Reich and Bunker $^{23}$) for the rare earth nuclei, was used in the preparation of this figure. In the ground state the 111 neutron is placed in the 1/2 (510) orbit for this deformation. The single-particle levels of $^{185}_{76}$W $^{24,25}$) $^{187}_{70}$Os $^{12}$), and
the currently studied $^{189}$Pt are shown in this figure. It is interesting to note that the 1161 keV positive parity (probably $1/2^+$) level, recently identified in $^{186}$Pt $^{11,20}$, can also be simply explained by the Nilsson model as a $1/2$ (651) state. However, this state is probably strongly mixed with

\{$p3/2(402), p11/2(505), n9/2(505)$\} $1/2$,

the three-quasi-particle state, as is indicated by the low log ft (≤ 5.5) to this level$^{11,20}$.

Within this description the character of the state depopulated by the 166 keV transition is, however, not explained. Its spin-parity is probably $5/2^-$ or $7/2^-$. It is tempting to explain the 166 keV E2 transition as a crossover ($I + 2 \rightarrow I$) transition within a rotational band built on a $I = 1/2^-$ or $I = 3/2^-$ single-particle state. The recent, preliminary value for the half-life of this transition$^{26}$ (0.2 - 0.3 nsec) seems to be in good agreement with this supposition. In this case the M1 or E2 (0.5 μsec) low-energy isomeric transition would be K-forbidden, which may explain its hindrance. However, the suspicious lack of a cascade de-excitation ($I + 2 \rightarrow I + 1$) as well as any de-excitation to the members of the second band puts doubt on this simple interpretation. (The interpretation of this level as $7/2$ (503) is even less justified within currently available data.)

3.2 Odd-A Ir nuclei

Our investigations$^{4}$ of the decay scheme of $^{187}$Pt extend the systematics of energy levels in odd-A Ir nuclei$^{27-33}$. Some of the established states in these nuclei are presented in the Fig. 4. More details concerning $^{189-193}$Ir can be found in a paper$^{34}$ presented during this conference. The description of the lowest levels as members of Coriolis mixed 3/2 (402) and 1/2 (400) rotational bands has been proposed$^{31,35}$, although the $1/2^+$ level certainly also has a strong vibrational (K=2) component$^{16,36}$. From the Nilsson model one also
expects the existence of low-lying 1/2 (541), 3/2 (532), 1/2 (660), and 3/2 (651) states. The identification in the present work of a number of negative parity states in $^{187}$Ir can probably be explained by excitation to the first two mentioned orbitals.

The extension of the energy levels systematics to $^{187}$Ir shows that the low-energy positive parity states follow well the trends observed in heavier Ir nuclei. Using the half-life value of the first excited state in $^{187}$Ir (and also our more exact values for $^{189}$Ir and $^{191}$Ir) and recent multipolarity mixing determinations, we confirm the systematic change with $A$ of the M1 and E2 transition probabilities from the 1/2$^+$ to 3/2$^+$ states in these nuclei$^{16}$ [see Fig. 5, based on our data and data quoted elsewhere$^{37-40}$].

The hindrance factor of the M1 part of this transition has a very high value in the case of $^{187}$Ir. The mechanism of this delay may perhaps be attributed to the M1 matrix element cancellation due to Coriolis mixing, $\Delta N = 2$ mixing (expected in this region), or vibrational mixing. Detailed calculations of the mixed wave function in these nuclei are necessary in order to find if this is the reason for the systematic change of the transition probabilities from the 1/2$^+$ to 3/2$^+$ states and for the unusually high retardation (in comparison with the estimate of Ref. 38) of the M1 component in $^{189}$Ir and $^{187}$Ir (respectively, $1.1 \times 10^5$ and $\geq 5.4 \times 10^5$, $S = 2$).

The 155 nsec isomer in $^{187}$Ir apparently has no analogue in heavier odd-mass Ir nuclei. Only one delayed transition (247 keV) was seen to be
associated with this isomer, indicating that its energy is not higher than \( \sim 300 \text{ keV} \) (or that this isomer decays to another much longer-lived isomeric state). The M1 multipolarity of the delayed transition can hardly explain the existence of this isomeric state. We are searching for a low-energy transition to account for the observed half-life of the metastable state.

3.3 Isomers in \( ^{185}\text{Os} \)

The 157 keV (measured multipolarity M1 from our data and Ref. 24) transition was observed to be associated with a 0.7 \( \mu \text{sec} \) isomeric state in \( ^{185}\text{Os} \). As in the case of \( ^{187m2}\text{Ir} \), the existence of this metastable state is difficult to understand on the basis of the observed multipolarity. The same isomer was recently discovered by the \((\alpha,\gamma n)\) reaction\(^{31}\).

We have evidence for a second isomeric state in this nucleus, having a half-life of about 3 \( \mu \text{sec} \). The existence of this isomer is, however, only tentative at present.

4. CONCLUSIONS

We have discovered a number of isomeric states in the \( 185 \leq A \leq 189 \) mass region. In some cases, taking advantage of supplementary information from decay studies, the Nilsson model characteristics were tentatively attributed to the observed isomeric levels.

However, supplementary experimental information or detailed theoretical calculations are necessary to account for the observed unusual retardation of some transition probabilities.
REFERENCES


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NUCLEAR SPECTROSCOPY OF IODINE ISOTOPES WITH A≥134


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Buenos Aires, Argentina.

1. INTRODUCTION

During our exploration of the heavy fission peak of 235U in search of iodine isotopes at the Buenos Aires ISOL xx) we were able to record significant activity from mass numbers 134, 135 and 136.

The available information for 134I at that time 1) 2) 3) 4) included the observation by Erten and Coryell 1) of a 273 keV gamma ray with a half-life of 3.8 ± 0.5 min, and we decided to investigate the nature of this transition and of the associated isomeric state. Erten and Coryell had concluded that this isomeric level belonged to the I level scheme on the basis of their observations of the growth and decay of the activity due to low-lying transitions in 134Xe. Once our measurements had started, Erten, Coryell and Walters 5) suggested that the 273 keV ray itself belonged to I and that it was the transition which depopulated the isomeric level. In a still later report 6) the same authors included a 44.3 keV transition in cascade with the 273 keV ray, which they had detected as having the proper half-life. After our experimental work on the isomeric transition had been completed, our

x) Member of the Scientific Research Career of the Consejo Nacional de Investigaciones Científicas y Técnicas
xx) Detailed description to be published.
attention was called to a paper by Carraz, Blachot, Monnand and Moussa 7) dealing with the study of the 272 keV transition. Our conclusions agree with their experimental results, but also allow us to put the assignment of this transition to the I level-scheme on a much firmer basis. Besides, our on-line electron spectra allowed us a definite determination of the multipolarity of the isomeric transition, as well as of other gamma rays following $^{134}\text{I}$ ground-state decay.

2. EXPERIMENTAL RESULTS

We made a considerable number of measurements, mainly on-line, obtaining single and coincidence spectra, both for gamma rays and for electrons. The detectors used were two Ge(Li) diodes of 35 and 45 cm$^3$ with a resolution of about 2.5 and 3.5 keV FWHM for the 1332 keV $^{60}\text{Co}$ peak. We also used a small Si(Li) detector which had a resolution of about 7 keV for the 976 keV conversion line of $^{207}\text{Bi}$. These detectors were coupled by means of high quality electronics (Nuclear Diodes preamplifiers and Tennelec amplifiers) to a multichannel analyzer system provided with two 4096 channel ADC's and several readout peripherals under control of a small computer. The software for this system was suitably modified for our requirements and in part developed at our laboratory.

In this way the energies and intensities for the most important transitions following the $^{134}\text{I}$ ground-state decay were determined (Table 1) and found to be in reasonable agreement with the data obtained by other groups 2) 4) in their study of chemically separated I, except in the case of some intensities at low energy.
For the energy of the short-lived transition we obtain
\[ E = 271.7 \pm 0.3 \text{ keV} \]
and by following the decay of this ray for several half-lives we also get
\[ T_{1/2} = 3.50 \pm 0.12 \text{ min} \]
for the half-life of the isomeric level.

Combining the intensity data from the gamma and electron spectra and using the well-known value for the conversion coefficient of the 279 keV transition in $^{203}\text{Hg}$ as a standard, we get for the 272 keV ray
\[ \alpha_K = 0.20 \pm 0.03 \]
\[ K/L+... = 2.6 \pm 0.3 \]  \hspace{1cm} (E3)

We thus conclude that the 272 keV gamma ray is the time-setting isomeric transition.

To confirm the suggested placing of this transition as part of the $^{134}\text{I}$ level-scheme, we measured the energy difference between the L- and K-electron conversion lines, obtaining $28.4 \pm 0.2 \text{ keV}$ which agrees well with the expected value of 28.3 keV for I, as compared to 29.5 keV for Xe. This conclusion also agrees with the fact that the 272 keV peak does not appear in an on-line coincidence spectrum gated by the prominent 847 and 884 keV transitions in $^{134}\text{Xe}$.

Further analysis of the on-line spectra obtained with the Si(Li) detector showed that there is conclusive evidence for the existence of a low-lying, highly converted transition in cascade with the 272 keV ray, which might be identical with the 44.3 keV ray reported in ref. 6). Due to experimental difficulties, mainly a very high neutron-induced background, we were unable to obtain more definite data on this transition.
By analyzing the decay of $^{134}$I activity for energies above 280 keV we conclude that there is a significant feeding of Xe levels from the isomeric state in I. Further work on the details of this feeding is in progress.

We next used our set-up to measure a number of conversion coefficients for mass 134; the resultant data are summarized in Table 2.

Our work at other masses has been limited so far to the recording of gamma and electron spectra and the calculation of some internal conversion coefficients. The results of these measurements are given in Tables 3 and 4.

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5) H.N. Erten, C.D. Coryell and W.B. Walters, BAPS 14, 1255 (1969) and private communication


### Table 1

$\gamma$-rays observed in the decay of $^{134}$I

<table>
<thead>
<tr>
<th>$E_\gamma$(keV)</th>
<th>$I_\gamma$</th>
<th>$E_\gamma$(keV)</th>
<th>$I_\gamma$</th>
</tr>
</thead>
<tbody>
<tr>
<td>134.4 ± 0.5</td>
<td>4.59 ± 0.14</td>
<td>948.3 ± 0.2</td>
<td>4.07 ± 0.09</td>
</tr>
<tr>
<td>161.7 ± 0.5</td>
<td>0.28 ± 0.01</td>
<td>966.8 ± 0.3</td>
<td>0.47 ± 0.05</td>
</tr>
<tr>
<td>186.8 ± 0.4</td>
<td>0.86 ± 0.09</td>
<td>974.8 ± 0.2</td>
<td>4.75 ± 0.10</td>
</tr>
<tr>
<td>217.8 ± 0.3</td>
<td>0.20 ± 0.02</td>
<td>1039.8 ± 0.2</td>
<td>1.90 ± 0.06</td>
</tr>
<tr>
<td>235.35 ± 0.07</td>
<td>2.16 ± 0.05</td>
<td>1072.2 ± 0.2</td>
<td>15.52 ± 0.22</td>
</tr>
<tr>
<td>279.1 ± 0.3</td>
<td>0.17 ± 0.09</td>
<td>1103.6 ± 0.2</td>
<td>0.88 ± 0.05</td>
</tr>
<tr>
<td>313.1 ± 0.3</td>
<td>0.14 ± 0.04</td>
<td>1136.3 ± 0.2</td>
<td>10.09 ± 0.17</td>
</tr>
<tr>
<td>319.8 ± 0.2</td>
<td>0.53 ± 0.02</td>
<td>1159.4 ± 0.2</td>
<td>0.37 ± 0.06</td>
</tr>
<tr>
<td>351.4 ± 0.2</td>
<td>0.37 ± 0.02</td>
<td>1191.3 ± 0.2</td>
<td>0.28 ± 0.08</td>
</tr>
<tr>
<td>405.7 ± 0.2</td>
<td>7.66 ± 0.10</td>
<td>1270.3 ± 0.2</td>
<td>0.61 ± 0.05</td>
</tr>
<tr>
<td>411.5 ± 0.2</td>
<td>0.58 ± 0.03</td>
<td>1323.0 ± 0.2</td>
<td>0.19 ± 0.12</td>
</tr>
<tr>
<td>434.1 ± 0.2</td>
<td>4.58 ± 0.07</td>
<td>1337.2 ± 0.2</td>
<td>0.14 ± 0.16</td>
</tr>
<tr>
<td>459.9 ± 0.2</td>
<td>1.38 ± 0.04</td>
<td>1354.4 ± 0.2</td>
<td>0.41 ± 0.07</td>
</tr>
<tr>
<td>489.7 ± 0.2</td>
<td>1.66 ± 0.03</td>
<td>1456.3 ± 0.2</td>
<td>2.10 ± 0.08</td>
</tr>
<tr>
<td>515.1 ± 0.2</td>
<td>2.71 ± 0.05</td>
<td>1470.8 ± 0.2</td>
<td>0.77 ± 0.06</td>
</tr>
<tr>
<td>541.7 ± 0.2</td>
<td>8.57 ± 0.12</td>
<td>1542.7 ± 0.2</td>
<td>0.48 ± 0.08</td>
</tr>
<tr>
<td>566.0 ± 0.2</td>
<td>1.26 ± 0.04</td>
<td>1614.2 ± 0.2</td>
<td>4.25 ± 0.11</td>
</tr>
<tr>
<td>570.8 ± 0.2</td>
<td>0.50 ± 0.15</td>
<td>1629.6 ± 0.3</td>
<td>0.14 ± 0.18</td>
</tr>
<tr>
<td>595.7 ± 0.2</td>
<td>11.62 ± 0.15</td>
<td>1645.1 ± 0.2</td>
<td>0.44 ± 0.07</td>
</tr>
<tr>
<td>622.4 ± 0.2</td>
<td>11.39 ± 0.16</td>
<td>1656.4 ± 0.3</td>
<td>0.16 ± 0.15</td>
</tr>
<tr>
<td>628.6 ± 0.2</td>
<td>2.38 ± 0.06</td>
<td>1741.7 ± 0.2</td>
<td>2.46 ± 0.09</td>
</tr>
<tr>
<td>677.8 ± 0.2</td>
<td>8.22 ± 0.19</td>
<td>1806.6 ± 0.2</td>
<td>5.52 ± 0.14</td>
</tr>
<tr>
<td>731.0 ± 0.2</td>
<td>1.94 ± 0.06</td>
<td>1925.6 ± 0.2</td>
<td>0.18 ± 0.13</td>
</tr>
<tr>
<td>739.3 ± 0.2</td>
<td>0.63 ± 0.04</td>
<td>2020.4 ± 0.2</td>
<td>0.17 ± 0.15</td>
</tr>
<tr>
<td>767.0 ± 0.2</td>
<td>4.30 ± 0.09</td>
<td>2157.1 ± 0.3</td>
<td>0.15 ± 0.16</td>
</tr>
<tr>
<td>817.2 ± 0.2</td>
<td>0.75 ± 0.04</td>
<td>2260.2 ± 0.3</td>
<td>0.12 ± 0.17</td>
</tr>
<tr>
<td>836.7 ± 0.2</td>
<td>0.65 ± 0.05</td>
<td>2310.2 ± 0.3</td>
<td>0.22 ± 0.13</td>
</tr>
<tr>
<td>847.3 ± 0.2</td>
<td>100.0 ± 1.1</td>
<td>2407.6 ± 0.3</td>
<td>0.07 ± 0.29</td>
</tr>
<tr>
<td>857.8 ± 0.2</td>
<td>7.88 ± 0.13</td>
<td>2450.1 ± 0.3</td>
<td>0.07 ± 0.29</td>
</tr>
<tr>
<td>884.2 ± 0.2</td>
<td>68.0 ± 0.8</td>
<td>2466.4 ± 0.3</td>
<td>0.12 ± 0.19</td>
</tr>
<tr>
<td>Source</td>
<td>$E_Y$(keV)</td>
<td>$a_K \times 10^3$</td>
<td>$K/L+$...</td>
</tr>
<tr>
<td>--------</td>
<td>------------</td>
<td>------------------</td>
<td>------------</td>
</tr>
<tr>
<td>$^{134m}$I</td>
<td>271.7</td>
<td>200 ± 30</td>
<td>2.6 ± 0.3</td>
</tr>
<tr>
<td>$^{134g}$I</td>
<td>134.4</td>
<td>260 ± 40</td>
<td></td>
</tr>
<tr>
<td></td>
<td>186.8</td>
<td>170 ± 30</td>
<td></td>
</tr>
<tr>
<td></td>
<td>235.4</td>
<td>70 ± 10</td>
<td>6.1 ± 0.7</td>
</tr>
<tr>
<td></td>
<td>405.7</td>
<td>11 ± 2</td>
<td></td>
</tr>
<tr>
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<td>434.1</td>
<td>15 ± 3</td>
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<td>847.3</td>
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<td>7.5 ± 0.6</td>
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<td></td>
<td>857.8</td>
<td>4 ± 2</td>
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</tr>
<tr>
<td></td>
<td>884.2</td>
<td>1.9 ± 0.3</td>
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</tr>
</tbody>
</table>

a) From the gamma intensity and a limit on the electron intensity.
Table 3

Properties of low-energy transitions following decay of $^{139}$Xe

<table>
<thead>
<tr>
<th>$E_\gamma$(keV)</th>
<th>$I_\gamma$</th>
<th>$a_\alpha \times 10^3$</th>
<th>$K/L+$...</th>
<th>QL</th>
</tr>
</thead>
<tbody>
<tr>
<td>104</td>
<td>6 ± 1</td>
<td>800 ± 100</td>
<td>0.7 ± 0.1</td>
<td>E2</td>
</tr>
<tr>
<td>175</td>
<td>372 ± 10</td>
<td>120 ± 10</td>
<td>1.9 ± 0.1</td>
<td>E2</td>
</tr>
<tr>
<td>219</td>
<td>1100 ± 30</td>
<td>63 ± 5</td>
<td>4.2 ± 0.1</td>
<td>E2</td>
</tr>
<tr>
<td>290</td>
<td>198 ± 6</td>
<td>28 ± 5</td>
<td>3.4 ± 0.5</td>
<td>E2(+M1)</td>
</tr>
<tr>
<td>297</td>
<td>424 ± 10</td>
<td>31 ± 5</td>
<td>4.5 ± 0.5</td>
<td>E2(+M1)</td>
</tr>
<tr>
<td>339</td>
<td>13 ± 1</td>
<td>6 ± 1</td>
<td>3.2 ± 1</td>
<td>E3</td>
</tr>
<tr>
<td>357</td>
<td>11 ± 1</td>
<td>75 ± 10</td>
<td>2.2 ± 1</td>
<td>E3</td>
</tr>
<tr>
<td>394</td>
<td>156 ± 5</td>
<td>10 ± 3</td>
<td>3.3 ± 0.5</td>
<td>E2-M1</td>
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Table 4
Conversion coefficients of transitions observed in the mass region 133-137

<table>
<thead>
<tr>
<th>Source</th>
<th>$E_{\gamma}$(keV)</th>
<th>$\sigma_K \times 10^3$</th>
<th>$K/L+$...</th>
<th>QL</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{133}\text{I}$</td>
<td>529</td>
<td>$8.7 \pm 1.4$</td>
<td></td>
<td>M1</td>
</tr>
<tr>
<td></td>
<td>707</td>
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<td></td>
<td>M1-E2-E1 a)</td>
</tr>
<tr>
<td></td>
<td>875</td>
<td></td>
<td></td>
<td>E2-E1 a)</td>
</tr>
<tr>
<td>$^{133}\text{I} +$</td>
<td>233</td>
<td>$7200 \pm 1300$</td>
<td>$2.45 \pm 0.15$</td>
<td>M4</td>
</tr>
<tr>
<td>$^{133}\text{mXe}$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{135}\text{I}$</td>
<td>857</td>
<td></td>
<td></td>
<td>E1 a)</td>
</tr>
<tr>
<td>$^{135}\text{I} +$</td>
<td>527</td>
<td>$200 \pm 10$</td>
<td></td>
<td>M4</td>
</tr>
<tr>
<td>$^{135}\text{mXe}$</td>
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<td></td>
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</tr>
<tr>
<td>$^{135}\text{Xe}$</td>
<td>250</td>
<td>$60 \pm 4$</td>
<td>$5.6 \pm 0.6$</td>
<td>M1 + E2</td>
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<tr>
<td>$^{136}\text{I}$</td>
<td>198</td>
<td>$74 \pm 15$</td>
<td>$6 \pm 2$</td>
<td>E2</td>
</tr>
<tr>
<td></td>
<td>382</td>
<td>$19 \pm 5$</td>
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<td>M1-E2</td>
</tr>
<tr>
<td>$^{137}\text{Xe}$</td>
<td>456</td>
<td>$9.8 \pm 0.7$</td>
<td>$4.2 \pm 0.8$</td>
<td>E2</td>
</tr>
</tbody>
</table>

a) From the gamma intensity and a limit on the electron intensity.
Fig. 4: Electron spectrum for mass 139
ENERGY LEVELS IN $^{114,116,118,120,122}$Cd
POPULATED IN THE DECAY OF AG ISOTOPES

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Studsvik, Nyköping, Sweden.

Abstract
The decays of Ag isotopes to even Cd isotopes with $114 \leq A \leq 122$ have been investigated. The following levels are populated (energies in keV):

$^{114}$Ag $\rightarrow$ $^{114}$Cd ($T_{1/2} < 10 \text{ s}$): 558 ($2^+$), 1134 ($0^+$).

$^{116}$Ag $\rightarrow$ $^{116}$Cd ($T_{1/2} = 3 \text{ min}$): 513.5 ($2^+$), 1154.0 ($0^+$), 1212.8 ($2^+$), 1922.

$^{116}$Ag $\rightarrow$ $^{116}$Cd ($T_{1/2} = 10 \text{ s}$): 513.5 ($2^+$), 1219 ($4^+$), 2026.

$^{118}$Ag $\rightarrow$ $^{118}$Cd ($T_{1/2} = 4.5 \text{ s}$): 487.1 ($2^+$), 1165, 1269 ($2^+$).

$^{120}$Ag $\rightarrow$ $^{120}$Cd ($T_{1/2} = 1.3 \text{ s}$): 505.5 ($2^+$), 1203, 1324 ($2^+$).

$^{122}$Ag $\rightarrow$ $^{122}$Cd (short): 569 ($2^+$) (tentative).

1. Introduction
The even Cd nuclei with mass-number $\leq 114$ have been extensively studied 1), and some data exist also on $^{116}$Cd 2, 3). The lowest excited states in these nuclei are generally interpreted in terms of quadrupole vibrations of a spherical nucleus. However, neutron-rich nuclei with $Z$ around 44 are expected to be deformed 4, 5). From this point of view it is of interest to investigate Cd nuclei around mass-number 120, which can be expected to be on the border between this deformed region and the spherical region around $Z = 50$.

The present paper is a preliminary report of a study of the levels in $^{118}$Cd and $^{120}$Cd from the decay of $^{118}$Ag and $^{120}$Ag. In addition we have studied the little known decays of $^{114}$Ag $\rightarrow$ $^{114}$Cd and $^{116}$Ag $\rightarrow$ $^{116}$Cd. Tentative evidence was also found for the first excited state in $^{122}$Cd.

2. Experiment
The Ag activities were produced as fission products and mass-separated in the OSIRIS facility 6) at Studsvik. In this arrangement a
uranium target is placed within the ion-source of the mass-separator close to the reactor core. The fission products are thermalized in graphite, which has a temperature of about 1500°C, thereby allowing the products to diffuse and evaporate rapidly. After mass-separation the fission products of the selected mass-number were collected on a movable aluminium tape for rapid transportation to a position 11 cm away from the ion beam, where the measurements were made. Gamma-ray spectra were measured with a 30 cc Ge(Li)-detector connected to a Nuclear Data 3300 4k analyzer. By varying the time of collection and measurement spectra associated with different half-lives were enhanced. The element(s) associated with each half-life was determined from similar measurements of X-rays and conversion electrons using Si(Li)-detectors. For $^{116}\text{Cd}$ and $^{120}\text{Cd}$ gamma-gamma coincidence measurements were performed with 25 cc and 30 cc Ge(Li) coaxial detectors using the digital gate system of the ND 3300 analyzer.

3. Results

The measurements are not yet complete, and in this paper only some preliminary results will be given. Table 1 gives the energies and intensities of the strongest gamma-rays in each of the decays studied. For $^{122}\text{Cd}$ only a tentative assignment of one gamma-ray could be made. This was based on the observation of a half-life of the gamma-ray, which is shorter than that of the decay of $^{122}\text{In} \rightarrow ^{122}\text{Sn}$, and the energy, which is too large to be likely to be a transition in $^{122}\text{In}$ but fits into the systematics of $2^+ \rightarrow 0^+$ transitions in even Cd nuclei. Tentative level schemes deduced from the strongest transitions are shown in fig. 1. In $^{114}\text{Cd}$ the two strongest transitions 558 and 576 keV were identified also in the (n,γ) reaction 7). Several additional transitions in $^{114}\text{Cd}$ were found, which have to await a coincidence experiment to be localized.

The energy of the strongest transition in $^{116}\text{Cd}$ agrees well with the energy of the first excited state earlier observed $^{2,3}$) in $^{116}\text{Cd}$. The remaining part of the level scheme of fig. 1 rests mainly on observed coincidence relationships. The two different half-lives found for $^{116}\text{Ag}$ were found mainly to feed different levels in $^{116}\text{Cd}$. The assignment of the 1213 keV state as $2^+$ is based on the fact that a cross-over to the
ground state was found and that no $1^+$ or $1^-$ state has been observed in this energy region in lighter Cd isotopes. The tentative $4^+$ assignment of the 1219 keV level is based on the earlier observation of a close doublet of $2^+$ and $4^+$ states at $\sim$1220 keV. The tentative $0^+$ assignment of the 1154 keV state is also based on systematics.

For $^{118}\text{Cd}$ no coincidence measurements have been made so far. The 487 and 1165 keV levels are based on the fact that the 487 and 678 keV transitions are the strongest and second strongest transitions observed. The 1269 keV $2^+$ state is based on the energy-fit of the transitions involved and the systematic occurrence of $2^+$ states in even Cd nuclei. The 129 keV transition observed to follow the 4.5 s decay was found to convert in Ag, implying an isomeric level in $^{118}\text{Ag}$. The fact that we observed no gamma-rays following a half-life different from 4.5 s implies that the ground state in $^{118}\text{Ag}$ either has $T_{1/2} < 4.5$ sec or that it not at all or only very weakly feeds excited states in $^{118}\text{Cd}$.

The levels in $^{120}\text{Cd}$ are based mainly on coincidence relationships. The arguments for the $2^+$ assignment of the 1324 keV state are the same as for the corresponding states in the other Cd nuclei.

As mentioned above the observation of the first excited state in $^{122}\text{Cd}$ is very tentative.

4. Comments

The lowest part of the level schemes of $^{116,118,120}\text{Cd}$ look very similar to the lighter Cd nuclei and give no evidence of any rotational structure. In table 2 we give the ratio of the reduced E2 transition probabilities from the second excited $2^+$ state. This is expected to be small (0) for a vibrational nucleus, while it is large for a rotational nucleus (0.70 for a pure K=2 state). We see that this ratio although it raises a factor of 3 in going from $^{114}\text{Cd}$ to $^{120}\text{Cd}$ still is fairly small, both for $^{118}\text{Cd}$ and $^{120}\text{Cd}$. This indicates that the excitation mechanism in even Cd nuclei is essentially unaltered when the neutron number changes.

We are indebted to the members of the OSIRIS collaboration for making the facility available to us and to Messrs. L Jacobsson and B Johnsson for technical assistance.
Table 1.

Energies and intensities of transitions in even Cd isotopes populated in the decay of Ag isotopes

<table>
<thead>
<tr>
<th>Decay</th>
<th>Gamma-ray energy keV&lt;sup&gt;a)&lt;/sup&gt;</th>
<th>Gamma-ray intensity&lt;sup&gt;b)&lt;/sup&gt;</th>
</tr>
</thead>
<tbody>
<tr>
<td>114&lt;sup&gt;Ag&lt;/sup&gt; → 114&lt;sup&gt;Cd&lt;/sup&gt; (&lt;10 s)</td>
<td>558.27&lt;sup&gt;c)&lt;/sup&gt; 575.93&lt;sup&gt;c)&lt;/sup&gt; 597</td>
<td>100 8.8 4.8</td>
</tr>
<tr>
<td>116&lt;sup&gt;Ag&lt;/sup&gt; → 116&lt;sup&gt;Cd&lt;/sup&gt; (10 s)</td>
<td>513.5 666.2 699.7 705.3 806.8 1029.8 1213.2</td>
<td>100 4.0 14.5 36.5 6.5 14.0 6.3</td>
</tr>
<tr>
<td>116&lt;sup&gt;Ag&lt;/sup&gt; → 116&lt;sup&gt;Cd&lt;/sup&gt; (3 min)</td>
<td>513.5 640.5 699 7 835 1213.2</td>
<td>100 3.4 13 4.7 7.6</td>
</tr>
<tr>
<td>118&lt;sup&gt;Ag&lt;/sup&gt; → 118&lt;sup&gt;Cd&lt;/sup&gt; (4.5 s)</td>
<td>487.1 677.9 770.6 781.8 797.8 1060.0 1269.7</td>
<td>100 43 5.5 5.5 7.9 7.5 3.1</td>
</tr>
<tr>
<td>120&lt;sup&gt;Ag&lt;/sup&gt; → 120&lt;sup&gt;Cd&lt;/sup&gt; (1.3 s)</td>
<td>203.2 505.5 697.6 818.4 1323.5</td>
<td>13.2 100 43 15 10.8</td>
</tr>
</tbody>
</table>

<sup>a)</sup> The uncertainty is estimated to ± 1 keV

<sup>b)</sup> Normalized to 100 for the strongest transition in each decay. The uncertainty is estimated to ± 20% for all values

<sup>c)</sup> From ref. 7)
Table 2.

The ratio of the reduced E2 transition probabilities from the second $2^+$ state in even Cd nuclei.

<table>
<thead>
<tr>
<th>A</th>
<th>$\frac{B (E2; \ 2^{+} \rightarrow 0^{+})}{B (E2; \ 2^{+} \rightarrow 2^{+})}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>114</td>
<td>0.018</td>
</tr>
<tr>
<td>116</td>
<td>0.028</td>
</tr>
<tr>
<td>118</td>
<td>0.050</td>
</tr>
<tr>
<td>120</td>
<td>0.065</td>
</tr>
</tbody>
</table>

a) Assuming the $2^+ \rightarrow 2^+$ transition to be pure E2. The Uncertainty can be estimated to about 25%.

References

1) C M Lederer, J M Hollander and I Perlman, Table of Isotopes, J Wiley and Sons, New York 1967.
5) S G Nilsson and I Ragnarsson, private communication.
Figure 1. Preliminary level schemes of even Cd isotopes as populated in the decay of Ag isotopes. Energy errors are given in parentheses. When no error is given the uncertainty is ±1 keV. The energies in $^{114}$Cd are taken from ref. 7).
DECAY STUDIES OF NEUTRON DEFICIENT Ga AND Zn ISOTOPES

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Abstract:

Spectroscopy is performed on neutron deficient Ga and Zn isotopes. The decay of 31 sec $^{63}$Ga was investigated with Ge(Li) detectors. A decay scheme is presented. Preliminary results are given from the study of the decay of 15 min $^{65}$Ga. The connection of an isotope separator to a cyclotron by means of a pneumatic transport system will be tested with 2.4 min $^{60}$Zn.

Introduction:

In the simple shell model description of odd-mass, neutron deficient Zn nuclei the two protons outside the doubly closed $^{56}$Ni core occupy the 2p 3/2 subshell, while for the extra neutrons the 2p 3/2, 1f 5/2 and 2p 1/2 orbits are available. Comparing the results of a shell model calculation of odd-mass Ni-isotopes with the experimental data indicates that excitations from the 1f 7/2 shell should be taken into account (cf. for instance ref. 1). Moreover, from experimental data it has to be concluded that collective effects play a role too.

Kisslinger et al. $^{2,3}$ gave a description of the low lying levels of odd-mass Ni- and Zn-isotopes. Kisslinger and Kumar $^{3}$ started from a phonon model in which anharmonic effects were incorporated. They took a pairing plus quadrupole force for the phonon-quasi particle interaction. Their results for $^{59}$Ni, $^{61}$Ni, $^{63}$Ni, $^{65}$Zn and $^{67}$Zn show low lying 1/2$^{-}$, 3/2$^{-}$ and 5/2$^{-}$ states, one of them being the ground state.

Weidinger et al. $^{4}$ compared their experimental level scheme of $^{65}$Zn with the results of a calculation in the weak-coupling model.

It can be concluded that for the Zn-isotopes no extensive theoretical work exist, like in the case of Ni and Cu isotopes.

Experimental data on the decay of Ga isotopes to neutron poor odd Zn nuclei were rather poor. For this reason we investigated the decay of $^{63}$Ga and $^{65}$Ga.
Experimental method:

Irradiated foils were transported in 6 sec from the cyclotron to the spectrometers with a pneumatic transport system. No chemical separation was performed. Gamma ray spectra were taken with Ge(Li) detectors of 40 cm$^3$ and 5 cm$^3$ (resolution 4 keV and 2.2 keV at 1.33 MeV). A Si(Li) detector with a resolution better than 1 keV at 13.96 keV was used for measuring the $\gamma$-ray spectrum below 0.12 MeV.

An anti-annihilation set-up was used to suppress the effect of annihilation-in-flight. Energy calibration was performed with standard sources, using energy values given by Marion. The spectra were analysed with a computer program giving energies and relative intensities of the peaks.

Decay of 31 sec $^{63}$Ga

Only the half life of this isotope was known from the work of Nurmia and Fink. We irradiated enriched $^{64}$Zn foils with 28 MeV protons, $^{63}$Ga being produced by a $(p,2n)$ reaction. The $\gamma$-ray spectrum and the results of the analysis are given elsewhere. To find absolute $\beta^+$-intensities corrections for impurities ($^{64}$Ga, $^{63}$Zn and $^{62}$Cu) had to be applied to the intensity of the $\gamma^\pm$. Results of decay curve unfolding were used for determining $\log ft$ values. The resulting decay scheme is presented in fig. 1. Until now no detailed theoretical description of the excited states of $^{63}$Zn has been given. In fig. 2 a comparison is given between the low-

![Decay Scheme for $^{63}$Ga](image)

Fig. 1 - Decay scheme for $^{63}$Ga.

All energies are given in MeV.
lying levels of $^{63}_{30}$Zn$_{33}$, $^{61}_{28}$Ni$_{33}$ and $^{59}_{26}$Fe$_{33}$. (Data for $^{61}$Ni and $^{59}$Fe were taken from reference 9). These nuclides all have a ground state with $J^m = 3/2^-$. The energy of the lowest $5/2^-$ level is much different. In $^{61}$Ni and $^{59}$Fe the lowest $1/2^-$-level lies at approximately the same energy. Further resemblance is poor. Interesting is the known M 1-character of the $5/2^- + 3/2^-$ transition in $^{61}$Ni and $^{63}$Zn. The 0.067 MeV transition in $^{61}$Ni is nearly 100% M 1 9). According to Menti 10) and Birstein et al. 11) the 0.193 MeV transition in $^{63}$Zn is = 90% M 1 + 10% E 2. In the shell model the ground state and the first excited state would have the configuration $[(\nu 2p 3/2)^3 (\nu 1f 5/2)^2, 3/2^-]$, respectively $[(\nu 2p 3/2)^4 (\nu 1f 5/2)^1, 5/2^-]$. The M 1-transition between these states is $\ell$-forbidden. Comparing in the $^{61}$Ni-case the partial M 1 life-time with the single particle value reveals a retardation of the order of 70. For the same type of transition in the Z = 33 nuclides $^{73}$As and $^{75}$As a retardation is also found. Laulainen and McDermott measured the quadrupole moment of $^{63}$Zn 12). The value of +0.29(3)b leads them to the conclusion that collective effects should be incorporated in the description of $^{63}$Zn. Assuming for the retardation of the M 1-part of the 0.193 MeV transition in $^{61}$Zn the same value as found for the comparable transition in $^{61}$Ni gives an enhancement of $\approx 13$ for the E 2-part. This would support the conclusion of Laulainen and McDermott.

Decay of 15 min $^{65}$Ga

This decay was recently studied by Li-Solchoz and Bakhru 13). They produced $^{65}$Ga by irradiating natural copper foils with 28 MeV α particles and were troubled by a lot of contamination in their sources. We obtained
rather clean sources by irradiating enriched $^{64}$Zn-foils with 5 MeV deuterons, $^{65}$Ga being produced by a (d,n) reaction. About 90 $\gamma$-rays are ascribed to this decay on the basis of the results of spectrum multiscaling (for comparison: Li Scholz and Bakhru found 15 $\gamma$-rays). In fig. 3 a preliminary decay scheme is presented. In recent

![Decay Scheme Diagram]

Fig. 3 - Preliminary decay scheme for $^{65}$Ga. All energies are given in MeV. (p,$\alpha$,$\gamma$) work 4) the level at 0.865 MeV appeared to be a doublet. Our Ge(Li) spectra show a broadened peak at 0.866 MeV, indicating that both members of the doublet are also reached (directly or indirectly) in the $\beta^+$-decay of $^{65}$Ga.

**Short-living Zn-isotopes**

The decay of Zn-isotopes will be studied using sources prepared with an isotope separator, which is directly connected to a cyclotron by means of a pneumatic transport system 5). Rabbits carrying irradiated foils are transported from the cyclotron to the separator. After removing (cutting) the foils from the rabbits they are sluiced into the high vacuum system of the isotope separator, which is in full operation. In less than 40 seconds after the end of the irradiation the activated foils are inside
the ion source, where the reaction products are diffused out of the foils. The whole procedure can be run automatically. With this system a study will be made of the decay of 2.4 min $^{60}$Zn. Details about the system will be presented at the International Conference on Electromagnetic Isotope separators and the techniques of their applications at Marburg (Germany).

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   J. Vervier, Nuclear Data Sheets for A = 59, Nucl. Data B2 - 5 - 1 (1968)
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13) A. Li-Scholz and H. Bakhru, Phys. Rev. 168 (1968) 1193
NUCLEAR SPECTROSCOPY OF SOME LIGHT Tl ISOTOPES

Joint Institute for Nuclear Research, Dubna, USSR.

1. INTRODUCTION

In our earlier work on Tl isotopes we have investigated the radiation of $^{188}\text{Tl}$ and $^{190}\text{Tl}$. With the aim of producing neutron deficient thallium isotopes, PbF$_2$ powder was irradiated with 660 MeV protons. The irradiated target was transported by a pneumatic rabbit system in the chemical laboratory. After rapid thermochromatographic separation the $\gamma$-spectrum of thallium isotopes was measured with a Ge/Li/spectrometer. New isotopes $^{188}\text{Tl}$ and $^{190}\text{Tl}$ were identified on the basis of known $\gamma$-transitions between levels of the daughter isotopes. The energy and relative intensity of $\gamma$-lines, which might belong, on the basis of the half-life, to the decays of $^{188}\text{Tl}$ and $^{190}\text{Tl}$ have been also investigated.

The aim of the present study was also the nuclear spectroscopic investigation of the neutron deficient, few-minutes-lived thallium isotopes, but in this case - in order to get stronger sources - the proton irradiation and thermochromatographic thallium separation were performed at the same time and also mass separation was used. As a result of investigations we identified the $^{191}\text{Tl}$ isotope, determined its half-life, measured the $\gamma$-spectrum of $^{190}\text{Tl}$, $^{191}\text{Tl}$, $^{193}\text{Tl}$, and the conversion electron spectrum of $^{193}\text{Tl}$.

2. EXPERIMENTAL PROCEDURE

Neutron deficient thallium isotopes were produced irradiating PbF$_2$ powder in the external beam of the Dubna synchro-cyclotron with protons of 660 MeV energy and $\approx 10^{12}$ p/sec intensity. The schematic drawing of the target system and thermochromatographic apparatus is shown in Fig.1.
The express separation of thallium from the target material was based on the volatility of thallium fluorides. The boiling point of PbF$_2$ is 1293 C$^0$, while that of thallium fluorides substantially lower /655 C$^0$ at TlF, and 927 C$^0$ at TlF$_3$/2/. During irradiation the target was kept at 660 C$^0$. The evaporating – from the target – thallium fluorides were carried by the nitrogen stream into a quartz tube, which was kept at approx. linearly decreasing temperature in the function of distance. According to the experiments, on the tube region of temperature 270 → 170 C$^0$ the Tl was deposited. After finishing irradiation the part of the tube, containing Tl, was cut out and transported into the ion source of isotope separator.

The PbF$_2$ powder target of 490 gr mass was loosely spreaded on shelves inside the oven. The nitrogen stream was purified from humidity, from oxygen, and was heated. The flow rate of the nitrogen stream was ≈100 ml/min. After the thermochromatographic apparatus the nitrogen was purified from the radioactive contamination by acidic, alkaline, NaF and water absorbents.

According to the previous measurements carried out with a Ge/Li/ detector, thallium was deposited on the quartz tube of 270 → 170 C$^0$ temperature region without substantial contamination by other radioactive products.

The thallium ions were separated according to their masses by an 1.5 m radius, 55° sector-field isotope separator. The description of the working basis of the isotope separator is to be found in refs. 3/, 4/, and the plans of the Dubna isotope separator system in 5/.

We transferred the thallium isotopes into the magnetron type ion source with the aid of a pneumatic remote control apparatus. The separated ions were collected on an aluminized plastic tape of ≈ 5 mg/cm$^2$ surface density in the vicinity of the focal plane. The quick removal of the sources from the collector chamber was enabled by an air-lock system.
The main characteristics of the isotope separation were as follows.

/i/ The collection efficiency on the tape of the thallium isotopes, transported into the ion source, was 2.5 - 5% after 5 min separation time, and 13% - after 15 min.

/ii/ The distance between the neighbouring thallium isotopes in the collector plane = 13.5 mm.

/iii/ The mass resolution at half-maximum was ≈1000 for lead mass marker lines, which were used during separation.

The analysis of γ-spectra was performed with 40 cm$^3$ sensitive volume Ge/Li/ detectors. The half-width of γ-lines was 4.5 keV at 609 keV. We used two 4096 channel analysers of AI-4096 /USSR/ type in the measurement. The second /buffer/ analyser served for intermediate storage of the information before its recording on the magnetic tape. The data processing was made with a MINSK-22 computer. The spectrometer calibration for energy and detection efficiency was performed with γ-lines of $^{226}$Ra + daughters 6/.

The investigation of the conversion electron spectra was made with 0.6 cm$^2$ sensitive surface Si/Li/ detectors. The half-width of e$^-$-lines was 4.7 keV at 625 keV, and at -120 °C temperature, at which the measurements was performed. The quick transport of the sources to the detector was made possible by an air-lock. We used $^{137}$Cs, and the known e$^-$-lines of Tl isotopes for spectrometer calibration.

3. RESULTS

The γ-spectrum of Tl isotopes was investigated within the energy interval from 60 keV to 3500 keV, the spectrum of conversion electrons - from 20 keV to 800 keV. The measuring time for the $^{190}$Tl and $^{191}$Tl sources was ≈40 min, and for the $^{193}$Tl sources ≈70 min.

Some characteristic parts of the spectra are shown in the following figures: Fig. 2 /$^{193}$Tl γ/, Fig. 3 /$^{193}$Tl e$^-$/, Fig. 4 /$^{191}$Tl, $^{190}$Tl γ/. 

Fig. 1 Production and separation of thallium isotopes. Schematic drawing of the target system and thermochromatographic apparatus.
Fig. 2 Characteristic γ-spectrum part of chemically purified and mass separated $^{193}\text{TI}$ source. $N_\gamma$ - counts per channel per 45 min.
Fig. 3 Characteristic $e^-$-spectrum part of chemically purified and mass separated $^{193}$Tl.
$N_{e^-}$ - $e^-$-counts per channel per 33 min.
Fig. 4 Three γ-spectrum in succession of chemically purified and mass separated $^{191+190}\text{Tl}$ source. The first measurement was started after 11.5 min after the end of the 10 min irradiation. During the second exposition only $^{190}\text{Tl}$ source was investigated.
The results of the γ-spectrum measurements are summarised in Table I, and the results of the conversion electron measurements — in Table II.

Table I. γ-lines of chemically purified and mass-separated Tl isotopes

<table>
<thead>
<tr>
<th>Isotope</th>
<th>Energy /keV/</th>
<th>Relative intensity</th>
<th>Half-life /min/</th>
</tr>
</thead>
<tbody>
<tr>
<td>193\textsuperscript{Tl}</td>
<td>207.8 ± 0.9</td>
<td>21 ± 4</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>284.6 ± 1.0</td>
<td>22 ± 4</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>324.2 ± 0.9</td>
<td>100</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>335.3 ± 1.0</td>
<td>27 ± 3</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>344.1 ± 1.0</td>
<td>44 ± 6</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>676.2 ± 0.9</td>
<td>41 ± 5</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td></td>
<td>1047.7 ± 0.9</td>
<td>54 ± 9</td>
<td>21.6 ± 0.0</td>
</tr>
<tr>
<td>191\textsuperscript{Tl}</td>
<td>215.7 ± 0.8</td>
<td>100</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td></td>
<td>264.7 ± 0.8</td>
<td>46 ± 5</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td></td>
<td>326.1 ± 1.0</td>
<td>64 ± 7</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td></td>
<td>336.5 ± 1.2</td>
<td>34 ± 4</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td></td>
<td>378.2 ± 1.5</td>
<td>27 ± 7</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td></td>
<td>563.4 ± 1.5</td>
<td>17 ± 4</td>
<td>5.2 ± 0.4</td>
</tr>
<tr>
<td>190\textsuperscript{Tl}</td>
<td>305.0 ± 1.5</td>
<td>16 ± 4</td>
<td>3.6 ± 0.5</td>
</tr>
<tr>
<td></td>
<td>416.8 ± 0.8</td>
<td>100</td>
<td>3.6 ± 0.5</td>
</tr>
<tr>
<td></td>
<td>625.5 ± 1.0</td>
<td>75 ± 7</td>
<td>3.6 ± 0.5</td>
</tr>
<tr>
<td></td>
<td>730.6 ± 1.2</td>
<td>32 ± 5</td>
<td>3.6 ± 0.5</td>
</tr>
<tr>
<td></td>
<td>859.5 ± 1.5</td>
<td>23 ± 5</td>
<td>3.6 ± 0.5</td>
</tr>
</tbody>
</table>

+\textsuperscript{1} See also our earlier publication\textsuperscript{1}.

The 247 keV /\approx 26\text{ min/} and 494 keV /\approx 25\text{ min/} γ-lines, observed in the spectrum of the 193\textsuperscript{Tl} source, may also belong to 193\textsuperscript{Tl}. Similarly, the 477.7 keV /\approx 4\text{ min/}, 535.7 keV /\approx 5\text{ min/}, 637.6 keV /\approx 6\text{ min/} weak lines, observed in the γ-spectrum of the 191\textsuperscript{Tl} source, may belong to 191\textsuperscript{Tl}, but also other assignment is possible. The 196.0 ± 0.8/, 224.5 ± 1.2/, 240.8 ± 1.0/, 252.5 ± 0.8/,
Table II. $e^-$-lines of chemically purified and mass-separated $^{193}$Tl

<table>
<thead>
<tr>
<th>$E_e$ keV</th>
<th>$E_{e^-}$ keV</th>
<th>$e^-$ rel. int.</th>
<th>$\alpha_k/\alpha_l$</th>
<th>Exp. $\alpha_k^+$</th>
<th>Multipolarity$^\Delta$</th>
<th>$E_e$ keV</th>
<th>$e^-$ rel. int.</th>
</tr>
</thead>
<tbody>
<tr>
<td>284.6</td>
<td></td>
<td></td>
<td></td>
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<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>K 202 30.4$^{\pm}$ 4.2</td>
<td>6.3$^{\pm}$ 1.1</td>
<td>4.8$^{\pm}$ 0.8</td>
<td>0.345</td>
<td>M1 + E2</td>
<td>270</td>
<td>$\sim$10</td>
</tr>
<tr>
<td></td>
<td>L 270</td>
<td></td>
<td></td>
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* The half-lives of $e^-$-lines, shown in the Table, are in agreement with the 21.6$^{\pm}$ 1 min period.

$^\Delta$ We used the Hager - Seltzer tables $^7$ for the $\alpha_k/\alpha_l$ rate calculations.

$^+$ At the calculation of $\alpha_k$ values the $e^-$ relative intensities were renormed so that we could get the theoretical $\chi_k/M1$, 324.2 keV/ = 0.25 value /see text/. 
From the experimental $\alpha_K/\alpha_L$ ratio the multipolarity of the 324.2 keV transition must be M1 or E1. The M1 assignment is more probable, because with this assumption the determination of multipolarities of the other transitions on the basis of $\alpha_K/\alpha_L$ and $\alpha_K$ values may be performed without contradiction, while the E1 assumption leads to difficulties.

4. DISCUSSION

$^{193}$Tl. The $\gamma$-spectrum of $^{193}$Tl has not been directly investigated till now in the literature. The conversion electron spectrum results, obtained by Andersson et al., are shown in Table II.

The obtained half-life is in agreement with the published 22.6 $\pm$ 1 min \(^8\) value.

$^{191}$Tl. K. F. Chackett and G. A. Chackett obtained for the half-life of $^{191}$Tl 10 $\pm$ 1 min \(^9\) in 1960. They investigated the neutron deficient thallium isotopes, obtained in heavy ion reactions, after chemical separation with GM and scintillation counters. According to the measurements of Andersson, Häller, Ringh \(^8\), the half-life of $^{191}$Tl is less than 10 min; probably $^{192}$Tl /11.4 min/ admixture disturbed in the \(^9\) measurements. In the present work for the first time succeeded to assign concrete $\gamma$-lines to $^{191}$Tl. On the basis of investigation of the 216, 265, 326, 337, 378 and 563 keV $\gamma$-lines, we have obtained 5.2 $\pm$ 0.4 min as the half-life of $^{191}$Tl.

$^{190}$Tl. The measurement, accomplished with the isotope separated sources, confirmed that the 417, 626 and 730 keV $\gamma$-lines really belong to $^{190}$Tl. In our earlier study \(^1\) the half-life determination of $^{190}$Tl new isotope was performed on the basis of the intensity decrease of these lines.
We are grateful to Dr. K. Ya. Gromov and Dr. V. I. Raiko for the support of the work, to Mrs. T. B. Vandlik for help in chemical separation, and to the Institute for Nuclear Research of the Hungarian Academy of Sciences /Debrecen/ for placing some measuring equipment at our disposal.

REFERENCES


NEUTRON-RICH NUCLIDES OBTAINED USING THE OSIRIS FACILITY

The OSIRIS Collaboration,

B. Grappengigsser, E. Lund and G. Rudstam,
Studsvik, Nyköping, Sweden.

Abstract

A survey of the short-lived nuclides produced in the isotope-separator-on-line facility at Studsvik is presented. Fission fragments recoiling out of a $^{235}\text{U}$-target are caught by a graphite cloth in the ion source. The temperature of the ion source is kept high enough to release eighteen of the elements produced by fission so fast that it is possible to detect their short-lived isotopes after electromagnetic mass separation. No chemical separation but for the element-dependent separator efficiency has been used.

During the course of the work several methods for the determination of the half-life of the radioactive nuclides have been employed. The decays have been counted by means of the following systems: a $4\pi\beta$-counter made of plastic scintillating material, a neutron detector consisting of BF$_3$-tubes, an anthracene crystal with beta pulses sorted according to energy, and Ge(Li)- and Si(Li)-detectors for following the decay of gamma-rays, X-rays, and conversion electrons. The results are collected in a table containing half-life determinations of 75 nuclides or isomers, 32 of them new or with half-lives significantly different from published data.

1. INTRODUCTION

When the OSIRIS facility$^1$ for mass separation of fission products came into operation in June 1968 it was considered of urgent interest to survey the nuclides obtained. This initiated a systematic investigation of the radioactive components appearing in the different mass chains. Various instrumentation for determining the half-life of the nuclides produced have been used. Results from this survey are presented in Table 1. As a target-ion source arrangement giving essen-
tially no element discrimination has been used in these experiments (see Ref. 1) the element assignments given in the table are mostly based on a comparison with half-lives published in the literature and reasonable "parent-daughter" arguments. In a few cases the assignments have been confirmed by observing K-X-rays and conversion electrons.

2. INSTRUMENTATION

2.1 Integral counting of beta particles and of neutrons (Experiment I)
(B. Grapengiesser, E. Lund and G. Rudstam)

Most of the results in Table 1 were obtained with a $4\pi$-detector connected to the magnetic tape transport system of the isotope separator. After a predetermined collection time the sample was rapidly moved to the detector for measurement, usually multisampling, during a fixed period of time. This cycle was then repeated automatically in order to reach sufficient statistics. The sample stays all the time in the same vacuum system, and the time loss during the transportation of the sample between the collection position and the measuring position - a distance of about 1 meter - could be kept as small as 2 sec. This means that the practical limit for measurements lies in between 0.5 and 1 sec. depending on the strength of the samples produced.

The most important features required for the beta detector are the following:

i) a high efficiency which is constant, reproduceable and little dependent on the beta energy;

ii) a low background to make possible measurement even of very weak samples;

iii) capability of measuring high as well as low counting rates with precision.

A plastic scintillator was considered to best fulfill the third requirement. It was therefore chosen for the construction shown in Fig. 1, which well responds to the first two requirements. Its volume
is only 2.5 cm$^3$ which means easy shielding and a low background rate, typically 10 cps at full reactor power, and its solid angle is almost $4\pi$. The electronic components (see Fig. 1) accept high counting rates. Most of the results were obtained using a Nuclear Data Model 2200 analyzer in the multiscaling mode, but for more long-lived nuclides a printing scaler can be connected to the detector.

The discriminator level was set to accept pulses corresponding to electrons of energy larger than 50 – 200 keV. Although roughly half of the electrons have to pass through the tape (5.6 mg/cm$^2$) the efficiency was generally over 90% in the measurements reported here.

The dead time was around 0.3 usec in experiments without the prescaler. The fast prescaler arrangement shown in Fig. 1 should considerably decrease the dead time. Corrections for dead time losses have usually been negligible.

The possibility that certain elements diffuse out of the tape thereby causing the decay curve to appear too steep must be considered. This can be checked for long-lived components with known half-lives. No such effect was observed.

The multiscaling data are recorded on punched paper and evaluated by means of a computer programme which resolves the decay curve into a given number of components using an iterative procedure. Systematic trends in the differences between measured points and best fit reveal components not taken into account, and a new evaluation with better input data can be carried out. Obviously, two components with similar half-lives might not always be resolved by this method. In such cases a specific detection method, i.e. following the decay of a gamma line or X-ray line, would be preferable to the non-specific beta measurement.

The errors given for the $4\pi$-data in Table 1 are purely statistical. Possible systematic errors caused by several components having closely the same half-life, or by contaminating activities of half-lives similar to that of the measured nuclide are not included.
The latter reason of systematic error can usually be ruled out for cases where duplicate measurements give consistent results as the degree of contamination and hence the effect would presumably vary from experiment to experiment. Cases with several determinations (the number of determinations is denoted in the table) generally show very good internal and external consistency indicating negligible effect of contaminating activities. Cases with only one determination should be considered as preliminary, however, until duplicate runs confirm the results.

For a few mass chains the decay was also followed using a neutron detector consisting of ten BF$_3$- tubes imbedded in paraffine and arranged in a semicircle around the sample position. The BF$_3$-tubes were coupled in parallel to a multiscaler for analysis of the decay of delayed neutron activities. The half-lives obtained in this way are also included in the table.

2.2 Anthracene detector for selective counting of beta particles

(Experiment II)


In order to be able to measure activities with very short half-lives (below one second) a selected ion beam is focused onto a second tape system where the detectors are arranged in such a way as to view the collection spot. Then the delay caused by the tape transport is avoided, and the moving tape is only used to remove long-lived activities. A special 1/2-inch tape with an aluminium coating was chosen because it can withstand the rather intense ion beams of stable krypton and xenon isotopes used as mass markers.

Both anthracene and Ge(Li)-detectors have been used to measure beta or gamma radiation. Usually the experimental sequence for the half-life determinations was supervised by a 8K PDP-9 computer. Such a sequence consisted of an accumulation on the tape for a predetermined time after which the ion beam was interrupted and the measurement was started. When using an anthracene detector, the output pulses were divided into four ranges, each corresponding to a 2 MeV energy interval. This enhances
short-lived activities with large Q-values in the higher energy gates. The PDP-9, run in a multiscaler mode, was used for storing the pulses from the different gates simultaneously. A similar procedure was adopted in some confirming experiments using the Ge(Li)-detector, now with gates set on different gamma lines.

All the data were analyzed on the PDP-9 using various fitting programmes. The resulting half-lives (averages of measurements using different energy gates) are given in Table 1.

2.3 Multispectrum scaling of gamma-ray, X-ray and conversion electron spectra from Ge(Li)- and Si(Li)-detectors (Experiment III)  
(B. Fogelberg, A. Bäcklin and G. Hedin)

Identification of several short-lived nuclides have been made in the mass-region A = 114 - 125 by multispectrum scaling experiments in which the tape transport system mentioned in subsection 2.2 has been used. Ge(Li)-detectors for gamma-ray measurements and high resolution Si(Li)-detectors for X-ray or conversion electron measurements were used. For background reasons the detectors were in some cases placed 11 cm away from the beam implying a time delay of 0.1 sec for the movement of the tape before the measurements could be started. After an appropriate collection time the beam was shut off, and successive spectra were recorded in a ND 3300 multichannel analyzer operated in the 4 x 1024 channel mode. Generally, one spectrum was measured during the collection time, and three quadrants of the analyzer memory were used to record the successive decays. The cycle was automatically repeated until sufficient statistical accuracy was reached. Half-lives determined in this way are not particularly accurate since only three points on the decay curves are obtained, but reliable identifications of elements could be made from X-ray energies and from the K to L line distances in the conversion electron spectra. In some cases elements could also be identified through the presence of gamma transitions known to take place in the daughter nuclei.

The results obtained (averages of measurements using different gamma gates) are included in Table 1.

3. EXPERIMENTAL RESULTS

The experimental results are presented in the table below
### Table 1

<table>
<thead>
<tr>
<th>Mass number</th>
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<th>Experiment</th>
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<th>Comment</th>
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<th>Ref</th>
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<td>120</td>
<td>2.87±0.05 s</td>
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<td>In</td>
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<td>- 5 s</td>
<td>III</td>
<td>1.171</td>
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<td>Element suggested</td>
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<td>Assignment Ref.</td>
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<td>0.1946 1.171</td>
<td>In</td>
<td>0.1946 MeV γ converts in Sn. Sn X-rays observed</td>
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<td>III</td>
<td>0.194 0.420 0.448</td>
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<tr>
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<td>129</td>
<td>0.8±0.3 s</td>
<td>II</td>
<td></td>
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<td>II</td>
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<tr>
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<td>I</td>
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<td>I</td>
<td>2)</td>
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<td>Cs</td>
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<td>Cs</td>
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<td>Ba</td>
<td>9)</td>
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<td>La</td>
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<td>La</td>
<td>9)</td>
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<td>Ce</td>
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<tr>
<td>146</td>
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<td>La, Ce</td>
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<td>Ce</td>
<td>2)</td>
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<td></td>
<td>La, Ce</td>
<td>65 s</td>
<td>Ce</td>
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</table>

a) Figure within brackets denotes the number of determinations.

b) Indium X-rays show growth corresponding to a ~ 5 sec. parent activity.
   This is more probably due to the decay of cadmium than of a very low-energy indium isomer.
REFERENCES


Fig. 1 4π beta plastic scintillator detector arrangement.
PRODUCTION AND DECAY PROPERTIES OF SOME SHORTLIVED COBALT NUCLEI

H.W. Jongema, J.C. de Lange, J.C. Boddendijk,
T. Kamermans and H. Verheul,
Natuurkundig Laboratorium, Vrije Universiteit,
Amsterdam, The Netherlands.

Abstract
Several new gamma transitions have been observed in the decay of
the following isotopes: $^{62}$Co (13.5 min), $^{62}$Co (1.4 min), $^{62}$Cu
(9.73 min) and $^{63}$Co (26.7 sec).

1. Introduction

The decay of some neutron rich shortlived cobalt nuclei, decaying
to nickel isotopes, was investigated. In the past few years the
results of several theoretical calculations on the excited states
of the Ni isotopes were reported $^{1,2,3,4,5,6}$. However the experi-
mental data about the decay of $^{62,62m,63}$ and $^{64}$Co were rather poor.
The 13.5 min isomeric state of $^{62}$Co decays to levels in $^{62}$Ni with
high spin. There is no consistency in the level schemes presented
by various authors $^{7,8,9,10,11}$, except for the position of the
$2^+$ and $4^+$ level in $^{62}$Ni.

$^{62}$Cu and the 1.4 min isomer of $^{62}$Co both decay to levels in
$^{62}$Ni with low spin. Decay properties of levels with low spin are
known from a recent study of the $^{61}$Ni (th n,$\gamma$)$^{62}$Ni reaction by
Fanger et.al. $^{12}$.

As pointed out in the N.D.S. $^{13}$, confusion exists in lit-
terature about the half-lives of $^{62}$Co and $^{64}$Co and whether or not
these nuclei have isomeric states. Only one gamma ray with an
energy between 80-100 MeV and a half-life of around 30 sec has been
observed $^{14,15,16,17}$ and assigned as well to the decay of $^{63}$Co as
to the decay of $^{64}$Co.

We measured a.o. single gamma-ray spectra. This was mainly
done with a 30 cm$^3$ Ge(Li)-detector. The amplified signals of these
detector were analyzed with 100 MC ADC's (4096 channels). The ADC's
were coupled to a 4K Nuclear Data 50/50 memory and via the PDP-8/L
computer to a CDC 1700 computer.
2. Decay to $^{62}$Ni

2.1 Decay of 13.5 min $^{62}$Co

$^{62}$Co was produced by the (16 MeV d,$^a$)-reaction on enriched $^{64}$Ni in the AVF cyclotron of the 'Vrije Universiteit'. The sources were in about 6 sec transported to the spectrometers with a pneumatical system $^{18}$). A chemical separation with a ion-exchange column was performed in about 1.5 min. Besides single gamma-ray spectra, two dimensional $\gamma$-$\gamma$- coincidence spectra were measured with Ge(Li) and Na I(Tl) detectors. The results of the analysis of the spectra and the construction of the decay-scheme is published elsewhere $^{19}$). The gamma transitions assigned to this decay are given in Table I.

Table I

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<tr>
<th>$E_\gamma$ (MeV)</th>
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<td>1.1291 3</td>
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<td>6.4 5</td>
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<tr>
<td>2.251x) 1</td>
<td>0.4 2</td>
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<td>2.3020 5</td>
<td>1.8 2</td>
</tr>
<tr>
<td>2.8826 5</td>
<td>1.1 1</td>
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</table>

x) Not placed in level scheme

2.2 Decay of 1.4 min $^{62}$Co

For a more detailed discussion of the results of the investigation of this decay, see ref. 19.
The results of the single gamma-ray measurements are given in Table II. Except of the gamma-rays of 1.8863 and 2.3458 MeV, these gamma-rays were also seen in the decay of $^{62}\text{Cu}$.

**Table II**

Energies and relative intensities of the gamma-rays of 1.4 min $^{62}\text{Co}$ and $^{62}\text{Cu}^x$)

<table>
<thead>
<tr>
<th>$E_\gamma$ (MeV)</th>
<th>Int. (1.4 min $^{62}\text{Co}$)</th>
<th>Int. ($^{62}\text{Cu}$)</th>
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<td>0.87571</td>
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<td>0.30</td>
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<tr>
<td>3.8617</td>
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<td>0.08</td>
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</tbody>
</table>

$x$) For the energies the mean values of $^{62}\text{Cu}$ and $^{62}\text{Co}$ were taken.

2.3 Decay of $^{62}\text{Cu}$

To investigate the feeding of the high excited states in $^{62}\text{Ni}$, found in the decay of $^{62}\text{Co}$, we reinvestigated this decay with an anti-annihilation set-up $^{20})$. This set-up is especially suitable for the study of this decay, because of the very strong annihilation radiation (99.5% of the beta-decay goes to the ground state). Due to count-rate and background reduction in this set-up, very weak gamma-rays (till 3.10 $^{-4}$%) could be found in the high energy part of the
spectrum. The energies and relative intensities of the other gamma-transitions were determined with much more accuracy than in our previous work 21). Also the accuracy of the experimental value for the intensity ratio I (1.173)/I (0.511) was improved. For this ratio we found $(1.72 \pm 0.09) \times 10^{-3}$ in good agreement with the value recently reported by Van Patter et.al. 22). In our investigation the $^{62}$Cu sources were produced with the (7 MeV p,n)-reaction on $^{64}$Ni-foils. Two foils were alternately irradiated and measured during 6 hrs.

2.4 Levels in $^{62}$Ni

In figures 1a and 1b the proposed decay schemes are given. Spin assignments are only based on the results of this investigation. For parity assignments the known odd values of $l_p$ and $l_n$ from particle transfer reactions 8) were used too.

Combining the spin assignments of Fanger et.al. with our results, we found that almost all levels till 4 MeV have spin $1^+$, $2^+$ or $4^+$. The first candidate for a $3^+$ level is the 3.52269 MeV level of Fanger ($2^+$, $3^+$). Besides the well established $0^+$ level of the two phonon triplet at 2.04873 MeV, a next $0^+$ level could exist at 2.8096 MeV. A $0^+$ level was found indeed at 2.85 MeV from a (p,t)-measurement by Davies et.al. 23). We have tried to establish the $0^+$ assignment by searching the conversion electrons of the possible E0 transition of the 2.8906 to the 2.04873 MeV level. However, no conversion line was found. So we could not exclude the assignment $2^+$ of Fanger for this level.
3. Decay of \(^{63,64}\)Co

Both isotopes were produced by irradiating enriched \(^{64}\)Ni (15 mg) with neutrons from the \((42\ \text{MeV} \ \text{^3He, n})\)-reaction on a beryllium target. In about 2 sec the source, packed in a polyethylene capsule, was pneumatically transported to the detector.

Some preliminary results of this study are mentioned below.

Table III

<table>
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<tr>
<th>(E\gamma ) (MeV)</th>
<th>(I\gamma )</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.090</td>
<td>15</td>
</tr>
<tr>
<td>0.157</td>
<td>15</td>
</tr>
<tr>
<td>0.527</td>
<td>15</td>
</tr>
</tbody>
</table>

\(5/2^-, 7/2^- \quad 26.7\ \text{sec} \)

\(^{63}\)Co

\(\frac{5/2^-, 7/2^-}{63}\)\text{Co}

\(\frac{5/2^-, 7/2^-}{63}\)\text{Ni}

\(\frac{3/2^-}{63}\)\text{Co}

\(\frac{3/2^-}{63}\)\text{Ni}

\(\frac{1/2^-}{63}\)\text{Ni}

\((\leq 6\%. \geq 5.8)\)

\((\leq 4.6\%. \geq 2.6)\)

\((90.4\%. 4.8)\)

\(5/2^-, 7/2^- \quad 37\)

\(5/2^-, 7/2^- \quad 527\)

\(5/2^-, 7/2^- \quad 157\)

\(5/2^-, 7/2^- \quad 90\)

\(5/2^-, 7/2^- \quad 0\)

Fig. 2 Preliminary decay scheme of \(^{63}\)Co

A strong gamma-ray with an energy of 0.090 MeV and a half-life of 26.7 \(\pm 0.8\) sec was observed (cf. figure 3). Two other short-lived gamma-rays were found at 0.157 MeV and 0.527 MeV. We tentatively assigned these 3 gamma-rays to the decay of \(^{63}\)Co, because their energies fit with the already known levels in \(^{63}\)Ni. Moreover, a half-life of 26.7 sec for \(^{63}\)Co would be in agreement with the value reported by Ehrlich\(^{15}\), Kiselev\(^{16}\) and Ward\(^{17}\).

No indication of \(^{64}\)Co gamma-rays was found until yet; this would be in agreement with a half-life of 0.4 sec as reported by Levkovski\(^{24}\) and Ward\(^{17}\).

This study is still in progress.
4. Discussion

Most of the theoretical calculations on the Ni-isotopes treat these isotopes as a double closed $^{56}\text{Ni}$-core plus additional neutrons in the next three single particle shell model states. A large number of states are calculated by Auerbach 1) using an effective matrix element method. For $^{62}\text{Ni}$ a comparison (figure 4) between theoretical and experimental levels is separately made for the different spins.

The spins and parities $2^+$ and $5^+$ for the $^{62}\text{Co}$-isomers are in agreement with the revised Nordheim rules 25) if we assume a configuration $p (f \frac{7}{2})^{-1} n (p \frac{3}{2})$ for these states. This is in analogy with the isomeric states in $^{58}\text{Co}$ and $^{60}\text{Co}$. The ground state spin and parity of $^{63}\text{Co}$ is $\frac{7}{2}^-$, $\frac{7}{2}^+$ 13). This is in agreement with the log ft = 4.8 for the feeding of the 90 KeV ($\frac{5}{2}^-$, $\frac{7}{2}^-$) level in $^{63}\text{Ni}$. The feeding of the 157 and 527 KeV levels has to proceed by unobserved gamma-transitions from higher levels, or by beta-decay. Allowed log ft values for these beta-transitions should mean spin and parity $\frac{5}{2}^-$ for the $^{63}\text{Co}$ ground state. This however is very unlikely within the framework of the nuclear shellmodel.
- 1112 -

Theoretical 1) and experimental levels in $^{62}\text{Ni}$

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LIFETIMES OF EXCITED NUCLEAR LEVELS

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The measurement of nuclear lifetimes plays a central role in the investigation of low energy nuclear structure. The obtained absolute transition probabilities give a direct information on the different excitation modes and can be used to test the validity of proposed nuclear models. It is therefore natural that it is of great interest to measure nuclear lifetimes in nuclei far from the beta stability line where the nature of the excited levels is unknown. For this purpose an electron-electron coincidence spectrometer has been set up in connection with the two on-line mass separator facilities ISOLDE (at CERN) and OSIRIS (at Studsvik). At CERN the radioactive products are produced by 600 MeV proton spallation in different targets giving a number of neutron deficient nuclei. Several of these activities have been studied by us over the last years and we have measured half-lives of excited levels in $^{119,120,121,122,123}_{\text{I}}$, $^{184,185,186,187,188,189}_{\text{Os}}$, $^{185,187,188,189,191}_{\text{Ir}}$, $^{184,186,187,188,189,190,191,193}_{\text{Pt}}$, $^{190,191,192,193,195}_{\text{Au}}$, $^{193,195}_{\text{Hg}}$, $^{205}_{\text{Po}}$ and $^{211}_{\text{At}}$. At Studsvik neutron rich radioactivities are produced by neutron irradiation of $^{235}_{\text{U}}$ giving rise to a lot of fission products. New determinations of half-lives of excited levels have here been made in $^{88}_{\text{Rb}}$, $^{91}_{\text{Rb}}$, $^{93}_{\text{Sr}}$, $^{93}_{\text{Rb}}$, $^{121}_{\text{Sn}}$, $^{128}_{\text{Sb}}$, $^{128,130}_{\text{Te}}$, $^{134,138}_{\text{Cs}}$ and $^{141}_{\text{La}}$. Many of the measurements mentioned above are very recent and the data has not been as yet completely analysed. We will therefore only give some details in a few specific cases below.

CERN. An extensive study of the products produced by the Pb (p, 3pxn) reaction has been performed. Many of the most neutron deficient nuclei produced by this reaction have not been studied at all up to now and at present in many cases only vague ideas of the level scheme are available. This makes measurements of nuclear lifetimes with the delayed coincidence method rather a delicate affair and although many delayed transitions are observed the true level half-lives can not be deduced.
until reliable level scheme are obtained. This is the reason why interpretations of the present data at the moment is only available for nuclei facing the beta stability line although data for nuclei further away from stability are available.

One of the advantages with the above mentioned activity production methods is that sequencies of nuclei with equal number of protons or neutrons are produced making systematic studies possible. Good examples of that are the Os, Ir, Pt and Au nuclei. As an example we can here mention the very recent measurements of the half-lives for the first excited $2^+$ level in $^{184}$, $^{186}$, $^{188}$ Os and $^{184}$, $^{186}$, $^{188}$, $^{190}$ Pt.

Another set of half-lives have been measured for the low energy positive parity levels in odd mass Ir-nuclei giving a direct evidence for the possible deformation of these nuclei. The half-lives for a set of $7/2^-$ states in odd mass Ir and Au nuclei have also recently been measured. An example of measured transition rates between single particle levels is given for $^{193}$, $^{195}$ Hg where the M1 and E2 strengths
between the $f_5/2 \rightarrow p_3/2$, $p_3/2 \rightarrow p_1/2$ and $f_5/2 \rightarrow p_1/2$ odd neutron states were observed [2]. Recently our data was used in a similar evaluation of transition rates between the low lying positive parity states in the odd proton gold nuclei [3]. The result is shown in fig 1 together with the theoretical calculations of Sorensen [4] and Reehal and Sorensen [5].

**STUDSVIK.** Several of the nuclei produced in fission have been investigated as regards to half-lives of excited levels. Besides the above mentioned electron coincidence spectrometer, a $\beta - \gamma$ scintillation coincidence system has recently been set to work or-line the OSIRIS mass separator, directly facing the spot where the selected ion beam hits a moving tape. In this way investigations can be done on very short lived decay products for instance illustrated by a measured half-life in $^{93}$Rb populated in the decay of $^{93}$Kr (1.2 sec.). It is expected that this new half-life is associated with a three quasi-particle state similar to the one recently discovered in $^{91}$Rb [6]. Other examples of recently measured half-lives in fission products are the 79.5 keV level in $^{134}$I ($T_{1/2} = 1.62 \pm 0.10$ nsec) [7] and the 190.3 keV level in $^{141}$La ($T_{1/2} = 1.27 \pm 0.08$ nsec) [8]. A lot of the data still has to be analysed and it is expected that the results will be published within the next year.

We are indebted to Dr John McDonald especially for his efforts put into the on-line coincidence system.
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SHORT-LIVED ISOMERS OF $^{136}$I AND $^{138}$Cs
DECAYING VIA METASTABLE STATES OF $^{136}$Xe AND $^{138}$Ba

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$^{136}$I ($T_{1/2} = 83$ s) and $^{138}$Cs ($T_{1/2} = 32$ min) are well known among the fission products $^1$. In the course of a study of charge distribution in fission $^2$ we have found another $^{136}$I isomer ($T_{1/2} = 48$ s), which was discovered meanwhile by Lundan and Siivola $^3$. This isomer has a higher spin than the 83 s state and decays through a $^{136}$Xe metastable state $^4$. Such states are known in various $N = 82$ nuclei $^4$ and this has led us to a search for homologues in the isobaric chain $A = 138$.

Fast chemical separation of cesium from the fission products after a short ($\simeq 30$ s) irradiation of $^{235}$U by thermal neutrons, gives a mixture of short-lived Cs nuclides. The gamma spectrum is analyzed with a Ge(Li) diode, and the decay of the 1436 keV gamma of $^{138}$Ba can be resolved in two components, the half-lives of which are 3.1 min and 32 min. The decay of the 463 keV gamma ray gives the same result. A gamma ray of 192 keV decays with only the 3.1 min half-life $^5$.

The complex spectrum of the fission sources suggests the utilization of $^{138}$Cs samples obtained by $^{138}$Ba (n,p) $^{138}$Cs (14 MeV neutrons). A chemical separation is still necessary, due to the strong interference of $^{137m}$Ba ($T_{1/2} = 2.6$ min). Such sources show mainly the two $^{138}$Cs isomers with negligible contributions of others cesium isotopes.

For $^{136m}$I $\longrightarrow$ $^{136}$Xe as for $^{138m}$Cs $\longrightarrow$ $^{138}$Ba the most conspicuous result is the absence of coincidences between the most intense component of the beta spectrum and the $4^+ \longrightarrow 2^+$ and $2^+ \longrightarrow 0^+$ gamma transitions in the daughter nucleus, for a short resolving time (100 ns). Consequently, we have studied the delayed coincidences and found a metastable state ($T_{1/2} = 2.8$ $\mu$s for $^{136}$Xe; $T_{1/2} = 0.8$ $\mu$s for $^{138}$Ba) which is strongly fed by the beta decay of the previously mentioned isomers.

In the decay scheme of fig. 1, the positions of the $2^+$ and $4^+$ levels are established by $\gamma - \gamma$ coincidences; they are confirmed by other measurements:
a) decays of $^{136}$I (83 s) $^{3, 4}$ and $^{138}$Cs (32 min) $^{1, 6}$;

b) nuclear reactions $^{136}$Xe ($p,p'$) $^{7}$, $^{137}$Ba ($n,\gamma$) $^{8}$.

The position of the isomeric level is given tentatively, because the observation of low energy $\gamma$ rays or conversion electrons is made very difficult by the intense and energetic beta spectrum. The nature of the isomeric transition is not completely ascertained; nevertheless some analogies can be observed with the isomeric states in $^{140}$Ce and $^{142}$Nd to which the spin $6^+$ is attributed. The isomeric transition E2 has then an hindrance factor, relative to single particle estimates, of about 60 for $^{136}$Xe and 20 for $^{138}$Ba. These values ought to be compared with 2 for $^{140}$Ce and 50 for $^{142}$Nd $^{9}$.
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ESTIMATION OF DECAY PROBABILITIES
OF SHORT-LIVED NUCLEI

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The prediction of partial half-lives for the different
modes of decay is very helpful in planning the search for
unknown nuclides. In addition, major deviations between
the calculated and experimental half-lives, are a good
indication for unusual nuclear properties. Therefore,
Viola and Seaborg (1) made an attempt to develop methods
for evaluating the systematics in the half-lives of heavy
nuclei. The accuracy of these calculated values is limited.
Thus, for instance, extrapolations to the region of super-
heavy nuclei are rather uncertain.

For this reason, we tried to improve the methods for
calculating partial half-lives of $\alpha$-decay, electron
capture and $\beta^+$-decay. The probability of spontaneous
fission was also taken into account (2).

For the calculation we used more general equations for
the decay probabilities. Further we assumed that the
transitions to the levels of the daughter nuclei can be
lumped together to a mean excitation energy. With these
assumptions it was possible to improve the estimation of
half-lives considerably. For instance, in the case of
Z-even, N-even actinide isotopes, the factors of error (standard deviation of the mean) are for \( \alpha \)-decay 1.14, for electron capture 1.9 and for \( \beta \)-decay 1.3. With our method, half-lives can be calculated more accurately by a factor of 4 to 10, as compared to Viola and Seaborg (1).

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Gamma-ray decay studies have been made for the decays of the following nuclei: $^{90}\text{Kr}$, $^{90}\text{Rb}$, $^{91}\text{Kr}$, $^{91}\text{Rb}$, $^{92}\text{Kr}$, $^{92}\text{Rb}$, $^{92}\text{Sr}$, $^{140}\text{Xe}$, $^{140}\text{Cs}$, $^{141}\text{Xe}$, $^{141}\text{Cs}$, $^{141}\text{Ba}$, $^{141}\text{La}$, $^{142}\text{Xe}$, $^{142}\text{Cs}$, $^{142}\text{Ba}$, and $^{142}\text{La}$. Although it is not possible to be comprehensive in describing all these cases, the following features are briefly summarized:

For the mass 90 and 91 decay chains, more than half of the observed gamma rays representing more than 90% of the gamma intensity have been placed into decay schemes, with the exception of the decay of $^{91}\text{Kr}$, where about 75% of the gamma intensity has been accommodated in a level scheme in which about 40% of 150 gamma rays have been placed. The other decays exhibited about 70 gamma rays each. The isomeric state at 106.9 keV in $^{90}\text{Rb}$ has been observed to populate several levels in $^{90}\text{Sr}$, but as yet no firm interpretation of the various levels observed has been made.

For the mass 92 decay chain, about 100 gamma rays are observed in both Kr and Rb decays, while the decay of Sr exhibits only 19 gamma transitions, with one intense transition and four moderately strong ones. Decay schemes have been developed which incorporate most of the gamma intensities, but are not yet interpreted.

At mass 140, the decay of Xe exhibits at least seven low-lying levels in the first 150 keV of excitation energy, although most of the beta strength is to a single level at 1427 keV which may have a possible
shell-model configuration of $\pi(h_{11/2})\nu(h_{9/2})$. The decay of Cs is observed to populate strongly a level at 1510 keV in Ba, but with other high-energy levels present. The even-even $^{140}$Ba nucleus has a first excited state at 602 keV and a second excited state (probably $2^+$) at 1130 keV.

The mass $^{141}$ nuclei show very complicated gamma-ray spectra, ranging from over 300 gamma rays in the decay of Xe to 30 transitions in the decay of La. Levels in $^{141}$Ce at 14 energies accommodate 22 of the 30 observed transitions, while for $^{141}$La, 27 levels have been assigned from 109 out of the 113 observed transitions. Only partial level schemes exist for Ba and Cs nuclei, due primarily to the great complexity of the information at hand.

At mass $^{142}$, the level schemes for Cs, La and Ce are rather well established, and a partial level scheme for Ba was determined without the benefit of extensive coincidence data. Again, the decays exhibit very complicated gamma-ray spectra.

The results for these studies at mass $^{140}$ and $^{142}$ can be compared with other recent studies at neighboring even masses to provide further evidence for the onset of substantial deformation in the transition neutron region of $N = 88$ to $N = 90$. The systematic level progressions of selected states (e.g., the first $2^+$, $4^+$ and $3^-$ states) show smooth trends toward rotation as neutron number progresses away from $N = 82$.

Decay schemes for these mass numbers will be made available upon request shortly before and during the conference. It should be mentioned that these decay schemes have been deduced from great quantities of singles and high-resolution coincidence data, and represent the combined efforts of no less than seven investigators during the past two years. The author
would like to recognize the efforts of these people at this time: W. C. Schick, Jr., J. W. Cook, J. T. Larsen, R. J. Olson, K. B. Nielsen, C. L. Duke, and J. R. McConnell. It is anticipated that publication of these results will commence soon under individual mass numbers.
CO-EXISTENCE OF DIFFERENT STRUCTURES IN VARIOUS DOMAINS OF THE PERIODIC TABLE

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I. Introduction. A general survey of empirical nuclear properties (in particular the various level schemes) exhibits a natural division of the periodic table in different domains. First of all one realizes immediately that there are special domains which correspond to what should be called a 'pure structure'. With this concept we characterize the following typical cases: 1) The ideal spherical shells for medium heavy nuclei and for the lead-region. 2) The deformed shells for the rare earthes and for the very heavy nuclei. 3) The few cases of ideal complex or \( \alpha \)-structures like He\( ^5 \), Be\( ^8 \) etc. (Perhaps one might add also the special, somehow hypothetical case of space and parity deformation for F\( ^{19} \)).

In view of these well-known properties we wish, in the framework of this short contribution, to draw special attention to the following fact: In between the domains mentioned above we have a large number of isotopes which correspond to neither of the ideal cases mentioned so far. It appears therefore quite natural to assume that the total state function in these new cases should be described in a first approximation by a linear superposition of the pure wave-functions considered before. From the viewpoint of general quantum mechanical principles it is clear that the relative strength of admixture will depend strongly on the position in the periodic table (\( A \)-dependence) as well as on the excitation energy of a given nucleus. This assumption can already be checked in a qualitative way with the help of special properties of the relevant level schemes.
II. Co-Existence from the empirical view-point. In the realm of light nuclei the following properties might be characteristic for such a behaviour: Whereas the ideal complex structure which contains $\alpha$-particles exhibit in their level scheme the characteristic groups of levels which are found at nearly the same position as in the free $\alpha$-particle, these groups are gradually displaced (to lower energies) in the neighbouring isotopes. This fact can easily be understood by an admixture of, let's say, a shell structure by which levels with the same assignment will mix and will, therefore, be displaced. (This displacement becomes finally so large, that this special group of levels can no longer be distinguished within the scheme). Another typical case represents $^{16}_0\mathrm{O}$ where spherical structure prevails for the ground state whereas a deformed wave-function was introduced for a typical band of excited levels. ¹)

In the domain of heavier nuclei there are two characteristic regions where the change-over (from spherical to deformed, $A$ about 160, and from deformed to spherical, $A$ about 200) occurs which might be understood by our gradual change of state functions. This assumption might easily explain the characteristic behaviour of the so-called softness parameter as pointed out by Mrs. Scharff-Goldhaber ²) as well as a few gradual changes of excitation energies near $A = 200$.

III. Mathematical model of co-existence. Under the circumstances mentioned above it might be interesting to see whether the state function of an exactly solvable mathematical model for a $N$-Fermion-problem would exhibit such a behaviour. For this reason we have studied a two-shell model containing two types of forces (of pairing and monopole-type) as proposed by H.J. Lipkin and worked out by the present authors with the help of the representation of
the underlying group $O_3^3$). The energy levels and the state function were determined as functions of a few characteristic parameters, i.e. relative strengths ($V$) of the 2 forces and the number of particles ($N$). In figure I we have plotted the generalized Fourier coefficients $C$ (corresponding to states of sharp seniority $S$) as function of the strength parameter $V$ (with $N = 18$). It is easily seen that the abscissa can be divided in 3 domains corresponding in turn to 1) spherical structure (only one slowly varying coefficient $\neq 0$), 2) admixed structure (rapid variation of all coefficients), 3) deformed structure (with a smooth behaviour of the coefficients). The figure II shows very much the same property using now an expansion according to state functions which are well adopted to the deformed state (quasi spin corresponding to isospin $j$) where, in domain III, only one coefficient is $\neq 0$. In figure III a few excited levels are plotted against the same strength parameter $V$ (cp. full lines) and a characteristic rapid variation is immediately realized in the middle domain (II). In the same time it is interesting to observe that the well-known RPA scheme (represented by the dotted lines) yields a rather good result in the domains I and III, whereas the correspondence is extremely bad in II. (The same holds for the HB-scheme $^4$). This again is evidence that the structure must be of a more general type as soon as we are in the transitional domain between two different structures. (Some level-curves of our diagram nearly correspond to the realistic behaviour of $2^+$- states considered as functions of $A$).

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ONE BOSON EXCHANGE POTENTIAL AND NUCLEAR MATTER PROPERTIES

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Previously\(^1\), we gave a momentum dependent one boson exchange potential (OBEP) containing central, \(\sigma_1\sigma_2\), tensor, \(\sigma\), \(\sigma_1\sigma_2\), and \(\sigma_1^L\sigma_2^L\) terms which fits all the nucleon-nucleon phase shifts up to 330 MeV as well as the two-nucleon bound state properties. (Recently we found that the OBEP gives also the tendency of most of the phase-shifts of a high-energy (up to 600 MeV) analysis done by a Dubna group\(^2\)). In addition to the pole parameters (meson coupling constants and masses) of the exchanged \(\pi\), \(\eta\), \(\omega\), \(\varphi\), \(\sigma_0(J^P=0^+, T=0)\) and \(\sigma(J^P=0^+, T=1)\) mesons a zero cut-off radius \(d\) (the potential is set equal to zero within \(d\leq 1/M\), \(M\) the nucleon mass) was introduced to regularize non-physical singularities at the origin. Because of the shielding effect of the short-ranged \(\varphi\) and \(\omega\) mesons the phase-shifts (except of \(\sigma(\sqrt{3}S_1)\)) are rather insensitive to variation of \(d\) in the range of \(d<1/M\).

With a total of ten free parameters (the coupling constants are in agreement with those determined from other sources) we found a least-squares value of about 2.7/datum for 89 pieces of data. This small number of free parameters necessary to describe the nucleon-nucleon interaction in the scheme of meson theory should be compared with a total of 30-50 free parameters used in phenomenological potentials\(^3-5\).
The next step in checking the idea of the one meson exchange in the potential theory is suitably made by the calculation of infinite nuclear matter properties (binding energy and saturation density) of the OBEP. We have done this using Bethe's reference spectrum method\(^6\) where we assumed the intermediate state single particle potential to vanish. The reference matrix \(G^R\) and its second order correction (Pauli term) have been calculated whereas the higher-order corrections are estimated applying the 'modified geometric approximation' proposed by Dahll et al.\(^7\). In agreement with Bhargava and Sprung\(^8\) we found that the potential energy contribution from first, second and third order are in the ratio 80 : 20 : 2. The binding energy per particle (see the figure) as a function of density, i.e. of the Fermi momentum \(k_F\), shows a minimum of \(-11.2\) MeV at \(k_F=1.55\) \(\text{F}^{-1}\).

![Graph](image)

\(\text{The OBEP binding energy due to two particle clusters as a function of density.}\)

The energy should be corrected for the 3-body contribution which, however, gives less than 1 MeV binding for quite different potentials as has been shown by Dahlblom\(^9\). Accepting this value for our OBEP too we are still over
3 MeV short of the semi-empirical value (-15.7 MeV). Thus the zero cut-off OBEP gives underbinding at a too large density (the observed value is \( k_F = 1.36 \, \text{F}^{-1} \)).

The relatively small deviation of our nuclear matter values from the observed ones (both energy and density) is probably induced by small effects neglected in the OBEP. Considering the results - too little binding at too large density - it seems to us that especially the ratio of repulsion (core region) to attraction (intermediate range) in the S-partial wave states of the OBEP must be improved.

This is under consideration using the methods of high-energy physics (Regge, Veneziano) for the repulsive core region and the correct treatment of the \( 2\pi \)-exchange contribution (thus leaving the pure one particle exchange model) for the attractive intermediate range.

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ON THE DECAY OF $^{90}$Nb, $^{89m}$Nb, $^{89}$Nb, $^{85m}$Y AND $^{85m}$Sr

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The sources of Nb, Y and Sr were obtained by irradiating Mo, SrCl$_2$ or Nb, and Y$_2$O$_3$, respectively, with 660 MeV or 120 MeV protons using the following chemical separation. The single gamma-spectra were investigated using 6 and 10 cm$^3$ Ge(Li) detectors.

About 34 $\gamma$-rays assigned to $^{89}$Nb; 34 $\gamma$-rays to $^{85m}$Y and 2 $\gamma$-rays to $^{85m}$Sr were observed. The $\gamma-\gamma$ coincidences of $^{90}$Nb were measured using the two-dimensional spectrometer with a pair of 10 cm$^3$ or 30 cm$^3$ Ge(Li)-crystals.

On the basis of obtained results, the following levels of $^{90}$Zr, $^{89}$Zr and $^{85}$Sr were proposed:

$^{90}$Zr: 0(0$^+$); 1760,7 (0$^+$); 2186,2 (2$^+$); 2318,6 (5$^-$);
2740,7 (3$^-$); 3077 (4$^+$); 3447,7 (6$^+$); 3589,2 (8$^+$);
3975 (5$^-$); 4232 (6 7$^-$); 4542 (7 6$^+$); 5060 (7 8$^+$);
5164 (8 7$^+$); 5288 (9 7$^-$); 5330 (7,8$^+$); 5377 (7,8$^+$);
5432 (8 7$^+$); 5465 (9 8$^-$); 5674 (9 9$^-$).

$^{89}$Zr: 0(9/2$^+$); 587,6 (1/2$^-$); 1094 (3/2$^-$); 1449,6 (5/2$^-$);
1511,5 (9/2$^+$); 1267,0 (5/2$^+$); 1833,1 (5/2$^+$);
2102 (7/2 9/2 11/2$^+$); 2128,1 (7/2 9/2 11/2$^+$), 2573 (7/2
9/2 11/2$^+$); 2612 (7/2 9/2 11/2$^+$), 2754,0 (11/2 9/2 7/2$^+$);
2926 (11/2 7/2 9/2$^+$), 2960 (7/2 9/2$^+$); 2983 (7/2 9/2
11/2$^+$), 3016 (7/2 9/2 11/2$^+$); 3092,8 (7/2 9/2$^+$);
3471 (7/2 9/2 11/2$^+$); 3511,3 (9/2 11/2$^+$); 3559 (7/2 9/2
11/2$^+$), 3576 (7/2 9/2$^+$), 3911 (7/2 9/2 11/2$^+$) and 3917.

$^{85}$Sr: 0(9/2$^+$); 232 (7/2$^+$); 239 (1/2$^-$); 742 (3/2$^-$);
768 (5/2$^+$); 1154 (3/2$^-$); 1364 (5/2$^-$); 1797 (5/2$^+$);
1935 (7/2 9/2 11/2$^+$); 2123 (7/2 9/2$^+$); 2173 (7/2 9/2$^+$);
2584 (7/2 9/2$^+$); 2745 (7/2 9/2 11/2$^+$); 2786 (7/2 9/2$^+$);
3266 (7/2 9/2 11/2$^+$).
\(^{167}\)Yb DECAY

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The level scheme of the deformed nucleus \(^{167}\)Tm has been investigated in the decay of \(^{167}\)Yb using germanium detectors and a magnetic beta spectrograph. The conversion electron spectrum, single gamma-ray spectra and prompt and delayed gamma-gamma-coincidences have been measured. About 45 new transition have been observed. Most of them have been introduced into a level scheme of \(^{167}\)Tm containing 13 new levels. In the low-energy region besides the rotational bands \(1/2^+\) \(^{[411]}\), \(7/2^+\) \(^{[404]}\) and \(7/2^-\) \(^{[523]}\) the 5/2 and 9/2 members of the \(1/2^-\) \(^{[541]}\) band could be identified. The possibilities of interpretation of the high-lying levels are discussed basing on the microscopic model.

*) Polytechnical Institute, Tashkent, USSR.
THREE-QUASIPARTICLE EXCITATION IN 158\[1/2]\textsuperscript{+}Pm

NEW SHORT-LIVED ISOTOPE 160\[1/2]\textsuperscript{+}Sm

R. Arlt, G. Beyer, G. Musiol, E.S. Ryndina,
S. Seidler and H. Strusny,

Joint Institute for Nuclear Research,
Dubna, USSR.

Using a rapid chemical separation method for rare earths \[1\] we have prepared sources of 141\textsuperscript{m}Sm by irradiating Dy-targets with 680 MeV protons. The decay of the earlier identified \[2\] high spin isomeric state 141\textsuperscript{m}Sm (T_{1/2} = 22.5 m, I = 11/2^-) has been investigated, using single \gamma and \gamma-\gamma-coincidence measurements with two large volume Ge(Li) counters (senst. vol. 30 cm\(^3\) and 33 cm\(^3\)). In this way we have found three levels with high excitation energies in the 141\textsuperscript{m}Pm nucleus (1982.6 keV, 2091.6 keV, 2119.0 keV). They are populated via B-decay with low logft values (5.6; 5.6; 5.4). We interpret this levels as the low spin members of a three-quasiparticle multiplet with the configuration

\[ p(d_{5/2})^{-3}, n_1(d_{3/2})^{-1}, n_2(h_{11/2})^{-1}(I^\pi = 9/2^-, 11/2^-, 13/2^-), \]

which are populated from the isomeric state in 141\textsuperscript{m}Sm with the structure \[ p(d_{5/2})^{-2}, n_1(d_{3/2})^{-2}, n_2(h_{11/2})^{-1} \] 11/2^-.

The halflife of the earlier unknown groudstate 141\textsuperscript{15}Sm (I^\pi = 1/2^+ or 3/2^+) has been determined to be 9.5 ± 0.5 m. This investigation is a part of the YASNAP programme of investigations of short-lived isotopes with the external beam of the 680 MeV Dubna synchrocyclotron.

\[2\] R. Arlt et al., preprint P6 - 3540, Dubna 1967.
NEW ISOTOPES $^{181}\text{Ir}$, $^{180}\text{Ir}$, $^{178}\text{Ir}$


Joint Institute for Nuclear Research, Dubna, USSR.

New neutron-deficient Ir isotopes have been produced by irradiating 4 mg/cm$^2$ Ho and Tm targets in the external beam of the Dubna U-300 cyclotron in the reactions $^{165}\text{Ho}(^{22}\text{Ne},6\text{n})^{181}\text{Ir}$, $^{169}\text{Tm}(^{160},4\text{n})^{181}\text{Ir}$, $^{165}\text{Ho}(^{22}\text{Ne},7\text{n})^{180}\text{Ir}$, $^{169}\text{Tm}(^{160},5\text{n})^{180}\text{Ir}$ and $^{169}\text{Tm}(^{160},7\text{n})^{178}\text{Ir}$.

The mass assignment of Ir isotopes have been established by gamma measurements with a Ge(Li) detector (13 cm$^3$ sensitive volume) and by measuring the excitation functions.

$^{181}\text{Ir}$: The $^{181}\text{Ir}$ isotope has been identified by studying the genetic relationship with 2.7 min $^{181}\text{Os}$ isotopes$^{1/1}$. Its half-life has been determined by measuring the intensity decrease of the $^{181}\text{Os}$ decay gamma transitions with energies of 118.0 and 144.7 KeV$^{1/1}$ and it is $(5.0 \pm 0.3)$ min. About 20 new transitions have been observed in the $^{181}\text{Ir}$ decay.

$^{180}\text{Ir}$: This isotope has been observed by measuring the intensity decrease of the gamma transitions with the energies of 132 and 276 KeV de-exciting the levels of the ground state rotational band $^{180}\text{Os}^{2/2}$. The half-life of $^{180}\text{Ir}$ is $(1.5 \pm 0.1)$ min.

$^{178}\text{Ir}$: From the observation of gamma-transitions with the energies of 132 and 266 KeV de-exciting the levels of the ground state rotational band $^{178}\text{Os}^{2/2}$, the isotope $^{178}\text{Ir}$ decaying with a half-life of $(0.5 \pm 0.3)$ min has been identified.

*) Polytechnical Institute, Tashkent, USSR.
References

/1/ R. Arlt et al. Preprint P6-4635, Dubna, 1969
/2/ R.M. Diamond, G. Burde, F. Stevens
BETA-TRANSITIONS OF THE $p(p_{3/2}) \rightarrow n(p_{1/2})$ TYPE IN SOME SPHERICAL NUCLEI

B. Kracik and Trần Thanh Minh,
Joint Institute for Nuclear Research,
Dubna, USSR.

The $\beta^+ - \epsilon$ -transition from $1/2^-$ states in some odd-$N$ spherical nuclei to the $3/2^-$ states in even-$N$ nuclei ($^{85m}$Sr$_{47}$ $(1/2^-)$ $^{87}$Rb$_{50}$ $(3/2^-)$; $^{87m}$Sr$_{47}$ $(1/2^-)$ $^{89Zr}_{49}$ $(1/2^-)$ $^{89}$Y$_{50}$ $(3/2^-)$) are in fact transitions of the $p(p_{3/2}) \rightarrow n(p_{1/2})$ type. They are characterized by very low values of logft (4.5, 4.25 and 4.3 respectively).

However, nuclei in the states characterized by $I^\pi = 9/2^+$ like $^{85}$Y$_{46}$ $(9/2^+)$ and $^{89}$Nb$_{48}$ $(9/2^+)$ can undergo $\beta^+ - \epsilon$ - transitions of this type as well. Such transitions would lead to the three-quasiparticle states of the $p(p_{3/2})^{-1} p(g_{9/2})^{1} n(p_{1/2})^{1}$ configurations. If we assume that the $9/2^+$ state in $^{85}$Y is a mixture of $p(g_{9/2})^{1}$ $n(p_{1/2})^{2}$ $n(g_{9/2})^{6}$ and $p(g_{9/2})^{1} n(p_{1/2})^{0} n(g_{9/2})^{8}$ configurations, and the $9/2^+$ state in $^{89}$Nb is a mixture of $p(g_{9/2})^{1} n(p_{1/2})^{2} n(g_{9/2})^{8}$ and $p(g_{9/2})^{1} n(p_{1/2})^{0} n(g_{9/2})^{10}$ configurations, then only the second member takes part in the $\beta$-transition in both cases. This will put the corresponding logft to higher values.

In $^{85}$Sr and $^{89}$Zr there are several high-excited states populated by $\beta^+ - \epsilon$ -decay of $^{85}$Y $(9/2^+)$ and $^{89}$Nb $(9/2^+)$ with the logft values of $\sim 5.5$ and $\sim 5$, respectively [1,2], which may be considered as three-quasiparticle states of the above mentioned type.

It is possible, too, that some of the high-excited states in $^{90}$Zr$_{50}$ [3] are the four-quasiparticle states of the $p_1(p_{3/2})^{-1} p_2(g_{9/2})^{1} n_1(p_{1/2})^{1} n_2(g_{9/2})^{1}$ configuration populated by $\beta^+ - \epsilon$ -decay of $^{90}$Nb$_{49}$ of the $p(p_{3/2}) \rightarrow n(p_{1/2})$ type. Corresponding logft values are rather high
(6.5), but these could be explained by the fact that only the \( n(g_{9/2})^{-1} n(p_{1/2})^{-2} n(d_{5/2})^2 \) admixture to the \(^{90}\text{Nb}\) ground state configuration, which should be expected to be quite small, takes part in the \( p(p_{3/2})\leftrightarrow n(p_{1/2}) \) type \( \beta \)-transition.

1 R. Arlt, B. Kracik, M.G. Loshchilov, G. Musiol, L.K. Peker, Tran Thanh Minh, N.G. Zaitseva
   Preprint JINR 6-5093 (1970)

2 R. Arlt, B. Kracik, G. Musiol, L.K. Peker, Tran Thanh Minh, N.G. Zaitseva
   Preprint JINR 6-5088 (1970)

3 B. Kracik, G. Musiol, L.K. Peker, W.I. Fominich, Tran Thanh Minh, N.G. Zaitseva
   Preprint JINR 6-5106 (1970)
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