# Effects of Pairing and Blocking on the Structure of the Octupole Septuplet of $^{209}$Bi†

J. C. HAFELE
*Washington University, Saint Louis, Missouri*

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The Pauli principle and the nuclear pairing interaction are shown to play a major role in the determination of the level structure of the 2.6-MeV octupole multiplet of $^{209}$Bi.

## I. INTRODUCTION

THE structure of the 2.6-MeV octupole multiplet[1] of $^{209}$Bi appears to give fundamental information about the microscopic excitations that form the octupole vibration. This nucleus represents the addition of a single proton to the doubly-magic $^{208}$Pb core and therefore its low-lying states are expected to have a particularly simple shell-model description. In addition to the single-particle states representing different excitations of the extracore proton, there should exist collective states with properties similar to those of the corresponding collective vibrations of the core. If there is only a weak residual particle-core interaction, a close multiplet is expected for each collective state of the core. This weak-coupling model[2] predicts that the well-known octupole excitation of $^{209}$Bi near 2.6 MeV should consist of seven levels with spins ranging from $\frac{3}{2}^+$ through $15/2^+$. It has been known for some time that the levels of this multiplet are quite close in energy[1,3–6]; they were adequately resolved only recently at the Los Alamos Scientific Laboratory during a high-resolution study of inelastic proton scattering from nuclei in the Pb region.[1,6] The data from this study are consistent with the predictions of the weak-coupling model, and the resulting level energies with their spin-parity assignments are shown in Fig. 1. The observation that the octupole strength of the $^{208}$Pb core is preserved within the seven members of the $^{209}$Bi multiplet to within a couple of percent experiment uncertainty[4–6] indicates the high purity of the assumed particle-core configuration (a proton in the $1h_{9/2}^-$ level coupled to the $3^-$ vibration of the core) and supports the view that there is no appreciable splitting of the levels from admixtures with other particle-core configurations, since any appreciable admixture would destroy this equality by spreading the strength among the admixed states.[7] Thus the analysis of at least the gross structure of the multiplet should require only the first order of perturbation theory. Finally, a striking feature of the observed multiplet structure is the way the residual particle-core interaction pushes the $15/2$ member up and the $\frac{3}{2}$ member down by about 125 keV and leaves the rest of the levels nearly degenerate with the energy of the core state (see Fig. 1). The purpose of this note is to show that precisely this type of gross structure is expected when account is taken of the partial blocking of the microscopic core excitations and of the nuclear pairing interaction that occur when an extracore particle is added to the system.

## II. THEORY

The remarkably successful random-phase approximation[8–10] (RPA) provides a description for the states of even-even nuclei in terms of the microscopic, unperturbed single-particle excitations from occupied shell-model levels below the Fermi level to higher energy, unoccupied levels above the Fermi level and of the mixing of these excitations by the residual interaction between the excited particle and the hole left in the core by the removal of this particle. The collective states come out of this formalism as a coherent superposition of many of these one-particle–one-hole ($1p$-$1h$) excitations with the "collectiveness" of the state increasing with the number of contributing $1p$-$1h$ excitations. Negative-parity states, such as the octupole vibrations, involve predominantly $1p$-$1h$ excitations from the highest filled or partially filled major shell to the next higher (unfilled) major shell and, particularly for the doubly-magic nuclei which have filled major shells, the oscillator spacing between the major shells functions as a large energy gap. The octupole states of doubly-magic nuclei therefore should be well described by assuming that in the ground state there is a sharp Fermi level with levels below completely filled and those above completely empty. The assumption of a sharp Fermi level allows avoidance of the additional complications of a description in terms of quasiparticles and associated occupation parameters. Furthermore, since we desire only a simple, qualitative description

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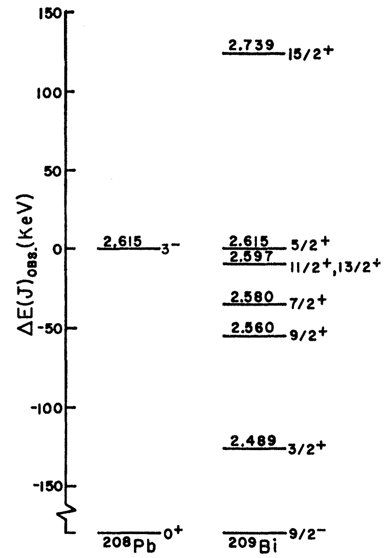


FIG. 1. Octupole states of $^{208}$Pb and $^{209}$Bi. The level energies and spin-parity assignments were taken from Ref. 1. Other known states for these nuclei are not shown and the distance to the ground states is not to scale. The energy splittings $\Delta E(J)_{\rm obs} \equiv E(J)_{\rm obs} - 2.615$ MeV.

of the effects of the addition of an extracore particle, we use the RPA degenerate schematic model which has been fully developed by Brown.[8] In this case all the $1p$-$1h$ excitation energies are taken equal, $\epsilon_{ph}=\epsilon$, and the energy of the core state turns out to be given by

$$E^2=\epsilon^2-2\epsilon\sum_{ph} D_{ph}^2\,, \qquad (1)$$

where $E$ is the energy of the state and $D_{ph}$ is the diagonal matrix element for the particle-hole interaction. The minus sign in Eq. (1) is correct for states with zero relative isotopic spin, which is the case for the octupole vibrations. The sum in Eq. (1) is over all $1p$-$1h$ excitations that can couple to form the spin and parity of the core state. Each term of the sum adds to the coherence of the state and pushes its energy *down* from the unperturbed $1p$-$1h$ energy $\epsilon$. The wave function for the core state is given approximately by

$$|j_c m_c\rangle=\sum_{ph} x_{ph}|j_p j_h j_c m_c\rangle\,, \qquad (2)$$

where the kets on the right represent $1p$-$1h$ states that couple to form the core state, and the square of the amplitude $x_{ph}$ is the probability for the occurrence of the corresponding $1p$-$1h$ excitation. The $1p$-$1h$ annihilation amplitudes that arise naturally in the general PRA treatment have not been included in the wave function of Eq. (1) because their inclusion does not appear to cause any significant change in the results given below. Finally, using Brown's formulation of the degenerate schematic model,[8] it is easy to show that the elements $D_{ph}$ are simply related to the amplitudes $x_{ph}$ by

$$D_{ph}^2=x_{ph}^2\frac{2E(\epsilon-E)}{\epsilon+E}\,. \qquad (3)$$

The addition to the even-even system of a single particle in the level $j_q$ will give rise to a multiplet of $N$ $[=2\,{\rm Min}(j_q,j_c)+1]$ levels with spin-parities $J^\pi$ given by the vector addition of $j_q$ and $j_c$ and $\pi=\pi_q\pi_c$. Any residual particle-core interaction will split the levels of this multiplet. If the residual interaction and hence the energy splittings are weak, a simple, elaborate theory[1,2] in terms of tensor operators of unspecified explicit form can be developed. This approach, however, leaves unspecified interaction parameters that can be determined only by fits to the data. A more fundamental but considerably more complicated approach would be to redo the RPA calculations for the odd-$A$ case[11–13]; although such "correct" calculations are highly desirable, they are quite complicated and tend to obscure the specific effects of the nuclear pairing force and of the Pauli principle on the $1p$-$1h$ excitations that occur unhindered when the extracore particle is not present. We wish to show in an "incorrect" but heuristic way that it is these two effects that give rise to the gross structure of the multiplet shown in Fig. 1.

Let there be an extracore particle in the level $j_q$ and assume that $1p$-$1h$ excitations into this level are the only ones that are affected by the presence of the extracore particle. Since the energy gained by adding a spin-zero correlated pair to the level $j$ is $G(j+\frac{1}{2})$, where $G$ is the pairing force constant,[14,15] we choose the following form for the pairing interaction Hamiltonian:

$$H_{\rm pair}=-G(j_q+\tfrac{1}{2})\delta(j_{pq},0)\delta(j_p,j_q)\,, \qquad (4)$$

where $|j_{pq}m_{pq}\rangle$ is the state formed by coupling the state of the excited particle $|j_p m_p\rangle$ to the state of the extracore particle $|j_q m_q\rangle$. The diagonal matrix elements of $H_{\rm pair}$ between the multiplet states $|j_q j_c JM\rangle$ give the expected energy splittings from the pairing interaction. These matrix elements are derived in the Appendix; they give

$$\Delta E(J)_{\rm pair}=-G\frac{2j_c+1}{2J+1}\sum_{j_h} x_{qh}^2\delta(J,j_h)\,, \qquad (5)$$

where $j_h$ is the spin of the hole level for the $1p$-$1h$ excitations with a particle in the level $j_q$. Thus, a level of the odd-$A$ multiplet is strongly affected by the pairing force only if there occurs in the free core system a strong $1p$-$1h$ excitation with the spin of the hole level equal

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to the spin of the multiplet level. The energy for any such level of the multiplet is pushed down by the pairing interaction.

The pairing interaction of Eq. (4), however, cannot push any of the multiplet levels up. The Pauli principle, on the other hand, can cause level energies to rise because any $1p$-$1h$ excitation with $j_{pq}$=odd is suppressed, and the partial suppression of a $1p$-$1h$ excitation is equivalent to the partial removal of the corresponding term from the coherent sum of Eq. (1) and thus to a smaller lowering of the energy from the unperturbed energy. Since we wish to resurrect those terms that would be naturally suppressed in a properly antisymmetrized, "correct" RPA calculation, it is necessary to use an unantisymmetrized wave function $|j_q j_c JM\rangle_{\rm NAS}$ because otherwise the wave function would vanish when $j_{pq}$=odd. Then, the recognition that the addition to the coherent sum of a particular $1p$-$1h$ excitation depresses the energy approximately by the value of the corresponding element $D_{ph}^2$ leads to the following choice for an effective interaction Hamiltonian for the blocking effect of the Pauli principle:

$$H_{\rm block}=+D_{ph}^2\delta(j_{pq},{\rm odd})\delta(j_p,j_q)\,. \qquad (6)$$

The diagonal matrix elements of $H_{\rm block}$ between the multiplet states $|j_q j_c JM\rangle_{\rm NAS}$ give the expected energy splittings. They are also worked out in the Appendix, where it is shown that

$$\Delta E(J)_{\rm block}=(2j_c+1)\sum_{j_h}D_{qh}^2\chi_{qh}^2\sum_{j_{pq}={\rm odd}}(2j_{pq}+1)\times\begin{Bmatrix}J & j_h & j_{pq}\\ j_q & j_q & j_c\end{Bmatrix}^2, \qquad (7)$$

where { } is a 6-$j$ symbol.[16–18] The sums are over all nonvanishing 6-$j$ symbols and over all nonvanishing amplitudes for $1p$-$1h$ excitations with the excited particle in the level $j_q$; they are more easily (but not necessarily) calculated with the aid of a computer.

## III. RESULTS AND CONCLUSIONS

Since these ideas apply equally well to both proton and neutron $1p$-$1h$ excitations, it is interesting to compare the calculated results for the octupole multiplet of $^{209}$Bi to those for the so-far unobserved octupole multiplet of $^{209}$Pb, a case which represents the addition to the $^{208}$Pb octupole vibration of a single neutron in the $2g_{9/2}^+$ level. Although for this case inelastic scattering studies are impossible because the $^{209}$Pb nucleus is unstable, it may be possible to study the collective states of this nucleus by using two-nucleon transfer reactions; for example, the reaction $^{207}$Pb$(t,p)^{209}$Pb*. For these reasons calculations of $\Delta E(J)_{\rm pair}$ and $\Delta E(J)_{\rm block}$ have been done for the 2.6-MeV octupole multiplets of both $^{209}$Bi and $^{209}$Pb. The amplitudes for proton excitations into the $1h_{9/2}^-$ level and neutron excitations into the $2g_{9/2}^+$ level were taken from the RPA wave function for the 2.6-MeV state of $^{208}$Pb calculated by Gillet *et al.*[19,20] These amplitudes are listed in Table I. The oscillator spacing $\epsilon=41$ MeV$/A^{1/3}$ $=6.92$ MeV as well as the pairing force constant $G=23$ MeV$/A=110$ keV were taken from the literature.[15] The elements $D_{ph}^2$ are related to the amplitudes $x_{ph}$ by Eq. (3), and they were calculated using the above value for $\epsilon$ and the value 2.615 MeV for $E$. The calculated results are listed in Table II and are illustrated graphically in Fig. 2. It can be seen that the pairing interaction concentrates almost entirely on pushing the $\frac{3}{2}$ member of the multiplet down and the effective blocking interaction concentrates almost entirely on pushing the 15/2 member up, while the other levels of the multiplet are hardly split by either of these interactions. This general result is in excellent agreement with the observed multiplet structure shown in Fig. 1.

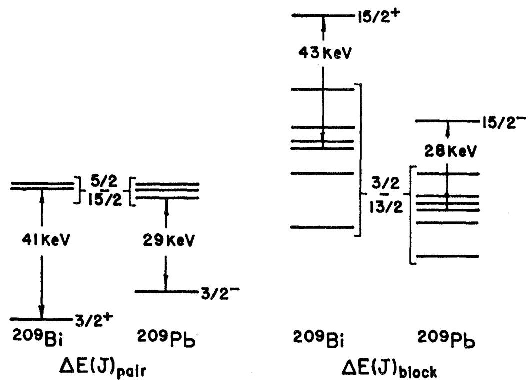


FIG. 2. Calculated energy splittings for the octupole multiplets near 2.6 MeV of $^{209}$Bi and $^{209}$Pb. These results suggest that the energy splittings in the Pb multiplet will not be as large as those in the Bi multiplet.

TABLE I. Particle-hole creation amplitudes for proton excitations into the $1h_{9/2}^-$ level and neutron excitations into the $2g_{9/2}^+$ level. These amplitudes are for the 2.6-MeV octupole state of $^{208}$Pb and were taken from the report of Gillet *et al.*[a]

| $j_h$ (protons) | $2d_{3/2}^+$ | $2d_{5/2}^+$ | $1g_{7/2}^+$ | $1g_{9/2}^+$ |
|---|---|---|---|---|
| $x_{9/2,h}$ | 0.47 | 0.09 | 0.14 | 0.03 |
| $j_h$ (neutrons) | $3p_{3/2}^-$ | $2f_{5/2}^-$ | $2f_{7/2}^-$ | $1h_{9/2}^-$ |
| $x_{9/2,h}$ | 0.42 | −0.19 | 0.15 | −0.03 |

[a] Reference 20.

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TABLE II. Calculated energy splittings from the pairing and effective blocking interactions for the 2.6-MeV octupole multiplets of $^{209}$Bi and $^{209}$Pb. The splittings for the $^{209}$Bi multiplet have been measured[a] and the observed values for this case are included in the table. All values are in keV.

| $J^\pi$ | $\Delta E(J)_{\rm pair}$ | $\Delta E(J)_{\rm block}$ | $\Delta E(J)_{\rm obs}$ |
|---|---|---|---|
| $^{209}$Bi | | | |
| 3/2+ | −42.5 | 58.7 | −126 |
| 5/2+ | −1.0 | 56.2 | 0 |
| 7/2+ | −1.9 | 63.2 | −35 |
| 9/2+ | −0.1 | 48.4 | −55 |
| 11/2+ | 0.0 | 75.0 | −18 |
| 13/2+ | 0.0 | 31.5 | −18 |
| 15/2+ | 0.0 | 98.6 | 124 |
| $^{209}$Pb | | | |
| 3/2− | −34.0 | 39.1 | |
| 5/2− | −4.6 | 37.6 | |
| 7/2− | −2.2 | 41.7 | |
| 9/2− | −0.1 | 33.1 | |
| 11/2− | 0.0 | 48.8 | |
| 13/2− | 0.0 | 22.4 | |
| 15/2− | 0.0 | 65.2 | |

[a] Reference 1.

The actual numerical magnitudes, however, fall somewhat short of being in good agreement with the observed values, but numerical agreement can hardly be expected from such a simple approach and depends, in any case, on the numerical values taken for various parameters. Nevertheless, it does appear that the general result is established and that the Pauli principle and the pairing interaction will play a vital role in any refined calculation of the level structure of these multiplets. Finally, an actual observation of the level structure of the $^{209}$Pb multiplet would give valuable information about the difference between the interaction of a proton and of a neutron with the vibrating core.

## ACKNOWLEDGMENTS

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## APPENDIX

One way to derive Eq. (5) for $\Delta E(J)_{\rm pair}$ and Eq. (7) for $\Delta E(J)_{\rm block}$ is to expand the wave function for the odd-$A$ multiplet $|j_qj_cJM\rangle$ into product wave functions of the form $|j_qm_q\rangle|j_pm_p\rangle|j_hm_h\rangle$, and then recouple $|j_qm_q\rangle|j_pm_p\rangle$ to form $|j_{pq}m_{pq}\rangle$. In the case where antisymmetrization is required, it can be accomplished by transforming

$$|j_qm_q\rangle|j_pm_p\rangle\rightarrow 1/2^{1/2}\times[|j_qm_q\rangle_q|j_pm_p\rangle_p-|j_qm_q\rangle_p|j_pm_p\rangle_q]$$

before recombining to form the coupled system. Since only terms with $j_p=j_q$ arise in the calculation, recombining the antisymmetrized product wave functions to form the coupled system gives[21]

$$\frac{1}{\sqrt{2}}\sum_{j_{pq},m_{pq}}(-1)^{m_{pq}}\left[\begin{pmatrix} j_q & j_q & j_{pq}\\ m_q & m_p & -m_{pq}\end{pmatrix}-\begin{pmatrix} j_q & j_q & j_{pq}\\ m_p & m_q & -m_{pq}\end{pmatrix}\right]|j_qj_qj_{pq}m_{pq}\rangle$$
$$=\sqrt{2}\sum_{j_{pq},m_{pq}}(-1)^{m_{pq}}\begin{pmatrix} j_q & j_q & j_{pq}\\ m_q & m_p & -m_{pq}\end{pmatrix}\times\tfrac{1}{2}[1+(-1)^{j_{pq}}]|j_qj_qj_{pq}m_{pq}\rangle,\quad\text{(A1)}$$

where the equality follows from the relation

$$\begin{pmatrix} j_q & j_q & j_{pq}\\ m_p & m_q & -m_{pq}\end{pmatrix}=(-1)^{2j_q+j_{pq}}\begin{pmatrix} j_q & j_q & j_{pq}\\ m_q & m_p & -m_{pq}\end{pmatrix}.$$

Thus, the antisymmetrized wave function vanishes for $j_{pq}$= odd and gives rise to a factor of 2 for $j_{pq}$= even. Although this procedure makes the origin of the factor 2 obvious and could be continued, it is more elegant and straightforward to generate the required wave functions with a "change of coupling transformation" and then multiply the result by 2.

To zero order the addition of an extracore particle in the level $j_q$ leaves the $1p$-$1h$ composition of the vibrating core unaltered and therefore the wave function for the odd-$A$ case is similar to the core wave function given by Eq. (2);

$$|j_qj_cJM\rangle=\sum_{ph}x_{ph}|j_pj_h(j_c)j_qJM\rangle,\quad\text{(A2)}$$

where the ket notation means that $j_c=j_p+j_h$ and then $J=j_c+j_q$. The transformation from this coupling scheme to the one where $j_{pq}=j_p+j_q$ and $J=j_{pq}+j_h$ involves 6-$j$ symbols[17];

$$|j_pj_h(j_c)j_qJM\rangle=\sum_{j_{pq}}(-1)^{j_q+j_h+j_c+j_{pq}}\times[(2j_c+1)(2j_{pq}+1)]^{1/2}\begin{Bmatrix} j_p & j_q & j_{pq}\\ J & j_h & j_c\end{Bmatrix}\times|j_pj_q(j_{pq})j_hJM\rangle.\quad\text{(A3)}$$

Substitution of Eq. (A3) into Eq. (A2) and then formation of the diagonal matrix elements of $H_{\rm pair}$ gives

$$\Delta E(J)_{\rm pair}=-G(2j_q+1)(2j_c+1)\times\sum_{j_h}x_{qh}^2\begin{Bmatrix} j_q & j_q & 0\\ J & j_h & j_c\end{Bmatrix}^2.\quad\text{(A4)}$$

[21] See Refs. 16–18 for the properties of the 3-$j$ and 6-$j$ symbols.

The factor 2 mentioned above has been applied in the derivation of Eq. (A4). Then, the relationship

$$\begin{Bmatrix} j_q & j_q & 0 \\ J & j_h & j_c \end{Bmatrix} = \frac{(-1)^{J+j_q+j_c}}{[(2J+1)(2j_q+1)]^{1/2}} \delta(J,j_h)$$

gives the result of Eq. (5).

The derivation of Eq. (7) for $\Delta E(J)_{\text{block}}$ is similar to that given above, except that the contributing terms of the wave function are restricted to those with $j_{pq}=$ odd. As is explained in the text, it is necessary to use an unantisymmetrized wave function because, on the one hand, the antisymmetrized form vanishes when $j_{pq}=$ odd [see Eq. (A1)] and, on the other hand, it is precisely these ordinarily vanishing terms that are desired in the derivation of the effective blocking interaction. Moreover, it seems that the factor of 2 mentioned above should not be applied in this case. Formation of the diagonal matrix elements of $H_{\text{block}}$ using Eqs. (A2) and (A3) gives

$$\Delta E(J)_{\text{block}} = {}_{\text{NAS}}\langle j_q j_c JM | H_{\text{block}} | j_q j_c JM \rangle_{\text{NAS}} = (2j_c+1) \sum_{j_h} D_{qh}^2 x_{qh}^2 \sum_{j_{pq}=\text{odd}} (2j_{pq}+1) \times \begin{Bmatrix} j_q & j_q & j_{pq} \\ J & j_h & j_c \end{Bmatrix}^2, \quad \text{(A5)}$$

which by virtue of the symmetry relations for the 6-$j$ symbols is equivalent to Eq. (7).

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# Decay Properties of Neutron-Deficient Osmium and Rhenium Isotopes. II. The $A=180$ Decay Chain*

K. J. HOFSTETTER† AND P. J. DALY
*Department of Chemistry, Purdue University, Lafayette, Indiana*
(Received 1 March 1967)

The decay properties of $Os^{180}$ and $Re^{180}$ have been investigated using Ge(Li), anthracene, and NaI(Tl) detectors in singles and coincidence measurements. A solvent extraction method for rapid milking of daughter rhenium activities from osmium was used to confirm the genetic relationship between 21.5-min $Os^{180}$ and 2.45-min $Re^{180}$. No radiations other than $K$ x rays were detected in $Os^{180}$ decay and it is concluded that this isotope decays by electron capture with a log $ft$ value in the range $4.7\pm0.3$ In $Re^{180}$ decay the following radiations were identified: $K$ x rays (105), $\gamma$ rays of energies $103.6\pm0.3$ (26), $232.6\pm1.0$ (1.2), $750.8\pm1.0$ (0.8), $826.4\pm0.8$ (11), and $902.2\pm0.5$ keV (100), and a single positron group of endpoint $1.76\pm0.4$ MeV (8). Relative intensities are given in parentheses. A $Re^{180}$ decay scheme which is consistent with all the observations, including $\gamma$-$\gamma$ coincidence measurements, is proposed; its most interesting feature is a two-quasiparticle state ($J^\pi=2^-$) at 1006 keV, which de-excites almost exclusively to the first excited state of $W^{180}$. The disintegration energy of $Re^{180}$ was determined to be $3.78\pm0.04$ MeV.

## I. INTRODUCTION

SINCE the discovery of 2.45-min $Re^{180}$ in 1955,[1] no investigation of the $W^{180}$ levels populated in its decay has been reported. This is in striking contrast with the large volume of published work describing attempts to locate and characterize the levels of the analogous nucleus $W^{182}$ populated in the $\beta$ decay of $Ta^{182}$ and $Re^{182}$. Extensive data on the ground-state rotational band of $W^{180}$ have resulted from $(\alpha,xn)$ reaction spectroscopy studies[2] and a 5.2-msec metastable state at 1525 keV has been recently characterized[3] as a $8^-$ state which de-excites by a strongly $K$-forbidden $E1$ transition to the $8^+$ member of the ground-state band. Graetzer *et al.*[4] have used conversion-electron spectroscopy to investigate the vibrational states of $W^{180}$ populated in the $Ta^{181}(p,2n)W^{180}$ reaction.

Recently, Hofstetter and Daly[5] have disproved the existence of the 20-hr positron-emitting isomer previously assigned to $Re^{180}$. Evidence for a new isotope, 21.5-min $Os^{180}$, has been independently reported by two groups[5,6] but no characteristic $Os^{180}$ radiations were identified in either case. This paper describes more detailed investigations of the radiations emitted in the decay of $Os^{180}$ and $Re^{180}$. Many of the experimental

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