HYPERFINE STRUCTURE IN LINE SPECTRA AND NUCLEAR SPIN

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by

S. TOLANSKY
D.Sc., Ph.D., D.I.C.

WITH 29 DIAGRAMS

METHUEN & CO. LTD. LONDON
36 Essex Street, Strand, W.C.2
PREFACE TO THE SECOND EDITION

During the twelve years that have passed since the writing of the first edition, considerable advances have been made both in the theoretical and experimental aspects of the study of hyperfine structure in line spectra. This has of necessity resulted in a good deal of re-writing in the preparing of a second edition. Chapters I, II, VIII, X, XI, and XII are largely new, the other chapters being modified mainly in details.

Nuclear spins are now known for 55 different atoms, including among them 79 isotopic species, and for 65 of these, nuclear magnetic moments have been determined. These relatively large numbers enable an approximate statistical approach to be made, but it cannot yet be said that theories of the origin of either nuclear spin or nuclear magnetic moment are out of the hypothetical stage. Although a number of such theories have of course been proposed, it is considered that as yet they are not sufficiently developed to be included in an introductory book such as this.

As far as possible, however, nuclear data have been brought up to date. References have been omitted in compiling the tables of moments, since the data have been compiled from a considerable number of original papers and a cumbersome list would be out of place.

A brief chapter has been added dealing with nuclear electrical quadrupole moment, since this promises to be of especial future importance. Isotope effects are treated more comprehensively than formerly.

S. TOLANSKY

MANCHESTER, January, 1946
PREFACE

This book is meant as an introduction to the study of fine structure in line spectra for those who are already familiar with the elements of spectroscopy. The study of fine structure has already led to many important facts concerning the constitution of atomic nuclei, and will certainly lead to more. It is hoped that this short survey will prove useful to the student who contemplates undertaking research in this branch of the subject. The experimental and theoretical material is accumulating at a very rapid rate, so that no claim to completeness can be made, but the more important facts known up to date have been included. It has not been considered advisable, because of lack of space, to treat the Zeeman effect in detail, so that only a brief summary is given in an appendix.

I take this opportunity of expressing my thanks to Dr. W. E. Williams of King's College, London, who has completely revised the manuscript and made many valuable and helpful suggestions.

I also wish to thank Miss E. Pincasovitch, who has kindly executed all the diagrams for me.

S. T.

Imperial College of Science,
London, February, 1935

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CHAPTER I

INTRODUCTION

It is now some fifty years since spectroscopists first found that some of the lines of the spectra of a number of elements were in reality complex close line groups. This complex close structure, revealed only by instruments of high dispersion and high resolution, was formerly called "fine structure" but for sound reasons this terminology has generally now been replaced by "hyperfine structure".

With improvements in technique, it has now been established that in many spectra the majority of the lines are complex. Such a complex usually consists of a small number of components, rarely exceeding perhaps a dozen, and often with no apparent easily recognised regularity within the complex pattern. These hyperfine structures, for long a puzzle, assumed considerable importance when it was recognised that atomic nuclear properties are responsible for their appearance and at once a new and powerful method for studying atomic nuclei was developed. It has been established beyond all doubt that these hyperfine structures are due either to (a) isotopic mixtures or (b) spin properties of the atomic nuclei.

As a result of extensive spectroscopic investigation, important contributions to our knowledge of atomic nuclei have been made. The physics of the nucleus has become a study of fundamental importance, and recent developments, such as the evolution of the atomic bomb, based on nuclear fission, point to the bigger and bigger place that nuclear physics will hold.

There are available perhaps some five broadly distinct experimental methods for studying nuclear properties. Taking them historically, these are (a) Radioactive, (b) Spectroscopic, (c) Magnetic, (d) Collision-disintegrative, (e) Cosmic rays.

Radioactive experiments first established the nuclear structure of the atom and revealed details about nuclear
structure, nuclear energy levels and nuclear constituent particles. Leaving for later the discussion of the spectroscopic methods, the magnetic technique, whilst limited in application, is of very high precision. It is found that a jet of atoms, when passing through a non-uniform magnetic field, suffers deflection which, when interpreted, yields the mechanical and magnetic spin properties of the nuclei of the atoms constituting the jet. This magnetic method correlates very closely with the spectroscopic method.

So important and well known are the collision-disintegration experiments made with both natural radioactive projectiles and with accelerated particles from machines like cyclotrons, that little need be said about the fundamental discoveries and contributions made by this study. Cosmic ray experiments have contributed in a very fundamental way to the study of the nucleus and one need but mention the discoveries of the positron and meson to realise this.

The purpose of this book is to discuss one aspect of the lesser known spectroscopic method of approach, namely that of hyperfine structure in line spectra and it will be shown that not only has a great deal already been found by this technique, but that there is considerable future scope in addition.

Two distinct spectroscopic methods for investigating problems of nuclear structure have been developed. The first method involves the difficult measurement of the alternations in intensity found in the rotation lines of many molecular band spectra. The simplest type of molecule is of the homogeneous kind consisting of two similar molecules such as \( \text{O}_2 \) or \( \text{N}_2 \). A homogeneous molecule is one in which the constituent atoms are of the same isotopic species, e.g. \( \text{O}_{16-16} \) (but not \( \text{O}_{16-17} \)). When such a homogeneous molecule emits an electronic band spectrum, alternate rotation lines of the band bear a fixed ratio to one another. It can be shown on theoretical grounds that the alternation of intensity arises because of a nuclear spin, the spin being characteristic for the particular constituent atoms forming the molecule. Each nucleus has a mechanical spin moment which is some small integral or half integral, multiple of \( h/2\pi \), \( h \) being Planck's constant. The intensity ratio depends, then, upon the number \( I \) (called the nuclear spin) where \( I = \frac{h}{2\pi} \) is the mechanical moment. This ratio is \( \frac{I+1}{I} \).

The intensity ratio arises as a purely statistical effect, alternate transitions being more probable than their neighbouring transitions by virtue of the mechanical moment.

It can be shown by various methods that the mechanical moment of the hydrogen nucleus (this is a single proton) is \( \frac{1}{2} \) \( \left( \frac{h}{2\pi} \right) \) and it is now considered probable that an \( I \) greater than \( \frac{1}{2} \) for stable atomic nuclei in unexcited states is not found. The relatively small values of \( I \), especially in heavy atoms containing some two hundred particles, is an indication that possibly only a small fraction of the nucleus is involved in the spin property.

Since the nucleus is electrically charged, one expects that the mechanical spin will produce an associated magnetic moment, which indeed it does. (It will be shown later that even uncharged spinning nuclear particles can produce also a magnetic moment.) This nuclear magnetic moment couples with any other electronic magnetic moment of the atom. Each valency electron, which is involved in the production of the line spectrum, has a magnetic moment and this couples with the nuclear magnetic moment. As a result, the spectroscopic levels characteristic of the valency electrons are split up into a number of sub-levels.

The magnetic moment of the proton, being inversely as the mass, is some two thousand times smaller than that of the electron. Now, as will be indicated later, the separation of lines of a multiplet in a spectrum depend largely upon the spin magnetic moment of the electron. The separations due to nuclear spin are thus considerably smaller than those due to electron spin, hence the use of the term "hyperfine structure" to distinguish this splitting from "gross multiplet structure", sometimes called "fine structure".
It is clear now why the study of hyperfine structure can be expected to reveal information both about nuclear mechanical and nuclear magnetic moments. But recent work has revealed much more. Interesting isotope effects have been discovered, the statistical distribution of nuclear spins has proved to be of value, and unsymmetry in nuclear electrical charge distribution has been demonstrated.

The measurement of hyperfine structure presents a number of particular experimental difficulties due in the first place to the minute scale of the structures. Most readers will be familiar with the close doublet gross structure multiplet separation of the sodium D lines, usually just separated by typical small dispersion instruments. The wave number separation for this doublet is of the order 25 cm.\(^{-1}\) (some 6 angstrom units). Taking this as a standard, the problems involved will be realised when it is noted that hyperfine structures of \(0.010\) cm.\(^{-1}\) are commonly met with and much smaller structures occur and are detectable only with the most refined techniques.

There are two experimental problems, namely the development of suitable line sources and the employment of high resolution instruments. This experimental aspect is being treated by the writer in a separate and much more comprehensive treatise. The salient features will be briefly reviewed in the next chapter, which will deal with line width, sources and instruments.

REFERENCE

\(^1\) Tolansky, "High Resolution Spectroscopy" (Methuen).
produced by any closely approaching neighbouring atoms. In absorption observations these are neutral, but in emission some will be charged. Highly complex effects result, leading to effective term broadening increasing proportional to the pressure and depending in a very marked way on the nature of the perturbing gas. Typical collision broadenings in atmospheric pressure arcs are of the order of 0.4 to 1.0 cm$^{-1}$. To reduce these effects, pressures below 1 mm. are desirable.

**Resonance broadening**

This is a particularly severe form of pressure broadening, which sets in when the perturbing atoms are of the same kind as the emitter, e.g. sodium in sodium vapour as distinct from sodium in, say, nitrogen gas. At a pressure of 0.2 mm. of sodium vapour the resonance width of the D lines is 0.6 cm$^{-1}$, which is one hundred times that produced by the same pressure of foreign gas.

**Stark effect broadening**

When emitting or absorbing atoms are in a strong uniform electric field, spectral levels are split up and the lines become complex narrow groups, the extent of the pattern, according to circumstances, being proportional to either the field strength or to its square. A non-uniform field leads to an effective line broadening. In bad cases a field of only 1000 volts/cm. can lead to line broadening of some cm$^{-1}$. The remedy is to reduce the effective field in the region of emission.

Interatomic stark broadening is a related effect produced by local fields set up by ion concentrations. It is clear that powerful currents and condensed discharges always lead to prohibitive line broadening and must therefore be avoided. This feature has been a principal factor in stimulating the development of special sources.

**Self reversal broadening**

When radiation from a hot gas passes through a cooler layer of the same vapour, the central portion of the emission
line is preferentially absorbed. The amount of absorption depends upon the absorption coefficient and on the square root of the total number of absorbing atoms in the line of sight. A 20 cm. length of path in sodium vapour at a pressure of 0.01 mm. (300° C.) leads to a line width of some 2 cm.\(^{-1}\).

When the central absorption becomes pronounced, a spurious doublet results, as shown in Fig. 1. In Fig. 1A, \(a, b, c\) show successive stages of reversal broadening. In Fig. 1B the further stages leading to reversal doubling are illustrated.

![Diagram](image)

**Fig. 1.**—A. Shows how a line is broadened by self-absorption. B. Shows how further self-absorption produces reversal.

**Development of special sources**

To combat the effects of line broadening special sources have been developed. The main essentials are that brilliance is required, yet the temperature, pressure and field strength must be kept low. A brief résumé of the more useful sources is as follows.

1. **Hollow cathode discharge**
   This is a glow discharge maintained at a pressure of about a millimetre in a rare gas. A direct current supply is used and the cathode, a hollow cylinder, is usually cooled, often with liquid air. Pressure, Stark and temperature broadening effects are all small and the brilliance is high.

2. **Low pressure arc**
   These have limited application but are of some use.

3. **Cooled low pressure Geissler discharge**
   A good deal can be done with such sources when carefully designed.

4. **High frequency electrodeless discharge**
   A high frequency discharge can be obtained by means of a thermionic valve oscillating on moderately short wavelengths (say 20 m.). The merit of this discharge is that it is quite bright at fairly low pressures and only small electric fields are involved, the result being very sharp line emission.

5. **Atomic beam**
   If a jet of atoms, moving in a high vacuum is viewed at right angles to the direction of travel of the atoms, then the resolved component of velocity in the line of sight is so small that Doppler width is very greatly reduced. It is possible by this means to reduce Doppler width to that which would result if a source were cooled to the abnormally low temperature of only 6° Abs.
   With this source the highest resolving powers can be used and the structures measurable are very small.

**High resolution instruments**

Not only must the lines be sharp, but the instruments employed must have high resolution and dispersion. Interferometers must generally be employed and the apertures and light losses are such with these instruments that brilliance in emission becomes a necessity. When it is remembered too that some components within a hyperfine structure are often very weak and that long exposures are experimentally difficult to obtain, the importance of source brightness is further emphasised.
The most important instrument in use is the variable gap Fabry-Perot interferometer, silvered for the visible and infra-red and aluminised for the ultra-violet. The compound Fabry-Perot (two different etalons in tandem) has found useful application also.

The quartz Lummer plate has contributed a good deal, particularly in the ultra-violet region. Echelons, both reflection and transmission, have been employed, but the application of all three instruments (the Fabry-Perot excluded) is restricted by the fixed dispersion. For the coarser structures, the higher orders of gratings have been utilised.

CHAPTER III

MULTIPLETILITY IN LINE SPECTRA

In order to be able to understand the origin of the hyperfine structures of spectra, it is necessary to have a clear understanding as to how gross structure multiplets arise. It is not possible to give a comprehensive account here, and for a full treatment the reader is referred elsewhere. The more important general features, in so far as they bear upon the subject of hyperfine structures, will, however, be discussed here. It will be assumed that the reader is familiar with the elementary Bohr theory of the hydrogen atom, and also understands how the Periodic Table has been built up by the process of completion of the electron shells. An excellent introductory account is given in Johnson’s “Spectra” in this series. The simplest type of spectrum, namely that arising from a hydrogen-like atom such as sodium, will first be discussed. Such an atom possesses only one valence electron, the transition movements of which give rise to the optical spectrum. The discussion will then be extended to two electron and later to more complex cases by the help of a vector model. Although the wave mechanics has shown that the vector model is not quite exact, yet it is extremely helpful and valuable as an aid to the solution of problems and sufficiently correct for our purposes.

One-electron spectra

A convenient example of a hydrogen-like atom is sodium, which consists of a nucleus and eleven outer electrons. In this atom, two inner electron shells, namely the K and L shells are completely filled, and in the third shell there is one electron, the valence or optical electron. For spectroscopic purposes the first shell is denoted by the number 1, the second by 2, etc. In the vector model of the atom,
each electron is considered to be spinning, the mechanical moment of the spin being the same for each electron and equal to \( s \frac{h}{2\pi} \), where \( s = \frac{1}{2} \) and \( h \) is Planck’s constant.\(^3\)

\( s \) is called the spin quantum number of the electron. Since, in addition, an electron rotates in its orbit, it also possesses an orbital angular momentum, and the value of this orbital momentum is \( l \frac{h}{2\pi} \) for each electron, where \( l \) is called the orbital quantum number. In contradistinction to the spin, different electrons may possess different \( l \) values; they may have any of the values 0, 1, 2, 3, 4, etc., according to the eccentricity of the orbit in which the electron moves.

**TABLE I**

<table>
<thead>
<tr>
<th>I value</th>
<th>0, 1, 2, 3, 4, ...</th>
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<tbody>
<tr>
<td>Electron designation</td>
<td>s, p, d, f, g, ...</td>
</tr>
<tr>
<td>Atomic state</td>
<td>S, P, D, F, G, ...</td>
</tr>
</tbody>
</table>

In general an electron moves in an elliptical and not in circular path around the nucleus, and orbits of the same major axis, but of different eccentricities, are possible for a given electron. The numerical value of \( l \) is a measure of the ellipticity of the orbit, the smaller the \( l \) value, the more eccentric is the path. An orbit of given eccentricity is designated by a letter, the letters corresponding to different \( l \) values. The particular letters which are employed are so used because of historical associations. Rows one and two of Table I show the electron designations for various \( l \) values.

It will be recalled that the first orbit is termed a \( 1s \) orbit, the second a \( 2s \) orbit, etc., and an electron in any orbit may have different \( l \) values. Thus an electron in the first orbit with \( l = 0 \) is called a \( 1s \) electron and similarly, for example, an electron in the second orbit with \( l = 1 \) is designated as a \( 2p \) electron. Using this notation the complete structure of the normal sodium atom, which possesses eleven electrons, is written as \( 1s^2 \), \( 2s^2 \), \( 2p^6 \), \( 3s \), meaning that there are two \( 1s \) electrons, two \( 2s \), six \( 2p \) and one \( 3s \) electron in the whole atom. The number of similar or “equivalent” electrons, as they are called, is indicated by the index. The optical spectrum is produced by movements of the \( 3s \) electron, for the inner shells, the first with two and the second with eight electrons, are completely filled, and only contribute to the spectrum in so far as they screen off the nuclear charge from the valence, or optical electron, as it is called.

**Spectral terms of a one-electron atom**

By “one-electron atom” is meant an atom with one valence electron, such as sodium or caesium. In such an atom only this one valence electron comes into consideration spectroscopically. The electron possesses both a spin and an orbital motion, and being a charged particle, these two angular momenta produce magnetic fields. The spin magnetic field and the orbital magnetic field interact with each other so that in effect the mechanical spin \( s \) \(^*\) and the mechanical orbital moment \( l \) combine together to form a resultant by virtue of their magnetic fields. This resultant, called the inner quantum number, is denoted by the letter \( j \) and is obtained by the vectorial addition of \( l \) and \( s \), that is \( j = l + s \). Since \( s \) is constant and equal to \( \frac{1}{2} \), and as \( l \) varies for different types of orbit, it follows that the inner quantum number also varies for orbits of different eccentricities. The sodium atom in the normal unexcited state has the optical electron in the \( 3s \) orbit and as seen from Table I the \( l \) value for an \( s \) orbit is zero, so that \( j = \frac{1}{2} + 0 \), i.e. \( j = \frac{1}{2} \). The same is true for \( 4s \), \( 5s \), \( 6s \), etc., orbits.

**Term notation**

So far only the notation for the optical electron itself has been considered, and a further symbol is required to

* It is usual to omit the \( \frac{h}{2\pi} \) factor since this is always understood.

The letters \( l, s, j \) represent vectors, and are printed in heavy type.
describe the energy of the whole atom when the optical electron is in a given orbit. When this electron has a given \( l \) value the atom itself is said to be in a given state, corresponding to the notation given in the third row of Table I, and it is seen that in the case of a one-electron atom, the letter designating the state is the same as that describing the electron, except that capital letters are used for atomic states and small letters for electron orbits. An atomic state is virtually a description of the energy and position of the optical electron, and the different states, or different energy values, which make up all the energy levels out of which the spectrum arises, are called terms. To make this clearer, it will be seen that, for example, a number of \( S \) states are possible for \( 3s \), \( 4s \), \( 5s \), etc., electrons, and each of these is called an \( S \) term.

In describing a term, the \( j \) value is written to the right. For example, in the term arising from a \( 3s \) electron, since \( j = \frac{1}{2} \) a \( ^2S_\frac{1}{2} \) term is formed. The left-hand upper subscript is called the term multiplicity, and indicates the maximum number of values \( j \) can have. For it will be remembered that \( j \) is the resultant of the vectorial combination of \( l \) and \( s \), and in general a number of resultants are possible. If \( l \) is greater than zero, \( j \) cannot have more than \( 2s + 1 \) values. This can be seen as follows. Since \( s \) adds vectorially to \( l \) the values of \( j \) vary between \( l + s \) and \( l - s \), each value diminishing by one, i.e.

\[
(1 - s), \quad (l - s + 1), \quad \ldots \quad 1, \quad (l), \quad 0, \quad (l + 1), \quad \ldots \quad (l + s - 1), \quad (l + s).
\]

There are therefore \( 2s + 1 \) values, so that the term multiplicity is given by \( 2s + 1 \). In the case considered above, as \( s = \frac{1}{2} \) the multiplicity is two, and the term is called a doublet term. In more complex atoms \( s \) can have other values, and when \( s = \frac{1}{2}, \frac{3}{2} \), \( 2, \) etc., the terms are called doublet, triplet, quartet, quintet, etc., the multiplicity subscripts being 2, 3, 4, 5, respectively.

Consider what happens when the sodium atom is excited so that the optical electron is forced out of its normal \( 3s \) orbit into the higher \( 3p \) orbit. Since the electron is now in a \( p \) orbit, the \( l \) value is 1 (see Table I) and the atom is in a \( P \) state. Two values for \( j \) are now possible, that in which the electron spin \( s \) is parallel with \( l \) and that in which it is anti-parallel. There are thus two values of \( j \) namely, \( j = l + \frac{1}{2}, \) i.e. \( j = 1 \pm \frac{1}{2} \), so that \( j \) may be equal to \( \frac{3}{2} \) or \( \frac{1}{2} \). It follows that two terms are formed, these being \( ^2P_{\frac{3}{2}} \) and \( ^2P_{\frac{1}{2}} \). The terms are doublet terms, as \( s = \frac{1}{2} \) for the single electron concerned, and, as just shown, multiplicity is \( 2s + 1 \). It is easy to see that a \( 3d \) electron (with \( l = 2 \)) will produce \( ^2D_{\frac{5}{2}} \) and \( ^2D_{\frac{3}{2}} \) terms in a similar manner. Since, say, a \( ^2P_{\frac{1}{2}} \) term may be produced by either a \( 3p \), \( 4p \), \( 5p \), etc., electron, it is customary to specify the terms more completely by writing the designation of the optical electron before the term. For instance, the \( ^2P_{\frac{1}{2}} \) terms produced by \( 3p \) and \( 4p \) electrons respectively, are described as \( 3p \cdot ^2P_{\frac{1}{2}} \) and \( 4p \cdot ^2P_{\frac{1}{2}} \). Other examples are, say, \( 4d \cdot ^2D_{\frac{5}{2}} \) and \( 5f \cdot ^2P_{\frac{1}{2}} \), etc. A line transition from a \( 4p \) to a \( 3s \) orbit would be written as \( 3s \cdot ^2S_{\frac{1}{2}} - 4p \cdot ^2P_{\frac{1}{2}} \) as a particular instance.

**Terms in alkali spectra**

The terms which arise in alkali spectra are shown in Fig. 2. All the terms other than \( S \) terms are double, since two values for \( j \) are permitted, whereas for the \( S \) terms \( j = \frac{1}{2} \).

In the two components of a doublet term, e.g. \( 4p \cdot ^2P_{\frac{1}{2}} \) and \( 4p \cdot ^2P_{\frac{1}{2}} \) that with the smaller \( j \) usually lies deepest, and such a spectrum is called a normal spectrum. In some atoms the terms are completely inverted, the biggest \( j \) value lying deepest and there exist regular intermediate stage gradations between normal and inverted spectra.

A group of terms such as \( 3p \cdot ^2P_{\frac{1}{2}}, 4p \cdot ^2P_{\frac{1}{2}}, 5p \cdot ^2P_{\frac{1}{2}}, \) etc., is called a sequence of terms, and each sequence tends to a common limit which represents the ionisation limit of the atom. When the optical electron is lifted to this limit it does not fall back, but leaves the atom, which is then charged, and is called an ion. The number 3, 4, 5, 6, etc., before the \( p \) in the term designation is called the current total quantum number \( n \), and Fig. 2 shows that the intervals
between the components of a doublet term steadily diminish as \( n \) increases.

When the electron is excited and lifted up to a high orbit, it tends to fall back to a lower orbit. Not all transitions are possible, and amongst those which do occur some are more probable than others. Transitions are limited by a selection principle which states that \( I \) must change by \( \pm 1 \) and \( J \) by \( \pm 1 \) or 0, the transition \( \Delta J = 0 \rightarrow 0 \) being excluded. Under special conditions of excitation these rules may be violated, but such departures need not concern us at present.

As \( \Delta I = \pm 1 \), transitions from \( S \) to \( P \) or \( P \) to \( D \) (or vice versa) are allowed, but an \( S \) to \( D \) transition is, for example, forbidden. All the transitions, from, say, a given \( ^2P \) to a given \( ^2D \) state, that are allowed by the \( J \) selection

---

**Fig. 3.** A typical doublet multiplet; note that three lines are produced but the doublet name is derived from the term multiplicity.

A \( ^2P - ^2D \) multiplet arises. Some transitions are more probable than others, resulting in lines of differing intensity, and in the illustration the heights of the lines are drawn proportional to the line intensities. The intensities of the components of a multiplet obey laws which are very similar to those obeyed by hyperfine structures, and since the latter have to be considered in detail, it will be sufficient to state the “sum rule” which governs the relative intensities of the components of a given multiplet. Simply stated, the rule is that the sum of the intensities of all the lines coming
to any given j level is proportional to 2j + 1. It is to be noted that a doublet multiplet can consist of more than two lines; the doublet appellation refers to the multiplicity of the terms and not to the number of lines in the multiplet patterns.

Two-electron spectra: Electron couplings

Having examined the alkali-like spectra, the next step is to discuss the spectra produced by atoms which have two valence electrons, e.g., Ca or Hg, etc. The two valence electrons interact with each other, and as the principles employed here are fundamentally the same as those for more complex electron systems, they will be discussed in detail. As in the previous case considered, each of the two valence electrons possesses a spin angular momentum \( s = \frac{1}{2} \), and an orbital angular momentum \( l \) which may have any of the values 0, 1, 2, 3, 4, 5, etc. There are four vectors to take into consideration, and these may combine in different ways to form a resultant \( J \) which will represent the total angular momentum of the atom. (In the vector notation small letters are used to represent specific angular momenta for individual electrons, and capital letters for resultant sums.) The vectors may combine in a variety of ways and one may distinguish two extreme types of coupling called LS* and \( jj \) respectively.

LS coupling

There are four vectors which will couple in a two-electron atom, namely, \( s_1 \) and \( s_2 \), the two electron spins, and \( l_1 \) and \( l_2 \), the two orbital angular momenta. In LS coupling, the binding between the spins \( s_1 \) and \( s_2 \) is so strong that these form a resultant spin \( S \). Correspondingly, the binding between the two orbital angular momenta \( l_1 \) and \( l_2 \) is strong and these add up to form a resultant orbital angular momentum \( L \). This means that the linkage of the mutual spins is so strong that it is not influenced by the orbital angular momenta; at the same time, the linkage of the orbital angular momenta is such that the electron spins cannot affect it. There are now two resultant vectors \( S \) and \( L \) and these in turn couple to form a resultant \( J \). As in the simple alkali case, the multiplicity is now 2\( S + 1 \).* Furthermore the term type is determined by the numerical value of the resultant \( L \), i.e., for \( L = 0, 1, 2, 3, 4, 5, \) etc., the terms are \( S, P, D, F, G, H, \) etc. It is important to notice that the term letter notation is not the same in capitals as the letter notation of the optical electron in small letters, as was the case with hydrogen-like spectra. The selection rules now become

\[
\Delta J = \pm 1 \text{ or } \Delta [o \rightarrow o \text{ excluded}]
\]

\[
\Delta L = \pm 1.
\]

*The type is often termed Russell-Saunders coupling, after its discoverers.

jj coupling

Whereas in LS coupling all the spins unite to form a resultant spin \( S \), and all the orbital momenta form a resultant \( L \), the case called \( jj \) coupling is different. The same four vectors \( s_1, s_2, l_1, l_2 \) are under consideration, but the mode of combination is altered. Here the coupling between the \( s_1 \) and \( l_1 \) of the one electron is so strong that it forms a resultant \( j_1 \). Similarly the spin and orbital momenta \( s_2, l_2 \) of the other valence electron form a resultant \( j_2 \). The other possible influences are not sufficient to break these pairings so that the two resultants \( j_1 \) and \( j_2 \) now couple to form a resultant \( J \). Summarising then, in LS coupling all the \( l \) values form a resultant \( L \) and the \( s \) values a resultant \( S \) which then combine to form \( J \), whilst in \( jj \) coupling the \( l \) and \( s \) of each electron combine to a resultant \( j \) for each electron, and then the \( j \) values add up to form \( J \).

Regular intermediate stages exist between LS and \( jj \) couplings and these are conveniently shown below.†

* This type is often termed Russell-Saunders coupling, after its discoverers.

† See Pauling and Goudsmit, "Structure of Line Spectra" (1931).
HYPERFINE STRUCTURE

\[
\begin{align*}
\{(l_1l_2)(s_1s_2)\} &= \{LS\} = J \\
\{(s_1l_1)(s_1l_2)\} &= \{j_1j_2\} = J \\
\{(s_1l_1s_2l_2)\} &= \{j_1s_2l_2\} = J \\
\{(s_1l_1l_2s_2)\} &= \{j_1l_2s_2\} = J
\end{align*}
\]

Two vectors within a bracket are directly coupled together. The first and second cases are the LS and \( jj \) types already mentioned.

In the third case \( s_1l_1 \) give a resultant \( j_1 \) which breaks up the link between \( s_1l_1 \) and first combines with \( s_2 \). The vector pair \( j_1s_2 \) then combines with \( l_2 \) to give \( J \). The fourth case is similar, except that \( j_1 \) first binds itself with \( l_2 \) instead of \( s_2 \).

It can be shown that no matter what manner of coupling may exist, the number and type of the terms which are produced by the vector combinations remain the same; this has considerably simplified analysis. The multiplicity has been defined as \( 2S + 1 \), but since it is only in pure LS coupling that a resultant \( S \) exists, the multiplicity for \( jj \) or intermediate stage coupling has, in this sense, no real meaning. As, however, the number and type of terms are independent of the nature of the coupling, it is convenient to use the same notation, for the multiplicity gives the number of terms, independent of the existence of \( S \). The greater number of actual spectra hitherto analysed are cases of the LS type of coupling, which will be illustrated by an example.

Graphical representation of coupling processes

As a typical example of a two-electron spectrum of the normal type, calcium will be considered. In the normal state calcium has two valence electrons, both of which are in \( 4s \) orbits and either of these may be excited, either separately or simultaneously. Let us first consider the case when only one of the valence electrons is excited, the other remaining permanently in its \( 4s \) orbit. The optical electron may go either into a \( ns, np, nd, \) etc., orbit according to the excitation energy given to it, and the two electrons with their different energies interact to give a resultant \( J \).
\[ I_1 + I_2 = 0 \] for both electrons are in \( s \) orbits. Hence as \( J = L + S \) a \( ^3S_1 \) term is formed. The optical electron being also capable of moving to higher \( s \) orbits, a whole sequence of \( ^3S_1 \) terms is similarly produced.

The graphical representation of Fig. 4 enables us to see what occurs when the electron is lifted into a \( p \) orbit. As before, the \( 4s \) electron has \( I_1 = 0 \) and \( S_1 = \frac{1}{2} \), while the \( p \) electron has \( I_2 = 1 \) and \( S_2 = \frac{1}{2} \). The two spins \( S_1 \) and \( S_2 \) have formed a resultant \( S = 1 \), and similarly \( I_1 + I_2 = 1 \). From this it is seen that \( J \) can have any of the three values \( 0, 1, 2 \), and the terms \( ^3P_0, ^3P_1, ^3P_2 \) are produced. From the graphical point of view, it is very useful if we start from the \( ^2S_{1/2} \) ion term. That is to say, consider the atom ionised, in which state the optical electron is completely removed, leaving the other valence electron in the \( 4s \) orbit.

The ionised atom will now be in a state identical with the normal state of the one-electron type of atom, i.e. in a \( ^2S_{1/2} \) state. Now, to produce a Ca I atom with one electron in a \( p \) orbit all that is necessary is to add the vector representing the \( p \) electron to the vector representing the \( ^2S_{1/2} \) state of the Ca II atom. This is shown in Fig. 4. The left, A, shows the \( ^2S_{1/2} \) ionic term of Ca II in which the only existing vector is the electron spin, \( S_1 = \frac{1}{2} \), of the \( 4s \) electron. This is shown vectorially by a thick hollow arrow in the vector picture at the right of the term. The direction of the resultant \( J \) is shown by means of a dotted arrow, and of necessity coincides with the direction of \( S_1 \) in the \( ^2S_{1/2} \) term. The effect of adding a \( 4p \) electron to the \( 4s \) electron is shown vectorially at B. The orbital angular momentum \( (I_2 = 1) \) of the \( 4p \) electron is shown by the long thin black arrow, and its spin \( (S_2 = \frac{1}{2}) \) by means of a short thick black arrow. [The reason for distinguishing between the spins of the \( 4s \) and \( 4p \) electrons, although both equal \( \frac{1}{2} \), will appear later.] The four terms which result from the singlet and triplet combinations of \( 4s \) and \( 4p \) are shown at the right. The first term is the \( ^1P_1 \) term previously discussed. In the second all the vectors act parallel, and so produce a \( ^3P_2 \) term. The third term shows how both the spins couple to a resultant \( S = 1 \), which then combines at an angle with \( L = 1 \) to give a resultant \( J = 1 \). In the fourth the two spins still act parallel, but their resultant opposes \( L \) and a \( ^3P_0 \) term is produced.

A schematic diagram of the more important terms in Ca I when one of the valence electrons remains in the \( 4s \) orbit and the other moves successively to \( s, p, d, f, g, \) etc., orbits is shown in Fig. 5. As in the hydrogen-like spectra, all these terms tend to a limit, which is the normal state in which the ion is left, in this case a \( 4s \) \( ^2S_{1/2} \) state. It is very convenient to imagine all the terms of an atom as being built up...
by the addition of an electron vector to the terms of the ion. For example, the addition of a 5s electron to the 4s \( ^3S_{1/2} \) state of Ca II produces the \( ^1S_0 \) and \( ^3S_1 \) terms of Ca I, and, as shown graphically, a 4p electron produces \( ^1P_1 \) and \( ^3P_0 \), \( ^3P_1 \), \( ^3P_2 \) terms. The addition of \( d, f, g \) electrons can be made in a similar way.

**Series limits**

It has so far been assumed that only one of the valence electrons has moved into a higher orbit, but it is possible for both to be displaced when excitation of the atom takes place. For example, the 3d orbit is higher than the 4s and one electron may be lifted up into this orbit and remain there, whilst the other may be lifted to an \( s, p, d, f \), etc., orbit. Or else the first electron is lifted to 4p and the second to some other orbit. In each case the terms will tend to different limits, since the ion is left in a different state. Thus, in the first case previously discussed, where one electron was permanently in a 4s orbit, the calcium ion is left in a 4s \( ^3S_{1/2} \) state when the optical electron is completely removed. In the second case mentioned above, the electron of the ion is left in a 3d orbit when the optical electron is removed, so that the ion may be left either in a 3d \( ^2D_{5/2} \) or a 3d \( ^2D_{3/2} \) state. It follows that all the terms of the calcium atom, which arise when one electron remains in a 3d orbit and the other moves to other orbits, tend to two limits, \( ^2D_{5/2} \) and \( ^2D_{7/2} \). These two terms lie close together, being members of a doublet group, but are both far away from the \( ^3S_{1/2} \) state which was the single limit reached by all the terms when only one electron was excited. In an identical manner the terms produced when one electron is in the...
4p orbit and the other is in some other orbit tend to the limits 4p, 3p 1/2, and 4p, 3p 3/2, which again lie close together but are both far removed from the 3D terms of the previous case. The limits reached by the terms of Ca I are shown in Fig. 6 and are seen to be widely separated.

In order to be able to build up term schemes correctly, it is necessary to know exactly what the series limits will be and to which particular limit a given sequence of terms will converge. It will be seen later that this has a very important bearing upon the problem of hyperfine structures, and for this reason series limits have been discussed here. Complete tables of series limits have been investigated by Hund 4 and the reader is referred to Hund’s work for particular cases. Hund 5 has given two sets of limits, one set being those to which normal terms tend; the other for inverted terms. An example of the application of the rules will be given later when series limits are applied to hyperfine structures.

Equivalent electrons

The Pauli exclusion principle strictly limits the number of electrons with the same n and l value in any atom. In a given orbit there cannot be more than a fixed number of, say, p or d of f electrons. Such electrons are called equivalent electrons. As an example of an atom which has many equivalent valence electrons, the bromine atom may be cited. In the normal state it has seven valence electrons, two of which are in 4s orbits and five in 4p orbits. The configuration is called a 4s 2.4p 5 electron configuration, using the notation already described. The two 4s electrons are equivalent, and the five 4p electrons are also equivalent. The 4s group is called a closed group, and takes practically no part in producing the spectrum, so that for present purposes it may be neglected, and only the 4p 5 group need be considered. The spectral terms of the atom are produced by one of these electrons moving into various orbits. In the normal state the five 4p electrons are arranged so that the resultant S = 1/2 and the resultant L = 1. This means that 2P terms are formed and these are the deepest lying terms, the 2P 1/2 term being deeper as the multiplets are inverted. The remaining important terms of the spectrum are produced by the moving optical electron.

The first group of terms arises when one of the 4p electrons is lifted up to the 5s orbit, so that the configuration becomes 4s 2.4p 4.5s. Another group is formed if the optical electron goes instead into a 5p orbit producing a 4s 2.4p 4.5p configuration, and so on. Each separate electron configuration forms a number of terms and these can all be conveniently built up by treating the 4s 2.4p 4 electron group as a core to which the optical electron is added in the 5s, 5p, etc., orbits. It is seen that the process is not very different from that involved in building up the two-electron spectra, for in the latter two electrons were coupled together, and at present one electron is to be coupled to a core.

The series limits must be known before any attempt to build up the terms can be made. All the terms of the bromine atom can be predicted by means of Hund’s theory. And the most important of these are shown in Table II.6

<table>
<thead>
<tr>
<th>Electron Configuration</th>
<th>3P</th>
<th>1D</th>
<th>1S</th>
</tr>
</thead>
<tbody>
<tr>
<td>4s 2.4p 5</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>4s 2.4p 4.5p</td>
<td>2P</td>
<td>4P</td>
<td></td>
</tr>
<tr>
<td>4s 2.4p 4.5p</td>
<td>2(SP)</td>
<td>4(SP)</td>
<td></td>
</tr>
<tr>
<td>4s 2.4p 4.4d</td>
<td>2(PDF)</td>
<td>4(PDF)</td>
<td></td>
</tr>
</tbody>
</table>

In Table II the electron configurations are shown in the first column and each configuration produces the group of terms shown in the remaining columns. It is usual to distinguish between two classes of terms called odd and even terms. Odd terms are those in which the sum of the l values of the individual electrons of the term configuration
is odd and even terms those in which the sum is even. For example, in the previously considered case of Ca a $4p^2$ term is even and a $4p \cdot 4d$ term is odd. In Br I $4s^2 \cdot 4p^6$ terms

necessary. As Table II shows, transitions such as $P \rightarrow P$ or $P \rightarrow F$ are possible providing they are odd-even transitions. In addition to this L restriction, transitions can only take place when $\Delta J = \pm 1$ or $0$ ($0 \rightarrow 0$ excluded).

In Table II the basic ion terms of Br II are shown in small type. These are produced by the configuration of the ion Br II and are the basic terms which produce the terms of Br I when the optical electron is added. They are also the series limits, that is, the limits to which the Br I terms converge. In order to avoid making the table too cumbersome some abbreviations are used. First, a term written as $3P$ or $4P$ means the whole sequence of allowable J values, in these cases, $3P_0, 3P_1, 3P_2$ and $4P_{1/2}, 4P_{3/2}, 4P_{5/2}$ respectively. Second, a group of terms $2S, 2P$ and $2D$ is written $2(PD)$ and similarly for other multiplicities. Each horizontal row shows the terms which arise from a given configuration, and the limits to which these terms tend is indicated by the column in which the terms are found.

The question of series limits is of importance, and since the limits for normal and inverted spectra are different, care must be taken to use the correct limits. The terms in Br I are inverted; a pictorial diagram of the limits is given in Fig. 7. The terms of Br II are also inverted and this appears in the diagram. The diagram illustrates two cases only, namely the important $4s^2 \cdot 4p^4 \cdot 5s$ and $4s^2 \cdot 4p^4 \cdot 5p$ terms which tend to the $3P$ limit. The limiting $3P$ terms are shown by the thick horizontal lines and form an inverted multiplet (the drawings are schematic and not to scale) as shown by the J values on the left. The effect of adding the $5s$ electron is shown by the vertical lines which indicate to which limit each term sequence tends. For instance, $3P_{1/2}$ tends to the limit $3P_0$ whilst $1P_{1/2}$ tends to $3P_1$. Conversely the $2P_{1/2}$ term of Br I is formed when the $5s$ electron is added to the $3P_0$ term of Br II, and the $4P_{1/2}$ term of Br I is formed by adding the electron to the $3P_1$ term of Br II. The method by which this is obtained will be shown later by means of a vector diagram. The lower diagram shows the effect of adding a $5p$ electron to the ion terms, giving in fact the terms and limits of the $4s^2 \cdot 4p^4 \cdot 5p$
For comparison purposes Fig. 8 shows how the series limits for normal and inverted multiplets are different, although similar multiplets are considered.

![Diagram showing normal and inverted multiplets](image)

**Fig. 8.**—Showing the difference between the series limits in normal and inverted spectra.

*Vector coupling in bromine*

The vector coupling in bromine is somewhat similar to that of a two-electron spectrum and is shown for the $4s^2 \cdot 4p^4 \cdot 5s$ configuration in Fig. 9. In effect, the optical electron $5s$ is to be coupled vectorially to the core $4s^2 \cdot 4p^4$. The core produces the $3P_0, 3P_1, 3P_2$ inverted multiplet terms shown at the left, which are the series limits, so that the core is equivalent to a vector with $\mathbf{L} = 1$ and $\mathbf{S} = 1$,

![Diagram showing vector coupling in bromine](image)

**Fig. 9.**—Vector coupling diagram showing how the $3P$ and $4P$ multiplet terms of the $4s^2 \cdot 4p^4 \cdot 5s$ configuration of Br I arise.

*In order to avoid complication it is assumed that strict LS coupling takes place. This is not true for bromine, but as it has only been chosen as an example because of its many valence electrons, the gloss is of no consequence here.*
shown to the right of the terms. The long thin black arrow gives the resultant \( L = 1 \) and the short thick black arrow the resultant \( S = 1 \) (compare Fig. 4). As the terms are inverted, the largest \( J \) value (2) lies deepest. As before, the direction of \( J \) is shown by the dotted arrow.

The result of coupling the 5s electron to the core is shown at the right and this should be studied carefully. The electron is added in \( LS \) coupling which means that the original coupling of the \( L \) and the \( S \) can be broken, but the individual \( L \) couplings and \( S \) couplings remain intact. The new \( S \) value of the added electron may combine either parallel or anti-parallel with the already existing \( S \) value of the core, to form a new resultant \( S \). The same is true for \( L \) (in this case the \( l \) of the added electron happens to be zero). The new \( L \) and \( S \) values then combine in the usual manner to give \( J \). The 5s electron is shown by means of a short thick hollow arrow and is always added in the same direction (either positively or negatively) as the original \( S \) value of the core.

In the two higher lying \( \text{Br I} \) terms, the spins oppose, the resultant spin quantum number is \( \frac{1}{2} \), producing an inverted \( ^2P \) doublet. In the three deeper terms the resultant spin of the core acts parallel to that of the 5s electron, so that the terms are \( ^4P \) \((S = \frac{3}{2}\) and multiplicity = \(2S + 1\)) and are also inverted.

The vector method just described can be used for all configurations if the terms predicted by Hund’s theory are taken as a guide and the correct series limits used.

**Interval rule**

In most of the diagrams previously given, the number and description of the terms of a multiplet has been given, but little attention has been paid to the relative positions of the terms such, for example, as the relative spacings between the terms of \( ^3P \) group. In a normal triplet, the terms lie in the order \( ^3P_0, ^3P_1, ^3P_2 \), the term with the smallest \( J \) value lying deepest. The interval between the first pair of terms is different from that of the second pair, because the interaction energy between the spins and angular orbital momenta differs for the various terms, and the position of a term depends upon this interaction energy. It can be shown that, in \( LS \) coupling, the interaction energy between the resultant spin \( S \) and the resultant angular orbital momentum \( L \) is given by \( E \), where

\[
E = A' \cdot SL \cos (SL),
\]

\( A' \) being called the interval factor for a reason which will appear later. The interaction energy between two vectors is proportional to the cosine of the angle between those vectors, and it can be shown that the above expression is equivalent to

\[
E = \frac{A'}{2} \cdot \{J(J + 1) - L(L + 1) - S(S + 1)\}
\]

This formula expresses what is known as the Landé Interval Rule, after its discoverer, and from it the ratio of the intervals in a multiplet can be calculated.

Since the position of a term is given by its interaction energy, the difference of energy between two terms will be proportional to their separations and the formula shows that for a multiplet group arising from a given \( S \) and \( L \), the energy differences of the different terms are proportional to the differences of the \( J(J + 1) \) values. With a \( ^4D \) multiplet group, for example, \( S = \frac{3}{2} \) and \( L = 2 \), giving four terms with the \( J \) values \( \frac{3}{2}, \frac{5}{2}, \frac{7}{2}, \frac{9}{2} \). According to the Landé rule the successive energy differences between adjacent pairs of terms, being proportional to \( J(J + 1) \), are in the ratio

\[
\{\frac{3}{2} \cdot \frac{3}{2} - \frac{5}{2} \cdot \frac{7}{2}\} : \{\frac{5}{2} \cdot \frac{5}{2} - \frac{7}{2} \cdot \frac{9}{2}\} : \{\frac{7}{2} \cdot \frac{7}{2} - \frac{9}{2} \cdot \frac{9}{2}\},
\]

that is, in the ratio 7:5:3, or that the ratios of the intervals of the terms of the multiplet are the same. It will be seen that these ratios are given directly by the ratios of the different \( J \) values when the smallest \( J \) value is neglected. This simplified form of the Landé rule is true for all cases in which strict \( LS \) coupling takes place. It is both true for normal and for completely inverted multiplets, but it will be obvious that when a multiplet is partially inverted
the rule completely fails. Whilst the ratio of the intervals between the terms is a function of $J$, the actual separations depend upon the value of the interval factor $A'$. This is determined by the distance of the electrons from the nucleus, the degree of penetration of the orbits, the screening effect of other electrons, and the interaction energy between the electrons themselves. The interval factor, therefore, differs for different multiplets, usually diminishing as the $n$ of the optical electron increases. For this reason the separations in, say, the $4s \cdot 4p \cdot 3p$ multiplet of Ca I are greater than those in the $4s \cdot 5p \cdot 3p$ multiplet, although the ratios of the intervals are the same in both cases (in effect, the difference is merely one of scale).

An example of the line pattern formed by the transitions between two multiplet terms is shown in Fig. 10, the case illustrated being that of transitions between two $5D$ multiplet term groups. The interval factor of the upper set of terms is made less than that of the lower, but the interval rule is obeyed by both. The intensities of the lines differ according to the intensity rule on page 18. The resulting line pattern is complex because of the different interval factors and the varying intensities.

The line pattern of a multiplet

The line pattern of a multiplet can vary in width from a fraction of a wave number to several thousand wave numbers, and various methods have been used to assist in sorting out members of a multiplet from a large group of lines. The most important way is the method of constant frequency differences which can be seen from Fig. 10. An examination of the transitions shows that the two hatched lines $ab$ have the same frequency difference as the two dotted lines $cd$, namely the interval between the upper terms with $J = 3$ and $J = 2$ respectively. It follows that a search among the lines suspected to belong to the multiplet will reveal this common difference, even when other foreign lines are present. It will be observed that a large number of common differences exist, and from these the multiplet can generally be built up, due regard being paid to the intensities which should obey the intensity law. There are, of course, many multiplets in which the intensity or interval rules break down and other means must be used to assist in analysis, such as Zeeman effect observations.

For further study in multiplet analysis the reader is referred to Hund’s "Linienspektren".

REFERENCES

2. Johnson, "Spectra" (Methuen, 1929).
CHAPTER IV
THE EFFECT OF A NUCLEAR SPIN

Nuclear spin

Two types of hyperfine structure occur in line spectra, and these arise from entirely different causes. One type, called the isotope hyperfine structure, is due to the fact that many elements are not simple but consist of a number of isotopes. Since the separate isotopes of an element have identical extra-nuclear electron configurations, the spectra emitted by each of these are almost the same, and the lines of the one isotope practically coincide with the lines of the other. There is, however, in some cases, a measurable displacement for certain lines, the displacement being due either to the different masses or varying electrical properties of the different isotopic nuclei or else to the motion of the nucleus. In favourable cases these displacements can be observed as a hyperfine structure effect. This will be discussed in a later chapter.

The second type of structure arises because the nuclei of certain atoms have the property of a spin. It was first suggested by Pauli that hyperfine structures (excluding isotopic effect) could be accounted for if it were assumed that the nucleus possessed a magnetic moment which perturbed the motions of the external electrons. This nuclear magnetic moment would couple vectorially with the magnetic moment of the valence electrons and behave as if a new quantum number had been introduced. If the nuclear and the electron magnetic moments couple in accordance with the usual vector rules, each electron term would be divided into a group of terms, the energies varying between the two limits formed when the nuclear moment is parallel and anti-parallel to the electron moment. In between these limiting terms further energy states may occur, the number depending on the quantum values of the respective moments. The expression "level" will be retained to describe any of these hyperfine structure energy levels and the word "term" confined to denote the gross structure energy state.

Associated with the magnetic moment is a mechanical moment, which being quantised, becomes some multiple of $\hbar/2\pi$. The mechanical moment is usually denoted by $I\hbar \over 2\pi$ and the coefficient $I$ is termed the nuclear spin. The existence of a nuclear spin had already been previously inferred, for the nuclei of certain atoms, from the study of band spectra. The investigation of the band spectrum of hydrogen showed definitely that each of the protons, which are the nuclei of the hydrogen atoms forming the hydrogen molecule, possesses a spin with $I = \frac{1}{2}$, i.e. the mechanical moment of the spinning proton is $\hbar \over 2\pi$, the same as in the case of the spinning electron. The ratio of the magnetic moment of the proton to that of the electron should be approximately in the inverse ratio of the two masses so that the magnetic moment of the proton should be about $184\pi$ that of the electron. This explains the minuteness of the scale of hyperfine structures, in comparison with that of the ordinary gross structure multiplet. A study of the intensities of the band spectra of homogeneous molecules seems to point to the rule that the nuclear spin of light even atomic weight elements of mass a multiple of four is zero, while the odd elements have half integral spin values. It can be shown that if the identical nuclei of the two atoms forming a homogeneous dipolar molecule have a nuclear spin $I$, the alternate lines in the rotational structure of the band spectrum have an intensity ratio $I + 1 \over I$. When $I = 0$, the ratio becomes infinite, so that alternate lines are absent. This occurs in the spectra of $\text{He}_2$, $\text{C}_2$, $\text{O}_2$, and $\text{S}_2$.

* Actually it is 2.8 times greater than the value expected on the simple theory (p. 97).
Nuclear spins are usually found to be half integral and vary in value from $\frac{1}{2}$ to $\frac{3}{2}$. A detailed discussion will be given later of nuclear building-up processes with reference to the experimental values of the spins.

**Interaction of nuclear and electron magnetic moments**

As we have shown in the previous chapter the resultant moment of the electron system of an atom is given by $\mathbf{J}$, so that if the atom possesses a nuclear spin $\mathbf{I}$ it is to be expected that the two vectors $\mathbf{I}$ and $\mathbf{J}$ should combine to form a resultant quantum number. This is termed the *hyperfine structure quantum number*, and is denoted by the letter $\mathbf{F}$, which is a vector quantity, giving the total angular momentum of the whole atom. We find that $\mathbf{I}$ and $\mathbf{J}$ can combine vectorially to give a number of values so that each term of the spectrum (henceforth called the gross structure term) will divide into a tiny multiplet group of hyperfine structure terms, which for purposes of differentiation will be referred to as *levels*. The values that $\mathbf{F}$ can have are given by the vectorial relation $\mathbf{F} = \mathbf{I} + \mathbf{J}$, so that by analogy with the case of multiplets previously discussed it will be seen that if $\mathbf{J}$ is larger than $\mathbf{I}$ the total number of combinations possible is $2\mathbf{J} + 1$, hence the gross structure term divides into $2\mathbf{J} + 1$ levels. If, on the other hand, $\mathbf{J}$ is less than $\mathbf{I}$ the number of arrangements which is possible is $2\mathbf{I} + 1$. Since many terms split up in this manner, a gross structure line will have the line structure of a multiplet on a minute scale, and general considerations show that a selection principle will operate in the production of this hyperfine structure multiplet. The predicted and observed selection principle is $\Delta \mathbf{F} = \pm 1$ or 0 ($0 \rightarrow 0$ being excluded), and hitherto no exceptions to this rule have been found.

The interaction energy between the two vectors $\mathbf{I}$ and $\mathbf{J}$ is dependent upon a number of different factors. First, it depends upon the value of the magnetic moment of the nucleus as distinct from its mechanical moment, for it is primarily the magnetic fields which couple. It is, of course, true that the magnetic fields arise because the charged particles in motion have mechanical moments, but since the coupling between the mechanical moments occurs through magnetic linkage, it is the value of the magnetic moment that determines the interaction energy of $\mathbf{I}$ and $\mathbf{J}$. The nuclear magnetic moment is usually written as $\mu$, where $\mu = I \cdot g(I)$, so that $g(I)$ is the ratio of the magnetic to the mechanical moment when quantum units are employed. $g(I)$ is a pure number, and as its value is determined by the internal structure of the nucleus, it may be very small at times. A second factor affecting the interaction energy is the nuclear charge, for the higher the degree of ionisation the greater is the coupling energy between the extra-nuclear electron moment $\mathbf{J}$ and the nuclear spin $\mathbf{I}$. A third factor which is extremely important depends upon the eccentricity of the orbit of the optical electron. If the orbit is very eccentric, a high degree of penetration of the inner shells takes place and the optical electron approaches closely to the nucleus so that the coupling becomes very strong. It will be apparent that the coupling energy of an $s$ electron will generally be much greater than that of a $p$ electron, etc.

Let us consider a one-electron system; it is convenient to imagine the coupling between $\mathbf{I}$ and $\mathbf{J}$ as consisting of two parts. $\mathbf{I}$ couples with the orbital angular moment of the electron, that is, $\mathbf{IL}$ coupling takes place, and $\mathbf{I}$ couples with the electron spin so that $\mathbf{IS}$ coupling also occurs. When these two effects are added the interaction energy between nucleus and electron is obtained. The final result shows that the energy $W$ is given by the expression

$$W = A \cdot IJ \cdot \cos (IJ)$$

which can be reduced to the form

$$W = \frac{A}{2} \{F(F + 1) - J(J + 1) - I(I + 1)\}.$$ 

Comparing this with the formula previously given for ordinary multiplets it will be seen that the two formulae are identical if $\mathbf{J}$, $\mathbf{L}$ and $\mathbf{S}$ are replaced by $\mathbf{F}$, $\mathbf{J}$ and $\mathbf{I}$. Since this is identical in form to the Landé formula for gross structure multiplets, the energies of the hyperfine structure
levels are also such that they will obey the Landé interval rule. The interval factor of the hyperfine structure multiplet \( \Lambda \) can be approximately calculated for the simple case of one electron spectra. It is found that \( \Lambda \) is proportional to
\[
g(I) \cdot \frac{Z_i z^2}{n^6 L(L + \frac{1}{2})(L + 1)}
\]
where \( Z_i \) and \( z \) are respectively the effective electric charges of the nucleus in the inner and outer portions of a partially penetrating orbit, and \( n \) is the effective Rydberg quantum number. This formula includes the conditions previously referred to concerning the value of the nuclear magnetic moment, the nuclear charge and the degree of electron penetration. Calculation shows that the interval factor for the gross structure multiplet, \( \Lambda' \) of a one-electron system differs from the above value \( \Lambda \) by the factor \( g(I)/Z_i \). From this it follows that the ratios of the scales of hyperfine structures to gross structure multiplets should be given approximately by \( g(I)/Z_i \), but the formula is only very approximately correct, and the agreement with experimental results is only partial.

The formulae for the hyperfine structure interval factors have only been developed for a few electron systems, but although the general theory is very incomplete, a number of important facts have already been deduced and experimentally verified. As previously stated the width of a structure in a term, that is to say, the total spread of the hyperfine structure multiplets, depends upon the coupling between \( I \) and \( J \). This, in turn, depends upon the closeness of approach of the optical electron to the nucleus. The more elliptic the orbit, the more likely will the electron penetrate the inner shells. It will not only approach the nucleus more closely, but, when it penetrates closed electron shells, the screening effect of these shells will be reduced. This is equivalent to a considerable increase in the effective nuclear charge which further increases the tightness of the coupling, as it is called. Since the smaller the \( I \) value of an orbit, the more eccentric is its shape, it will be expected that electrons in \( s \) orbits will produce

much wider structure than electrons in \( p \) orbits. Correspondingly, the width of a \( 4s \) electron term should be greater than that of a \( 5s \) electron. Special attention may be directed to the case of \( p \) orbits. There are two distinct kinds, depending upon whether the electron spin assists or opposes the orbital angular momentum. In particular, in \( jj \) coupling the \( p \) electron can have \( j \) values \( \frac{3}{2} \) or \( \frac{1}{2} \) [for \( l = \frac{1}{2} \) and \( s = \frac{1}{2} \)]. These different orbits are designated as \( p_{1/2} \) and \( p_{3/2} \) respectively. The \( p_{1/2} \) electron is in the more elliptical and more penetrating orbit and should produce wider fine structures than the \( p_{3/2} \) electron. This has been verified experimentally.

**Hyperfine structure in one-electron spectra**

The vector model previously employed to illustrate gross structure multiplets can also be applied to study hyperfine structure multiplets.

The first new consideration is the question of the numerical value of the nuclear spin which introduces a complication. For example, two atoms may give spectra in which the gross structure multiplets are identical, except that their scales differ, but if the atoms have different nuclear spins, the hyperfine structures of the two spectra will be totally different. This can be made clear by specifying a certain value for \( I \) and it will then be shown how a variation in \( I \) will affect the levels of the terms.

As an example of the types of structures which are to be expected in one-electron spectra, a case like that of sodium will be considered. Referring to Fig. 2 we see that the deepest lying term in sodium involves the most penetrating electron and is the \( 3s \cdot ^2S_1/2 \) term. Next to it come the two terms \( 4p \cdot ^2P_{1/2} \) and \( 4p \cdot ^2P_{3/2} \) which involve \( p \) electrons. We will examine the structures of the two lines \( 35s \cdot ^2S_{1/2} - 4p \cdot ^2P_{1/2} \) and \( 35s \cdot ^2S_{1/2} - 4p \cdot ^2P_{3/2} \). These two lines are the well-known yellow D lines of sodium. For simplicity we shall assume the nuclear spin \( I \) of sodium to be equal to \( \frac{1}{2} \). This spin is to be coupled with each of the three terms concerned in producing the lines, in order to see what hyperfine structure levels will arise. The \( 3s \cdot ^2S_{1/2} \) term has the
J value $\frac{1}{2}$ and as the spin I is $\frac{1}{2}$, it is obvious that I and J can only combine in two ways, to give $F = 0$ or $F = 1$ (it will be remembered that the maximum number of levels is $2J + 1$ or $2I + 1$ according to whether J or I is the smallest). The F value 0 is produced when the nuclear spin and the electron spin oppose each other, and the F value 1 arises when the two spins act parallel.* It follows therefore that the $3s \cdot ^2S_{1/2}$ term is double.

In this particular case $I = \frac{1}{2}$ and as all the J values in sodium are $\geq \frac{1}{2}$ every term of the spectrum must also be double. The $4p \cdot ^2P_{1/2}$ term will have the two F values 0 and 1 as in the previous case, but in the $4p \cdot ^2P_{3/2}$ term

* The question as to true spin direction of an electron and of a spinning nucleus is somewhat modified by the wave mechanics, but for simplicity the vector model is adhered to here.

The F values will be 1 and 2, for J equals $\frac{3}{2}$ and I equals $\frac{1}{2}$, hence $F = 1$ or 2. Since both the lower and upper term of each of the two lines are double, each line will be a tiny multiplet. The structure of the common lower $3s \cdot ^2S_{1/2}$ term will be much wider than that of the upper terms, since an s electron is here involved. The structure of $4p \cdot ^2P_{1/2}$ will be somewhat greater than that of $4p \cdot ^2P_{3/2}$.

Remembering the selection principle, $\Delta F = \pm 1$ or 0 (0 $\rightarrow$ 0 excluded) it is seen from Fig. 11 that each of the D lines will have three components. If the structures in the p terms are very small, and in sodium this is the case, the separations of the upper terms cannot be easily experimentally resolved and both lines appear to have the same structure, namely two components, the intensity ratio and frequency separation being the same for both of the D lines. It has been shown by Fermi that in the s terms of one-electron spectra, the nuclear spin can be calculated from the intensity ratio of these apparent doublets, when the upper terms are unresolved. It can be shown that the quantum weight of a hyperfine structure level is $2F + 1$.

By this is meant that the sum of the intensities of all the lines coming to a level of given F value is proportional to $2F + 1$. If I is the nuclear spin, the F values of the two levels of the $^2S_{1/2}$ term are $I + \frac{1}{2}$ and $I - \frac{1}{2}$ respectively. The intensity ratio of the two components of the doublet is therefore

$$\frac{2(I + \frac{1}{2}) + 1}{2(I - \frac{1}{2}) + 1} = \frac{I + 1}{I - 1},$$

so that the accurate measurement of the intensity ratio will give the spin, I. Experimentally this is very difficult to determine, because the lines involved are resonance lines, which usually suffer self-absorption and this gives rise to false intensity measurements. In the case under consideration, the intensity ratio of the components in both the D lines would be $3 : 1$. It is easy to see by making I equal to zero, that the formula given above is self-consistent. The intensity ratio in the hyperfine structure doublet then becomes $1 : 0$, meaning that one line is missing. In other
words, the doublet becomes a single line when the atom has no nuclear spin, which is what one expects. Measurements have been carried out upon intensity ratios of doublets of alkali resonance lines, but the results are usually disturbed by reversal.

**Effect of varying nuclear spins**

The calculated structure of the D lines becomes very different for different values of \( J \), although the observed structure may not appreciably change. The latter effect is, however, only accidental and is due to the fact that the structure in the \( s \) electron term is so much wider than in that of the \( p \) electron. The pattern formed by a larger spin such as \( \frac{3}{2} \) is shown in Fig. 12. The \( ^3S_{\frac{3}{2}} \) term again divides into two levels, but as \( J = \frac{1}{2} \) and \( I = \frac{1}{2} \) the \( F \) values now become 2 and 3 respectively. The \( 4p \cdot ^2P_{\frac{1}{2}} \) term behaves in the same manner, except that the structure is on a smaller scale. As a result \( ^3S_{\frac{3}{2}} \rightarrow 4p \cdot ^2P_{\frac{1}{2}} \) has three components, two of which are grouped closely together.

With a nuclear spin of \( \frac{3}{2} \) the \( 4p \cdot ^2P_{\frac{1}{2}} \) term is no longer double but becomes more complex. Since \( J < I \), \( 2J + 1 \) components are formed, so that with \( J = \frac{3}{2} \) the term splits up into four components. We can see from Fig. 13 how these four levels can occur. The \( J \) value of the level is represented by the long thin arrow and the nuclear spin by the short thick arrow. In the highest level the two vectors are parallel, and in the lowest they are anti-parallel. In between these values, the vectors take up intermediate...

---

**Fig. 12.**—Hypothetical hyperfine structure of sodium D lines with the assumed nuclear spin of \( \frac{3}{2} \). Two doublets are formed, the separation being 0.06 cm.\(^{-1}\), and the intensity ratio 5:7.

**Fig. 13.**—Illustrating the vector coupling between a nuclear spin \( I = \frac{3}{2} \) with a \( J = \frac{3}{2} \) to form four hyperfine structure levels.
positions that satisfy the cosine law. Expressed simply, this means that the angles between the vectors are such that the resultant vector F can have all integral values between 4 and 1. As shown in the diagram we obtain four levels with F values 1, 2, 3, 4, the smallest F value lying deepest. These levels will be so spaced as to obey the Landé interval rule, although this is not shown in the schematic vector diagram, but it is so drawn in Fig. 12.

It is seen in this figure that the line 3s \( {^2S_{1/2}} \rightarrow {^4P_{1/2}} \) has six components in its structure pattern. The components are grouped into two pairs of threes, owing to the smallness of the upper interval factor. If the interval factors of the upper terms are both very small compared with that of the lower term, the patterns cannot be completely resolved, and the structures of both the D lines will be doublets with the intensity ratio 5:7 and of approximately the same separation as was obtained when I was taken to be \( \frac{3}{2} \). The above values, \( \frac{5}{2}, \frac{7}{2} \), are taken for illustration only. The nuclear spin of sodium is now known to be \( \frac{3}{2} \). This apparent similarity shows that the prime importance of experimentally resolving all the components of a line before a reliable estimate of the nuclear spin of the atom can be made, and when I \( \geq J \) of checking the results by the examination of other terms.

When the nuclear spin of an atom is small, the term structures are simple, and the line patterns not very complicated, but when the spin is large, the terms with higher J values become quite complex and the line patterns correspondingly more so. Fortunately the interval factor in one of the terms of a line is often very much smaller than that in the other term and this introduces a considerable simplification into the line pattern which is formed.

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1. Pauli, Naturwiss., 12, 741, 1924.
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structure. For all the other terms I and J compound in the usual manner to form multiplets. Suppose, for example, the nuclear spin is $\frac{3}{2}$, then the $^3P_0$ term will be single, the $^3P_1$ term will have three components, and the $^3P_2$ term will have four, as shown in Fig. 14.

$$\begin{array}{c}
^3P_2 \\
\frac{5}{2} \\
\frac{3}{2} \\
\frac{1}{2} \\
^3P_1 \\
\frac{5}{2} \\
\frac{3}{2} \\
\frac{1}{2} \\
^3P_0 \\
\frac{5}{2} \\
\end{array}$$

$I < J$

$2I + 1 = 4$

$J < I$

$2J + 1 = 3$

$J < I$

$2J + 1 = 1$

Fig. 14.—The hyperfine structures of $^3P$ terms when the nuclear spin is $\frac{3}{2}$. The number is $2I + 1$ or $2J + 1$, according to which is the smaller.

Since all the terms which go to the $4s\cdot 2S_{1/2}$ limit have a $4s$ electron in their configurations, the structures of all will be wide (excluding terms with $J = 0$). It is seen from this that a line transition such as $4s\cdot 4p \cdot ^3P_2 \to 4s\cdot 4d \cdot ^3D_2$

will have structures of the same order in both the upper and lower terms. The result is that the line pattern is complex, as, for example, in the case shown in Fig. 10. Most of the terms which go to the next limit, $3d\cdot 2D_{5/2}$, of the calcium atom, do not possess a penetrating $4s$ electron in their configurations, so that practically all the terms going to this limit have narrow structures, and the same is true of the terms which go to the next limit. Line transitions between one of the terms which go to the $4s\cdot 2S_{1/2}$ limit and one of the terms which go to another limit have their patterns much simplified, as was the case with the sodium lines in Fig. 12. The lines which arise from transitions between two terms, neither of which possess unpaired s electrons in their configurations, often appear single, because the structure is so minute. Generally speaking, the appearance of the structures is very dependent on the specific electrons in the configurations.

**General vector method for many-electron spectra**

A very convenient method for predicting the structures which arise in atoms with complex electron configurations has been developed by White.\(^1\) This is an extension of the vector method which has already been used to show how gross structure multiplets arise. It was shown in Chapter III that the best way to see how the gross structure multiplets arise was by adding the vector representing the optical electron to the ion terms. The hyperfine structures are obtained by the simple extension of adding the nuclear spin vector to the terms already obtained in this manner. In a large number of cases only the configurations which possess a penetrating $s$ electron show any structure in the terms, so that this type of term will be considered first.

If a configuration possesses a penetrating $s$ electron, it can be shown that the coupling between the nuclear spin and the electron configuration is proportional to $\cos (I\mathbf{s})$ * and not, as in general, to $\cos (IJ)$. The direction of the

* $s$ is the spin vector of the $s$ electron. It is assumed throughout that LS coupling is used.
coupling between the nuclear spin and the $s$ electron becomes of great importance, so that it is necessary to know exactly in which direction the $s$ electron is pointing with reference to the other electrons of the configuration. For this reason the series limits must be known and the multiplet structure built up from the ion, in order that the coupling of the $s$ electron can be followed carefully from term to term.

For the purpose of explaining the vector method, recourse will be made to Fig. 9, which shows how the $^3P$ and $^4P$ multiplets of the $4s^2 \cdot 4p^4 \cdot 5s$ configuration of Br I are formed. The problem is to see how the addition of the nuclear spin affects the terms. The terms are split up according to the $J$ values, but the interval factors, on which the coarseness of the structures depend, are determined by the value of $\cos (I_s)$. White assumes that the complex electron group $4s^{2} \cdot 4p^{4} \cdot 5s$ behaves as a unit and couples with $I$ in the usual vector manner, the degree of coupling depending on the orientation of $I$ and $s$. The relative values of the interval factors, in a group of levels belonging to a gross structure term, are proportional to $\cos (I_s)$ when $\cos (IJ) = 1$. That is to say, when the nuclear spin is set parallel to $J$ the angle which $I$ then makes with the spin $s$ of the $5s$ electron determines the coarseness of the structure. From the properties of vectors we have

$$\cos (I_s) = \cos (I J) \times \cos (J_s),$$

and as $\cos (I J) = 1$ when $I$ and $J$ are parallel, therefore $\cos (I_s) = \cos (J_s)$.

As the coarseness of the structures is proportional to $\cos (I_s)$ it must therefore also be proportional to $\cos (J_s)$, hence the problem is to find the relative orientation of the spin of the $s$ electron, with respect to the resultant $J$ value of the term, and the angle between these vectors determines the coarseness of the structures.

As in the two electron spectra the number of hyperfine structure levels in any gross structure term is determined by the values of $I$ and $J$. There will be $2I + 1$ or $2J + 1$ levels according to which is the smaller value. From the

\begin{figure}  
\centering  
\includegraphics[width=\textwidth]{fig15.png}  
\caption{Vector coupling diagram showing how the hyperfine structures in the $4s^2 \cdot 4p^4 \cdot 5s$ configuration are produced by a nuclear spin of $\frac{3}{2}$.}  
\end{figure}
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Landé formula it is easy to see by subtraction of two levels that the actual separation between two levels with quantum numbers \( F + 1 \) and \( F \) is \( A(F + 1) \), \( A \) being the interval factor for the term in question. When \( I \) and \( s \) are parallel \( \cos(1s) = 1 \) so that \( A \) is positive. When \( I \) and \( s \) are anti-parallel \( \cos(1s) = -1 \) so that \( A \) is negative. A negative value for \( A \) means that the hyperfine structures are inverted, the highest \( F \) value lying deepest. The vector method will now be applied to bromine in order to trace the various possible couplings.

Fig. 15 shows the vector combinations. The notation is exactly the same as in Fig. 9 and in it A and B of Fig. 15 are identical with those of Fig. 9. A shows the \( 3P \) terms of \( Br \) II based upon the \( 4s^2 \cdot 4p^4 \) configuration and B shows the \( 5s \cdot 3P \) and \( 5s \cdot 4P \) terms for \( Br \) I which are formed by the addition of the \( 5s \) electron. The vector process up to this point has already been discussed in detail. The experimental value for the nuclear spin of bromine is \( \frac{3}{2} \) and C shows the hyperfine structure levels (dotted) that should arise when a spin of \( \frac{3}{2} \) is added to the doublet and quartet multiplets. In order to simplify the complicated vector diagrams, only the coupling between the nuclear spin and the \( 5s \) electron is shown at the extreme right. As before, the thin dotted arrow in B shows the direction of \( J \). I is set parallel to this, so as to conform to the rule that \( \cos(1) = 1 \) and in this position the interval factor is proportional to the cosine of the angle between \( I \) and \( s \). The spin of the \( s \) electron is shown, as before, by the hollow arrow, and the nuclear spin by the long thick dotted arrow. The small curved arrow at the head of this implies that it is to be compounded vectorially.

Considering each term in detail it is seen that the \( 5s \cdot 3P_{1/2} \) term has \( I \) parallel to \( s \), and as \( J = \frac{1}{2} \) the hyperfine structure is therefore a wide normal doublet; wide because the two spin vectors assist, and normal since they act parallel. In the next term \( 5s \cdot 3P_{3/2} \), the \( I \) and \( s \) oppose, so that the term is inverted. As \( J = \frac{3}{2} \) a quartet term is formed, and the structures are still comparatively large. The \( 5s \cdot 4P_{1/2} \) term is again a wide normal doublet. The next term is \( 5s \cdot 4P_{3/2} \) and the nuclear spin acts at an angle with the \( s \) electron spin so that a medium width quartet is produced. Finally the \( 5s \cdot 4P_{5/2} \) term has all the vectors parallel, so that a wide quartet is formed.

The methods used to describe the coupling in the case of a penetrating \( s \) electron can be applied to the case of a non-penetrating \( p \) electron. For instance, it is possible to couple the \( 5p \) electron to the \( 4s^2 \cdot 4p^4 \) core of \( Br \) II, but it is found that the agreement with experiment is not very good. The supporting evidence for the \( s \) electron case has been very well established in manganese, but different investigations indicate that such a simple treatment is not possible for \( p \) electrons. It is quite certain that the \( p \) electrons couple vectorially to give the correct number of levels that theory demands.

* There is reasonably good experimental verification for bromine, the deviations being due to the incomplete \( LS \) coupling. In manganese, which shows full \( LS \) coupling, the vector method is completely verified. Bromine has only been selected as an example of a many-electron atom, for illustration purposes.

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CHAPTER VI

ANALYSIS OF THE HYPERFINE STRUCTURES

Intensities of hyperfine structure components

Although it is true that severe perturbations may occur in hyperfine structure levels most of the spectra hitherto investigated seem to show the interval rule to be reasonably valid. The perturbations have been found in the structures of mercury, bismuth and aluminium, etc., and will be discussed later, but for the present purpose all that need be said is that in certain cases the structure in a term can be perturbed in such a manner that the interval rule fails. The intensities of the hyperfine structure components of a line are just as important as the intervals, and are of equal help in the analysis. Perturbations have been found in intensities, but so far these have not been very serious. The intensity disturbances are of two kinds. In one case they are due to the source being possibly too hot, or the pressure too high, with resulting self reversal effects. The other type of intensity perturbation is not understood, but will be discussed in the chapter dealing with perturbation.

The intensity rules which hyperfine structure components obey are important. These have been given by Hill and since they are used so frequently in analysis, it has been considered worth while to reproduce them here. They form the basis of the graphical method of analysing structures, due to Fisher and Goudsmit, which is useful in practical work. In the formulae given below the following abbreviations are used: In. (F . F + 1) means the intensity of a line transition from an upper hyperfine structure F level to a lower term hyperfine structure level F + 1. The meanings of In. (F . F) and In. (F . F - 1) are then self-evident. In all cases the first F value refers to the upper term, and the second to the lower term. In the

formulæ themselves the following contractions are used for the sake of brevity.

\[
\begin{align*}
\text{P.} & \quad (a + b)(a + b + 1) - c(1 + 1), \\
\text{Q.} & \quad (c + 1) - (a - b)(a - b + 1), \\
\text{R.} & \quad a(a + 1) + b(b + 1) - c(1 + 1).
\end{align*}
\]

All the transitions are those for which \( \Delta L = \pm 1 \), and there are three sets of formulæ according to whether \( \Delta J \) is equal to +1, 0, or −1. These are set out in table form below. A, B, and C are constants determining the relative intensities of the gross structure lines but do not here concern us, since all the components for any one given line (i.e. single J transition) involve the same constant so that the intensity ratios in a hyperfine structure pattern are not dependent on the value of the constant.

On substituting in these formulæ one finds that the transition \( \Delta F = 0 \rightarrow 0 \) always has zero intensity, and also that the total intensity of all the components to a given F level is proportional to \( 2F + 1 \). The formulæ appear rather cumbersome, but are actually simple to apply, and check themselves automatically, so that errors in calculation are avoided. Numerical intensity tables derived from these formulæ are now available.

Irregular and regular line patterns

As a result of the intensity and interval rules, two types of line pattern are produced, and these may conveniently be called irregular and regular patterns. An irregular line pattern arises when both the upper and lower interval factors are of the same order. The line pattern becomes particularly complex when the hyperfine structure multiplicities are high. The irregular structures can only be analysed by seeking for frequency differences, as in the case of gross structure multiplets. When the hyperfine structure multiplet has been built up the F values of the terms are easy to find if \( J > I \), but if \( J < I \) the F values can usually only be obtained from the interval rule. If both I and J are already known, the analysis of a line is much simpler. However, when the F values have been found I
can always be calculated. The best way to analyse a spectrum with many lines is to select the lines which involve terms wherein the J value is greater than the suspected I value. In all of these terms the multiplicity is 2I + 1, so that I is determined when the line pattern has been analysed, and the interval rule can be applied as a check. It is often quite impossible to analyse the pattern of a line completely in this manner, because many structure components crowd together and cannot be resolved. The method has, however, been successfully used in very many cases, e.g. Cd, Tl, Pb, Hg etc. An example of an irregular pattern is shown in Fig. 16. The line pattern illustrated is that of 5s. $^4P_{1/2} - 5p. \; ^4D_{3/2}$ in Br I. The nuclear spin is $\frac{3}{2}$ and for the purpose of illustration it is assumed that the interval factor of the upper term is $\frac{5}{2}$ that of the lower. The resulting pattern is complex, and no obvious apparent regularities exist in it.

In the irregular pattern described above, the irregularity arises from the fact that the two interval factors have nearly equal values. It often happens that only one term involves a penetrating s electron, and so produces a structure which is very much coarser than the structure in the other term. When this is the case the line pattern becomes regular and has the appearance of a regular degraded series. For

**Irregular**

\[
\text{Interval factor ratio } 5:4
\]

**Regular**

\[
\text{Interval factor ratio } 20:1
\]

Fig. 16.—Dependence of the appearance of a fine structure line pattern on the fine structure interval ratio of the upper and lower terms.
example, the same line as above has the pattern shown in Fig. 16 when the interval factor of the upper term is only \( \frac{3}{2} \) that of the lower. Owing to the natural widths of the components, and partly because of imperfections in the instruments used, the closely packed components cannot often be separated, so that a regular quartet pattern will be seen in the present example. It is clear that a regular pattern will usually show the approximate structure of one of the terms only (exceptions occur). In the present instance the structure is that of the lower term, which involves the s electron.* This term is therefore quartet in structure, and as the J value is \( \frac{5}{2} \), then I must be equal to \( \frac{3}{2} \), because a greater nuclear spin would have produced more components. This line shows how a nuclear spin can often be very easily obtained from one regular pattern. The scale of the structures may be gathered from the fact that the total width of the structure is about 0.4 cm\(^{-1}\).

**Graphical method of analysis**

In the regular “pseudo quartet” just described there are actually ten components, and the appearance of the line pattern becomes more and more complicated as the relative value of the smaller interval factor increases, with respect to the larger. In order to deal with complex patterns which are incompletely resolved, a very valuable graphical method of analysis has been proposed by Fisher and Goudsmit.\(^2\) For this graphical method I and the J of both terms must be known for the line in question. It is often possible to get I from the other data, e.g. from the structure of favourable lines or from band spectra, etc., so that the method is generally applicable. The same line, namely

\[ 5s^2 \cdot {}^4P_{3/2} \rightarrow 5p \cdot {}^4D_{5/2} \]

of Br I, will be taken for an example of the use of this graphical method.

The graph is shown in Fig. 17. Firstly the interval

\[ \Delta J = \frac{1}{2} \rightarrow \frac{1}{2} \]

\[ I = \frac{3}{2} \]

Fig. 17.—Fisher-Goudsmit line complex for the transition \( J = \frac{5}{2} \rightarrow \frac{3}{2} \). The nuclear spin is \( \frac{3}{2} \) in this example.

**transitions as calculated from Hill’s formula.** For the purpose of calculation, thin lines can be used, the actual intensity ratios being written in. The points which represent the upper limit of the whole graph, along AB, show the line pattern given by the two terms in question, when the interval factors of the two terms are identical; in other words, when the interval factor ratio is 1. Similarly, the points at the lower limit of the graph show the line pattern when the ratio of upper to lower interval factors is 1. The horizontal line CD cuts the graph at the centre. Each line is cut at a point, and each point represents a component of the line pattern. This particular centre pattern is that
formed when the upper interval factor is equal to zero [since CD lies midway between the lines AB and EF which represent ratios of $+ \frac{1}{3}$ and $- \frac{1}{3}$ respectively]. For any given ratio of interval factors, the line pattern is determined by drawing a horizontal line through the graph at a point between AE corresponding to this ratio. Thus, for example, the line patterns of Fig. 16 were obtained by drawing horizontals at points corresponding to interval ratios of $\frac{1}{3}$ and $\frac{1}{3}$ respectively, the upper interval being the smaller in both cases. When a line pattern has been measured all that need be done is to fit it unambiguously into the graph, and the values of both the upper and lower interval factors are at once determined, although the resolution in the line pattern may be incomplete.

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4 See papers by Schüler and co-workers, Zeit. f. Phys., 1928.

CHAPTER VII

TERM INTERVAL FACTORS

Coupling of two $s$ electrons

From the calculated interval factors information can be obtained concerning the coupling processes of the individual electrons, and perturbation effects can be detected. Extensive theoretical calculations of the ratios of the interval factors of the different terms of a spectrum have been made by Goudsmit. From the experimental side, although data are lacking, enough has been found to show that the agreement between theory and observation is only partial. Discrepancies arise from the fact that the actual coupling process between a nucleus and an electron is insufficiently understood. Fermi, Casimir, Breit, and Goudsmit have all made different attempts to solve this mathematical problem, which presents grave difficulties.

There is general agreement for the case of the atom which possesses one valence electron, such as the alkali metals or ionised alkali earths. Theory shows that the interaction energy between the nucleus and the optical electron is of the form $E = A \cdot \mathbf{1J} \cdot \cos (\mathbf{1J})$. This leads to the Landé interval rule for hyperfine structures, and experimental observations support this strongly. Casimir showed that in the particular case of an $s$ electron, the interaction energy is given by $E = A \cdot \mathbf{ls} \cdot \cos (\mathbf{ls})$ and this fact has already been utilised. Following Casimir, Goudsmit has calculated the interaction energy for an atom which has two valence $s$ electrons. For example, all the $^{3}s_1$ terms of Ca I are terms of an atom which has two $s$ electrons in its configuration. The interaction between the nucleus and the spins of the two electrons, $s_1$ and $s_2$, is

$$E = a_1 \mathbf{ls}_1 \cdot \cos (\mathbf{ls}_1) + a_2 \mathbf{ls}_2 \cdot \cos (\mathbf{ls}_2).$$

The energy is treated as the sum of two interaction energies,
one for each electron. Each electron obeys the law deduced by Casimir, and in the equation, \( a_1 \) and \( a_2 \) are the numerical proportionality factors for each electron.* The cosines are to be averaged over a complete electron revolution and when this is done the formula reduces to

\[
E = IJ \cdot \cos (IJ) \cdot \left[ \frac{a_1 s_1}{J} \cdot \cos (s_1 J) + \frac{a_2 s_2}{J} \cdot \cos (s_2 J) \right],
\]

\( J \) having the usual significance, being here the resultant of \( s_1 \) and \( s_2 \). Comparing this with the previous general equation, namely,

\[
E = A \cdot IJ \cdot \cos (IJ),
\]

it is evident that the two \( s \) electrons obey the same general law as one electron, if the interval factor \( A \) is replaced by the expression in square brackets. When the quantum values of the cosines are substituted in the expression for \( A \) we get

\[
A = a_1 \cdot \frac{J(J + 1) + s_1(s_1 + 1) - s_2(s_2 + 1)}{2J(J + 1)} + a_2 \frac{J(J + 1) + s_2(s_2 + 1) - s_1(s_1 + 1)}{2J(J + 1)},
\]

since \( s_1 = s_2 = \frac{1}{2} \), this reduces to

\[
A = \frac{1}{2} a_1 + \frac{1}{2} a_2.
\]

Thus, the term interval factor \( A \) is half the sum of the proportionality factors of the two electrons. Experimental verification of this has been possible in a few cases. Thallium I has been examined and the interval factors for the structures of the terms of the \( 6s \), \( 7s \), \( 6s \cdot 8s \), and \( 6s \cdot 9s \), etc., configurations determined. From these it is possible to calculate the value of \( a_1 \) for the \( 6s \) electron and of \( a_2 \), \( a_3 \), \( a_4 \), etc., for the \( 7s \), \( 8s \), \( 9s \), etc., electrons. The value of \( a_1 \) increases slightly as the second electron approaches the ionising limit because the screening effect of the second electron diminishes. \( a_1 \) tends to a limiting value which can be calculated. In the Tl II atom, which possesses one electron in the \( 6s \) orbit, the value of \( a_1 \) should be identical with the limiting value of \( a_1 \) calculated from the Tl I spectrum. Allowance has to be made for the fact that the interval factor is proportional to \( z^2 \), where \( z \) is the degree of ionisation. More strictly the interval factor is proportional to \( Z_i z^2 \), but in a heavy atom \( Z_i \) is sensibly the same for the normal and the ionised atoms so that the coarseness of the structure is very nearly proportional to \( z^2 \).* The values of \( a_1 \) calculated from Tl I and Tl II agree within experimental error, so that the theory is supported.

### Interval factors for the general case

The general theory for the hyperfine structure interval factors in spectra, which show \( LS \) or \( jj \) coupling, has been worked out by Goudsmit,* and since it is not practicable to give any details of his calculations here, only a few of his results will be quoted. The calculations lead to different results in \( LS \) coupling and in \( jj \) coupling, the results being more satisfactory for the latter type. In \( LS \) coupling it is only possible to give the sum of the interval factors for all the terms of a multiplet group which have the same \( J \) value. For example, the \( ^1P \) and \( ^3P \) terms which arise from an \( s \cdot p \) electron configuration are \( ^1P_1, ^3P_0, ^3P_1, ^3P_2 \). The interval factor for \( ^3P_0 \) is zero, that of \( ^3P_2 \) can be calculated, but only the sum of the interval factors for the two terms \( ^1P_1 \) and \( ^3P_1 \) is obtained from the formulæ, for they have identical \( J \) values.

The following table gives some of the more important results. \( a \) is the proportionality constant for the \( p \) electrons and \( b \) the constant for the \( s \) electrons in any configuration. The value of \( a \) or \( b \) is the same for all equivalent electrons.

In the table only a selected group of results has been shown and for full details the original paper must be consulted.

The calculated values for the interval factors of spectra which show \( jj \) type of coupling are more satisfactory than

* This is obviously an important practical point, for the higher the degree of ionisation the coarser are the structures and therefore the easier to observe.
those for LS type, because the interval factor for every term is given separately. The s electron has the same proportionality constant b as in the LS coupling, but the p electrons are different. As jj coupling implies that each

**TABLE III**

*Interval factors in LS coupling*

<table>
<thead>
<tr>
<th>Configuration</th>
<th>J</th>
<th>Interval Factor = A.</th>
</tr>
</thead>
<tbody>
<tr>
<td>p and p²</td>
<td>½</td>
<td>A = ½a</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>A = 2²a</td>
</tr>
<tr>
<td>p² and p⁴</td>
<td>2</td>
<td>Σa = 1²a</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>A = 2²a</td>
</tr>
<tr>
<td>p³</td>
<td>2</td>
<td>A = ⅛a</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>Σa = 1₂³a</td>
</tr>
<tr>
<td></td>
<td>½</td>
<td>Σa = ½²a</td>
</tr>
<tr>
<td>p . s and p⁵ . s</td>
<td>2</td>
<td>A = ²a + ½b</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>Σa = 2a + ½b</td>
</tr>
<tr>
<td>p⁵ . s and p⁴ . s</td>
<td>2</td>
<td>Σa = 1²³a + ⅛b</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>Σa = ½²a - ⅛b</td>
</tr>
<tr>
<td>p⁵ . s</td>
<td>3</td>
<td>A = ⅛a + ⅛b</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>Σa = 2a + ⅛b</td>
</tr>
<tr>
<td></td>
<td>½</td>
<td>Σa = 2½a - ⅛b</td>
</tr>
</tbody>
</table>

The formulae for each. The constants will be called a“ for the p₁/₂ electron and a' for the p₃/₂ electron. According to theory these constants are not independent, but are related to the constant a of the p electron in LS coupling.

Theory gives

\[ a' = \frac{3}{4}a, \]

\[ a'' = \frac{3}{2}a, \]

therefore

\[ a' = \frac{1}{8}a''. \]

As an example of the interval factors calculated for jj coupling, see Table IV. The first column in this table gives the electron configurations, and the second column the J values of the individual electrons. The J value of the s electron is omitted, since it is always equal to ½. In the third column the final resulting possible J values are shown, and opposite each J value (in the fourth column) is the calculated interval factor of the term. The terms with J = 0 are omitted as the interval factor for such a term is zero.

**TABLE IV**

*Interval factors in jj coupling*

<table>
<thead>
<tr>
<th>Configuration</th>
<th>J</th>
<th>J</th>
<th>Interval factor.</th>
</tr>
</thead>
<tbody>
<tr>
<td>p² and p⁴</td>
<td>½</td>
<td>2</td>
<td>a' + ½a</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>1</td>
<td>a'' - ⅛a</td>
</tr>
<tr>
<td>p² . s and p⁴ . s</td>
<td>½</td>
<td>2</td>
<td>a' + ½b</td>
</tr>
<tr>
<td></td>
<td>⅛</td>
<td>1</td>
<td>a'' - ⅛b</td>
</tr>
<tr>
<td>p⁵ . s and p⁴ . s</td>
<td>½</td>
<td>2</td>
<td>a' + ½a + ⅛b</td>
</tr>
<tr>
<td></td>
<td>⅛</td>
<td>1</td>
<td>a'' - ⅛b + ⅛b</td>
</tr>
<tr>
<td>p⁵ . s</td>
<td>3</td>
<td>2</td>
<td>a' + ⅛b</td>
</tr>
<tr>
<td></td>
<td>⅛</td>
<td>1</td>
<td>a'' - ⅛b</td>
</tr>
</tbody>
</table>
Interval factors in multiply ionised spectra

So far there is not a great deal of experimental evidence which has a bearing on the verification of the Goudsmit formula, since only few spectra have been thoroughly examined. Of interest are the results which have been obtained in the spectra of bismuth.\(^4\) In bismuth the Bi I, Bi II, Bi III and Bi V spectra have all been studied. It has been possible to calculate the proportionality constants of the separate electrons when different neighbouring electrons are present. Thus, in Bi I the 6p\(^1\) and the 6p\(^2\).7s configurations give the values of a\(^"\), a\(^'\) and b\(^'\) for the 6p\(_{1/2}\), 6p\(_{3/2}\) and 7s electrons respectively.* In Bi II the configurations examined were 6s.7s, 6p.6d and 6p.7p, so that the constants for the same electrons could be calculated, as well as those of 7p\(_{1/2}\) and 6d. The Bi III spectrum gave data for the 6s.6p\(^5\) configuration, so that in addition the constant b for 6s was determined. Finally, the 6s configuration of Bi V also gave a value for the constant b. A convenient summary of the results is shown in Tables V and VI:

**TABLE V**

<table>
<thead>
<tr>
<th>Electron</th>
<th>Constant</th>
<th>Bi I</th>
<th>Bi II</th>
<th>Bi III</th>
<th>Bi V</th>
</tr>
</thead>
<tbody>
<tr>
<td>6s</td>
<td>b</td>
<td>1.6</td>
<td>1.6</td>
<td>1.8</td>
<td>2.6</td>
</tr>
<tr>
<td>7s</td>
<td>b'</td>
<td>0.116</td>
<td>0.352</td>
<td>0.472</td>
<td>—</td>
</tr>
</tbody>
</table>

For both the 6s and 7s electrons the value of the constant increases regularly with the increase in the degree of ionisation. This is due to the removal of screening outer electrons. Table VI gives a summary of the observations made on the p electrons:

It was possible to calculate a\(^'\) and a\(^"\) from more than one term, and as seen in Table VI there is good approximate agreement between these. The general conclusions to be made from the tables are as follows:

1. The constant for a given electron is bigger the higher the degree of ionisation, owing to the removal of screening effect as other electrons are removed from the configurations.
2. The constant for an electron falls rapidly in value when the total quantum number, \(n\), of the electron increases.
3. The constant for the p\(_{1/2}\) electron is very much greater than that of the corresponding p\(_{3/2}\) electron, the ratios of these constants being altogether much greater than the calculated value, 5:1.

The fact that a\(^"\) is very much greater than 5a\(^'\), which the simple theory at first predicts, is interesting. There

* The exact value of the p\(_{3/2}\) electron constant was difficult to determine, but the data showed unambiguously that it was relatively quite small.
are two possible explanations, the second of which is much more likely than the first. It is possible that, owing to insufficient knowledge of the interaction processes, the equations are wrong, and the calculated values of \(a'\) and \(a''\) are not genuine. However, against this is the fact that the methods used to deduce the formulae have given correct results when applied to gross structure multiplet data, and also to other hyperfine structure data. This explanation is therefore not very probable. A much more likely explanation has been suggested by Breit,\(^6\) who has shown by means of a wave mechanics treatment, that in heavy atoms a \(p_{1/2}\) electron behaves with respect to the nucleus as if it were a deeply penetrating electron. The effect is only marked in the heavier atoms, as it arises from second order corrections which are negligible when the atomic weight is small. If the \(p_{1/2}\) electron behaves as if it were penetrating, it will therefore have associated with it a very large interval factor, comparable, in fact, with that of the \(s\) electron. As seen from the data in bismuth this is the case, so that Breit's theory offers a reasonable explanation of the relatively great value of \(a''\). It is, of course, of practical importance, as the penetrating property of the \(p_{1/2}\) electron results in the production of structures far wider than at first expected on the simpler theory.

REFERENCES

1 Pauling and Goudsmit, "Structure of Line Spectra", p. 205.

CHAPTER VIII

ISOTOPE EFFECT

Isotope effect in bands

Up to the present any hyperfine structure effect arising from the existence of isotopes has been neglected in the discussion. The isotope effects have become very important from both the experimental and theoretical aspects, so they will be treated here in some detail. The first attempt to detect any isotope effect in line spectra was made by Aronberg,\(^1\) who found that the wave-length of the lead line \(\lambda_{4508}\) was slightly less in ordinary lead than in radium lead. He was not able to observe the true structure of the line, partly because the line was broadened in the source he used, and partly because the resolving power he employed was not very high. It was only later and more refined work that enabled the true isotopic effect in lead to be discovered.

The isotope effects in line spectra are always small (the effect in deuterium apart), but the isotope displacements which occur in some band spectra are very large indeed, and are easy to recognise. This arises because the positions of the vibration heads of an electronic band system depend upon the moment of inertia of the nuclei of the molecule, and are very sensitive to changes in the value of this moment of inertia. For example, boron has the two isotopes \(B_{10}\) and \(B_{11}\) so that the band spectrum of the oxide consists of two sets of bands, one due to \(B_{10}\) and the other to \(B_{11}\). The different boron atoms have different moments of inertia because of their different weights. The band constants differ and as a result the band heads of the one system are displaced from those of the other by large amounts. Displacement of the order of 20 cm.\(^{-1}\) are not unusual in band spectra, from which it is seen that experimental observation of isotope effect in band spectra is much
The frequency of a line is the difference of energy between two terms divided by Planck's constant, so that \( \nu = R(A) \), where \( \nu \) is the frequency, \( R \) the Rydberg constant, and \( A \) is of the form

\[
A = \frac{z^2}{\hbar} \left( \frac{1}{n_1 + a + \frac{b}{n_1^2}} \right)^{\frac{1}{2}} \left( \frac{1}{n_2 + a + \frac{b}{n_2^2}} \right).
\]

For the isotopic constituents of a line \( A \) is constant, since it does not involve the nuclear mass. The only variable factor is \( R \). It follows then, that

\[
\nu_{20} = R_{20}(A), \quad \nu_{22} = R_{22}(A),
\]

hence

\[
\nu_{22} - \nu_{20} = (R_{22} - R_{20})(A).
\]

But \( \nu_{22} - \nu_{20} \) is the difference between the two isotopic components of the line, i.e. the isotopic separation \( \Delta \nu \), hence

\[
\Delta \nu = R_{22} \cdot (A) \left[ \frac{1}{R_{22}} - \frac{1}{R_{20}} \right] = \nu_{22} \left[ \frac{1}{R_{22}} - \frac{1}{R_{20}} \right].
\]

But as the Rydberg constant is proportional to \( \frac{mM}{m + M} \), the reduced mass of the electron, it follows that

\[
R_{20} = \frac{m \cdot M_{20}}{m + M_{20}} \times \frac{m + M_{22}}{m \cdot M_{22}},
\]

\[
R_{22} = \frac{m}{M_{22}} = 1 + m \left( \frac{1}{M_{22}} - \frac{1}{M_{20}} \right)
\]

i.e.

\[
R_{22} = I + \frac{m}{M_{20}} = I + m \left( \frac{1}{M_{22}} - \frac{1}{M_{20}} \right)
\]

to a first approximation. If \( M \) is the mass of the hydrogen atom the masses of the neon isotopes are \( 20M \) and \( 22M \) respectively, so that

\[
R_{21} = I - m \left( \frac{1}{M_{220}} \right).
\]

Therefore \( \Delta \nu = \frac{m}{M_{220}} \nu \) (\( \nu \) being an approximate mean
value of \( \nu_{20} \) and \( \nu_{22} \). It is sufficiently accurate for the purpose to take the mean frequency of the line without consideration of hyperfine structure.

The ratio \( \frac{m}{M} \) is the ratio of the mass of an electron to that of a proton and is equal to \( \frac{1}{1836} \), so that

\[
\Delta \nu = \nu \cdot \frac{1}{1836} \cdot \frac{1}{220} = \frac{\nu}{400,000}
\]

in approximate round numbers. For a line whose wavelength is 5000 A the displacement would therefore be about 0.05 cm\(^{-1}\). The isotopic displacement can be calculated for any other line in a similar manner.

The mass isotope displacement effect has been thoroughly investigated in neon and it has been shown that \( \frac{\Delta \nu}{\nu} \) is nearly constant for lines of the same series. The value of the displacement is consistently a little greater than that predicted by the simple theory, but the motion of the nucleus accounts for this.

There is no room for doubt that, for each line, the two components are produced by the different isotopes, as Hertz\(^2\) has clearly shown. In ordinary neon the isotopes are present in the ratio 9 : 1, the lighter atom being the more abundant. The hyperfine structure components should have the intensity ratio 9 : 1, which is observed in practice. Hertz succeeded in separating the isotopes of neon by means of a continuous diffusion process so that in the two end samples obtained, the ratio of \( \text{Ne}_{20} : \text{Ne}_{22} \) was 1 : 1 and 100 : 1 respectively. The spectrum of the first sample showed each line as a doublet, both components having the same intensity. On the other hand, the second sample showed each line to be quite single, as the other component was too faint. This affords a direct proof of the isotopic origin of the hyperfine structures in neon.

The simple mass isotopic effect is usually very difficult to observe, as it decreases rapidly with increase in atomic weight. For instance, in the bromine line 6122 A the dis-

placement should only be 0.004 cm\(^{-1}\). The mass isotope effect has been detected in a number of atoms including \( \text{H}^3, \text{Li}^6, \text{Ne}^8, \text{Cl}^3, \text{K}^4, \text{Br}^7 \) and possibly \( \text{Br}^8 \) and more recently boron. The most important of these is the exceptional case of hydrogen, for the spectroscopic detection of the isotope displacement effect was the original method whereby the existence of deuterium was demonstrated.

The mass effect in hydrogen is abnormally large in relation to that shown by other isotope mixtures because of the favourable mass ratio. It will be seen from the previous treatment that \( \Delta \nu = \frac{m}{2M} \), and for \( \text{H}_2 \) this amounts to as much as some 4 cm\(^{-1}\), an amount easily resolved even by prism spectrographs.

**Mixed odd and even isotopes**

In the heavy and moderately heavy atoms the simple mass isotope displacement is usually far too minute to be observed. A new effect arises when the atom consists of a mixture of isotopes of even and odd atomic weight. This effect was first observed by Schüller and Brück\(^8\) in the cadmium spectrum. Cadmium consists of six isotopes, two of which have odd and the other four even atomic weights.

The two odd atomic weight isotopes have each a nuclear spin but no mass displacement can be detected in any of the lines, as the effect is far too small. At the optical centre of gravity of the structure of each line lies a single very strong component. Intensity measurements reveal the fact that the intensity ratio of this single component, to the summed intensities of all the other components in the line, is the same as the abundance ratio of the even to the odd isotopes. It was found possible to account for all the components excluding the central strong one by assuming \( I = \frac{1}{4} \). There seemed to be no doubt about the interpretation of the weak lines, since the number of components and their intensities fitted a spin of \( \frac{1}{2} \). It was, therefore, suggested by Schüller and Brück that the even isotopes possess no nuclear spin, so that each contributes a single line, the lines all falling together at the optical centre of gravity of the spin hyperfine structure pattern.
HYPERFINE STRUCTURE

and appearing as one. The even isotopes form a single line because the mass displacements are too small to be measured.

odd       even

\[
\begin{array}{c}
\text{F} \\
\frac{1}{2} \\
\frac{3}{2} \\
\frac{5}{2} \\
\text{7p} \cdot 3s \\
\frac{1}{2} \cdot \text{P}_1 \\
\text{7p} \cdot 1s \\
\frac{3}{2} \\
\end{array}
\]

FIG. 18.—Undisplaced isotope line due to the even isotopes in the cadmium hyperfine structures. The thin lines show the doublet multiplet due to the odd isotopes, and the thick line represents the unaffected lines from the even isotopes all falling at the centre of gravity of the hyperfine structure pattern.

Obviously the intensity ratio of the central compound component to that of the total sum of the others must be the abundance ratio of the even to the odd isotopes.

The way in which the cadmium hyperfine structures arise is shown in Fig. 18. The line illustrated is \(7p \cdot 1s - 7s \cdot 3s\). Attention may be drawn to the fact that the structures in cadmium are inverted, which means that the biggest F values are deepest. This inversion is not to be confused with the term inversion which was previously discussed when \(\cos (1s) = -1\). The whole of the hyperfine structures are inverted, as if the nuclear magnetic moment arose from a negatively charged nucleus. Referring again to the diagram, the doublet multiplet arising from the split terms of the odd atoms is shown. The quantum weight of a hyperfine structure level is \(2F + 1\), so that the dotted lines indicate the centres of gravity of the split terms. The dotted lines are the positions the terms occupy in the even atoms, for the effect of a nuclear spin is to split up a term about its centre of gravity. The thick vertical line represents the transitions in the even atoms, the final resulting line pattern being shown below. It is to be observed that the isotope effect just described only arises from the fact that different isotopes have different nuclear spins. There is no displacement of the lines of isotopes which have the same spin.

Isotope displacement effect

The so-called isotope displacement effect is of quite a different nature and was first found in the thallium structures by Schüler and Keystone. Thallium consists of two isotopes, 205 and 203, the lighter of which is the more abundant in the ratio 2:3:1. Both isotopes have the nuclear spin of \(\frac{1}{2}\). Fig. 19, showing the analysis of the line \(6s \cdot 7s \cdot 3s_1 - 6s \cdot 7p \cdot 1s_1\), reveals the displacement effect. The line transitions due to the more abundant isotope, 203, are shown to the left at A, and those due to the less abundant, 205, are shown dotted at B. The observed line pattern is drawn, as usual, below. The analysis was carried
out by the use of frequency differences and the careful observation of intensities. Each transition in A is 2.3 times as strong as the corresponding transition in B, so that the analysis is quite certain. Furthermore, there is no doubt about the value of the nuclear spin which has been obtained from other lines which show no isotope effect. The transitions B should fall into the same system as A, but have been displaced to the right in order to clarify the diagram.

FIG. 19.—Isotope displacement effect in thallium.

A study of the line shows just what the displacement effect is. The structure of the $^1P_1$ term is the same in both isotopes, but there is a small measurable difference in the widths of the $^3S_1$ structures. The most marked effect is the bodily displacement of the whole of the $^1P_1$ term in one of the isotopes, that is to say, the centre of gravity of the term has been displaced without much appreciable alteration of the size of the structure. The result of this is a bodily shifting of the whole line pattern of the one isotope with respect to the other. The displacement is far greater than that calculated for the simple mass effect and must be ascribed to some other cause.

Displacement in odd and even isotopes

Displacement effects in atomic species containing a mixed group of odd and even isotopes have been observed in mercury, zinc, platinum and others. Most of the salient features appear in the platinum arc spectrum studied by Tolansky and Lee. The platinum isotopes revealed by the mass spectrograph are 192, 194, 195, 196, 198, of which 192 is very weak and will be disregarded. The analysis of the structure of a typical line ($\lambda 5390.8$) is shown in Fig. 19 (a) 1. The analysis in Fig. 19 (a) 2 is derived from a comprehensive term scheme including the numerous line transitions involving thirty-three separate terms. The dotted levels and transitions refer to the odd isotope 195 which has a nuclear spin $\frac{1}{2}$. The unbroken levels and transitions refer to the even isotopes 194, 196, 198. The two patterns are separated for clarity, and the scale for the even isotopes increased for clarity in Fig. 19 (a) 2.

It will be seen that these do not fall together, as in cadmium, but suffer a small displacement (about 0.044 cm.\(^{-1}\)), the amounts for 194-6 and 196-8 being the same, within limits of error of observation (some 2 units).

The intensity ratios of the even isotope lines agree with the mass spectrograph abundance ratios. A point of considerable interest, as yet not explained, is that the “centre of gravity” of the odd isotope pattern (195) does not fall exactly half-way between the 194 and 196 even isotope lines. The extent of this effect is shown in Fig. 19 (b) in which the “centre of gravity” of isotope 195 is shown dotted. This effect appears in other platinum lines, such as $\lambda 5360.0$ also shown in Fig. 19 (b).

The case of mercury is also interesting. There are six isotopes of measurable intensity, namely, 198, 199, 200, 201, 202, 204. Both the odd isotopes have nuclear spins, but the point of interest is that the spins differ, that of 199
being $\frac{1}{2}$ and that of 201 being $\frac{3}{2}$. (Other atoms also have different spins for different isotopes.) The hyperfine struc-

ture terms in 201 are all inverted relative to those of 199.

Some of the terms of mercury illustrate isotope displacement of a more complex nature than that encountered in thallium. Fig. 20 shows an analysis of 4358 A in Hg I. The four even isotopes do not coincide, but are relatively displaced (see hatched lines). The observed intensities of
these even isotope lines are in excellent agreement with the abundance ratios. The two odd isotopes each show the pattern expected from their respective spins, but in addition, the centre of gravity of each pattern is displaced. In the figure the straight lines belong to 199 and the dotted lines to 201. The hatched lines are 198, 200, 202, 204, respectively.

\[ \text{Hg}_{199} \text{ I-}\frac{1}{2} \qquad \text{Hg}_{201} \text{ I-}\frac{1}{2} \]

\[ \begin{array}{c}
\frac{1}{2} \\
\frac{1}{2} \\
\frac{1}{2} \\
\frac{1}{2}
\end{array} \quad \begin{array}{c}
\frac{1}{2} \\
\frac{1}{2} \\
\frac{1}{2} \\
\frac{1}{2}
\end{array} \]

\[ 7^2S_1 \quad 6^3P_1 \]

Fig. 20.—Hyperfine structure of the mercury line 4358 A, showing the isotope displacements. The hatched lines arise from the even isotopes, the dotted lines from 201, and the straight lines from 199.

The isotopic displacement effect implies that the binding energies of the electrons are different in different isotopes. Fig. 21 shows the displacements which have so far been found in the Hg I spectrum.\(^{10}\) The \(^1S_0\) series has big displacements in the deepest term, and these rapidly diminish up the term sequence so that by \(^8^1S_0\) the displacements can no longer be measured. The \(^7^3S_1\) term also shows a small displacement which diminishes so rapidly up the term sequence that it can no longer be measured in higher terms. Two other terms show a displacement effect, namely the complex \(6\ ^1P_1\) term arising from the \(5d^9 \cdot 6s^2 \cdot mp\) configuration, and the \(8\ ^1P_1\) term of the ordinary configuration. The complex term is the term shown at the extreme left of the diagram. It follows, therefore, that regular displacements, falling off up a term sequence, only occur in the \(^1S_0\) and \(^3S_1\) terms. The \(8\ ^1P_1\) term shows displacement because it is perturbed. This will be discussed in a later chapter.

Fig. 21.—Isotope displacements in the terms of the Hg I spectrum. The displacements in all the terms are in the same order as those numbered in \(6\ ^1S_0\).

The relative displacement of the individual isotopes in one term is also of interest. Fig. 22 shows the displacements arising in the complex \(6\ ^1P_1\) term. The displacements in all the terms are similar, except that the scales are different. The deepest term is that of 198, and the even isotopes are approximately the same distance apart, in the order of increasing atomic weights. These are shown
above the line in Fig. 22. Below the line the centres of gravity of the structures of the odd isotopes are drawn, and it is seen that although the distance between the centres of gravity is practically identical with the distance between a pair of even isotopes, both the even isotopes are displaced towards the next lightest atom.

**Dependence of isotope displacement effect on electron configuration**

Isotope displacements have now been measured in a large number of spectra including Hg I,14 Hg II, TI I, TI II, Pb I, Pb II,12 Cu II,12 Zn II,15 Wo I,16 Pt I,17,18

The most comprehensive data are those given by Tolansky and Lee for the Pt I spectrum.

Table VII shows the configurations in some different spectra in which even isotope displacement has been detected and measured.

Before discussing this table it should be noted that perturbed terms in mercury and in lead show anomalously large isotope displacement and as the lead spectra are suspected to be severely perturbed, they are best left out in discussion.

---

**Table VII**

<table>
<thead>
<tr>
<th>Pb I</th>
<th>6s², 6p², mp</th>
<th>6s², 6p, md</th>
<th>6s², 6p, mp</th>
</tr>
</thead>
<tbody>
<tr>
<td>Hg I</td>
<td>6s, ms</td>
<td>6s, 6s</td>
<td>6s, 6d</td>
</tr>
<tr>
<td>Hg II</td>
<td>5d⁹, 6a²</td>
<td>5d⁸, 6a²</td>
<td>5d⁹, 6a²</td>
</tr>
<tr>
<td>TI I</td>
<td>6s², mp</td>
<td>6s², 6s</td>
<td>6s², 6s</td>
</tr>
<tr>
<td>TI II</td>
<td>5d⁹, 6a², 6p</td>
<td>5d⁹, 6a²</td>
<td>5d⁹, 6a²</td>
</tr>
<tr>
<td>Cu II</td>
<td>3d³, 4s²</td>
<td>3d³, 4s²</td>
<td>3d³, 4s²</td>
</tr>
<tr>
<td>Wo I</td>
<td>5d⁴, 6a²</td>
<td>5d⁴, 6a²</td>
<td>5d⁴, 6a²</td>
</tr>
<tr>
<td>Zn II</td>
<td>3d⁹, 4s²</td>
<td>3d⁹, 4s²</td>
<td>3d⁹, 4s²</td>
</tr>
<tr>
<td>Pt I</td>
<td>5d⁹, 6s</td>
<td>5d⁹, 6s, 7s</td>
<td>5d⁹, 6s, 7s</td>
</tr>
<tr>
<td></td>
<td>5d⁹, 6p</td>
<td>5d⁹, 6p, 7s</td>
<td>5d⁹, 6p</td>
</tr>
</tbody>
</table>

Two additional features emerge, (1) displacements diminish usually on going up a sequence, (2) a configuration with an s² group shows a very large effect.

---

The details for the platinum spectrum given by Tolansky and Lee deserve further discussion and a small selected group of the results is shown above. The displacements are in units of cm⁻¹ x 10⁻³.
HYPERFINE STRUCTURE

It is assumed (although this is not absolutely certain) that the completely closed shell $5d^{10}$ has virtually no displacement effect (if it has, the displacement in this term is to be added to that of the others).

The largest displacements (203) are found in the $^9F$ multiplet involving the $5d^8 \cdot 6s^2$ group. The displacements in the groups involving only one $6s$ electron are about half of these. It seems clear that the effects of the two $6s$ electrons in $6s^2$ are additive and do not compensate in the manner usually characteristic of gross structure multiplicity.

As in other spectra, the higher sequence members show much smaller effects, e.g., 117 displacement for $5d^9 \cdot 6s \cdot ^3D_2$ and 12 for $5d^9 \cdot 7s \cdot ^3D_2$.

Within a multiplet the displacements do not appear to be at all dependent on $J$. This is shown by a comparison of the nuclear spin hyperfine structure interval factors ($A$) and the corresponding isotope displacements ($D$) shown below for the $^9F$ group.

<table>
<thead>
<tr>
<th>Term</th>
<th>$A$</th>
<th>$D$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$5d^8 \cdot 6s^2$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^3F_2$</td>
<td>80</td>
<td>200</td>
</tr>
<tr>
<td>$^3F_3$</td>
<td>52</td>
<td>203</td>
</tr>
<tr>
<td>$^3F_4$</td>
<td>31</td>
<td>203</td>
</tr>
</tbody>
</table>

The manner in which $A$ varies with $J$ whilst $D$ remains constant is clearly demonstrated.

The analysis of the isotope displacement effects is considerably complicated by the effect of local term perturbation and much work remains to be done before a detailed theory can be developed.

It seems quite established that the Bohr mass effect cannot be invoked to explain the large displacements found in the heavier atoms. The mass effect for a mass $M$ increasing to $M + 2$ is proportional to $1/M^2$ which excludes simple mass modification of Rydberg constant as a possible explanation.\(^{13}\)

ISOTOPE EFFECT

Displacement due to different isotope magnetic moments

A relatively infrequent isotope effect, which produces displacement in the line pattern but not displacement in the terms, is that found in antimony by Tolansky.\(^{19}\) A typical case is illustrated in Fig. 22 (a).

![Fig. 22 (a)—Isotope effect in antimony.](image)

There are two antimony isotopes, 121, 123, with abundance ratio almost 5:4. The spin of one isotope is $\frac{5}{2}$, that of the second being $\frac{3}{2}$ (or $\frac{1}{2}$?). The isotopes have quite different nuclear magnetic moments, that of 121 being 1.37 times that of 123. Thus the pattern of every line consists of a hyperfine structure multiplet for isotope 121 and superimposed within this is a pattern from isotope 123, having about $\frac{3}{8}$ the intensity and shrunk down to $\frac{3}{8}$ of the scale of the former. As a result, the patterns for lines of even moderate $J$ value are exceedingly complex. For example, $a\Delta J = 3 \rightarrow 2$ line has twenty-eight components and these form so complex a pattern that in a specific case but ten could be resolved from the general overlap. In Fig. 22 (a), the continuous lines represent isotope 121 and the dotted lines isotope 123.

But simple lines, such as $\Delta J = 0 \rightarrow 1$, reveal the isotope effect in a clear manner. A typical case is shown in Fig. 22 (a). This is the line 6806, involving wide structure arising from a $J = 1$ term.
The pattern for each isotope is a regular triplet, that for 123 being shown dotted.

The difference in magnetic moment is responsible for the separation of the two patterns, whose centres of gravity almost coincide so that there is no term displacement effect in the sense already discussed formerly.

It has been established by calculations that all isotope displacement effects cannot be ascribed to nuclear motion. Such an explanation can account only for the displacements in the lighter atoms. Displacements are in effect due to deviations from Coulomb's law. A physical interpretation is possible by describing this in terms of an effective nuclear radius. The difference between two typical isotopes of mass M and M + 2 is that two extra neutrons are packed into the nucleus and clearly the nuclear volume might be expected to increase. The increased volume leads to an alteration in nuclear-electron interaction and this produces the isotope displacement. Schüler and Schmidt suggest that the nucleus with the larger volume has less electrostatic energy than that with smaller volume, since both have identical charge, and this leads to the displacement. 20

REFERENCES

3 Urey, Crickwedde and Murphy, *Phys. Rev.*, 40, 1, 1932.
14 See note 6.

CHAPTER IX

PERTURBATIONS IN HYPERFINE STRUCTURES

In the present chapter an account will be given of the perturbation effects which have now been found in many structures. Well marked perturbations have been observed in Hg I, 1 Al II, 2 Bi IV, 3 etc. The suggestion was first made by Schüler and Jones that those of Hg I are analogous to the gross structure series perturbations which were first explained by Shenstone and Russell. 4 From a study of two electron spectra, the latter investigators conclude that in a spectrum with two terms, produced by different electron configurations, but with identical L S and J values, these terms perturb each other, even although they may be widely separated. The perturbation shows itself as an apparent repulsion between the terms. Paschen 2 and Goudsmit and Bacher 5 have gone further in the explanation of the effect. It is found that two types of perturbation exist. One is probably similar to the Shenstone-Russell effect, but the other perturbation apparently takes place when the gross structure multiplet separation is of a similar order to that of the hyperfine structure separation. Each will be discussed in turn.

Total widths of the structures

Let us first consider the total widths of the structures which arise from two isotopes of an atom. In most cases when an atom has two odd isotopes, both have the same nuclear spin and the same nuclear magnetic moment μ. In Hg the isotopes 199 and 201 have different spins, namely 7/2 and 7/2, the spin being negative in the latter. It can be shown, 6 that the total splitting of a term, namely the distance
between the extreme $F$ levels for any given gross structure term is $\Delta W$, where

$$\Delta W = \mu \cdot H_0 \cdot \frac{2J + 1}{J} \text{ when } J \gg I, \text{ and}$$

$$\Delta W = \mu \cdot H_0 \cdot \frac{2I + 1}{I} \text{ when } I \gg J.$$

In these formulæ $H_0$ is the magnetic field of the optical electron at the nucleus, and only depends on the term, whilst $\mu$ is a nuclear function. The ratio of the magnetic moments of 199 and 201, namely $\mu_{199}$ and $\mu_{201}$, can be obtained by comparing the total widths of the structures of the two isotopes in a given term. In a particular case, when $J = 1$, for 199,

$$\Delta W_{199} = \frac{3\mu_{199} \cdot H_0}{\Delta W_{201}} = \frac{3}{8} \mu_{201} \cdot H_0.$$

Therefore,

$$\mu_{199} = \frac{8}{9} \Delta W_{199} \Delta W_{201}.$$

From five terms with $J = 1$ it was found that the total widths of the structures are identical, but inverted relative to each other, i.e.

$$\frac{\Delta W_{199}}{\Delta W_{201}} = -1,$$

which gives

$$\frac{\mu_{199}}{\mu_{201}} = \frac{8}{9} = -0.89.$$

For terms with $J = 2$ it also follows that

$$\frac{\Delta W_{199}}{\Delta W_{201}} = \frac{\mu_{199}}{\mu_{201}} = -0.89.$$

These numerical relations hold well in all the terms examined, except in the 6 $^1D_2$ term, where $\frac{\Delta W_{199}}{\Delta W_{201}} = -0.74$ instead of $-1$, and in the 6 $^3D_2$ term, where $\frac{\Delta W_{199}}{\Delta W_{201}} = -1.08$ instead of $-0.89$. It is of interest to note that the next member of the series of the latter term, namely 7 $^3D_2$, has a normal ratio for the spread of the structures in the two isotopes.

The centres of gravity of the structures of lines involving these terms are displaced and this must be connected in some way with the above effect. In lines which do not show any appreciably large isotope displacement, the centres of gravity of the structures due to the two odd isotopes practically coincide with the null line formed by all the even isotopes. However, in lines which involve the 6 $^1D_2$ term, the centres of gravity of the structures of the two isotopes practically coincide, but both are displaced from the null line by an amount roughly equal to 0.04 cm.$^{-1}$, the displacement being to the red. It will be remembered that $\frac{\Delta W_{199}}{\Delta W_{201}}$, calculated from this term, is too large. On the other hand, in lines involving 6 $^3D_1$ the centres of gravity of the structures due to the two odd isotopes again coincide, but this time are displaced from the null line to the violet by about 0.06 cm.$^{-1}$.

This term gives too small a value for $\frac{\Delta W_{199}}{\Delta W_{201}}$. Both these anomalies arise from a perturbation which will now be discussed.

**Perturbation of $F$ levels**

The perturbation in this case arises from the accidental fact that the two terms 6 $^1D_2$ and 6 $^3D_1$ are very close to one another. The structures in these terms are approximately 0.8 cm.$^{-1}$ wide and the distance between the centres of gravity of these structures is only 3 cm.$^{-1}$. The whole of the perturbations can be completely explained if it is assumed that *hyperfine structure levels with the same $F$ value repel one another*. The way this perturbation occurs in the 199 isotope is shown in Fig. 23. A very convenient fixed point exists from which the perturbations can be measured, this is the position of the even isotopes. The $F$ value of the even isotope terms in 6 $^1D_2$ is 2 (since $I = 0$) and the $F$ value of the even isotope terms in 6 $^3D_1$ is 1, and as these
F values are different, there is no repulsion effect between the even isotope components. Further there is no measurable isotope displacement in both terms, so that the even

<table>
<thead>
<tr>
<th>normal</th>
<th>perturbed</th>
</tr>
</thead>
<tbody>
<tr>
<td>F</td>
<td>F</td>
</tr>
<tr>
<td>( \frac{1}{2} )</td>
<td>( \frac{1}{2} )</td>
</tr>
<tr>
<td>CG</td>
<td>( \frac{5}{2} )</td>
</tr>
<tr>
<td>108</td>
<td>385</td>
</tr>
</tbody>
</table>

\( 6^3D_1 \) \{ \\
\( 6^1D_2 \) \} \\

\( Hg_{199} \) \( I = \frac{1}{2} \)

**Fig. 23.**—Perturbation effect in the 199 isotope of mercury.

isotopes in both terms form single fixed points. The fixed points are shown in the diagram by the thin dotted lines which are marked C.G., since they are the centres of gravity of the unperturbed terms. Considering now the 199 isotope, the two hyperfine structure levels which have the same F value, in this case \( \frac{3}{2} \), repel each other from the zero points by the same amount. The thick lines at the left of the figure show the normal levels, and the dotted lines at the right the perturbed levels. The drawing is to scale, the units* being millimicrometres^{-1}. The thick vertical arrow shows the two levels which repel each other, and it will be observed that, as the upper levels are inverted, the upper structure becomes narrower, and the lower structure wider. Since the levels of \( 6^3D_1 \) shrink together because of the perturbations, the centre of gravity of the hyperfine structure pattern due to \( Hg_{199} \) is moved up 72 mcm.^{-1}. Correspondingly, as the lower term is widened, the centre of gravity of the \( Hg_{199} \) pattern is moved downwards, in this case 43 units, so that a net displacement towards the violet, is obtained.

The perturbations in the structures of isotope 201 are a little more complex because the nuclear spin is greater and thus the hyperfine structure multiplicities are also higher. The perturbations are shown in Fig. 24, the notation being the same as in the previous figure. For clarity the displacements of the centres of gravity have been omitted. Only the levels with F equal to \( \frac{1}{2}, \frac{3}{2}, \frac{5}{2} \) find their counterparts in both the upper and lower structures, so that the repulsion only takes place between these levels. As in the previous case, the levels with the same F values are both repelled from the unperturbed position by the same amount. For F = \( \frac{1}{2} \) the upper level moves up, and the lower down, by a distance equal to \( a = 42 \) mcm.^{-1}. The F = \( \frac{3}{2} \) levels displace by \( b = 56 \) mcm.^{-1} and the F = \( \frac{5}{2} \) levels by \( c = 78 \) mcm.^{-1}. The value of the displacements is such that the interval rule still holds in the upper term structure; owing to the fact that the lower F = \( \frac{3}{2} \) level is not perturbed the

* A convenient unit to adopt is the millimicrometre^{-1}, which can be written as mcm.^{-1}. 
interval rule necessarily breaks down for the lower term as a whole.

\[
\begin{array}{c|c}
\text{normal} & \text{perturbed} \\
\hline
F & F \\
\frac{1}{2} & \frac{3}{2} \\
\frac{3}{2} & \frac{5}{2} \\
\hline
6^3D_1 & \\
\hline
\frac{1}{2} & \frac{1}{2} \\
\frac{3}{2} & \frac{3}{2} \\
\frac{5}{2} & \frac{5}{2} \\
\hline
6^1D_2 & \\
\hline
\frac{1}{2} & \frac{1}{2} \\
\frac{3}{2} & \frac{3}{2} \\
\frac{5}{2} & \frac{5}{2} \\
\hline
\end{array}
\]

\( \text{Hg}_{201} I = \frac{3}{2} \)

**Fig. 24.**—Perturbation effect in the 201 isotope of mercury, showing how the interval rule breaks down in the perturbed 6^1D_2 term.

It is at once apparent from the values of a, b, c that the perturbation is some function of F, and as the perturbation is different in Hg_{199} the effect must also be dependent on the value of I or \( g(I) \). Fig. 24 shows that the \( F = \frac{1}{2} \) levels are much closer together than the \( F = \frac{3}{2} \) levels, yet the perturbation in the latter pair is greater than in the former. Furthermore, the value of the displacement in the \( F = \frac{3}{2} \) levels of Hg_{199} is about twice that of the displacement in the \( F = \frac{5}{2} \) levels of Hg_{201}, although the latter are nearer together.

The perturbation effect just described explains all the anomalies in the term structures. The values of the total splittings of the terms are so affected that the ratios of \( \Delta W_{199} \) are no longer normal, but the observed values agree with the calculated values when the perturbation is taken into account. The displacements of the centres of gravity of the structures of the odd isotopes are completely accounted for, and the break-down of the interval rule in the 6^1D_2 term of Hg_{201} isotope is thus explained. Perturbation only seems to take place between levels with the same I and F values.

**Isotopic displacement perturbation**

Reference to Fig. 21 reveals the interesting fact that the 8^1P_1 term in Hg I has an anomalously large isotope displacement effect in all the isotopes. This is doubtless associated with the fact that this term is perturbed by the 5^2D, 6^3D, 6^3P_1 term shown at the left of Fig. 21. The whole of the 1P_1 series of Hg I is perturbed from the 8^1P_1 term upwards, because of the presence of the 6^1P_1 term from the complex configuration. As both terms have the same J value, the even isotopes have the same F value and so perturb each other. The perturbation is proportional to the mass of the atom so that the even isotope lines are regularly displaced. It is possible that some observed isotope displacements are perhaps related to perturbations.
Theory of perturbations

Extensive perturbations in Al II have been observed, and explained by Paschen. According to Paschen, when the coupling between I and J (which determines the total widths of the hyperfine structure pattern of a term) is of the same order as the coupling between L and S (which determines the gross structure multiplet width) then a perturbation sets in. Effectively it is as if the nuclear magnetic field were sufficiently strong to affect the coupling between L and S, and vice versa, the electron magnetic field perturbs the IJ coupling. (More correctly, it is as if the Paschen-Back effect sets in on both the structures.)

The theory is beautifully verified by Paschen's observations in Al II. By examining successive members of a series, he showed that the perturbations begin to occur when the expected hyperfine structure pattern becomes wider than the multiplet separation. In the term series $4\,^1F_3 - nG$, the lower series members up to $8G$ show no trace of structure, but beyond this, the terms $8G, 9G, 10G, 11G$, show a hyperfine structure practically independent of $n$. The terms change in character and partake of both singlet and triplet properties. Furthermore the J values are completely undetermined between the values 5, 4, 3. The constancy of the structure is due to the fact that the g electron contributes very little for the high $n$ values, and in effect the structures depend only on the value of the structure of the $3s\,^2S_1$ term of Al III, i.e. on the structure of the ion of Al II.

If the above theory is correct, then it can be shown that forbidden lines should appear, i.e. the J selection principle should be violated, and this is found to be the case. Paschen also explains the perturbation in Hg I as a Zeeman effect. The perturbing terms here are $6\,^1D_2$ and $6\,^3D_1$, and as long ago as 1902 Runge and Paschen found that these terms showed asymmetrical Zeeman patterns. Parchen was of the opinion that the repulsion effects reported by Schüler and Jones are not a resonance phenomenon, but are just what would take place if the perturbation was a Zeeman effect. Goudsmit and Bacher have given a more extended theoretical treatment of the effect. The theory predicts the existence of forbidden transitions when perturbation occurs. Where only two levels of the same F value are apparently repelled, then if $\delta$ is the displacement from the null position and if $\Delta$ is the actual separation of the two perturbing levels, the ratio of the intensity of forbidden transitions I to allowed transitions $I_0$ is given by the formula $\frac{I}{I_0} = \frac{\delta}{\Delta + \delta}$. In Hg I this is only true for the odd isotope terms, for there is no perturbation in the even isotopes. It follows from the theory that transitions should take place from $6\,^1D_2$ to terms with $J = 0$, and from $6\,^3D_1$ to terms with $J = 3$. It is of particular interest that the line $6d\,^1D_2 - 7p\,^3P_0$ actually occurs and when examined for hyperfine structure it is found that the even isotopes are entirely missing. This is a valuable proof of the theory which demands forbidden transition in the odd isotopes only.

Deviations from the interval rule which are due to the nucleus are discussed in Chapter XII.

Intensity perturbations

The intensities in hyperfine structure components can, of course, be modified by self absorption, but this cannot be regarded as a true perturbation effect. Intensity perturbations have only been found in a hollow cathode discharge in cadmium when the exciting currents are high. In certain lines the transitions which involve the highest F value of the upper term are relatively strengthened. This effect is not due to self absorption, and has not yet been explained.

REFERENCES

1 Schüler and Jones, Zeit. f. Phys., 77, 801, 1932.
NUCLEAR SPINS

Nuclear spins have now been determined for a large number of isotopic species. The results are given in Table VIII, A, B, C, D below. These are collected from an extensive literature.

Atoms can be divided into four classes according to their nuclear mass and charge, thus:—

A, odd atomic weight, odd charge.
B, odd atomic weight, even charge.
C, even atomic weight, odd charge.
D, even atomic weight, even charge.

Class A atoms possess an odd proton.
Class B atoms possess an odd neutron.
Class C atoms possess an individual single neutron-proton pair.
Class D atoms possess mostly a mass divisible by 4, i.e., alpha particle mass as unit.

The various tables will now be discussed. In no case has a spin exceeding \( \frac{3}{2} \) been recorded. The last column shows the nuclear magnetic moments in Bohr magnetons. These will be discussed in a succeeding chapter. The statistical distribution of the spins is of interest and will also be discussed later.

**Class A atoms**

In all these odd atomic weights observed spins are half integral multiples of \( \frac{\hbar}{2\pi} \) and spins vary from \( \frac{1}{2} \) to \( \frac{3}{2} \) with no apparent regularity.

**Class B atoms**

Again all observed spins are half integral and in the same range as Class A atoms. In quite a considerable proportion

---

**Table VIII (A)**

<table>
<thead>
<tr>
<th>Mass.</th>
<th>Charge.</th>
<th>Symbol</th>
<th>Spin ( \frac{\hbar}{2\pi} )</th>
<th>Mag. Mom. (Bohr Magneton)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>1</td>
<td>H</td>
<td>( \frac{1}{2} )</td>
<td>2.7896</td>
</tr>
<tr>
<td>7</td>
<td>3</td>
<td>Li</td>
<td>( \frac{1}{2} )</td>
<td>3.2532</td>
</tr>
<tr>
<td>11</td>
<td>5</td>
<td>B</td>
<td>( \frac{1}{2} )</td>
<td>2.686</td>
</tr>
<tr>
<td>15</td>
<td>7</td>
<td>N</td>
<td>( \frac{1}{2} )</td>
<td>0.280</td>
</tr>
<tr>
<td>19</td>
<td>9</td>
<td>F</td>
<td>( \frac{1}{2} )</td>
<td>2.625</td>
</tr>
<tr>
<td>23</td>
<td>11</td>
<td>Na</td>
<td>( \frac{1}{2} )</td>
<td>2.215</td>
</tr>
<tr>
<td>27</td>
<td>13</td>
<td>Al</td>
<td>( \frac{1}{2} )</td>
<td>3.63</td>
</tr>
<tr>
<td>31</td>
<td>15</td>
<td>P</td>
<td>( \frac{1}{2} )</td>
<td>—</td>
</tr>
<tr>
<td>35</td>
<td>17</td>
<td>Cl</td>
<td>( \frac{1}{2} )</td>
<td>1.368</td>
</tr>
<tr>
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<td>17</td>
<td>Cl</td>
<td>( \frac{1}{2} )</td>
<td>1.316</td>
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<tr>
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<td>Sc</td>
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<tr>
<td>51</td>
<td>23</td>
<td>V</td>
<td>( \frac{1}{2} )</td>
<td>—</td>
</tr>
<tr>
<td>55</td>
<td>25</td>
<td>Mn</td>
<td>( \frac{1}{2} )</td>
<td>—</td>
</tr>
<tr>
<td>59</td>
<td>27</td>
<td>Co</td>
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</tr>
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<td>29</td>
<td>Cu</td>
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</tr>
<tr>
<td>65</td>
<td>29</td>
<td>Cu</td>
<td>( \frac{1}{2} )</td>
<td>2.6</td>
</tr>
<tr>
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TABLE VIII (B)

B. Odd Mass and Even Charge, i.e. Odd Neutron

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TABLE VIII (C)

C. Even Mass and Odd Charge (Neutron-Proton Pair)

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TABLE VIII (D)

D. Even Mass (Multiple of 4) and Even Charge

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<td>32</td>
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<td>S</td>
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</table>

of cases the magnetic moment is negative. This means that the nuclear-electron coupling is such that the hyperfine structure multiplets are completely inverted, the lowest F value usually lying deepest.

Class C atoms

This is a limited group. In each case but the last the spin is 1 and the magnetic moment small. It is reasonable to conclude that these lighter elements are built up from alpha particles with an odd neutron-proton pair added in each case. The magnetic moment is small and as a result hyperfine structures are too small to be measured. Both mechanical
and magnetic moments have been derived by methods other than those using hyperfine structures, in particular the magnetic deflection method has been very successful here.

**Class D atoms**

The atoms in this group have mass divisible by four. The spins are known to be zero with certainty from the fact that alternate lines are missing in the band spectra. Absence of hyperfine structure is not evidence of zero spin, but only of small nuclear magnetic moment; an obvious fact, not universally recognised.

These spin values appear to indicate that alpha particles are basic building stones in nuclear construction.

An exceptional case of considerable interest is that of the potassium isotope mass 40 and charge 19, whose spin is given by magnetic deflection experiments as 4. This isotope is radioactive, with a very long decay period and beta-decay theories predict in such a case a large nuclear spin.

*It is clear how incorrect is the frequently made assumption of zero spin for even mass.* One might in fact state categorically instead that even mass has integral spin which may or may not be zero and indeed, as far as the evidence goes, appears only to be zero when the mass is divisible by 4 and the charge is even.

**Statistical distribution of nuclear spin**

Tolansky has reported a striking statistical difference in the distribution of nuclear spin of class A and class B atoms. This is shown in Fig. 25, which is a histogram showing for the two classes the spins of nuclei plotted against the number having those spins.

For the *odd proton class* above the distribution of spins is roughly that of a probability distribution curve.

For the *odd neutron class* below the distribution is strikingly different. It is clear that low spin values are favoured. This fact has been used to support a theory of nuclear spin which attributes the spin to an odd nuclear particle (theory of Landé). It was shown by Inglis that in such a case low spin values will arise in the odd neutron atoms.

![Graph](image-url)
CHAPTER XI

NUCLEAR MAGNETIC MOMENTS

Magnetic and mechanical moments

We have already seen that the number of hyperfine structure levels in a term depends upon the value of \( I \) (provided that \( I > J \)). The width of the structure is determined, not by the mechanical moment, but by the magnetic moment of the nucleus. The effect of the nuclear mechanical spin is to create a nuclear magnetic field, which, coupling with the magnetic field produced by electron spin and orbital moment, gives the observed hyperfine structure pattern. The coupling theory has been dealt with in earlier chapters, but it is apparent that the size of the structure is dependent on the value of the nuclear magnetic moment. It will be shown later that in favourable cases this moment can be calculated from the observed patterns.

It will be recalled that the phenomenon of alternating intensities in band spectra arises from the existence of the mechanical moment \( I \), whilst the widths of the hyperfine structures in the line spectra depend (amongst other things) upon the magnetic moment, i.e. on \( g(I) \). Cases occur where an atom has a definite value of \( I \) as revealed by the band spectrum, yet no structures are observable, even in terms involving penetrating electrons. The conclusion must be made that the nuclear magnetic moment is small, so that no measurable structures exist. Associated with this is the problem of the reported zero spins for so many even isotopes. All that has really been observed, in line spectra, is that the even isotope lines show no structure. Band spectrum observations in \( \text{He}^4 \), \( \text{C}^{12} \), \( \text{O}^{16} \), \( \text{S}^{32} \), prove that these four even atoms have zero spin, but attention may be drawn to the fact that these atomic weights are divisible by four. The assumption has been made that because these four even atoms are known to possess a zero spin, therefore the absence of structure in other even isotope lines is also an indication of zero spin. This may not be true. It is quite possible that some even isotopes may have a mechanical spin but have extremely small nuclear magnetic moments. This is certainly true in some cases. This point cannot be over-emphasised, for unfortunately a tendency has grown up in which workers have unquestioningly assumed that absence of structure in even isotopes means zero spin.

Magnetic moment of the proton

The value of the magnetic moment of the proton is of particular interest, for theories of nuclear spin have been proposed which attribute the nuclear magnetic properties to protons and neutrons. It is apparent from the relatively small values of the nuclear spins (\( \frac{3}{2} \) is the greatest observed with certainty) that only a very limited number of the nuclear protons or neutrons contribute to the spin. It may be possible, therefore, to correlate the value of the proton or neutron magnetic moment with nuclear magnetic moments in general, and indeed attempts to do this have been made. A value of the magnetic moment of the proton can be deduced from very simple considerations. Both electron spin and proton spin are \( \frac{1}{2} \), but as the protonic mass is 1836 times as great as the electronic mass, the magnetic moment of the spinning proton should be \( \frac{1}{1836} \) that of the electron, i.e. \( \frac{1}{1836} \) of a Bohr magneton. It has been pointed out before that the doublet separation in gross structure multiplets is due to the electron magnetic moment, so that on this basis, to a first approximation, hyperfine structures should be 2000 times smaller than the gross structure width.

The magnetic moment of the proton was first determined experimentally by Estermann and Stern by means of a refined improvement of the Stern-Gerlach experiment. In this experiment a beam of atoms or molecules is passed through an intense non-homogeneous magnetic field, and the deflections that the beam experiences enable the magnetic moment of the particles of the beam to be calculated. Molecular hydrogen exists in two states, para and ortho,
in one of which the electron spins oppose and in the other assist. A beam of hydrogen was used in which the electron spins cancel out, so that the magnetic moment due to the protons could be observed, and it was found that the proton magnetic moment was some 2.8 times that expected, i.e. \(2.8^{1836}\). Thus if the magnetic moment is \(\mu\) and the nuclear spin \(I\), then \(\mu = I \cdot g(I)\), providing the correct units are employed. \(\mu\) is expressed in theoretical proton magnetons \(\frac{eh}{4\pi Mc}\), that is to say, the Bohr electron magneton divided by the nuclear mass \(M\). By analogy with the electron \(g\) value, one expects for a proton that \(I = \frac{1}{2}\) and \(\mu = 1\), so that \(g(I)\) would be 2. Experiment, however, shows that \(\mu\) for the proton is 2.79, so that \(g(I) = 5.58\). (Precision values are discussed later.)

The magnetic deflection method of calculating the magnetic moments has been used in a number of cases. It is very important, since it is a direct method, and it is to be hoped that it will be extended in the near future.

**Calculation of nuclear magnetic moments**

The calculation of nuclear magnetic moments from the data of hyperfine structure has been given by Goudsmit. As we have already seen, the coupling energy between electron and nucleus determines the total splitting, and this is given by

\[
W = A \cdot J \cdot \cos (IJ) = \frac{3}{2} A [F(F + 1) - J(J + 1) - I(I + 1)].
\]

When an alkali-like spectrum (a single optical electron) is under consideration, then the interval factor \(A\) is that of the optical electron \(a\).

Coupling theory shows that

\[
a = \frac{Rz^2z^3}{n^3(1 + \frac{1}{2})j(j + 1)} \cdot g(I) \cdot 1836 \text{ cm}^{-1}.
\]

In this formula \(R\) is Rydberg's constant, \(z^2\) is the spectroscopic doublet constant, \(z\) is the nuclear charge, \(n\) the total quantum number, \(I\) and \(j\) have the usual significance, and \(g(I)\) is the numerical ratio of the magnetic to mechanical moments.

When the case of a penetrating electron is considered, the formula is altered and a relativity correction has to be applied. The final formula for the \(g(I)\) factor in the case of a penetrating \(s\) electron is

\[
g(I) = \frac{a \cdot n_0^3 \cdot 1836}{8Rx^2 \cdot Z_iz_0^2 \cdot K(jZ_i)}.
\]

It will be seen that only two changes have really been made. \(n_0\) is the effective total quantum number of the term, which is easily obtained if the absolute term value is known. \(z^2\) has now been replaced by \(Z_iZ_0^2\), where \(Z_0\) is the effective nuclear charge of the inner portion of the electronic orbit and \(Z_i\) is the average effective nuclear charge when the electron is passing through the inner penetrating portion of the orbit. For a neutral atom \(z_0 = 1\), for a singly ionised atom \(z_0 = 2\), etc. In heavy atoms and deeply penetrating orbits the electron approaches so closely to the nucleus that practically all the screening effect of the outer electrons is reduced to zero. The effective \(Z_i\) may thus be taken as approximately equal to the atomic number. The relativity correction \(K(jZ_i)\) depends upon \(j\) and \(Z_i\), and a table of values for \(K\) has been calculated by Goudsmit.

As an example of applying the formula, the following calculation of the \(g(I)\) factor of tin will now be given.

The nuclear spin of tin is \(\frac{1}{2}\) and the 6s \(^{2}S_{1/2}\) term (odd isotopes) is a doublet 0.207 cm\(^{-1}\). The \(F\) values of the hyperfine structure levels are \(1\) and \(0\), so that the separation is the interval factor. Hence

\[
A = a = - 0.207 \text{ cm}^{-1}.
\]

The \(n_0\) for 6s \(^{2}S_{1/2}\) in tin is 2.685, and \(R = 109,737\). The effective inner charge \(Z_i\) is 50, which is the atomic number of tin, and as the \(Sn\ II\) spectrum is under consideration,
$z_0 = 2$. The relativity correction for $j = \frac{1}{2}$ and $Z_i = 50$ is 1.3, and the spectroscopic doublet constant has the value $5.305 \times 10^{-5}$. On substituting these data, one obtains $g(I) = -1.78$. Since $I = \frac{1}{2}$, and as $\mu = 1\cdot g(I)$, the nuclear magnetic moment is $-0.89$ Bohr proton magnetons, or $\frac{20.89}{10}$ that of the electron.

The formula utilized above is only to be used with $s$ electrons, and when other relatively non-penetrating orbits are studied a different formula is required, and has been deduced by Goudsmit as follows:

$$g(I) = \frac{aZ_i j(j + 1)(1 + \frac{1}{2})}{\Delta \nu \cdot I(I + 1)} \cdot \frac{\lambda(IZ_i)}{K(jZ_i)} \cdot 1836.$$ 

$g(I)$, $a$, $Z_i$, $j$, $I$, $K$, have their previous meanings; $\lambda(IZ_i)$ is a relativity correction depending upon the $I$ value, i.e. on whether $p$, $d$, or $f$ electrons are involved. $\Delta \nu$ is the spin doublet separation. Thus, if a $6p$ electron is under review, then $\Delta \nu$ is the separation between $6p \cdot 2P_\frac{1}{2}$ and $6p \cdot 2P_\frac{3}{2}$, and similarly for other electrons. With $p$ electrons, $Z_i$ is not equal to the atomic number $Z$, but throughout the periodic table it has always been found that

$$Z_i \approx Z - 4.$$ 

As this is always a little uncertain, the $g(I)$ values calculated from $p$ electrons may be too small.

When the terms arise from complex electron configurations, the separate electron interval constants must be obtained before the $g(I)$ factor can be calculated. Uncertainty often arises here, so that many calculated $g(I)$ factors are only approximate.

**Nuclear magnetic moments observed**

The known nuclear magnetic moments are shown in the preceding chapter in Table VIII which give the known nuclear spins. A comprehensive list of data has thus now been accumulated. Magnetic moments have been derived in many cases by two independent methods, namely (1) magnetic deflection of molecular and atomic beams, (2) hyperfine structure measurements. Where results are obtained by both methods, agreement is effectively complete. The precision and certainty in the magnetic deflection experiments are far superior to those in the spectroscopic observations. The information given refers to the normal state only, but suffices to give both mechanical and magnetic moments.

There are broad distinctions between the nuclear magnetic moments of the various classes A, B, C, D, of Table VIII. In class A the moments distribute from 2.780 to downwards. Of particular importance is the large value for the proton. This has been explained theoretically in terms of the Yukawa theory of the meson. According to this, a proton can be regarded as partially dissociated for a small fraction of its life-time into a neutron and positive meson. The latter having a mass of less than one-tenth that of the proton has a correspondingly larger magnetic moment. Thus the mean magnetic moment of the combination exceeds the Bohr magneton.

In class B the most striking feature is the existence of so many negative magnetic moments. The only possible explanation is that the magnetic moment of the neutron is itself negative. This again is explicable on the Yukawa theory as shown in discussing class C. Class C is of considerable interest.

The $D^2$ nucleus has a spin of 1 and as there is considerable independent evidence that proton and neutron have each a spin of $\frac{1}{2}$, clearly the two spins add in parallel. But the magnetic moment of $D^2$ is much less than that of $H$ (0.8565 compared with 2.780).3 One must conclude that the magnetic moment of the neutron is negative, i.e. that it acts in an opposite direction relative to that of the mechanical spin in the case of $H$.

Since a neutron is uncharged, the spin of such a particle should apparently have no magnetic moment. The Yukawa dissociation theory avoids the difficulty. A neutron is partially dissociated for a fraction of its life-time into a proton and a negative meson and as the magnetic moment of the latter is greater than that of the proton, there remains
a residual mean negative magnetic moment for the neutron as a whole.

Clearly of some considerable importance is the fact that Li$^6$, B$^{10}$, N$^{14}$, each probably built out of alpha particles and an added proton-neutron pair, show a regular decrease in magnetic moment.\(^6\)

In class D, nuclear spin is zero and there appears to be no reason for supposing that there exists any magnetic moment.

Various theories of nuclear magnetic moment have been proposed. One theory attributes the moment largely to that of single (or odd) nuclear protons or neutrons. The theories, however, are as yet not sufficiently developed to justify their inclusion in a book of this nature.

**REFERENCES**


**CHAPTER XII**

**NUCLEAR ELECTRIC QUADRUPOLE MOMENT**

**Deviations from interval rule**

It has been found that deviations from the interval rule may arise through electrostatic interaction between a nucleus and the outer electrons. In 1932 Tolansky\(^1\) reported a marked break-down in the cosine interaction law for a term in the As II spectrum. Other cases were found, and in 1935 Schüler and Schmidt\(^2\) correctly interpreted such deviations from the interval rule as showing the existence of an asymmetrical nuclear electric charge distribution. This method of studying such nuclear effects is clearly one of much importance to future nuclear theories, and will be dealt with in the following.

Schüler and Schmidt pointed out that the interval rule, represented by $E = A_0 + A' \cos IJ$ fails and must be replaced by an expression which is

$$E = A_0 + A' \cos IJ + B' \cos 2IJ$$

This becomes equivalent when substitutions are made to $E = a_0' + \frac{1}{2}a'C + b'C(C + 1)$ in which

$$C = F(F + 1) - I(I + 1) - J(J + 1).$$

In this expression, $a_0'$, $a'$ and $b'$ are constants, $a_0'$ being the displacement of the centre of gravity from the position the term would occupy if the spin were zero, $a'$ is the hyperfine structure interval factor and $b'$ is a measure of the deviation from the interval rule and is also a measure of the nuclear quadrupole moment.

The separations produced by the electrostatic interaction of nucleus and electrons can be shown to be of the form

$$- P \cdot Q \cdot \frac{(\frac{3}{2})C(C + 1) - \frac{3}{2}IJ(I + 1)(J + 1)}{IJ(2I - 1)(2J + 1)}.$$
The quantity $P$ is a measure of the deviation from spherical symmetry of the charge density of the electrons. It is zero for a spherically symmetrical distribution. (This is characteristic of any $S$ state.) It is positive for a prolate distribution and negative for an oblate distribution in the direction of angular momentum.

In like manner $Q$ is an average over the charge density of the nucleus. It is positive or negative according as the nucleus is prolate or oblate in the direction of its angular momentum. $Q$ is called the electric quadrupole moment of the nucleus.

In conformity with the preceding formulae, it is found by observation that the perturbation of a hyperfine structure level depends upon the $F$ value of that level. To discuss further the nuclear-electron interaction, it is necessary to consider very briefly the wave-mechanics model of electron distribution.

**Wave-mechanics model**

It is now firmly established that the concept of the electron as a small spherical unit of electricity is only a rough approximation and it has been found necessary to replace the over-simplified idea of Bohr orbits by "eigen-functions". These can effectively be considered as expressions describing the distribution of electrical density representing the equivalent of the electron. Ultimately it is as if the small "billiard ball" electron model were replaced by a smeared out or blurred electrical charge distribution which has maxima of probability at distances from the nucleus corresponding to what were formerly the Bohr orbital radii.

The equivalent electrical distributions for $S$, $P$, $D$, etc., states have been evaluated and it appears that only $S$ terms have a spherically symmetrical charge distribution, whilst $P$, $D$, etc., terms are marked by unsymmetrical distribution.

Suppose there is interaction between a non-spherically symmetrical nucleus and a non-spherical electron charge distribution, then, as Schüler and Schmidt first pointed out, there will be a break-down in the interval rule, as shown by Fig. 26 which illustrates two cases.

**FIG. 26.—Origin of deviation from the interval rule.**

Suppose we postulate that the nucleus (shaded in B and C) is not spherically symmetrical also and let it be in the form of a rotation ellipsoid whose axis of rotation is the direction of the nuclear spin I. Then Figs. I (B) and II (B) show the situation when I and J are parallel (F is a maximum) and I (C) and II (C) when I and J are perpendicular. When I and J are anti-parallel, the effect is identical with that when I and J are parallel.

From energy considerations the levels for which the overlap of the electron and nuclear distributions is greatest will be the higher. Thus I (B) lies higher than I (C) and II (B)
lies deeper than II (C). B then corresponds to both the highest and lowest levels in a hyperfine structure pattern and C to an intermediate level. It follows that in Case I the intermediate levels displace downwards relative to the outer levels and in Case II upwards. This is illustrated by the levels at the extreme right.

**Break-down of interval rule**

Clearly, such displacements lead to the break-down of the interval rule. They can often be fitted into a quadratic formula of the type described above, but there are discrepancies in many cases. When the formula is reasonably fitted the constants \(a'\) and \(b'\) can be derived. Casimir has shown that the constant \(b'\) is related to the quadrupole moment \(Q\) such that

\[
Q = \frac{b'8I(2I - 1)J(2J - 1)}{3e^2 P}
\]

where \(e\) is the electronic charge, and the quantity \(P\) is calculable, depending on the particular term. Thus

\[
P = \frac{Z_i Z_0 K}{n^3 L(L + \frac{1}{2})(L + 1)}
\]

where \(K\) is a calculable constant, e.g. \(-\frac{1}{3}\) for \(^2D_1\) or \(-\frac{1}{7}\) for \(^2D_3\), etc. The symbols \(Z_i, Z_0, n^3\) are as used in earlier chapters.

Numerical substitution from results show that \(Q\) can vary from 0 to \(5 \times 10^{-24}\) cm\(^2\) from which it appears that the deviation from spherical symmetry is a matter of at most a few per cent. (Casimir).

There exists some doubt about many of the reported values for nuclear electric quadrupole moment. It is clear that many causes can lead to spurious values. Thus bad measurement, differential absorption effects, incipient electric and magnetic field perturbation, neighbouring level perturbation, and even photographic developer effects, can all give rise to a spurious break-down in the interval rule, and make a false quadrupole moment appear. It is perhaps significant that often two terms in a spectrum give different values of the quadrupole moment and indeed one case had been reported in which the apparent nuclear quadrupole moment is positive from one term and negative from another, clearly a serious inconsistency and strong evidence of the existence of external perturbation as the cause, and not quadrupole moment.

The quadrupole moments reported so far are given in the table below, the units being \(10^{-24}\) cm\(^2\). They have been collected from a number of sources.

**Electrical nuclear quadrupole moments**

**Table IX**

<table>
<thead>
<tr>
<th>Atom</th>
<th>Mass</th>
<th>Charge</th>
<th>Spin</th>
<th>Magnetic Moment</th>
<th>Quadrupole Moment $\times 10^{-24}$ cm(^2)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Cu</td>
<td>63</td>
<td>29</td>
<td></td>
<td>$+2.5$</td>
<td>$-0.1$</td>
</tr>
<tr>
<td>Cu</td>
<td>65</td>
<td>29</td>
<td></td>
<td>$+2.6$</td>
<td>$-0.1$</td>
</tr>
<tr>
<td>Ga</td>
<td>69</td>
<td>31</td>
<td></td>
<td>$+2.0$</td>
<td>$+1.0$</td>
</tr>
<tr>
<td>Ga</td>
<td>71</td>
<td>31</td>
<td></td>
<td>$+2.5$</td>
<td>$+0.5$</td>
</tr>
<tr>
<td>As</td>
<td>75</td>
<td>33</td>
<td></td>
<td>$+1.5$</td>
<td>$+0.3$</td>
</tr>
<tr>
<td>Kr</td>
<td>83</td>
<td>36</td>
<td></td>
<td>$-1.0$</td>
<td>$+0.5$</td>
</tr>
<tr>
<td>In</td>
<td>115</td>
<td>49</td>
<td></td>
<td>$+5.5$</td>
<td>$+0.8$</td>
</tr>
<tr>
<td>I</td>
<td>127</td>
<td>53</td>
<td></td>
<td>$+2.8$</td>
<td>$+0.5$</td>
</tr>
<tr>
<td>Xe</td>
<td>131</td>
<td>54</td>
<td></td>
<td>$-0.8$</td>
<td>$0 \pm 0.1$</td>
</tr>
<tr>
<td>Eu</td>
<td>151</td>
<td>63</td>
<td></td>
<td>$+3.4$</td>
<td>$+2.5$</td>
</tr>
<tr>
<td>Eu</td>
<td>153</td>
<td>63</td>
<td></td>
<td>$+1.5$</td>
<td>$+2.5$</td>
</tr>
<tr>
<td>Yb</td>
<td>171</td>
<td>70</td>
<td></td>
<td>$+0.4$</td>
<td></td>
</tr>
<tr>
<td>Yb</td>
<td>173</td>
<td>70</td>
<td></td>
<td>$-0.7$</td>
<td>$+3.9$</td>
</tr>
<tr>
<td>Cq</td>
<td>175</td>
<td>71</td>
<td></td>
<td>$+2.6$</td>
<td>$+5.9$</td>
</tr>
<tr>
<td>Re</td>
<td>185</td>
<td>75</td>
<td></td>
<td>$+3.3$</td>
<td>$+2.8$</td>
</tr>
<tr>
<td>Re</td>
<td>187</td>
<td>75</td>
<td></td>
<td>$+3.3$</td>
<td>$+2.6$</td>
</tr>
<tr>
<td>Hg</td>
<td>201</td>
<td>80</td>
<td></td>
<td>$-0.6$</td>
<td>$+0.5$</td>
</tr>
<tr>
<td>Bi</td>
<td>209</td>
<td>83</td>
<td></td>
<td>$+3.6$</td>
<td>$-0.4$</td>
</tr>
</tbody>
</table>

**REFERENCES**

APPENDIX 1

HYPERFINE STRUCTURE—ZEEMAN EFFECT

This will be only described very briefly, more detail being available elsewhere.1 The complete Paschen-Back effect is the simplest to consider. In this, the external magnetic field is so strong that its coupling with \( I \) and \( J \) is great enough to break the mutual interaction of \( I \) and \( J \). It can be shown that, to a first approximation, the Zeeman pattern is the same as that of an atom when no nuclear spin is considered, that is to say, the line will be split up into its usual Zeeman components. The introduction of a nuclear spin affects each of these Zeeman components slightly and this can sometimes be detected. The total interaction energy of the atom with the field \( H \) is

\[
v = O[m_1 \cdot g(I) + m_J g] + A \cdot IJ \cdot \cos (IJ),
\]

where \( O \) is the Lorentz unit \( \frac{eH}{4\pi m_e c} \), \( g \) is the ordinary electronic \( g \) factor, \( g(I) \) the nuclear factor, and \( m_1 \) and \( m_J \) the projections of \( I \) and \( J \) along \( H \). Averaging the cosine term reduces the expression to

\[
m_J(Og + Am_1).
\]

\( m_1 \) can assume any of the \( 2I + 1 \) integral or half integral values between \( + I \) and \( - I \). Hence, each Zeeman level, which without the spin is single, now splits up into \( 2I + 1 \) components equally spaced by the distance \( Am_J \). The only transitions allowed are those in which \( m_1 \) does not change, so that each gross Zeeman line splits into \( 2I + 1 \) hyperfine structure Zeeman lines. For two terms \( J \) and \( J' \) the separations will be \( Am_J - A'm_J \). The \( A \) and \( A' \) values can be calculated from the hyperfine structures which arise when there is no magnetic field, and \( m_1 \) is \( + \frac{1}{2} \) or \( - \frac{1}{2} \) for both states when parallel components are observed. Hence, the separation can be predicted. The number of components gives the simplest and most satisfactory way of obtaining the nuclear spin. Thus, in \( Bi \), where the Zeeman theory was beautifully verified, each gross Zeeman component consisted of 10 lines. As this is \( 2I + 1 \), then \( I = \frac{9}{2} \).

If incomplete Paschen-Back effects exist, the pattern becomes extremely complicated, but the experimental evidence seems to agree very well with theory, even to the extent that the normally forbidden transitions occur as is to be expected when external fields are applied.

REFERENCES

APPENDIX II

LANDÉ'S PARTICLE THEORY OF NUCLEAR SPIN

Landé has proposed that the nuclear spin of an odd weight atom may be due to the single odd particle, respectively a proton and neutron in types A and B of Table VIII (see before). Consider first odd-proton nuclei. The unbound odd proton is considered to have a spin \( s = \frac{1}{2} \) and an orbital angular moment \( I \). These combine vectorially to a resultant \( I \) which is the nuclear spin. Thus \( I \) may be either \( I + \frac{1}{2} \) or \( I - \frac{1}{2} \). From this Landé derives a formula for \( g(I) \) which is

\[
g(I) = g_t \left[ \frac{I(I + 1) + I(I + 1) - s(s + 1)}{2I(I + 1)} \right] + g_s \left[ \frac{s(s + 1)}{2I(I + 1)} \right]
\]

in which \( s = \frac{1}{2} \), the proton spin, \( g_t = 1 \) and \( g_s \) the magnetic factor of the proton. By taking \( g_s = 4 \), Landé was able to obtain an approximate fit between observed and calculated \( g(I) \) values, for \( I \) can only arise from one of two adjacent \( I \) values and the \( g(I) \) nearest to that observed can be adopted. However, it is now clear that the approximate fits obtained are probably largely accidental, since \( g_s \) is known with certainty to be much more than 4, in fact 5.58 (see Table VIII, magnetic moment of proton).

The above theory was extended to include the odd-neutron case and \( g_t \) was postulated zero, since the neutron was considered to have no magnetic moment due to orbital motion. Observations were best fitted by adopting \( g_s = -1.2 \), giving the neutron a negative magnetic moment of \(-0.6\) magnetons. This value leads to both positive and negative \( g(I) \) factors, as indeed are found experimentally. Agreement is, however, but partial, and Tamm and Altschuler consider that sometimes a pair of neutrons does not form a closed shell of zero spin, but add up to a spin of 1. This spin combines with the remaining odd neutron (or proton) and if the former, the combined spin will be \( \frac{3}{2} \). This assumption doubles the possible number of \( g(I) \) factors that can be calculated for a single nuclear spin and thus introduces further uncertainty.

Thus, whilst the theory that the spin is due to one particle (or at most three particles) is clearly of considerable interest, accounting as it does for the small observed spins, yet it cannot be considered to have established itself as yet.

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