Backscatter cross sections of H, He, C and N of KeV energies from gold atoms in a solid.

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BACKSCATTER CROSS SECTIONS OF H, He, C AND N OF keV ENERGIES
FROM GOLD ATOMS IN A SOLID

by

Jagdish S. Patel

A Thesis
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ABSTRACT

The backscatter cross sections for the scattering of keV projectiles of H⁺, He⁺, C⁺ and N⁺ through 160° from thin gold films have been measured. For the energy range 5 < E < 120 keV, the He cross sections lie 5% above Thomas Fermi cross sections, and for 20 < E < 120 keV, the cross sections for C and N lie 20-30% above Thomas Fermi cross sections, consistent with results obtained by Andersen et al. The discrepancy between theory and experiment is not understood.
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TABLE OF CONTENTS

ABSTRACT ........................................................................................................ iii

ACKNOWLEDGEMENTS ................................................................................. iv

LIST OF TABLES ............................................................................................. vi

LIST OF FIGURES .......................................................................................... vii

CHAPTER

I. INTRODUCTION ......................................................................................... 1

II. THEORY ........................................................................................................ 6
   II.a Validity of classical mechanics .......................................................... 6
   II.b Two body radial force problem ......................................................... 7
   II.c Rutherford scattering cross section ............................................... 9
   II.d Thomas Fermi theory ...................................................................... 10
   II.e Thomas Fermi potential between atoms and scattering cross section 12
   II.f Reduced Thomas Fermi cross section ............................................ 14
   II.g Conversion to laboratory coordinates from C.M. ......................... 15
   II.h Energy loss and multiple scattering ............................................. 15

III. EXPERIMENTAL ...................................................................................... 19
   III.a Overall plan ................................................................................... 19
   III.b Source and accelerator tube ......................................................... 19
   III.c Magnetic analyser ......................................................................... 22
   III.d Collimator and beam monitoring device ..................................... 23
   III.e Target .............................................................................................. 25
   III.f Electrostatic analyser and energy measurement ......................... 26
   III.g Estimate of target thickness ........................................................ 30
   III.h Note on film thickness variations ................................................. 30
   III.i Channeltron detector and detecting efficiency ............................ 34

IV. RESULTS AND CONCLUSIONS .............................................................. 38
   IV Results .................................................................................................. 38
   IV.a Conclusions .................................................................................... 45

REFERENCES ............................................................................................... 46

VITA AUCTORIS .......................................................................................... 47
LIST OF TABLES.

TABLE I  Some typical backscatter yield measurements for $H^+$, $He^+$, $C^+$ and $N^+$. $Y_o$ is the yield from the 13 Å Au film and the Be substrate. $Y_{Be}$ is the substrate yield, and $Y_{Au} = Y_o - Y_{Be}$ ...

TABLE II  Some typical backscatter yield measurements for $H^+$, $He^+$, $C^+$ and $N^+$. $Y_o$ is the yield from the 6.5 Å Au film and the Be substrate. $Y_{Be}$ is the substrate yield, and $Y_{Au} = Y_o - Y_{Be}$ ...
LIST OF FIGURES

Fig. 1 The orbit r(θ) for the scattering of a projectile by a fixed centre of force. p is the impact parameter, and θ is the deflection angle after scattering. .......................... 7

Fig. 2 The solid angle dΩ = 2πsinθdθ .................................. 10

Fig. 3 Thomas Fermi screening function F(x) .......................... 13

Fig. 4 Energy dependence in reduced units of the ratio between Rutherford and T.F. cross sections. The triangles refer to experimental points obtained by Andersen et al.3 from whose work, this figure has been reproduced ..................... 16

Fig. 5 The velocity dependence of the electronic stopping cross section S_e and the nuclear stopping cross section S_n. v_o = e^2/m_e ........................................... 17

Fig. 6 Schematic of experimental set up .................................... 20

Fig. 7 Source and accelerator tube (not to scale) ......................... 21

Fig. 8 Beam monitor and Faraday cup ..................................... 24

Fig. 9 Target sample showing evaporated gold deposits on a beryllium substrate ........................................... 25

Fig. 10 Electrostatic analyser for observing backscatter energy spectra for a laboratory angle θ_L of 154° ......................... 26

Fig. 11 Higher energy portion of backscatter spectrum for the scattering of 95 keV He^+ from gold. Points are experimental, and the solid line has been drawn to determine the maximum backscatter energy, E^+ = 88.81 keV. The error bar is the error in counting statistics ........................................... 28

Fig. 12 Backscatter energy spectrum of H^+ on gold. The width of the curve at half-height, ΔE, is proportional to film thickness. ........................................... 31

Fig. 13 Backscatter energy spectrum for 100 keV N^+ from a very thin gold film. The error bars show error in counting statistics ........................................... 33

Fig. 14 Apparatus for channeltron efficiency measurement. The diameters of S_e and S_n are (1.727 ± 0.003) mm and (0.085 ± 0.003) mm respectively .......................... 35

Fig. 15 The energy dependence of the detection efficiency for various ions ........................................... 37

Fig. 16 The energy dependence of the backscatter yield for H^+, He^+, and C from a 40 Å gold film ........................................... 41

Fig. 17 Backscatter yield x (energy)^2 for a 40 Å gold film (closed circles) and 13 Å gold film (open circles). Experimental points are normalised to the T.F. cross section shown as a solid line. The Rutherford cross sections are shown for comparison (dashed lines) ......................... 42
Fig. 18 The energy dependence of the backscatter cross sections of H, He, C and N from gold for a laboratory angle \( \theta_L = 159.95^\circ \). The open circles, crosses and triangles are measurements for a 13 Å, 6 Å and 2 Å thick gold film respectively (see text). The solid lines are T.F. cross sections and the dashed lines are Rutherford cross sections.
CHAPTER I

INTRODUCTION

Backscattering of MeV atomic projectiles from solid targets is important in several areas of physics. It is used in industry as a tool for investigating impurities in solids. Rutherford Backscatter Spectrometry (RBS) allows both the concentration levels and the depth profiles of the various elements in a solid to be measured. The latter is of importance in determining regions of dopants in a semiconductor.

The method consists in bombarding a solid target under investigation with a monoenergetic beam of light ions, usually He at 1-5 MeV. Although most of the ions come to a rest in the medium, a small fraction undergo nearly head-on collisions, scattering back into the vacuum. The target particle recoils with a recoil energy dependent on its mass. The backscattered projectile emerges with a reduced energy. The mass dependence of the elastic energy loss of the projectile, the Z dependence of the scattering yield and the slowing down process of the projectile by inelastic collisions to electrons, inside the target medium, all affect the energy distribution of the backscattered projectiles. Its analysis therefore yields high resolution depth information about the target.

Rutherford backscatter spectrometry makes the following assumptions:

1) An atomic projectile travels in a straight line until it approaches a large Z target atom, where it may undergo a single, violent back-scattering collision. Here, most of the scattering takes place at very small internuclear separations, between bare nuclei. After the collision, the projectile travels again in a straight line to re-enter vacuum, toward the detector. 2) The scattering event is an elastic collision,
between bare nuclei, permitting the use of Rutherford cross sections for
analysis.

These assumptions, however, cannot be fully justified. Since
scattering at small angles is far more probable than at larger angles,
a backscattering event in a solid target is most likely to be preceded
and followed by smaller angle scattering events (multiple scattering).
The first assumption is therefore contradictory. In assuming Rutherford
cross sections, the interatomic potential is taken to be the Coulombic
potential between the bare nuclear charges of the two particles. In
practice, the nuclear charges are screened by orbital electrons of the
two atoms. As the projectile energy decreases, the minimum distance of
approach increases, and the screening of the nuclear charges by electrons
becomes more important. It follows that the deviation from Rutherford
scattering is more pronounced at lower projectile energies.

However, even at higher energies, screening is important. Thus,
L'Ecuyer et al.\textsuperscript{1} find that the backscatter cross section for 1 MeV He
on Bi lies 4\% below the Rutherford cross section. There is a demand in
industry for accurate backscatter cross sections, and hence, the effect
of screening on backscatter cross sections needs to be understood. The
most widely used theory for cross sections that takes screening into
account is the Thomas Fermi (T.F.) theory. The work of van Wijngaarden
and Baylis\textsuperscript{2} on the scattering cross sections of H, He, B and N from Hg,
using a mercury gas beam target, for scattering angles below 40\textdegree is
consistent with the T.F. theory. In contrast, the cross sections of
Andersen et al.\textsuperscript{3} for H, He and Li projectiles in passing through thin
solid gold foils at scattering angles 3\textdegree <\theta < 15\textdegree generally lie above
the T.F. cross sections, by as much as 35\%. Data which significantly
tests screening effects are scanty, and needs to be supplemented for a variety of scattering conditions.

It appears\textsuperscript{1} that experimental measurements of backscatter yields from thick solid targets cannot give accurate backscatter cross sections since the backscatter yield is affected both by the electronic screening of nuclear charges and by multiple scattering effects. An ideal experiment for the purpose of measuring true backscatter cross sections consists of the study of single collisions between projectiles and isolated target atoms. From a measurement of the angular and energy dependences of the scattering yields, the cross section $\sigma(\theta, E)$ can be found. This in turn yields information about the nature of the interatomic potential $V(r)$. Such an experiment can be done by crossing a beam of projectiles with a target consisting of a gas beam of a sufficiently low pressure to eliminate multiple scattering events. Measurements with a gaseous mercury target beam at a pressure of about $3 \times 10^{-5}$ Torr have been carried out by van Wijngaarden et al.\textsuperscript{2,4} The scattering angles ranged from $2.8^\circ$ to $40^\circ$, the scattering yields for larger angles being too small to measure.

An order of magnitude calculation for the expected backscatter yield gives an appreciation for this limitation. Suppose one uses a $1$ cm thick mercury gas beam of $10^{-4}$ Torr pressure, equivalent to a number density of $3.4 \times 10^{12}$ cm$^{-3}$. Although in principle this number can be increased using higher gas pressures, in practice, the beam explodes at gas pressures above $10^{-4}$ Torr, and the scattering region becomes undefined. Using a beam of projectiles, of $1$ mm$^2$ cross section, the number of target particles present in the interaction region is $3.4 \times 10^{10}$. The Rutherford cross section for a $50$ keV $\text{He}^+$ beam backscattered at $160^\circ$ is

$$d\sigma/d\Omega = 1.4 \times 10^{-20} \text{cm}^2 \text{sr}^{-1}.$$
Assuming a very large current of $1.6 \times 10^{-6}$ Ampere, the number of collisions per second into unit solid angle is $4.8 \times 10^3$ s$^{-1}$. In a typical set-up, a detector would span at most $10^{-3}$ sr from the target, yielding a count rate of only $4.8$ sec$^{-1}$. This low count rate makes an accurate measurement at large backscatter angles very difficult.

Next, we compare this target with a hypothetical thin film of solid mercury, consisting of just a monolayer of atoms. We would now have approximately $10^{16}$ target atoms per cm$^2$, and the count rate for this film would be $1.4 \times 10^6$ sec$^{-1}$. In addition, the zone from which scattering occurs is precisely defined, compared with the diffuse nature of the gaseous target. Thus, scattering from a thin film of a monolayer (or less) has great advantages over scattering from a gas beam, where the number density of the target is too small. The much higher scattering yield allows a quick determination of the energy and angular dependences of the backscattering process.

It is not possible however to fabricate self-supporting thin films of thicknesses much below 100 Å. At these thicknesses, multiple scattering effects are large, but become smaller with thinner films. An approximation to a thinner self-supporting film can in principle be made by depositing a large Z scattering element onto a flat, low Z, backing material. Since the backscatter yield is proportional to the square of the atomic number, the backscatter yield from the supporting material is relatively small. The present experiment used thin gold films (Z=79) on a beryllium backing (Z=4). The scattering yield from the gold is then the difference between that from gold on beryllium and beryllium alone. By measuring the backscatter yields from varying thicknesses of targets, ranging from 2 Å to 40 Å, the results were essentially extrapolated to
zero thickness, thereby eliminating multiple scattering effects.
CHAPTER II

THEORY

II.a Validity of classical mechanics

The use of classical mechanics for describing a scattering process is valid if the de Broglie wavelength of a projectile
\[ \lambda = \frac{h}{p} \]
is smaller than any significant dimension of the scattering centre.

We consider an example where a 25 keV He\(^+\) projectile collides with a stationary Au atom (Z = 79). The de Broglie wavelength is
\[ \lambda = 9.0 \times 10^{-11} \text{ cm}. \]

The Thomas-Fermi scattering size of the collision (section II.e) is
\[ a = a_0 \left( \frac{Z_1}{3} + \frac{Z_2}{3} \right)^{1/2} \]
\[ = 1.0 \times 10^{-9} \text{ cm}, \]
where \( a_0 \) is the Bohr radius. The minimum distance of approach for a head-on collision is
\[ b = \frac{Z_1 Z_2 e^2}{1/2 \mu \nu} \]
\[ = 9.1 \times 10^{-10} \text{ cm}. \]

Since the screening length and the minimum distance of approach are larger than the de Broglie wavelength of the He projectile, interference effects are small. Hence backscattering of keV projectiles can be treated classically.
II.b Two-body radial force problem

\[ r = \frac{m_1 m_2}{m_1 + m_2} \]  \hspace{1cm} (2.1)

Fig.1. The orbit \( r(\theta) \) for the scattering of a projectile by a fixed centre of force. \( p \) is the impact parameter, and \( \theta \) is the deflection angle after scattering.

It is well known from classical mechanics that the two-body radial force problem can be reduced to that of a single particle of reduced mass

by a fixed centre of force, the centre of mass (CM). Fig.1 shows the orbit of a projectile approaching the centre of force at an impact parameter \( p \), with an asymptotic velocity \( v_0 \). The polar axis passes through the centre of mass at \( O \) and the point of minimum distance of approach at \( A \). For a radial force, the torque

\[ \vec{r} = \vec{r} \times \vec{F} = 0, \]

and hence the angular momentum \( \ell \)

\[ \ell = \mu r^2 \theta \]  \hspace{1cm} (2.2)
is conserved. Thus
\[ l = \mu v_o p \]
where the initial velocity \( v_o \) is related to the projectile energy as
\[ E = \frac{1}{2} \mu v_o^2. \]
The angular momentum \( l \) is then
\[ l = (2\mu E)^{\frac{1}{2}} p. \]
The total energy at an arbitrary distance \( r \) is
\[ E = \frac{1}{2} \mu v^2 + V(r) \quad \text{or} \quad E = \frac{1}{2} \mu v^2 + \frac{1}{2} \mu (r \phi)^2 + V(r). \]
Substituting 2.2 into 2.6 yields the time derivative of \( r \)
\[ \frac{dr}{dt} = \frac{2}{\mu} \frac{(E - V)}{r^2} - \frac{l^2}{\mu^2 r^2}. \]
The orbit \( \phi(r) \) is given by
\[ \frac{d\phi}{dr} = \frac{d\phi}{dt} = \frac{l}{\mu r^2} \frac{1}{dr/dt}. \]
Substituting for \( dr/dt \) (Eq. 2.7), we get
\[ \frac{d\phi}{dr} = \frac{l}{\mu r^2 \sqrt{\frac{2}{\mu}(E - V) - \frac{l^2}{r^2}}}. \]
Since the right hand side only depends on \( r \), the polar axis is a symmetry axis for the orbit. At the minimum distance of approach, \( r = \rho \), \( dr/d\phi = 0 \), hence \( \rho \) is given by equating
\[ \sqrt{\frac{2}{\mu}(E - V) - \frac{l^2}{\mu \rho^2}} \]
to zero. From inspection of Fig. 1, it can be seen that the scattering angle \( \theta \) is given by
\[ \theta + 2\phi_o = \pi \quad \text{or} \quad \frac{\rho^2}{2} = \frac{\pi}{2} - \phi_o \]
where
\[ \phi_o = \int_{r=\rho}^{\infty} \frac{d\phi}{dr} \, dr. \]
Substituting for $d\Phi/dr$, we get

$$\Phi = \int_{r=0}^{\infty} \frac{2}{\mu r^2} \sqrt{2\frac{(E-V)}{\mu}} - \frac{Z_1^2}{\mu} \frac{1}{r} \, dr$$

Thus the scattering angle $\theta$ for any arbitrary $V(r)$ is given by

$$\theta = \pi - 2 \int_{r=0}^{\infty} \frac{2}{\mu r^2} \sqrt{2\frac{(E-V)}{\mu}} - \frac{Z_1^2}{\mu} \frac{1}{r} \, dr$$

For Rutherford scattering, $V(r) = \frac{Z_1 Z_2 e^2}{r}$. Substituting this potential and the expression for $\Phi$ in terms of $p$ (Eq. 2.3) into 2.10 yields

$$\Phi = \int_{r=0}^{\infty} \frac{2}{\mu r^2} \sqrt{2\frac{(E-Z_1 Z_2 e^2/r)}{\mu}} - \frac{2Ze^2}{\mu} \frac{1}{r} \, dr$$

on integration, the relationship between the impact parameter $p$ and the scattering angle $\theta$ is found to be

$$p = \frac{Z_1 Z_2 e^2}{2E} \cot \frac{\theta}{2}$$

II.c Rutherford scattering cross section

Suppose the target particle is subjected to a stream of particles, with a current density $N \text{ cm}^{-2} \text{ sec}^{-1}$. All projectiles approaching the target with impact parameters between $p$ and $p+dp$ are scattered at a large distance from the scatterer into the solid angle $d\Omega = 2\pi \sin \theta d\theta$ (see Fig. 2). The scattering cross section $d\sigma(\theta)$, by definition is the fraction of particles scattered into angles $\theta$ to $\theta + d\theta$. Hence,

$$d\sigma = \frac{d\Omega}{N} = 2\pi p \, dp$$

The differential cross section $d\sigma/d\Omega(\theta)$ is the fraction of particles scattered into unit solid angle centred about a $\hat{\theta}$ direction. Hence,

$$\int \frac{d\sigma}{d\Omega} = \frac{2\pi p \, dp}{2\pi \sin \theta \, d\theta}$$
Substituting for $p$ from equation 2.13 gives the Rutherford differential scattering cross section:

$$\frac{d\sigma}{d\Omega} = \left(\frac{Z_1 Z_2 e^2}{4E}\right)^2 \sin^{-4} \frac{\theta}{2}.$$  \hspace{1cm} (2.15)

An identical result is obtained using quantum mechanics. Rutherford scattering cross sections for interatomic collisions are valid at high energies, for large scattering angles, when most of the significant scattering takes place at small internuclear distances. At lower energies, when screening becomes important, the Rutherford cross sections have to be replaced by Thomas Fermi cross sections.

II. d Thomas Fermi theory

The theory bases itself on treating the electrons of an atom as a Fermi gas. The atom is considered as a series of spherical shells, having radii $r_i$ to $r_i + dr_i$, in each of which the potential energy $V(r_i)$ is considered constant. $V(r_i)$ is then regarded as the well depth for the $i$th shell. Each shell is filled with electrons to the top of the well, similar to that for conduction electrons in a metal. Thus

$$E_F(r_i) = - V(r_i).$$  \hspace{1cm} (2.16)
where $E_f$ is the Fermi energy of the electron gas in the $i$th shell. The electron density $\rho$ for a Fermi gas in a well is given by
\[
E_f = \frac{2}{3} \pi^2 \frac{h^2}{2m} \rho \frac{2}{3}
\]
2.17

The density of the electron gas for a shell at a distance $r$ is now related to the potential energy:
\[
\rho(r) = \frac{2\frac{m}{3h^2}2}{3} \left[V(r)\right]^{\frac{3}{2}}
\]
2.18

To find the electric field, $E$ at a distance $r$, we use Gauss' law
\[
4\pi \rho^2 = 4\pi \left[Z_e + \int_0^r \rho(r) 4\pi r^2 dr\right]
\]
2.19

$Z_e$ denotes the nuclear charge, and the integral gives the total electronic charge. The potential energy for an electron, $V(r)$ is related to the electric field by
\[
-eE = \text{Force} = \frac{dV(r)}{dr}
\]
2.20

Substituting for $E$ in equation 2.19, we get
\[
\frac{2}{r_c} \frac{d}{dr} \left[V(r)\right] = Z_e \frac{2}{4\pi e^2} \int_0^r \rho(r) r^2 dr
\]
2.21

Differentiating with respect to $r$,
\[
\frac{dr}{dr} \left[\frac{2}{r} \frac{d}{dr} \left[-V(r)\right]\right] = 4\pi e^2 \rho(r) r^2.
\]

Substituting for $\rho(r)$ from 2.18 yields
\[
\frac{1}{r^2} \frac{d}{dr} \left[r^2 \frac{d}{dr} \left[-V(r)\right]\right] = \frac{4\pi e^2 \left[-V(r)\right]^{\frac{3}{2}}}{3 m^3} \frac{2}{3} \frac{2}{3}
\]
or
\[
\frac{2}{r} \frac{d}{dr} \left[-V(r)\right] + \frac{d^2}{dr^2} \left[-V(r)\right] = \left[-V(r)\right]^{\frac{3}{2}} \frac{2}{3} \frac{2}{3} \frac{2}{3} \frac{e^2 m^3}{3}
\]
2.22

To solve this equation, we write
\[
V(r) = \frac{2}{r} X(r)
\]
2.23

where the screening function, $X(r)$ describes the shielding of the
nuclear potential. As

\[ r \to 0, \quad \chi(r) \to 1 \]

2.24

giving a Coulombic interaction. At large distances from the nucleus, the
electrons shield the nuclear charges completely. Hence as

\[ r \to \infty, \quad \chi(r) \to 0 \]

2.25

Simplifying equation 2.22 yields

\[ x^{\frac{1}{2}} \frac{d}{dx} \chi(x) = \chi(x) \frac{1}{2} \]

2.26

where \( x = r/a \), a dimensionless parameter of distance, and

\[ a = \frac{1}{2} \frac{3\pi}{4} \frac{2}{h^2} \frac{1}{m e^2} \sqrt{\frac{2}{Z}} \approx 0.885 a_o \frac{1}{\sqrt{3}} \]

The differential equation 2.26, subject to the boundary conditions

2.24 and 2.25, is solvable by numerical methods. The behaviour of \( \chi(x) \) is
shown in Fig. 3. It is noted that the screening function \( \chi(x) = \chi(r/a) \).
only has appreciable values for \( r \approx a \). Thus, \( a = 0.885 a_o / \sqrt{3} \) is called
the screening length.

II. e Thomas Fermi potential between atoms and scattering cross section

In the theory of interatomic collisions, the potential is usually
written as

\[ V(r) = \frac{Z_1 Z_2 e^2}{r} \chi(r/a) \]

2.27

similar to equation 2.23 for the potential between an electron and its
parent atom. Since two atoms are involved, the concept of a screening
length is now somewhat uncertain. The screening length

\[ a = 0.8853 a_o \left( \frac{Z_1^2}{2} + \frac{Z_2^2}{2} \right) \frac{1}{\sqrt{3}} \]

2.28

has been used by Lindhard et al., in their paper on approximation
methods of scattering by screened Coulomb potentials. An alternative to
this form was also proposed:
Fig. 3. Thomas Fermi screening function $X(x)$. 
\[ a = 0.885 \lambda^2 \left( \frac{Z_1}{2} - \frac{Z_2}{2} \right)^2 \]  

Cross sections can be found numerically by substituting the Thomas Fermi potential into equation 2.28. These, for the screening length of Eq. 2.28 have been obtained by W. E. Baylis for comparisons. Thomas Fermi cross sections have previously been found to be consistent with experiment, for the scattering of \( \text{He}^+ \) on \( \text{Li} \) by Aberch and Lorents, and for the scattering of light atomic projectiles from a gaseous mercury target at small scattering angles by van Wijngaarden et al. There is also excellent theoretical agreement between Thomas Fermi cross sections for \( \text{H} \) on \( \text{Au} \) and \( \text{C} \) on \( \text{Au} \), and the corresponding electron gas model cross sections. The latter have been obtained by W. E. Baylis from a relativistic Hartree-Fock calculation.

II. Reduced Thomas Fermi cross section

Lindhard et al. found that after introducing the dimensionless parameter for energy

\[ \varepsilon = \frac{m_1 m_2 \nu^2}{2(m_1 + m_2)(Z_1 Z_2 e^2/a)} \]  

(c.g.s.), the differential cross section can be expressed in terms of a single parameter,

\[ r = \varepsilon^2 \sin^2 \frac{\theta}{2} \]  

as

\[ \frac{d\sigma}{d\Omega} = \frac{a^2}{8} \varepsilon^2 \frac{f(r^{1/2})}{2^{1/2}}. \]  

Here, \( f(r^{1/2}) \) is a universal function, which has been tabulated, and the scattering cross section therefore covers all combinations of \( m_1, m_2, Z_1, Z_2, E \) and \( \theta \). Equation 2.32 is exact for Rutherford scattering when

\[ f(r^{1/2}) = \frac{1}{2r^{1/2}}. \]
but is only approximately correct, to a precision of a few percent, for T.F. cross sections. In terms of the reduced parameters, the ratio between the Rutherford and Thomas Fermi cross sections is shown in Fig. 4. Experimental data have frequently been expressed using reduced parameters. For example, the triangles in the figure show the experimental results of Andersen et al. for H, He and Li projectiles in passing through thin gold films for scattering angles of 3 – 15°.

II.8 Conversion to laboratory coordinates from C.M.

In laboratory coordinates, the target particle is not fixed, but moves from recoil. The relation between the laboratory angle \( \theta_L \) and the centre of mass angle \( \theta \) is

\[
\tan \theta_L = \frac{\sin \theta}{\cos \theta + \frac{m_1}{m_2}}
\]

with \( m_1 \) being the projectile mass. The number of particles scattered into a given element of solid angle does not depend on the choice of coordinates. Hence,

\[
2\pi I \sigma(\theta) \sin \theta d\theta = 2\pi I \sigma'(\theta_L) \sin \theta_L d\theta_L
\]

where \( I \) is the incident beam current. Then,

\[
\sigma'(\theta_L) = \sigma(\theta) \frac{\sin \theta}{\sin \theta_L} d\theta_L
\]

and employing 2.34, we get

\[
\frac{d\sigma}{d\eta/L} = \frac{d\sigma}{d\eta}_{CM} \left[ 1 + \left(\frac{m_1}{m_2}\right)^2 + \frac{m_1}{m_2} \cos \theta \right]^{3/2}
\]

II.9 Energy loss and multiple scattering

A particle travelling through matter loses energy by both electronic excitation and elastic nuclear collisions. Thus the total nuclear stopping
Fig. 4. Energy dependence in reduced units of the ratio between Rutherford and T.F. cross sections. The triangles refer to experimental points obtained by Andersen et al. from whose work, this figure has been reproduced.
The cross section is

\[ S = S_e + S_n \]  \hspace{1cm} (2.36)

The stopping power \( \frac{dE}{dx} \) is related to \( S \) by

\[ -\frac{dE}{dx} = NS \]  \hspace{1cm} (2.37)

Here, \( S \) represents the average energy loss per collision, averaged over all impact parameters, and \( N \) is the atomic number density of the stopping medium. The magnitudes of \( S_e \) and \( S_n \) are primarily dependent on the projectile velocity rather than its energy, as sketched in Fig. 5.

Fig. 5. The velocity dependence of the electronic stopping cross section \( S_e \) and the nuclear stopping cross section \( S_n \). \( v_0 = e^2/\hbar \).

This figure emphasizes that at low projectile velocities, the electronic stopping cross section, \( S_e \), is proportional to \( v \), reaches a maximum, and then falls off as \( v^{-2} \ln v^2 \). Electronic excitations by the projectile produces only negligible angular deflections of the latter. In a nuclear collision, however, the projectile nucleus interacts with the target nucleus. After the collision, the target nucleus, together with its orbital electrons recoil, and the projectile is deflected. The scattering probability has a \( \sin^{-4}(\theta/2) \) dependence for unscreened
nuclear interactions, and approximately $\sin^{-3}(\theta/2)$ dependence\textsuperscript{2} for screened Coulomb interactions.

The effect of multiple scattering on backscatter yields has been described by van Wijngaarden et al.\textsuperscript{3}, quoted here:

'As a consequence of multiple small-angle scattering, an ion beam in matter gradually diverges about the direction of the incident beam. In any infinitesimal time interval, $dt$ there is a certain probability that a projectile will be scattered backward in a violent collision. The backscattered particles travelling towards the surface will also suffer small deflections. The net effect of the smaller deflections is that there exists a range of individual scattering angles $\theta$ for which the particles can be detected by a detector placed at a fixed laboratory angle $\varphi$, with respect to the incident ion beam. Because the angular distribution of scattered particles is strongly peaked in the forward direction more projectiles are multiply scattered into the solid angle of the detector than out of it.'

Hence, for solid targets of finite thickness, multiple scattering results in increased yields for backscatter angles. The additional contribution to the yield generally decreases with decreasing film thickness $t$. For very thin films, it is known\textsuperscript{3} that the contribution vanishes as $t \rightarrow 0$.

The angular divergence resulting from multiple scattering increases with the nuclear stopping cross section, $S_n$. Thus the relative importance of multiple scattering at a given energy increases with increasing mass and for a given mass increases with decreasing energies (Fig. 5).
CHAPTER III
EXPERIMENTAL

III.a Overall plan

Monoenergetic ion beams of keV energies are obtained from a magnetic analyser consisting of an accelerator tube and a π/2 magnetic sector field. (Fig. 5). The exit slit S₃ is placed at the magnetic focus of the entrance slit S₁. A movable beam limiting slit S₂, mounted near S₃, can be moved into or out of the beam path to reduce the current if necessary. The beam current is monitored by a movable Faraday cup, before entering the collimator, for alignment purposes. A beam monitoring device, placed after slit S₄, allows a relative measurement of the beam current during a backscatter measurement. The absolute incident beam current is measured by a movable Faraday cup following the beam monitor.

An electrostatic analyser is used for accurate measurements of the incident ion energies by obtaining Rutherford backscatter spectra. The collimated beam impinges onto the target perpendicularly, and backscattered particles at a laboratory angle of 159.95° pass through slit S₅ of known area. The backscattered particles are counted by a channeltron.

III.b Source and accelerator tube

Fig. 7 shows a radio frequency gas source. A mechanical leak controls the gas flow into the source bottle, and the gas pressure is adjusted for an optimum current density and stability. The gas is ionised by a RF field oscillating at about 50 MHz. A static axial magnetic field of about 20 Gauss, produced by a solenoid surrounding the source bottle,
confines the resulting plasma axially. The field also causes the plasma electrons to spiral about the axis of the magnetic field, thereby ionising more neutral atoms. A large density of ions is directed towards the exit channel by the positive extractor electrode. The glass shield shown in the figure prevents negative ions in the discharge from travelling backward to the extractor.

A quartz sleeve placed over the metal surface of the base prevents surface recombination of ions, and also provides a focussing field for the plasma exiting from the aluminium channel into the gap lens. The gap lens and the first few electrodes of the accelerator are supplied with a variable, negative high voltage potential of 0 - 10 kV. This produces focussing of the plasma prior to entering the accelerator tube.

The accelerator tube consists of 18 electrodes, equispaced in an evacuated tube. The electrodes 5 - 18 are connected to a resistor chain, so that adjacent electrodes have the same potential difference. The field lines between the electrodes are such as to maintain the beam on-axis. The source and accelerator region is evacuated by a mercury diffusion pump rather than an oil diffusion pump, to avoid insulating layers of decomposed oil vapours from settling on the electrodes.

III.c Magnetic analyser

The purpose of the magnetic analyser is twofold. Firstly, it acts as a mass filter. This is necessary since the ion beam from the accelerator tube contains a variety of masses. Secondly, it strongly narrows the momentum interval amongst ions entering the collimator, and the energy of the projectile beam becomes well defined.

An electromagnet provides a variable magnetic field, with a maximum
in excess of 4 kGauss. A Hall probe, model HTB-060B is placed at the
centre of the poles to accurately measure the magnetic field for energy
determination. The monoenergetic beam enters the collimator through a slit
$S_3$ of diameter 2.8 mm.

III. d Collimator and beam monitoring device

The beam leaving the collimator (Fig. 8) is limited to an angular
divergence of 0.6°. To monitor the current at $S_4$, a very small portion of
the beam is intercepted by a grid of fine, stainless steel wires. The
resulting secondary electron current from the wires is measured on a
collector plate, biased at +45V, by a Keithley 416 picoammeter, whose
output is fed to a chart recorder.

The Faraday cup used for the absolute current measurements is made
long enough to prevent multiple scattering of ions to the outside. A
repeller plate, kept at -90V stops secondary electrons from leaving the
cup. The Faraday cup current is measured with a Cary 401 vibrating reed
electrometer with a precision of 0.5 %. The electrometer was calibrated
with the aid of a Keithley 225 current source, a variable resistor and a
Fluke 8350A digital voltmeter. The calibration was done before and after
the entire experiment, to check on the stability of the calibration.

Fluctuations in the beam intensity, produced at the source, were
such that a greater error arose from estimating the average current than
the precision available for measuring it with the Faraday cup. During a
backscatter measurement, the Faraday cup is removed out of the beam path.
To take the current fluctuations into account, the level of the beam
current is continuously recorded from the beam monitor output. From the
recorded trace, an average current is obtained, to a precision of about
Fig. 8. Beam monitor and Faraday cup.
1 percent.

III. Target

The target sample consists of a smoothly polished beryllium backing plate (25mm x 75mm), coated with two gold films. (Fig 9).

Fig. 9. Target sample showing evaporated gold deposits on a beryllium substrate.

The thick gold film is used for energy measurements with the e.s.a. (section III.f) and the thin gold film for the backscatter measurements. The uncoated areas on either side of thin gold film are required for background measurements (section IV).

Since the beryllium plate has small variations in the concentration of high Z impurity atoms across the surface, we define the Be scattering yield to be the average of that from the two beryllium regions on either side of the thin gold film.

The beryllium plate is mounted on an externally adjustable holder, which brings any section of the plate into the beam path. The target region is evacuated by an ion pump to pressures of $10^{-8}$ Torr.

The entrance slit $S_5$ (Fig. 6) of the detector has a diameter of $5.66 \pm 0.03$ mm, and subtends a solid angle to the target of

$$\Delta \Omega = (2.21 \pm 0.02) \times 10^{-4} \text{ sr}.$$
centred about the laboratory scattering angle

\[ \theta_L = 159.93 \pm 0.06 \text{ degrees.} \]

III.f Electrostatic analyser and energy measurement

Energy measurements are facilitated by simply finding the magnetic field in the magnetic analyser with the Hall probe, connected to a Bell model 640 incremental Gaussmeter. An external digital voltmeter reads the output from the Gaussmeter, for accurate values. The magnetic analyser is calibrated with the aid of an electrostatic analyser (e.s.a.).

The electrostatic analyser (Fig. 10) consists of a \( \pi/2 \) cylindrical sector field, with a central radius of \( R_e = 19.05 \text{ cm} \). The field is produced between two electrodes, placed at a distance \( d = (0.3748 \pm 0.0025) \text{ cm} \) apart. The rectangular entrance and exit slits are both 0.64 mm wide. The electrostatic analyser views particles that are scattered from the target, through a laboratory angle of 154°.

Fig. 10. Electrostatic analyser for observing backscatter energy spectra for a laboratory angle \( \theta_L \) of 154°.
The potential difference across the plates is measured to a precision of 0.1% or better using a potential divider and a Keithly 190 digital voltmeter. The particles leaving the e.s.a. are detected by a Bendix windowless photomultiplier.

**Energy measurement**

A particle of energy \( eV_a = \frac{1}{2}mv^2 \) passing the entrance slit \( S_6 \) and exit slit \( S_7 \) of the electrostatic analyser will obey

\[
\frac{mv^2}{R} = eE
\]

where \( E \) is the electric field between the e.s.a. plates. Since their distance of separation is small compared with the central radius \( R_e \), the electric field for a potential difference \( V_e \) is given nearly by

\[
E = \frac{V_e}{d}
\]

We therefore have

\[
V_a = \left( \frac{R}{2d} \right) V_e \text{ in eV}
\]

\[
= (25.41 \pm 0.17) V_e
\]

The error of 0.7% in the absolute value of the energy primarily arises from the uncertainty in the \( d \) value. Because of the finite widths of the entrance and exit slits \( S_6 \) and \( S_7 \), particles with a given energy can pass through the e.s.a. over a small range of potential difference \( \delta V_e \). The relative range of \( \delta V_e \), or the energy resolution

\[
\frac{\delta V_e}{V_e} = \left( \frac{1}{S_6} + \frac{1}{S_7} \right) / R_e \approx \frac{1}{150}
\]

A part of a typical Rutherford backscatter spectrum used for energy measurements is shown in Fig. 11, where a beam of about 95 keV \( \text{He}^+ \) on a thick gold target film was used. The yield in energies above 89.0 keV arises from multiple scattering effects between the plates of the e.s.a. This is caused by the small separation of the plates, which is required for high electric fields. The step in the backscatter yield
Fig. 11. Higher energy portion of backscatter spectrum for the scattering of 95 keV He$^+$ from gold. Points are experimental, and the solid line has been drawn to determine the maximum backscatter energy, $E_{st} = 88.61$ keV. The error bar is the error in counting statistics.
is not sharp, but has a relative width of approximately

\[(88.8 - 88.4)/88.6 = \frac{1}{200}\]

Any significant spread in the energy of the incident beam would be manifested by an increase in the step width, on top of the finite resolution of the e.s.a. However, the observed relative width compares favourably with the energy resolution of the e.s.a. Hence the energy spread \(\delta V / V_a\) in the incident beam from the magnetic analyser is much smaller than the relative width of the step. The energy \(E_{st} (= 88.61\, \text{keV})\) at half-height of the step in the RBS is the energy of the particles scattered back from the surface layer of the target. It is lower than the incident energy \(E_0\) as a consequence of recoil of the target particle, and is given by

\[E_{st} = E_0 \left(1 - \frac{4m_1 m_2}{(m_1 + m_2)^2} \sin^2 \frac{\theta}{2}\right)\]

Particles scattered from deeper layers appear at lower energies in the spectrum due to energy loss in the target. Using Eq. 2.34 to convert \(\theta = 154^\circ\) from lab to CM coordinates, and putting \(E_{st} = 88.61\, \text{keV}\) yields the incident energy of the beam \(E_0 = 95.7\, \text{keV}\).

For a given monoenergetic beam, the magnetic field of the magnetic analyser should be proportional to the square root of the beam energy. This was verified for the present apparatus to a high degree of precision. This allowed us to determine the beam energy solely by a measurement of the magnetic field, with a relative error of 0.1% and an absolute error of 0.7%. The magnetic analyser was calibrated periodically to follow drifts in the instruments and to allow for changes in the alignment of the apparatus.
III.g Estimate of target thickness

The thickness of a particular gold film was measured with the aid of the RBS shown in Fig. 12. A 90 keV H\textsuperscript{+} beam was used. The sharp rise at high energies corresponds to scattering from the vacuum exposed gold surface. The sharp fall off at a lower energy corresponds to scattering from the gold surface in contact with the beryllium substrate. Here, the lower energy particles have travelled a distance essentially twice the film thickness. The particles cover a maximum distance \( t_0 \) into the target, and a distance \( t_0 / \cos \theta = 1.1 t_0 \) out of the target before entering the e.s.a. The total traversed gold thickness is 2.1\( t_0 \). The energy loss \( \Delta E \), given by the width of the curve at half-height is

\[
\Delta E = (dE/dx) \times 2.1 t_0
\]

where \( dE/dx \) is the stopping power for 85 keV H\textsuperscript{+} in gold. For this energy, \( dE/dx \approx 19.5 \text{ eV Å}^{-1} \), and the width \( \Delta E \) in Fig. 12 is 6.1 keV, which yields \( t_0 \approx 150 \text{ Å} \).

The backscatter yield of H from this film was also measured with the channeltron detector. Since the yield is proportional to the film thickness, at least for thin films, other film thicknesses were found by measuring the corresponding yields \( Y(t') \) for H at the same energy \( E_0 \). Then

\[
t' = \frac{Y(t')}{Y(t_0)} \times t_0
\]

This equation is also valid for films thin enough to be discontinuous, consisting of small 'islands'. But then, \( t' \) must be interpreted as the average material thickness of the target film.

III.h Note on film thickness variations

It is known that under vacuum evaporation onto a substrate, the atoms first settle at nucleation sites\textsuperscript{10}. The sites then continue to...
Fig. 12. Backscatter energy spectrum of $\text{H}^+$ on gold. The width of the curve at half-height, $\Delta E$ is proportional to film thickness.
grow in size, until finally, a continuous film is obtained. The question arises whether or not discontinuous films of very small average thickness, such as those used in the present experiment, possess many 'islands' of thicknesses much larger than the average material thickness. If they do, then there is the possibility that multiple scattering effects for our thinnest films are not small.

To investigate this, one of our target samples with a gold film of the order of 5 Å material thickness was studied with an electron microscope. Unfortunately, no difference in the micrographs of the 'pure' beryllium substrate and the Au evaporated area could be discerned.

The shape of the lower portion of an RBS however is very sensitive to the presence of thick islands in a discontinuous film. Since the backscatter count rate for the e.s.a. was too small to obtain a good RBS with H or He from a very thin target, the spectrum was obtained with N\textsuperscript{+} ions at a primary energy E\textsubscript{o} = 100 keV. (Fig. 13). The spectrum has a nearly continuous background of about 0.1 x 10\textsuperscript{3} counts. It will be noted that at lower energies, the count rate falls off rather sharply with decreasing energy. This occurs in spite of the fact that the straggling in the energy loss for heavier projectiles is large (Section II.h). The above implies that the individual island thicknesses do not deviate much from the thickness determined by the ΔE value, the width of the spectrum at half-height. The thickness \( \bar{t} \) is

\[ \bar{t} = \Delta E / 2.11 \frac{dE}{dx} \]

Here (see Fig. 13) \( \Delta E = 2 \) keV. An experimental \( \frac{dE}{dx} \) value for N in Au could not be found, but it was inferred as follows. The ratio of the experimental stopping powers of 400 keV N\textsuperscript{14} and He\textsuperscript{4} in Ar\textsuperscript{40}, tabulated by Whaling\textsuperscript{12} is 1.7:1, in agreement with the Lindhard theory\textsuperscript{13} for the
Fig. 13. Backscatter energy spectrum for 100 keV N\(^+\) from a very thin gold film. The error bars show error in counting statistics.
electronic stopping power. The experimental stopping power for 75 keV \( \text{He} \) in \( \text{Au} \) is \( 31.9 \text{ ev } \AA^{-1} \). Using Lindhard's theory now to calculate the ratio of \( \text{N} \) to \( \text{He} \) stopping powers in \( \text{Au} \), the value for the stopping power of \( \text{N} \) was found to be \( 64 \text{ ev } \AA^{-1} \). This value does not include nuclear stopping of the order of \( (\text{dE/dx})_n \approx 15 \text{ ev } \AA^{-1} \), and is therefore a lower limit. Substitution into Eq. 3.8 gives \( t \leq 15 \AA \). Hence, our evaporated gold films of thicknesses as low as 15 \( \AA \) are still semi-continuous: being nearly free of island sizes that are significantly larger than \( t \).

III.1 Channeltron detector and detecting efficiency

The channeltron used was Ceratron-E model EMS-6061B, operating at a potential difference between the collector (positive) and the input (grounded) of 3.5-4 kV and was connected to standard counting equipment. This channeltron was chosen for its very low UV photon detection efficiency. Both the amplifier gain setting and the operating voltage of the channeltron were adjusted for stable detection efficiency.

To find absolute backscatter yields, the counting efficiency of a detector needs to be known. To measure the efficiency, one can in principle allow an ion beam of known current to enter the detector, and compare the count rate with the incident particle current. This is not practical however for the present counting equipment for the following reasons.

The maximum allowable counting rate for the channeltron and counting electronics is of the order of \( 10^4 \) counts sec\(^{-1}\). This corresponds to a current for a singly charged particle beam of \( 1.6 \times 10^{-15} \) Ampere. Such a small current could not be measured accurately with our instruments. Since the ion beam had short term stability, we could only measure
currents larger than $10^{-13}$ Ampere with a precision of a few percent. This corresponds to two orders of magnitude above the maximum allowable counting rate for the detector.

To limit the counting rate of the channeltron, only a small, but known fraction of the ion beam was allowed to enter it. This was accomplished by the set-up shown in Fig. 14.

![Diagram](image_url)

**Fig. 14.** Apparatus for channeltron efficiency measurement. The diameters of $S_a$ and $S_b$ are $(1.727 \pm 0.003)$ mm and $(0.085 \pm 0.003)$ mm respectively.

Here the beam enters through the slit $S_a$ of diameter $(1.727 \pm 0.003)$ mm and the current is measured with the Faraday cup. Since the current density across the slit $S_a$ must be uniform, this size could not be increased.

The end of the Faraday cup was designed to hold a thin metal disc with a very fine hole, $S_b$ of diameter $(0.085 \pm 0.003)$ mm. The fraction

$$\frac{\pi \times \left(\frac{0.085}{2}\right)^2}{\pi \times \left(\frac{1.727}{2}\right)^2} = (0.00242 \pm 0.00017)$$

of the measured current then 'leaks' into the channeltron. The thin film served two purposes: Firstly, it stops low energy particles created by the interaction of the fast ion beam with the residual gas in the target chamber, and slits scattered particles heading towards the detector via
wall scattering, etc. Secondly, the film serves as a secondary electron emitter. Upon emerging from the film, these electrons are detected, thus increasing the detection efficiency for the incident particle.

Fig. 15A shows the absolute efficiency in detecting H, He and N at various energies. The detection efficiency for all three species is the same, and within experimental error it is 100%. The scatter in the points arises from difficulties in measuring small, unstable currents, and from the random variations in the beam profile across the slit $S_a$. The fall-off at lower energies arises from projectiles that come to a rest in the film. Since the stopping cross section for N is larger than that for either H or He, the 100% detection level for N is only reached at about 35 keV. Fig. 15B shows the detection efficiency for He and N with a thinner aluminium film. The fall-off now appears at a lower energy, as expected. The combination of this thinner Al film and channeltron served as the detector in the backscatter experiment. (Fig. 6). The detector has a nominal efficiency of 100% for H at energies above 5 keV, for He at energies above 8 keV, and for N at energies above 35 keV.
Fig. 15. The energy dependence of the detection efficiency for various ions.
CHAPTER IV
RESULTS AND CONCLUSIONS

The backscatter yield for a given projectile was measured for several energies. We define the backscatter yield $Y$ as the counting rate $N$ for the solid angle of our channeltron detector per incident projectile. Hence,

$$ Y = \frac{N e}{I\text{(Amp)}} $$

where $I$ is the beam current for singly charged ions. The total observed backscatter yield $Y_o$ from a thin Au film consists of the scattering from the Au film and from the Be substrate underneath:

$$ Y_o = Y_{Au} + Y_{Be} \quad 4.2 $$

The prime emphasises that the Be scattered projectile first pass the Au film before re-entering the vacuum.

The yield $Y_{Be}$ from 'pure' Be was found separately by measuring the backscattering from the Be substrate, not coated with Au. (Section III.e). For thin Au films ($t \leq 40 \text{ Å}$), the $Y_{Be}$ is nearly the same as $Y_{Be}'$. The reason for this is that the majority of Be scattered projectiles of energies above a few keV do not come to a rest inside the Au film. Their energy loss $\Delta E$ in traversing the Au film is $\Delta E \ll 1 \text{ keV}$. Hence, Eq. 4.2 may be re-written as

$$ Y_{Au} = Y_o - Y_{Be} \quad 4.3 $$

It is clear that accurate backscatter yields can only be obtained for very thin films. However, an arbitrarily small thickness cannot be used. We observed that the backscatter yield for beryllium was not uniform across the uncoated areas, and it also varied slightly from projectile to projectile, with a maximum variation of 1-2% for the
heaviest projectile. Although the averaging process for the $Y_{Be}$ measurement (section III.e) minimises errors due to this variation, for our thinnest films, where $Y_{Be} > Y_{Au}$ (see tables I and II), the error in the $Y_{Au}$ measurement could not be kept below 5%. Other errors contributing to the total error in the backscatter yield are much smaller.

The backscatter yield versus projectile energy from a 40 Å gold target is shown in Fig. 16. The sudden fall-off of the backscatter yield for C at energies below 25 keV is due to the lowering in the detection efficiency of the channeltron. The slopes of the H and He curves for energies above 20 keV are both 2.0, corresponding to a $E^{-2.0}$ variation of the yield for a large energy range. This is better seen in a plot of the yield $x$ (energy)$^2$ versus energy shown in Fig. 17.

This figure also shows results for a 13 Å film. For the purpose of comparison, the $Y_{Au} x E^2$ values for both films have been normalised to the Thomas Fermi (T.F.) cross sections (solid curves), for H at about 80 keV. Since Rutherford cross sections vary as $E^{-2}$ with energy, these appear as straight lines on the graph (shown dashed). The difference between the theoretical T.F. cross section for H at 80 keV and the corresponding Rutherford cross section is only about 15%, and decreases at higher energies. The 40 Å film gives normalised $Y_{Au} x E^2$ values substantially higher than those for the 13 Å film at energies $\leq$ 60 keV for H and at all energies for He and C. This effect is attributed to multiple scattering in the target film, described in section II.h, and has previously been observed$^3,8$. Van Wijngaarden et al. have used light atomic projectiles on thin gold targets, the minimum thickness being 75 Å.

Since we observe very large multiple scattering effects even for the 40 Å film, they may not be negligible for the 13 Å film. It was therefore
### Table I

Some typical backscatter yield measurements for H⁺, He⁺, C⁺, and N⁺. 
Yₒ is the yield from the 13 Å Au film and the Be substrate. 
Yₘₐₖ is the substrate yield, and Yₐₚ = Yₒ - Yₘₐₖ.

<table>
<thead>
<tr>
<th>Energy(keV)</th>
<th>Projectile</th>
<th>Yₒ(total)</th>
<th>Yₘₐₖ</th>
<th>Yₐₚ</th>
</tr>
</thead>
<tbody>
<tr>
<td>37.85</td>
<td>H⁺</td>
<td>1.77 x 10⁻⁸</td>
<td>1.03 x 10⁻⁸</td>
<td>0.742 x 10⁻⁸</td>
</tr>
<tr>
<td>92.05</td>
<td>H⁺</td>
<td>7.02 x 10⁻⁹</td>
<td>5.60 x 10⁻⁹</td>
<td>1.42 x 10⁻⁹</td>
</tr>
<tr>
<td>35.79</td>
<td>He⁺</td>
<td>3.40 x 10⁻⁸</td>
<td>0.696 x 10⁻⁸</td>
<td>2.70 x 10⁻⁸</td>
</tr>
<tr>
<td>40.85</td>
<td>C⁺</td>
<td>12.4 x 10⁻⁹</td>
<td>0.630 x 10⁻⁸</td>
<td>11.7 x 10⁻⁸</td>
</tr>
<tr>
<td>40.99</td>
<td>N⁺</td>
<td>15.1 x 10⁻⁹</td>
<td>0.620 x 10⁻⁸</td>
<td>14.5 x 10⁻⁸</td>
</tr>
</tbody>
</table>

### Table II

Some typical backscatter yield measurements for H⁺, He⁺, C⁺, and N⁺. 
Yₒ is the yield from the 6 Å Au film and the Be substrate. 
Yₘₐₖ is the substrate yield, and Yₐₚ = Yₒ - Yₘₐₖ.

<table>
<thead>
<tr>
<th>Energy(keV)</th>
<th>Projectile</th>
<th>Yₒ(total)</th>
<th>Yₘₐₖ</th>
<th>Yₐₚ</th>
</tr>
</thead>
<tbody>
<tr>
<td>40.79</td>
<td>H⁺</td>
<td>13.2 x 10⁻⁹</td>
<td>10.4 x 10⁻⁹</td>
<td>2.79 x 10⁻⁹</td>
</tr>
<tr>
<td>40.98</td>
<td>He⁺</td>
<td>18.2 x 10⁻⁹</td>
<td>8.52 x 10⁻⁹</td>
<td>9.69 x 10⁻⁹</td>
</tr>
<tr>
<td>40.96</td>
<td>C⁺</td>
<td>6.94 x 10⁻⁹</td>
<td>0.959 x 10⁻⁸</td>
<td>5.98 x 10⁻⁸</td>
</tr>
<tr>
<td>41.09</td>
<td>N⁺</td>
<td>7.93 x 10⁻⁸</td>
<td>0.776 x 10⁻⁸</td>
<td>7.16 x 10⁻⁸</td>
</tr>
</tbody>
</table>
Fig. 16. The energy dependence of the backscatter yield for H, He and C from a 40 Å gold film.
Fig. 17. Backscatter yield $x$ (energy)$^2$ for 40 Å gold film (closed circles) and 13 Å gold film (open circles). Experimental points are normalised to the T.F. cross section shown as a solid line. The Rutherford cross sections are shown for comparison (dashed lines).
necessary to study thinner targets in order to obtain backscattering yields free from multiple scattering.

Fig. 18 shows results for three targets, of nominal thicknesses 13 Å, 6 Å, and 2 Å. All \( Y_{\text{Au}} \times E^2 \) values for the 13 Å film have been normalised to the T.F. curve for H at 40 keV. Similarly, the results for the 6 Å film have been normalised to the T.F. curve for H at 30 keV. The backscatter yield of H from the 2 Å was difficult to measure due to the relatively high background from the beryllium substrate. The \( Y_{\text{Au}} \times E^2 \) values for the 2 Å film have been normalised instead to the experimental \( Y_{\text{Au}} \times E^2 \) value of 50 keV He on the 13 Å film.

The normalisation of our experimental values to the Thomas Fermi cross sections at \( E = 40 \text{ keV} \), corresponding to \( \frac{1}{r_0} = 3.65 \) (section II.f) is valid because it is known that for \( \frac{1}{r_0} \) values in the range \( r \geq 3 \), the T.F. cross sections are accurate\(^3\).

That the scatter in the experimental yields for the C and N projectiles for the various films is larger than that for H or He arises from the difficulty in measuring the background. It was found that the variation in the beryllium backscatter yield across the surface of the target sample was greater for C and N than for H or He. This is not surprising because the range in Be of the heavier projectiles compared to the lighter projectiles is much shorter, and therefore, the backscatter yield of heavier ions is more sensitive (than it is for lighter ions) to high Z surface impurities in the target sample.

We note that in passing from the 13 Å to the 2 Å films, the shape of the curves for a given projectile is independent of film thickness. We also note that the relative magnitudes of the normalised \( Y_{\text{Au}} \times E^2 \) values for the different projectiles is independent of the film
Fig. 18. The energy dependence of the backscatter cross sections of H, He, C and N from gold for a laboratory angle $\theta_s = 159.95^\circ$. The open circles, crosses and triangles are measurements for a 13 Å, 6 Å and 2 Å thick gold film respectively (see text). The solid lines are T.F. cross sections and the dashed lines are Rutherford cross sections.
thickness, within experimental error (± 5% for C and N). We interpret this to mean that the scattering yields for these thinner films are nearly free from multiple scattering effects. This, together with the valid normalisation procedure to T.F. theory for H at higher energies elevates our yield measurements to measurements of cross sections.

The deviation from T.F. theory of the observed H values at energies lower than the normalisation point is minimal. However, all observed cross sections for He are greater than the T.F. cross sections by about 5% on average. Experimental cross sections for C and N lie above the T.F. predictions by about 25%. For these projectiles, the values lie in the range $0.3 \leq t \leq 1.6$, where Andersen et al. also find similar deviations from the T.F. theory (see Fig. 4).

**IV.a CONCLUSIONS**

By a justified process of normalisation of experimental backscatter yields of H to calculated Thomas Fermi cross sections, the backscatter cross sections for H, He, C and N from Au at a scattering angle of 160° have been obtained. The observed cross sections for He are 5% larger than the T.F. ones, while the cross sections for C and N for all energies investigated are higher by 20-30%. This is consistent with the results of Andersen et al. for $t \leq 1.0$, where screening is large.

The Thomas Fermi theory appears to overestimate electronic screening for atoms in a solid. In fact, the cross sections are quite sensitive to the screening length. For example, if we employ the smaller screening parameter of Eq. 2.29, instead of the one in Eq. 2.28, the theoretical cross sections for N relative to those for H reduce by about 7%. The reason for the discrepancy between theory and the results of this experiment are not understood.
REFERENCES


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