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Heavy-ion Elastic Collisions and the Nuclear Surface Diffuseness

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Abstract

Taking the magnitude of the elements of the S matrix and the nuclear phases to be of Woods–Saxon form in the variable l with width parameter Δ, we find that the form of the scattering amplitude components $f^+(\theta)$ and $f^-(\theta)$ is specified by the parameters Δ^+ and Δ^- such that Δ is closely the geometric mean of Δ^+ and Δ^-. Many angular distributions have been analysed into $f^+(\theta)$ and $f^-(\theta)$ to obtain Δ^+ and Δ^- and hence Δ, the angular momentum diffuseness, from which the nuclear surface diffuseness is obtained, so reducing an ambiguity in the nuclear potential. The case of light ions incident on heavy ions has also been investigated.

1. Introduction

Most calculations of elastic scattering of heavy ions have employed the optical model, but this approach gives little insight into the collision problem and is also subject to ambiguity in the nuclear parameters so derived. More general approaches have involved approximate diffraction formulae. The most fruitful method has been to take the magnitude of the scattering matrix $S(l)$ to be of Woods–Saxon form, to replace the term $P_l(\cos\theta)$ in the partial wave series for the scattering amplitude $f(\theta)$ by its asymptotic form containing the expression $\sin\{(l+\frac{1}{2})\theta+\frac{1}{4}\pi\}$ and to split the sine into two exponentials. Thus $f(\theta)$ is broken up into additive components $f^+(\theta)$ and $f^-(\theta)$, and these vary uniformly in magnitude as θ increases, while being progressively in and out of phase through the factors $\exp(\pm\mathrm{i}l_c\theta)$, where $l_c \approx k(R_1+R_2)$ is the critical value of l for which $S(l) = \frac{1}{2}$. Thus $f(\theta)$ has maxima and minima with angular separation π/l_c, the oscillations being large when f^+ and f^- are nearly equal in magnitude, and small or negligible when they are very unequal.

In the strong absorption model, Frahn and Venter (1963) neglect the nuclear phase component of S_l to obtain a simple expression for $f^\pm$ by taking the value for a sharp cutoff below $l = l_c$ and multiplying it by a 'form factor'. The approximations involved have been briefly summarized by Bassichis and Dar (1966). However, the result is best derived by taking the Fourier transform of $S(l)$ with respect to the variable $(l-l_c)$ (Friedman *et al.* 1974) to give $f^\pm(\theta)$ in terms of the conjugate variable $(\theta\pm\theta_c)$, where θ_c is the quarter-point angle given by $2\arctan(n/l_c)$, with $n = Z_1 Z_2 e^2/\hbar v$. Then for $\theta > \theta_c$ we have

$$\sin^{\frac{1}{2}}\theta\, f^\pm(\theta) \propto \pi\Delta/\sinh\{\pi\Delta(\theta\pm\theta_c)\}\,. \tag{1}$$

For values of θ not too close to θ_c we have approximately (2)

$$\sin^{\frac{1}{2}}\theta\, f^{\pm}(\theta) \propto \exp\{-\pi\Delta(\theta \pm \theta_c)\},$$

and so

$$f^{+}(\theta)/f^{-}(\theta) \approx \exp(-2\pi\Delta\theta_c). \tag{3}$$

Hence for θ_c large, as is the case for energies near the Coulomb barrier, we have $f^{+} \ll f^{-}$, so there are no oscillations in the angular distribution and the curve of $\log\{\sin\theta\,\sigma(\theta)\}$ versus θ falls almost linearly past θ_c, as is found experimentally. Using the principle of stationary phase, Frahn (1972) has shown that, for $\theta = \theta_c$, the component f^{-} is approximately equal to $\frac{1}{2}f_C$, where f_C is the Coulomb value of f, provided θ_c is not too small. Similarly, if f^{+} is extended back to negative angles, we expect to have $f^{+} \approx \frac{1}{2}f_C$ at $\theta = -\theta_c$.

Now refraction takes place in the nucleus as well as absorption, and more recently attempts have been made to incorporate the nuclear phases δ_l into the theory, using the parameterization $\delta_l = \delta_0(1-S_l)$ and the Poisson summation formula. Hartmann (1976) has taken the $m = 0$ term in the formula to obtain expressions valid for δ_l small which are used for angles close to, and on both sides of, θ_c. Rowley and Marty (1976) have used saddle points in the complex angular momentum plane, and from the $m = 0$ terms obtain

$$\sin^{\frac{1}{2}}\theta\, f^{\pm}(\theta) \propto \exp\{-\pi\Delta^{\pm}(\theta \pm \theta_c)\}, \tag{4}$$

where Δ^{+} and Δ^{-} decrease and increase respectively from Δ by equal amounts as δ_0 increases from 0, so that

$$\Delta^{+} + \Delta^{-} = 2\Delta. \tag{5}$$

We shall investigate here the accuracy of the relations (4) and (5).

2. Analysis of Experimental Angular Distributions

As θ increases, f^{+} and f^{-} undergo equal and opposite phase changes $l_c\theta$, so that when they are in phase $f_{\max} = f^{+}+f^{-}$ and when out of phase $f_{\min} = |f^{+}-f^{-}|$. Hence by drawing smooth curves through the maxima and minima of the graph of $\sigma(\theta) = |f|^2$, it is possible to deduce the curves for f^{+} and f^{-}. The latter curves cross near the lowest minimum of $\sigma(\theta)$, the depth of this minimum depending on how nearly f^{+} and f^{-} are of opposite phase when they are equal in magnitude.

Fig. 1 shows the results of the analysis of angular distributions calculated by Goldberg and Smith (1974) from an optical model; these theoretical distributions avoid the imperfections due to experimental error and limited angular resolution, and in fact the presence of two distinct components in these angular distributions can be seen immediately without detailed calculation. The lines for f^{+} and f^{-} in Fig. 1 are nearly straight and pass nearly through the value $\sin^{\frac{1}{2}}\theta_c \frac{1}{2}f_C$ at $\theta = -\theta_c$ and $\theta = \theta_c$ respectively as expected. The lines A, B, C and D for the same energy (Fig. 1*a*) have slopes Δ^{+} and Δ^{-} that are nearly independent of the target nucleus; the lines E, F, G and H for the same pair of ions (Fig. 1*b*) show that $\Delta^{-} \to \Delta^{+}$ as $E \to E_B$, the barrier energy. A similar analysis of experimental angular distributions with their imperfections and often small angular ranges gives less satisfactory results; nevertheless approximate values for Δ^{+} and Δ^{-} have been obtained in this study for many pairs of heavy ions.

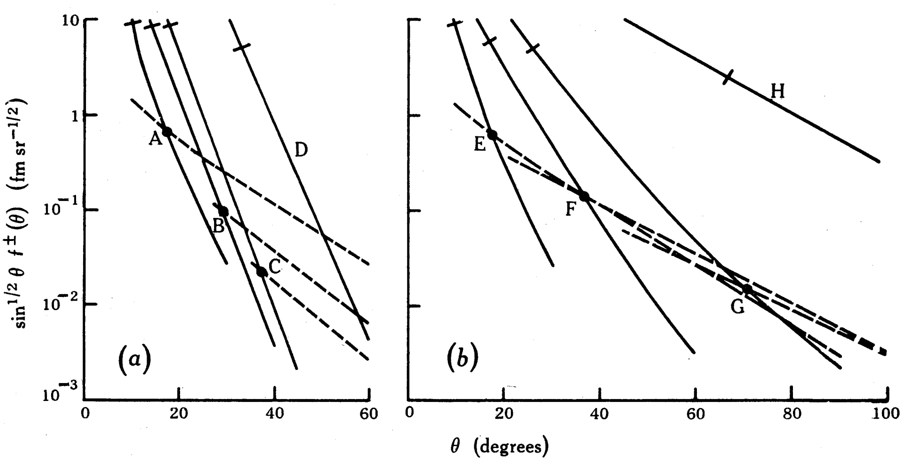

Fig. 1. Forms deduced for $\sin^{\frac{1}{2}}\theta\, f^{\pm}(\theta)$ as a function of θ, for (*a*) 200 MeV ^{16}O ions on ^{28}Si (A lines), ^{58}Ni (B), ^{90}Zr (C), and ^{208}Pb (D); and for (*b*) ^{16}O ions of energies 200 (E), 120 (F), 80 (G) and 40 MeV (H) on ^{28}Si. The full and dashed lines show the components f^- and f^+ respectively. The bar on each f^- line indicates the quarter-point angle θ_c, and each dot the crossing of lines for corresponding f^- and f^+.

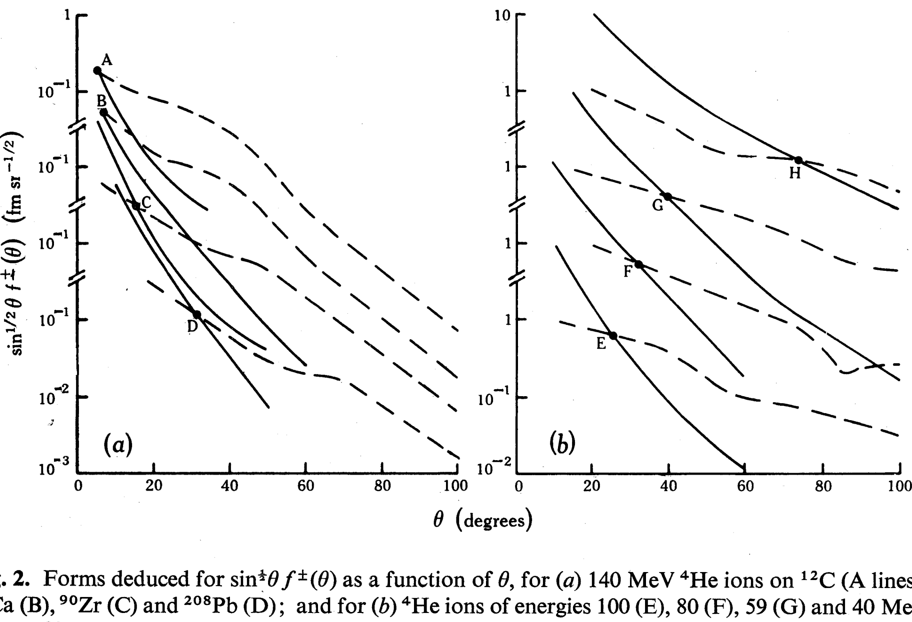

Fig. 2. Forms deduced for $\sin^{\frac{1}{2}}\theta\, f^{\pm}(\theta)$ as a function of θ, for (*a*) 140 MeV ^{4}He ions on ^{12}C (A lines), ^{40}Ca (B), ^{90}Zr (C) and ^{208}Pb (D); and for (*b*) ^{4}He ions of energies 100 (E), 80 (F), 59 (G) and 40 MeV (H) on ^{90}Zr. The full and dashed lines show the components f^- and f^+ respectively. Each dot indicates the crossing of lines for corresponding f^- and f^+.

Angular distributions for the scattering of fast α particles (Goldberg *et al.* 1974) and helions (Hyakutake *et al.* 1978) by heavy nuclei show the same feature of an intermediate angular range over which oscillations are more pronounced than at other angles, so it is of interest to carry out a separation of f^+ and f^- for incident light ions. Fig. 2 shows the results for collisions of ^{4}He, to be compared with Fig. 1. Here the lines for f^- show a more definite curvature, and in fact they finally curve upwards to meet the value of f^+ near 180° at the lower energies for which the oscillations in $\sigma(\theta)$ increase in amplitude towards 180°. The lines for f^+ are almost straight at the smaller and the larger angles, but show a bulge at intermediate angles, due no doubt to the behaviour of a few nuclear phases which do not conform to the parameterized form used in the theory, for l_c is smaller for light than for heavy ions incident at the same energy. Also l_c is too small for safe use of the approximation used to obtain the relation (1), whereby a summation over l is replaced by an integral.

The deviations of the curves for f^+ and f^- from linearity are greatest at the larger angles, and determinations of Δ^+ and Δ^- must be made from their slopes near their crossing point. There is some justification for this in that the lines for f^+ and f^- pass close to the value $\frac{1}{2}f_C$ at $\theta = -\theta_c$ and θ_c respectively. Although the determination of Δ^+ and Δ^- is thus subject to greater error for incident light ions than for heavy ions, the extension of the heavy-ion treatment to light ions is worth a trial.

3. Effect of Refraction by Graphical Analysis

Having checked the accuracy of the relation (4), we now check (5). It has so far proved difficult to do this by algebraic methods. As for numerical methods, the calculation of f is carried out using the Coulomb value f_C for the partial wave series, and correcting the lower order terms for the inclusion of the nuclear phases with the Coulomb phases, but f^+ and f^- cannot be calculated in this way because their Coulomb values are not known. It may be done, however, by using a regularity in the phase–amplitude diagram for the partial wave series (Mohr 1976).

Thus the terms in the series for f with $S(l) = 1$ for all l (pure Coulomb field) are represented by vectors in the complex plane, giving a diagram (Fig. 3) resembling the Cornu spiral, which first turns around O ($l = 0$) and finally around O′ ($l = \infty$). As θ increases, the spiral about O unwinds while the spiral about O′ winds up, and the point P at the end of the vector representing a particular term l moves from the former spiral to the latter, passing close to the midpoint of OO′ when $\theta = \theta_c$. For a sharp cutoff below $l = l_c$, f is given by O′P. We now apply the method to the diagrams for f^+ and f^- for $S(l) = 1$, and these are seen to take simpler forms, the exterior angle between the successive segments being $\theta^{\pm} = \theta \pm 2\rho_l$, where $\rho_l = \arctan(n/l)$ is the difference between the Coulomb phase shifts of order l and $l-1$. For f^+ there is only a single spiral; for f^- there is a double spiral, but our concern is with $\theta > \theta_c$ and hence with the spiral about O′.

For a sharp cutoff, $f^{\pm} = \mathrm{O'P}^{\pm} = \frac{1}{2}A/\sin\frac{1}{2}\theta^{\pm}$ from the geometry of Fig. 3, where A is the length of the segment of the spiral representing the term of order l in the partial wave series for $f^{\pm}$. This corresponds to the value given by Frahn and Venter (1963). For a smoothed cutoff, the segments of the spiral are shortened, as shown by the dashed line in the figure, and the value of f^- is now represented by O′T; similarly for f^+. As θ increases, $\theta^{\pm}$ increases and the spiral winds up more tightly, so that O′T becomes smaller, and this decreased value of $f^{\pm}$ is to be compared with the value given by the relation (1). Many calculations based on this graphical model were carried out and they gave results in fair agreement with (1) provided θ was not too large, for then the value of $f^{\pm}$ is small and therefore sensitive to small differences in $S(l)$ from the Woods–Saxon form.

Having dealt with the effect of absorption, we now introduce the additional effect of refraction. The exterior angle between successive segments of the spiral now becomes $\theta^{\pm} = \theta \pm 2(\rho_l - \delta_l)$, so that the spiral for f^+ winds up less tightly and f^+ is decreased by a smaller amount; conversely f^- is decreased by a larger amount.

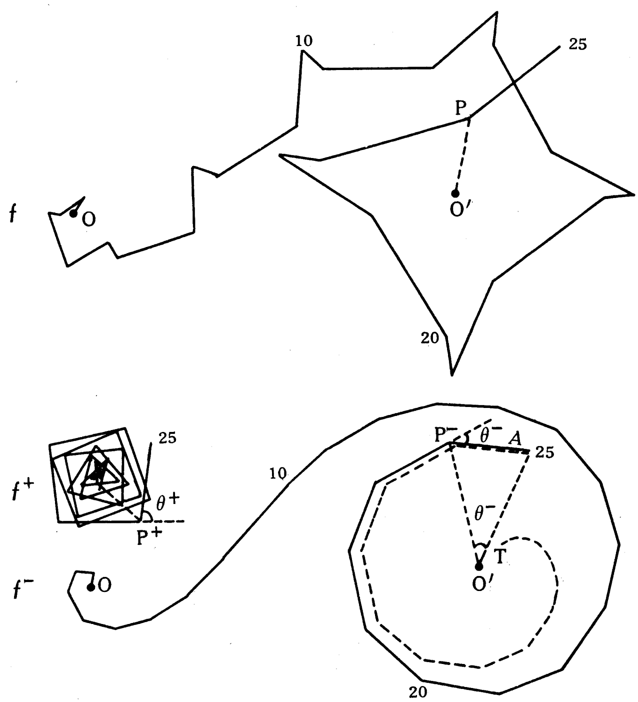

Fig. 3. Phase–amplitude diagrams for scattering by a Coulomb potential through an angle of 60° for $n = 5$. The successive terms in the partial wave series, starting from $l = 0$ at the point O, are shown for f, f^+ and f^-. For a sharp cutoff below $l_c = 24$, f is given by O′P, f^+ by O′P$^+$ and f^- by O′P$^-$. The effect of a smoothed cutoff of Woods–Saxon form with width parameter $\Delta = 2 \cdot 5$ is shown by the dashed line for f^-, which is then given by O′T.

Thus the effect of refraction is to decrease the slope of the line for f^+ from Δ to Δ^+ (say), and to increase that for f^- from Δ to Δ^-. The two changes in Δ are opposite in direction but not quite equal in magnitude: our detailed calculations show that Δ is nearer to the geometric mean of Δ^+ and Δ^- than to the arithmetic mean. The difference becomes significant for the largest values of Δ^-/Δ^+ which are close to 3. For $\theta \approx \theta_c$, f^- is given by the relation (1) with Δ^- in place of Δ, instead of by the relation (4), while f^+ is still given by (4).

4. Determination of Surface Diffuseness

We expect the nuclear surface diffuseness to be related to the angular momentum diffuseness, the latter being specified by the parameter Δ and obtained by analysing angular distributions as in Section 2 to obtain Δ^+ and Δ^- and hence Δ. Now the nuclear surface diffuseness may be specified in terms of the radial variation of either the nuclear density ρ or the nuclear potential V, both being approximately of Woods–Saxon form with width parameters d and a respectively. Because of the finite range of the force between nucleons we expect a to be a little greater than d, while d is related to Δ according to a semiclassical argument due to Frahn and Venter (1963), which we now give in slightly simplified form. For a sharp cutoff we have $l_c \approx kR$, where R is the combined radius R_1+R_2, and so for a smoothed cutoff the change Δ in l_c and a change d in R are related by $\Delta \approx kd$, where d is the combined diffuseness d_1+d_2.

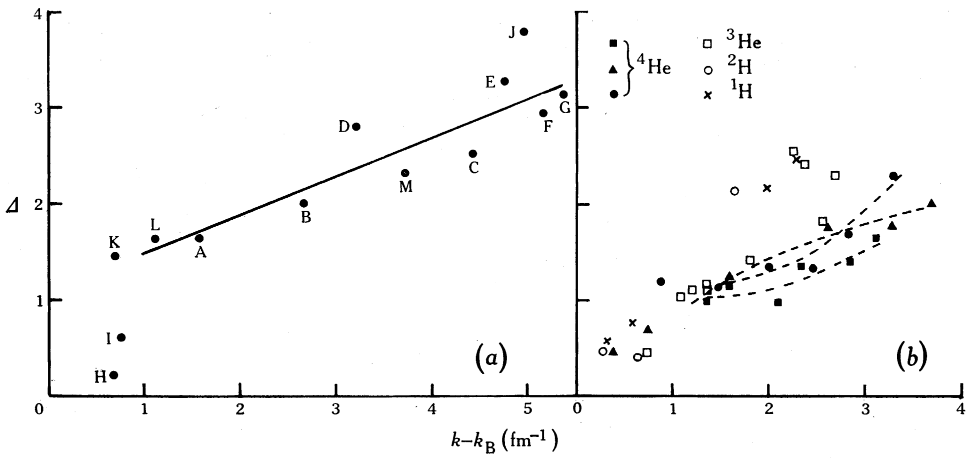

Fig. 4. Values of the width parameter Δ for various values of $k-k_B$ for the collision of (*a*) pairs of heavy ions and (*b*) light ions with heavy ions. Details of all the data plotted, with references, are given in Section 4.

The relation $\Delta \approx kd$, which is asymptotic, fails for energies near the Coulomb barrier, when there is a need for a detailed calculation of barrier penetration. This suggests that we consider a plot of Δ as a function of $k-k_B$, where k_B is the value of k at the barrier energy. Such a plot is shown in Fig. 4, where we see that Δ increases fairly uniformly with $k-k_B$, except near the origin where Δ has much smaller values due to the barrier. We therefore ignore the region near the origin, and take the slope of the mean line through the rest of the points as giving the value of d. The small displacements of the points from the mean line in Fig. 4*a* are due to a possible dependence of d on the pair of nuclei concerned, and also to inaccuracy in analysis of the experimental angular distributions, which often have too few oscillations to permit a fully satisfactory separation of f^+ and f^-. It should be noted that it is not possible to make use of angular distributions without oscillations, which is the case for the larger values of n/l_c, for then we have $f^+ \ll f^-$ and we can find only f^- and Δ^- and cannot determine Δ: the energy is too low for probing the nuclear surface. Actually there are few experimental studies which extend over a wide energy range for a particular pair of heavy ions.

Fig. 4*a* shows the values of Δ obtained for the collision of pairs of heavy ions. The points A, B and C are for 80, 120 and 200 MeV ^{16}O ions on ^{28}Si (Goldberg and Smith 1974), using the theoretical angular distributions referred to in Section 2 above, while D and E are for 141·5 and 215 MeV ^{16}O ions on ^{28}Si (Satchler *et al.* 1978), using experimental angular distributions; the latter two points are seen to lie a little above the former three points. The points F and G are for 200 MeV ^{16}O ions on ^{58}Ni and ^{90}Zr (Goldberg and Smith 1974), again using theoretical distributions. The rest of the points are from experimental distributions as follows: H, for 15 MeV ^{12}C on ^{14}C (Delic 1975); I for 52 MeV ^{18}O on ^{16}O (Reisdorf *et al.* 1975); J for 186 MeV ^{12}C on ^{28}Si (DeVries *et al.* 1977); K for 30 MeV ^{6}Li on ^{40}Ca (Bohn *et al.* 1977); L for 34 MeV ^{6}Li on ^{48}Ca (Cutler *et al.* 1977); M for 135 MeV ^{6}Li on ^{28}Si (DeVries *et al.* 1977). Apart from the three low energy points H, I and K nearest the Coulomb barrier, all the points lie close to a straight line whose slope corresponds to a value for d of 0·4 fm, with points away from the line by amounts corresponding to a variation in d of 0·1 fm. Our result $d = 0{\cdot}4 \pm 0{\cdot}1$ fm is to be compared with the result $a = 0{\cdot}8 \pm 0{\cdot}2$ fm obtained by optical model fits to angular distributions. We now make a similar comparison for the collision of light ions with heavy ions.

Fig. 4*b* shows the values of Δ obtained for the collision of light ions with heavy ions. The solid squares are for 40–120 MeV ^{4}He on ^{24}Mg (Singh *et al.* 1975), the solid triangles for 32·3 MeV ^{4}He (Cowley *et al.* 1974) and 23·4–166 MeV ^{4}He (Singh *et al.* 1975) on ^{58}Ni, and the solid circles for 40–118 MeV ^{4}He (Singh *et al.* 1975) and 140 MeV ^{4}He (Goldberg *et al.* 1974) on ^{90}Zr. The three dashed lines pass close to the points for Mg, Ni and Zr respectively, and, while the fluctuations from linearity are due partly to inaccuracies in analysis, there seems to be a small increase in Δ on going from the lightest to the heaviest of the three target ions. Taking the mean slope of all three curves for ^{4}He at the larger values of $k - k_{\rm B}$ we obtain a mean value for d of about 0·5 fm, slightly more than the mean value of 0·4 fm for incident heavy ions.

The open squares in Fig. 4*b* are for 20·5 and 32·6 MeV ^{3}He on ^{12}C (Karban *et al.* 1977), 46 MeV ^{3}He on ^{27}Al and ^{28}Si (Fulmer *et al.* 1978), 71 MeV ^{3}He on ^{60}Ni (Fulmer *et al.* 1973), 71 MeV ^{3}He on ^{209}Bi (Fulmer *et al.* 1975) and 109·2 MeV ^{3}He on ^{40}Ca, ^{58}Ni, ^{90}Zr and ^{116}Sn (Hyakutake *et al.* 1978). The open circles are for 11·8 and 21·6 MeV ^{2}H on Cu (Perey and Perey 1963) and 80 MeV ^{2}H on ^{208}Pb (Duhamel *et al.* 1971). The crosses are for 30·8, 45, 160 and 185 MeV ^{1}H on ^{208}Pb (van Oers *et al.* 1974). There are too few points for the ions ^{3}He, ^{2}H and ^{1}H for firm conclusions, but the indications are that Δ is significantly greater than for ^{4}He. The result that Δ is greater for incident ^{3}He than for ^{4}He is also shown by comparing the form of the S matrix for 217 MeV ^{3}He and 166 MeV ^{4}He incident on each of four very different heavy ions (Willis *et al.* 1973). The present comparison of values of Δ indicates that the penetration of incident heavy ions into target heavy ions is nearly the same for different pairs of heavy ions, but that incident light ions penetrate a little deeper, a result to be expected.

The less regular behaviour of incident ^{4}He ions as compared with incident heavy ions, as shown by comparing Fig. 2 with Fig. 1, and Fig. 4*b* with Fig. 4*a*, may be understood in terms of the degree of departure of the S matrix from Woods–Saxon form, a form on which the theory is based. Thus the form for 100 MeV ^{18}O on ^{120}Sn (Glendenning 1975) and 100 MeV ^{32}S on ^{27}Al (Garrett *et al.* 1975) is closely symmetric about the midpoint at $l = l_{\rm c}$, while the form for 166 MeV ^{4}He and 217 MeV ^{3}He on ^{12}C, ^{40}Ca, ^{120}Sn and ^{208}Pb (Willis *et al.* 1973) is markedly asymmetric. Another reason is that $l_{\rm c}$ is smaller for incident light ions than for heavy ions of the same energy, and the theory is valid only for sufficiently large $l_{\rm c}$.

In spite of the limitations of the theory when applied to incident light ions, it is worth comparing the values of d and a for this case. The distribution of density in most nuclei is approximately of Woods–Saxon form with width parameter $0\cdot5$ fm (Hodgson 1971). Taking this as d_2 and folding in a small extra width d_1 for the incident light ion, we obtain a value of rather more than $0\cdot5$ fm for d, as indicated by Fig. 4*b*. This is to be compared with the value for a (Hodgson 1971) of $0\cdot7$ fm for ^{1}H, $0\cdot85$ fm for ^{2}H (larger because ^{2}H is loosely bound), $0\cdot7$ fm for ^{3}He, and $0\cdot6\pm0\cdot1$ fm for ^{4}He (largest Igo ambiguity is for ^{4}He). Now our value of d is a little smaller for incident heavy ions than for light ions, and we expect the same tendency for the value of a, indicating a value for a of less than $0\cdot7$ fm for incident heavy ions, but just a little above our values of d of between $0\cdot6$ and $0\cdot4$ fm. It seems therefore that the values of a of between $0\cdot7$ and $1\cdot0$ fm used in some of the optical model calculations for pairs of heavy ions listed in the references above are too large, and that we have been able to reduce the ambiguity.

5. Conclusions

We now have a more complete understanding of the collision of heavy ions. This study suggests values for the nuclear surface diffuseness which are smaller, and lie within a smaller range, than the values tried in some optical model calculations.

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