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Density-dependent effect on $\alpha-^{12}C$ elastic scattering

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Abstract. An analytical expression for the real part of the folding potential, using a density-dependent effective interaction, is derived by taking the density of the ^{12}C nucleus according to Brink's α-particle model. The present model has succeeded, particularly at large angles, in describing the rainbow scattering observed in the elastic scattering of α-particles by ^{12}C at the energies 104, 139 and 166 MeV respectively. The density-dependent parameter β is found to be energy independent.

1. Introduction

The elastic scattering of α-particles at energies high enough ($E_\alpha > 100$ MeV) for refractive rainbow scattering at angles larger than the rainbow angle to be observed, resolves the discrete ambiguities of the optical potential that occur in the analyses of data at lower energies (Goldberg 1975, Kobos *et al* 1982). Another advantage of analysing the data at higher energies is that, since no potential exists for grazing l values, the cross section is almost exclusively due to barrier-scattering amplitudes with no interference from the internal wave amplitudes (Brink and Takigawa 1977). This leads to reduced sensitivity of the cross section to the details of the complex potential. For these reasons, the elastic scattering of α-particles at high energies, probing further into the target nucleus, can be used to assess the applicability of different folding models for calculating the real potential.

Single and double folding models of the real optical potential for nucleon, light and heavy ion scattering have been successfully applied to the analysis of the elastic and inelastic scattering cross sections (Jackson and Kembhari 1969, Rebel *et al* 1974, Petrovich *et al* 1977, Majka *et al* 1978a,b Satchler and Love 1979). The different folding models provide a relevant microscopic interaction for the α-particle scattering (Gils 1984). In many cases, the nuclear structure information (size and shape of nuclear densities) needed for the folding procedure has been taken from independent sources while the effective projectile–nucleon interaction has been the subject of several studies concerning the importance of including explicitly the density dependence in the folding model approach (Friedman *et al* 1978, Bernstein 1969, Gils *et al* 1980, Srivastava 1982, 1983a,b, 1984).

The effective interaction is weakened as the density of the nuclear medium, in which the two nucleons are embedded, is increased. This is a reflection of the saturation property of nuclear matter (Von Geramb 1979). This density dependence arises partly from the interchange of the two nucleons (antisymmetrization) and partly from the Pauli-forbidden intermediate states in the Brueckner reaction matrix (Jeukenne and Mahaux 1981).

Nuclear rainbow scattering, first observed in α-particle elastic scattering at energies $E_\alpha \geqslant 100$ MeV, has revealed the importance of the density dependence, in particular, when probing deeper into the nucleus (Gils *et al* 1984, Khoa 1988, El Azab Farid and Satchler 1985a,b 1988, Gils 1984, 1987, Majka *et al* 1978a,b).

Several analytical forms for the density-dependent (DD) effective interaction have been proposed (Srivastava 1984, Petrovich *et al* 1977, Friedman *et al* 1978, Majka *et al* 1978a,b, Jeukenne and Mahaux 1981, Gils 1987, Thompson *et al* 1986) and yet the problem of using the local density has not been solved in a unique way.

In a previous paper (Esmael *et al* 1990), calculations of the elastic scattering cross section of α–^{12}C, performed with Brink's α-particle folding model for the real part of the optical potential, for a density-independent (DI) interaction, and a phenomenological imaginary Woods–Saxon potential, have given, in general, a better agreement with the experimental data than those calculated with shell model or Hartree–Fock distributions (Tatischeff and Brissaud 1970). This has encouraged us to use, in the present work, the same model for the description of the ^{12}C nucleus but with the introduction of a density-dependent effective interaction in order to improve the fit to the experimental data of the elastic scattering of α–^{12}C at the three different energies 104, 139 and 166 MeV. Using the simple form of DD effective interaction given by Morgan and Jackson (1969), Hodgson (1978) and Feng *et al* (1976), an analytical form of the real optical folding potential is obtained.

The theory is formulated in section 2. The results and discussion are given in section 3 and the conclusion is presented in section 4.

2. Theory

The double folded real optical potential is written as (Nagata *et al* 1985):

$$V(R) = \int \mathrm{d}\boldsymbol{r}_1\, \mathrm{d}\boldsymbol{r}_2 \rho_1(\boldsymbol{r}_1)\rho_2(\boldsymbol{r}_2)V(s, \rho) \tag{1}$$

where $\boldsymbol{s} = \boldsymbol{R} + \boldsymbol{r}_1 - \boldsymbol{r}_2$.

$\rho_i(\boldsymbol{r}_i)$, $i = 1, 2$ is the nucleon point density distribution of the nuclei, $\boldsymbol{R}$ is the separation between the centres of mass of the two colliding nuclei and V is the DD effective NN interaction which may be taken as (Petrovich *et al* 1977, Gils 1984):

$$V = V_1(r) f(\rho) \tag{2}$$

where ρ is the local density, which may be assumed to consist of the sum of the two 'frozen' nuclear densities (Srivastava 1983a,b). This is normally taken as:

$$\rho = \rho_1(|\boldsymbol{r}_1 - \tfrac{1}{2}\boldsymbol{s}|) + \rho_2(|\boldsymbol{r}_2 + \tfrac{1}{2}\boldsymbol{s}|). \tag{3}$$

Further, introducing a short-range assumption (Srivastava 1983a,b) for the convenience of calculations, then, equation (3) may be written as:

$$\rho = \rho_1(\boldsymbol{r}_1) + \rho_2(\boldsymbol{r}_2). \tag{4}$$

Although this assumption is known to slightly underestimate the nuclear potential in the interior (Goldfarb and Nagel 1980), it is generally used as it considerably simplifies the computation of the double folded potential (Satchler and Love 1979), Petrovich 1975).

A simple linear dependence on ρ is assumed (Morgan and Jackson 1969, Hodgson 1978, Feng *et al* 1976)

$$f(\rho) = (1 - \beta\rho) \tag{5}$$

where β is the density-dependent parameter. For the radial part of NN effective interaction the set (1) (Knyazkov and Hefter 1981), which has given a reasonable agreement with the experimental data for DI calculations (Esmael *et al* 1990), is used:

$$V_1(r) = V_0 e^{-r^2/a^2} \tag{6}$$

where $a = 1.47$ fm and $V_0 = -20.97$ MeV.

Substituting equations (2), (5) and (6) into equation (1) gives:

$$V(R) = V_0 \int \exp[-(|\boldsymbol{R} + \boldsymbol{r}_1 - \boldsymbol{r}_2|)^2/a^2][\rho_1(r_1)\rho_2(r_2) \\ - \beta(\rho_2^2(r_1)\rho_2(r_2) + \rho_1(r_1)\rho_2^2(r_2)]\,\mathrm{d}\boldsymbol{r}_1\,\mathrm{d}\boldsymbol{r}_2. \tag{7}$$

According to Brink's model (Wadia and Moharram 1975), the density of the ^{12}C target nucleus may be obtained by expanding it in a series of spherical harmonics, by means of Sonine's expansion formula (Erdelyi *et al* 1953) in terms of the modified Bessel functions I and by applying the addition theorem for spherical harmonics to get the spherical part of the density in the form:

$$M_0(r_1) = \frac{r_1^{-1/2}}{6b^2\pi\sqrt{R_1}}\left[HI_{1/2}\left(\frac{2r_1R_1}{b^2}\right) + FI_{1/2}\left(\frac{r_1R_1}{b^2}\right)\right]\exp[-(r_1^2 + R_1^2)/b^2] \tag{8}$$

where $R_1 = 1.4$ fm, $b = 1.36$ fm, $H = 3N_1 + 12N_3$, $F = 2^{3/2}(3N_1 - 6N_3)$, $N_1 = 1/(3 + 6\eta)$, $N_3 = 1/(6 - 6\eta)$, $\eta = \exp(-\frac{3}{4}R_1^2/b^2)$ and $I_{1/2}(x)$ is the modified Bessel function of order $\frac{1}{2}$.

Choosing the density of the α-particle in the form (Knyazkov and Hefter 1981)

$$\rho_2(r_2) = \rho_{0\alpha}\exp(-\gamma r_2^2) \tag{9}$$

where $\rho_{0\alpha} = 4(\gamma/\pi)^{3/2}$ and $\gamma = 0.514\ \mathrm{fm}^{-2}$, which reproduces the experimental $\langle r^2 \rangle$ for the α-particles fairly well.

In the present calculations, the spherical part of $\rho_1^2(r_1)$ is derived following the same procedure used to obtain $M_0(r_1)$ given by equation (8).

Substituting equations (8) and (9) into equation (7) and carrying out the integrations over $\boldsymbol{r}_1$ and $\boldsymbol{r}_2$ yields for the real analytical potential:

$$V(R) = V_0 I_1(R) - \beta V_0 I_2(R) \tag{10}$$

where

$$I_1(R) = \sqrt{\pi}\left(\frac{2\gamma a^2}{\gamma(b^2 + a^2) + 1}\right)^{3/2}\exp\left(-\frac{\gamma(R^2 + R_1^2)}{\gamma(a^2 + b^2) + 1}\right)\left[M_1 I_{1/2}\left(\frac{2RR_1\gamma}{\gamma(a^2 + b^2) + 1}\right) \\ - M_2\exp\left\{-\tfrac{3}{4}\left[(R_1/b)^2\left(\frac{1 + \gamma a^2}{\gamma(b^2 + a^2) + 1}\right)\right]\right\}I_2\left(\frac{\gamma RR_1}{\gamma(a^2 + b^2) + 1}\right)\right]$$

$$M_1 = 4N_3^2 + N_1^2 \qquad M_2 = 4N_3^2 - 2N_1^2$$

and

$$I_2(R)=\frac{1}{3\pi}\left(\frac{2\gamma}{\gamma(2a^2+b^2)+2}\right)^{3/2}\left(\frac{a}{b}\right)^3\exp\left[-\left(\frac{2\gamma(R_1^2+R^2)}{\gamma(2a^2+b^2)+2}\right)\right]$$
$$\times\left[\!\!\left[M_1^2 I_{1/2}\left(\frac{4\gamma RR_1}{\gamma(2a^2+b^2)+2}\right)+(2M_1^2+M_2^2)\right.\right.$$
$$\times\exp\left\{-3\left[(R_1/b)^2\left(\frac{\gamma a^2+1}{\gamma(2a^2+b^2)+2}\right)\right]\right\}$$
$$\times I_{1/2}\left(\frac{2\gamma RR_1}{\gamma(2a^2+b^2)+2}\right)+2M_2(M_2-M_1)$$
$$\times\exp\left\{-\frac{15}{4}\left[(R_1/b)^2\left(\frac{\gamma a^2+1}{\gamma(2a^2+b^2)+2}\right)\right]\right\}I_{1/2}\left(\frac{\gamma RR_1}{\gamma(2a^2+b^2)+2}\right)-4M_1M_2$$
$$\left.\left.\times\exp\left\{-\frac{9}{4}\left[\left(\frac{R_1}{b}\right)^2\left(\frac{\gamma a^2+1}{\gamma(2a^2+b^2)+2}\right)\right]\right\}I_{1/2}\left(\frac{\sqrt{7}\,\gamma RR_1}{\gamma(2a^2+b^2)+2}\right)\right]\!\!\right]$$
$$+\frac{2^{7/2}a^3\gamma^3}{\pi[2\gamma(a^2+b^2)+1]^{3/2}}\exp\left(-\frac{2\gamma(R^2+R_1^2)}{2\gamma(a^2+b^2)+1}\right)\left[\!\!\left[M_1 I_{1/2}\left(\frac{4\gamma RR_1}{2\gamma(a^2+b^2)+1}\right)\right.\right.$$
$$\left.\left.-M_2\exp\left\{-\frac{3}{4}\left[\left(\frac{R_1}{b}\right)^2\left(\frac{1+2\gamma a^2}{2\gamma(a^2+b^2)+1}\right)\right]\right\}I_{1/2}\left(\frac{2\gamma RR_1}{2\gamma(a^2+b^2)+1}\right)\right]\!\!\right].$$

The total optical potential used may be written in the form (Cook *et al* 1982):

$$U(R)=N_F V(R)+\mathrm{i}W(R)+V_C(R) \tag{11}$$

where $V(R)$ is the real folded DD potential defined in equation (10), N_F is a normalization parameter and $V_C(R)$ is the Coulomb potential. The imaginary potential $W(R)$ is taken in two different forms: namely the Woods–Saxon form

$$W(R)=-W_0\left[1+\exp\left(\frac{R-r_I A_T^{1/3}}{a_I}\right)\right]^{-1}$$

and in the form (Burov *et al* 1983)

$$W(R)=\alpha V(R)$$

where α is a fitting parameter.

3. Results and discussion

The density-dependent effects have been taken into account in the present analysis of elastic scattering of α-particles by ^{12}C, considered according to Brink's α-particle model, at the three different energies 104, 139 and 166 MeV.

The computed optical potential using the present density-dependent (DD) model F_{1d} (given in equation (10)) for the value $\beta=0.6\,\mathrm{fm}^3$ obtained from best fit calculations of elastic scattering is compared with the density-independent (DI) potential F_1 (given in equation (10) for $\beta=0$) in figure 1. Both folding model potentials (F_1 and F_{1d}) are compared with the Woods–Saxon potentials at the three considered α-particle incident energies (Hauser *et al* 1969, Smith *et al* 1973, Tatischeff and Brissaud 1970).

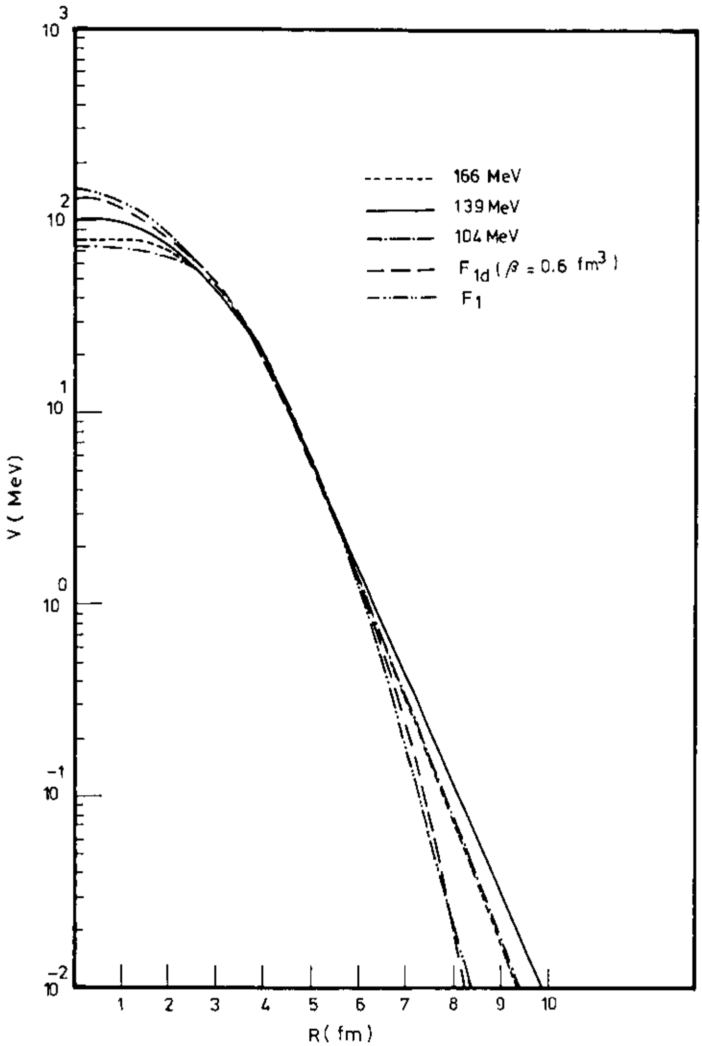

Figure 1. Comparison of the folding model potentials F_1(DI) and F_{1d}(DD) with the WS potentials at 104 MeV (chain curve, Hauser *et al* (1969)), 139 MeV (full curve, Smith *et al* (1973)) and 166 MeV (dotted curve, (Tatischeff and Brissaud (1970)).

The introduction of the density dependence in the effective interaction causes some reduction in the central depth in the nuclear interior. This result is consistent with previous calculations taking the DD into consideration (Kobos *et al* 1982). The two folding potentials are close to each other near the surface. The folding models and the Woods–Saxon potentials (WS) cross each other close to the strong absorption radius ($R_{sa} = 5.8$ fm). The real parts of the folding model potentials fall off more rapidly than the WS potential, which is in agreement with previous calculations (Kobos *et al* 1982).

In the present calculations, the effect of varying the parameter β between zero and unity is studied in order to reduce the contribution from the inner region of the nucleus which leads to an improved fit at large angles (Morgan and Jackson 1969).

As the value of β increases from zero (DI) to 0.6, a satisfactory fit with the data is

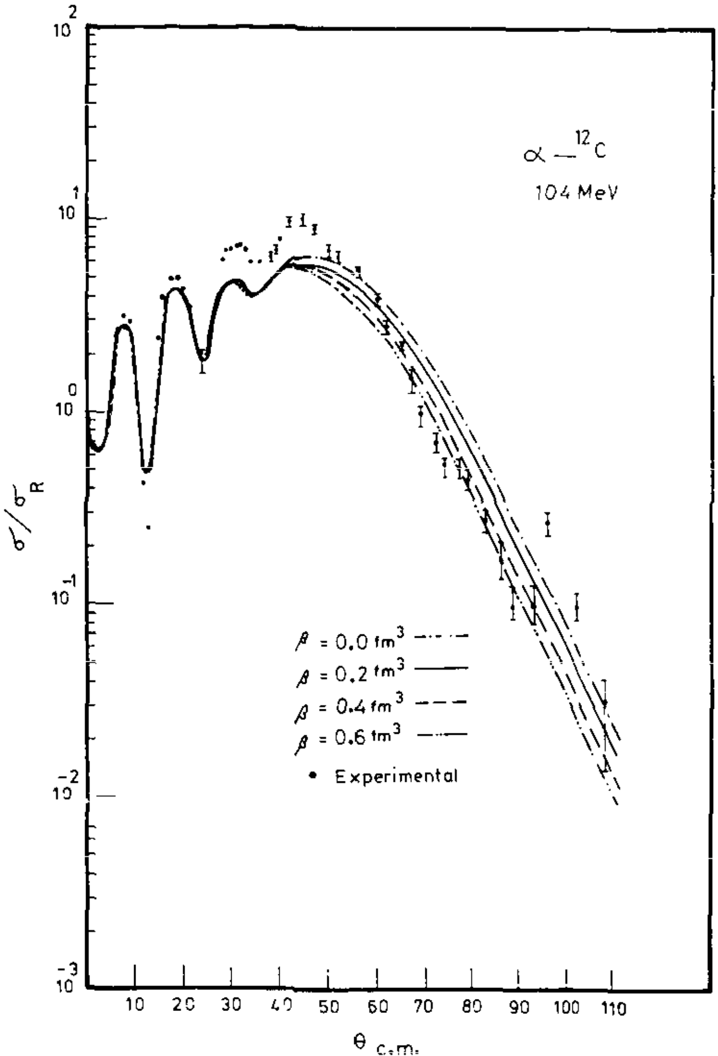

Figure 2. The differential cross section for the elastic scattering of α–^{12}C at 104 MeV calculated for different values of DD parameter β and for WS imaginary optical potential (OP) (Hauser *et al* 1969).

obtained at large angles (figure 2). There are no appreciable changes in the quality of the fit up to 40° for the different values of β. The best fit with experimental data, at large angles, is obtained for $\beta = 0.6\,\mathrm{fm}^3$.

Figure 3 shows the same calculations but for a different set of the optical model parameters of the imaginary part of WS potential (Hammad 1989). It can be seen that, although the fit, at forward angles, has been improved, the fit at large angles is less satisfactory than the calculations with the previous set of parameters (Hauser *et al* 1969). Indeed, the parameters of the imaginary part of the WS potential play a role in the quality of the fit with the experimental data.

Figure 4 displays the fit with the experimental data at energy 139 MeV. It may be noticed that there are no great differences between the fits obtained for the different

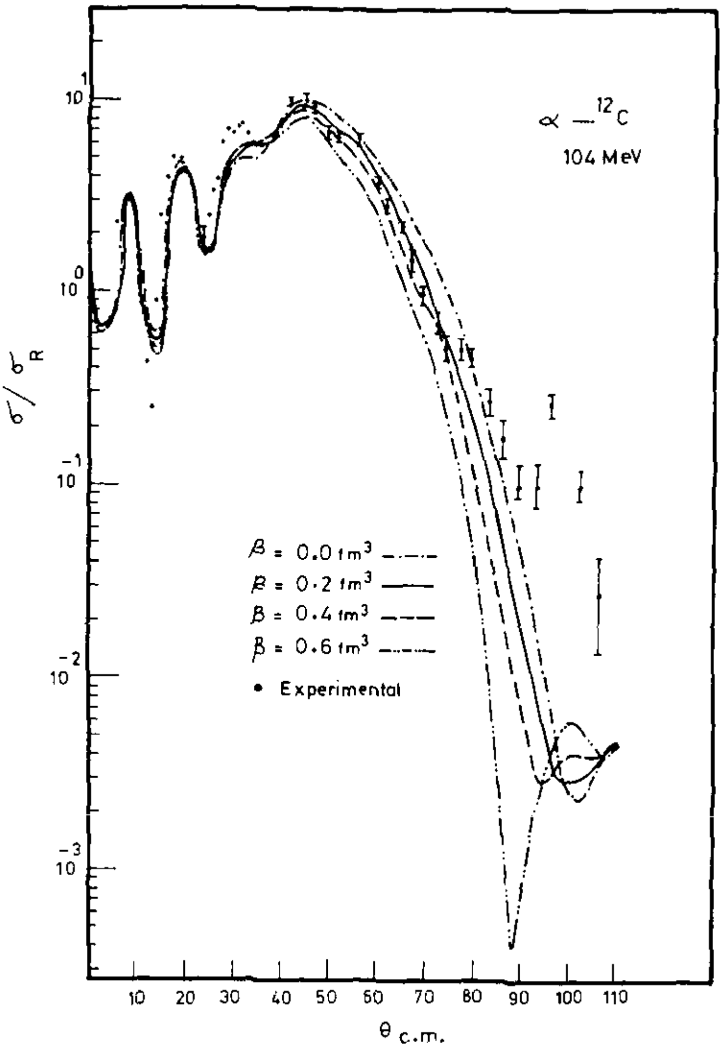

Figure 3. The differential cross section for the elastic scattering of α–^{12}C at 104 MeV calculated for different values of DD parameter β and for WS imaginary optical potential (OP) (Hammad 1989).

values of β up to 40°. There is an obvious improvement of the fit at large angles as β increases from zero up to 0.6 giving the best fit, as has been noticed in figure 2.

The present model gives better results than those of Khoa (1988), using only the direct part of the M3Y interaction without any exchange term, and also gives (for $\beta = 0.6\,\mathrm{fm}^3$) similar results to those of Khoa (1988) considering the finite range exchange term (M3Y/FRE). The present model has two advantages over that of Khoa (1988): first, the normalization coefficient is taken equal to unity while that of Khoa is 0.731, secondly, Khoa has varied the parameters of the imaginary part of the optical potential in order to obtain the best fit, while in the present calculations the parameters are kept unchanged. This means that the ^{12}C nucleus may be preferably represented by Brink's α-particle model rather than by the Fermi distribution as used by Khoa (1988).

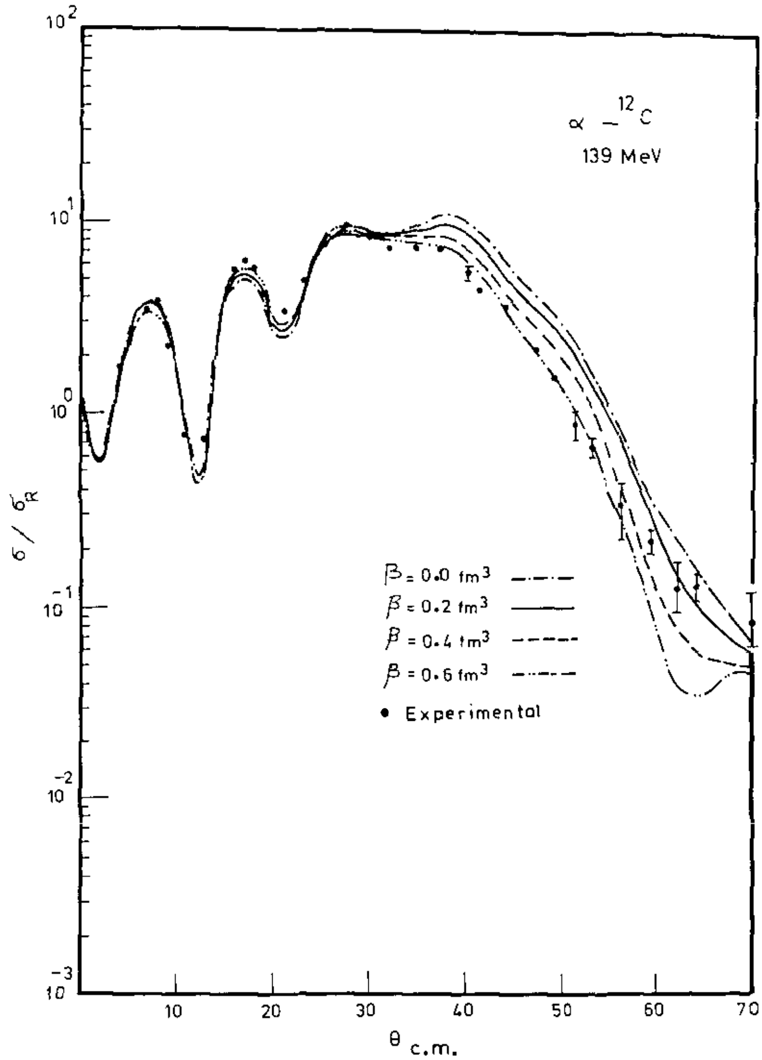

Figure 4. The differential cross section for the elastic scattering of α–^{12}C at 139 MeV calculated for different values of DD parameter β and for WS imaginary OP (Smith *et al* 1973).

The fit of the calculated elastic scattering of α-particles by ^{12}C at 166 MeV with experimental data (Tatischeff and Brissaud 1970) is shown in figure 5. It can be seen that the DD calculations for $\beta = 0.6\,\mathrm{fm}^3$ give a better fit with the experimental data than those for DI calculations, especially at large angles. It may be noticed that the effect of increasing the value of β is the same as that seen in figures 2 and 4, i.e. improving the fit at large angles.

In general, the reduction in the nuclear interior that occurs in the real part of the folding potential, due to the introduction of the density dependence in the effective nucleon–nucleon interaction, is responsible for the improvement of the fit in the description of the rainbow scattering observed in the elastic scattering of α on ^{12}C at energies $E_\alpha \geqslant 100$ MeV especially at large angles.

The calculations have been repeated considering the imaginary part of the

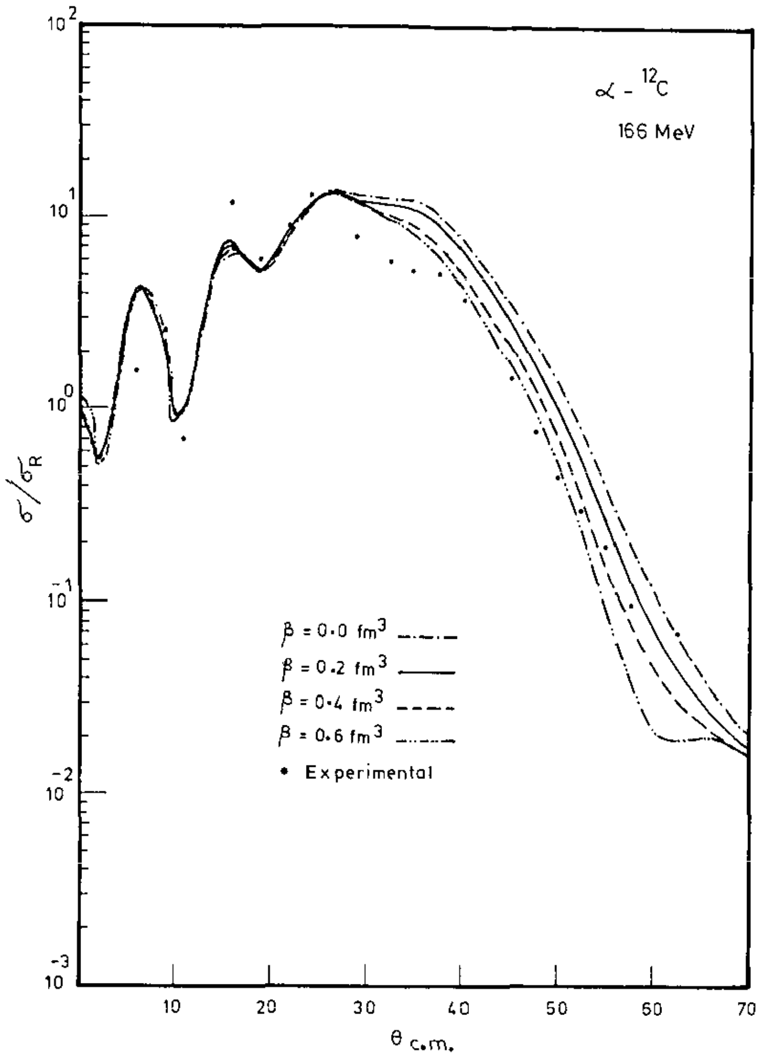

Figure 5. The differential cross section for the elastic scattering of α–^{12}C at 166 MeV calculated for different values of DD parameter β and for WS imaginary OP (Tatischeff and Brissaud 1970).

potential as a fraction of the real part. The results are shown in figures 6, 7 and 8 for 104, 139 and 166 MeV respectively. It may be seen that, for the smaller values of the adjusted parameter α, the fit is good only for small angles but is poor for large angles, while for the larger values of α, the fit becomes good for large angles and poor for small angles. The fit can be considered reasonable for the values of $\alpha = 0.4$, 0.46 and 0.5 at the energies 104, 139 and 166 MeV respectively over the whole angular distribution; these values are the same as those obtained for DI interaction in a previous paper (Esmael *et al* 1990), i.e. the density dependence effect is hidden by using this form of the imaginary potential.

The present model gives a reasonable fit to the experimental data at the energies 139 and 166 MeV rather than those obtained by Burov *et al* (1983) and Tatischeff

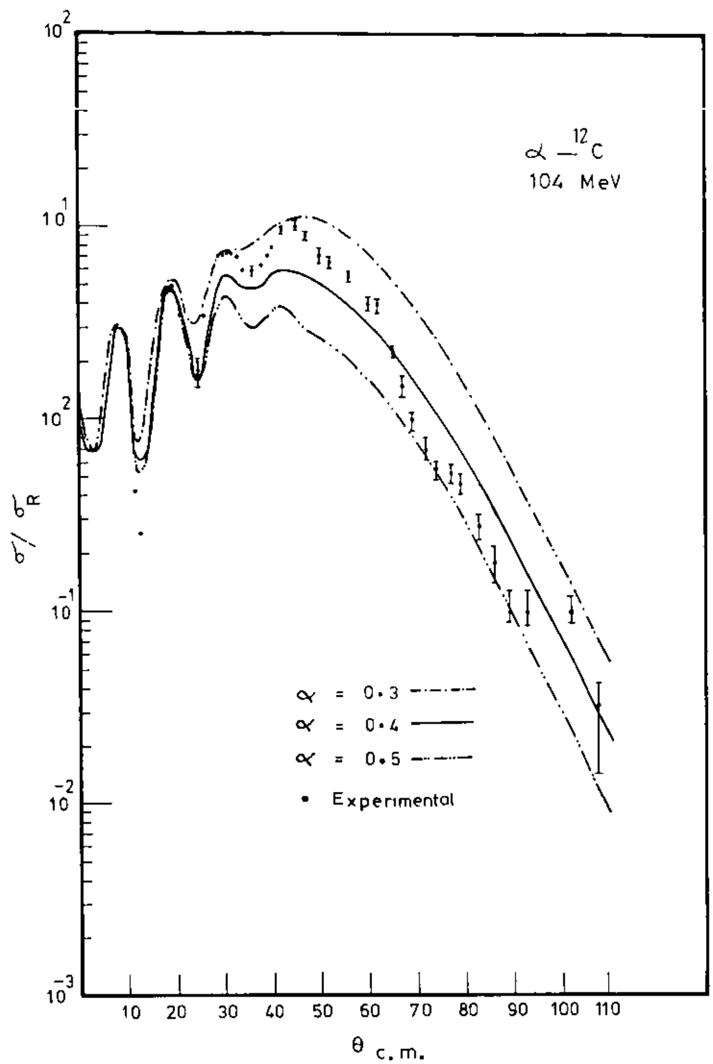

Figure 6. The differential cross section for the elastic scattering of α–^{12}C at 104 MeV calculated for different values of adjusted parameter α.

and Brissaud (1970) using densities of matter calculated in the hyperspherical method or shell model and Hartree–Fock respectively to describe the density of ^{12}C. Thus Brink's α-cluster model allows a reasonable representation of ^{12}C.

The values N_F/α are approximately 2.5, 2 and 2 for the energies 104, 139 and 166 MeV respectively, which are larger than the corresponding ratios W/V of the Woods–Saxon potential. This may be explained by the fact that, in the microscopic calculations, the same geometry for the real and imaginary terms are chosen, while in the WS case, the automatic search gives $R_I > R_R$ (Tatischeff and Brissaud 1970). The shapes of the potential are, however, similar at the nuclear surface.

Comparing between the calculations using the imaginary parts of the optical potential as a Woods–Saxon form and as a fraction of the real part (αV_R), it can be seen that the first model gives a better fit with the experimental data than the second one. This is consistent with the interpretation given by El Azab Farid and Satchler

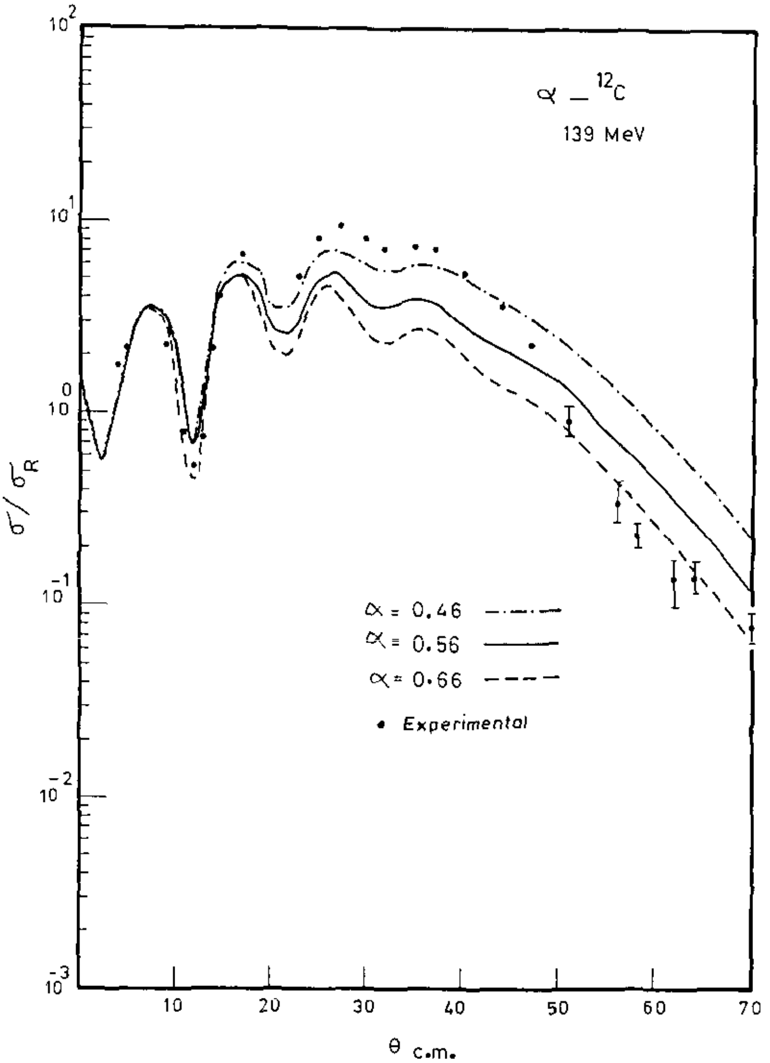

Figure 7. The differential cross section for the elastic scattering of α–^{12}C at 139 MeV calculated for different values of the adjusted parameter α.

(1988) in which they state that the values of α that give the correct peripheral absorption for α-particles then provide much too strong an absorption in the interior. However, the imaginary interaction is known less reliably, even for nucleon scattering (Jeukenne and Mahaux 1981, Jeukenne *et al* 1977) because it tends to be more strongly influenced by surface phenomena such as collective oscillations and transfer reactions. Because of this, a phenomenological form is frequently chosen for the imaginary optical potential.

4. Conclusion

From the above results and discussion, one may conclude that the cluster folding model potential with DD interaction gives a good description of the differential cross

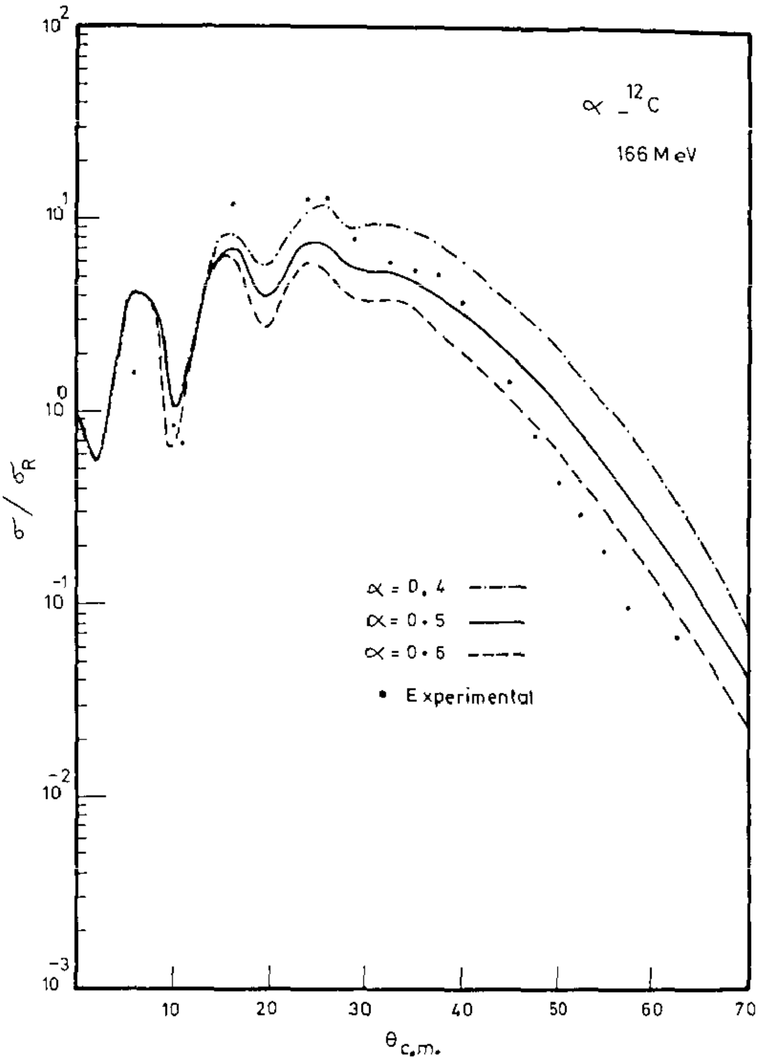

Figure 8. The differential cross section for the elastic scattering of α–^{12}C at 166 MeV calculated for different values of the adjusted parameter α.

sections showing rainbow scattering in the elastic scattering of α-particles by ^{12}C at the three energies 104, 139, 166 MeV specially at large angles. These results agree with previous calculations carried out for different nuclei by Kobos *et al* (1982), Gils *et al* (1984) and Chaudhuri (1986a,b).

The choice of the imaginary part in the Woods–Saxon form is more appropriate than the other choice (as a fraction of the real part). This result agrees with the conclusions of Kobos *et al* (1982), El Azab Farid and Satchler (1988), Burov *et al* (1983) and Katori *et al* (1988).

The effect of varying the parameter β between zero (DI) and unity has been studied in order to reduce the contribution from the inner region of the nucleus. The value of $\beta = 0.6\,\text{fm}^3$ gives the best fit with the experimental data which is consistent with that found by Morgan and Jackson (1969) and Hodgson (1978). Moreover, the value of β does not depend on the energy which confirms the results found by Chaudhuri (1986a,b) in the energy region 20–43 MeV per nucleon.

The discrepancies between the calculated differential cross sections and the experimental data may be reduced by taking into acccount the effects of the available non-elastic channels of the composite projectile (El Azab Farid and Satchler 1988, Katori *et al* 1988).

For the deformed light nucleus of ^{12}C, a perceptible effect due to target excitation is recognized on the elastic channel in comparison with the double folding and coupled channel calculations (Katori *et al* 1988, Faessler *et al* 1984).

The role of the exchange effects is most important in the refractive rainbow observed in α-scattering at large angles, where both elastic and inelastic scattering data can be reproduced only by propertly taking into account the Pauli principle (Khoa 1988). Also taking into account nucleon–nucleon correlations (Khoa *et al* 1989) is of essential importance in the description of the effects of nuclear rainbow scattering.

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