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FOREWORD

The Summary Report of the Physics Division of the Argonne National Laboratory is issued monthly for the information of the members of the Division and a limited number of other persons interested in the progress of the work. Each active project reports about once in 3 months, on the average. Those not reported in a particular issue are listed separately in the Table of Contents with a reference to the last issue in which each appeared.

This is merely an informal progress report. The results and data therefore must be understood to be preliminary and tentative.

The issuance of these reports is not intended to constitute publication in any sense of the word. Final results either will be submitted for publication in regular professional journals or, in special cases, will be presented in ANL Topical Reports.

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I. EXPERIMENTAL NUCLEAR PHYSICSI-18-16 Neutron Polarization and Differential Cross Sections (51210-01)

J. E. Monahan and A. J. Elwyn

THE LOCAL APPROXIMATION TO A
NONLOCAL OPTICAL-MODEL POTENTIAL

Recently Perey and Buck¹ have investigated the consequences of using a nonlocal potential in optical-model calculations. The empirical fact that the results of calculations based on the nonlocal model could always be reproduced by an equivalent local model led these authors to derive an approximate relation between their nonlocal potential $U_N(\vec{r})$ and an equivalent local potential $U_L(\vec{r})$, namely,

$$U_L(\vec{r}) \exp \left[\frac{M\beta^2}{2\hbar^2} [E - U_L(\vec{r})] \right] = U_N(\vec{r}). \quad (1)$$

The length β is a measure of the range of nonlocality of $U_N(\vec{r})$. In view of the assumptions employed in its derivation, this relation was found to be surprisingly accurate. However, the transcendental nature of Eq. (1) is inconvenient for numerical calculations and, more importantly, for obtaining a correspondence between the $U_N(\vec{r})$ defined by Eq. (1) and the nonlocal potentials which have been used previously in optical-model calculations. It is the purpose of this note to examine the error in the approximation (1) and, in particular, to investigate the possibility of obtaining a more convenient functional representation for an equivalent local potential.

Briefly, Eq. (1) is derived in reference (1) as follows. For the particular nonlocal potential used in these calculations, the operator form of the nonlocal wave equation is

¹F. G. J. Perey and B. Buck, Nuclear Phys. 32, 353 (1962).

$$\left[\frac{\hbar^2}{2M} \nabla^2 + E \right] \psi_N(\vec{r}) = \exp \left[\frac{\beta^2}{4} \left(\frac{1}{2} \nabla_1 + \nabla_2 \right)^2 \right] U_N(\vec{r}) \psi_N(\vec{r}), \quad (2)$$

where ∇_1 operates only on the potential form factor $U_N(\vec{r})$ and ∇_2 operates only on the nonlocal wave function $\psi_N(\vec{r})$. The first approximation in the derivation of Eq. (1) is the neglect of the contribution of the operator ∇_1 in Eq. (2). The error $\epsilon_1(\vec{r})$ associated with this approximation is

$$\epsilon_1(\vec{r}) = \left\{ \exp \left[\frac{1}{4} \beta^2 \left(\frac{1}{2} \nabla_1 + \nabla_2 \right)^2 \right] - \exp \left[\frac{1}{4} \beta^2 \nabla_2^2 \right] \right\} U_N(\vec{r}) \psi_N(\vec{r}). \quad (3)$$

An equivalent local potential $U_L(\vec{r})$, if it exists, must satisfy the relation

$$\nabla_2^2 \psi_N(\vec{r}) = f_L(\vec{r}) \psi_N(\vec{r}), \quad (4)$$

where

$$f_L(\vec{r}) = \frac{2M}{\hbar^2} \left[U_L(\vec{r}) - E \right]. \quad (5)$$

Finally a second approximation is made, namely that the operator $\exp \left[\frac{1}{4} \beta^2 \nabla_2^2 \right]$ in Eq. (3) is replaced by the operator $\exp \left[\frac{1}{4} \beta^2 f_L(\vec{r}) \right]$. This, of course, is equivalent to the assumption that $U_L(\vec{r})$ in Eq. (5) is independent of $\vec{r}$, just as the approximation (3) is equivalent to the assumption that $U_N(\vec{r})$ is independent of $\vec{r}$. The error $\epsilon_2(\vec{r})$ associated with this second approximation is

$$\epsilon_2(\vec{r}) = \left\{ \exp \left[\frac{1}{4} \beta^2 \nabla_2^2 \right] - \exp \left[\frac{1}{4} \beta^2 f_L(F) \right] \right\} U_N(\vec{r}) \psi_N(\vec{r}). \quad (6)$$

Written out explicitly, the operator equations (3) and (6) are respectively

$$\epsilon_1(\vec{r}) = \sum_{n=2}^{\infty} \frac{1}{(n-1)!} \left(\frac{\beta^2}{4} \right)^{n-1} \left\{ \left(\frac{1}{2} \nabla_1 + \nabla_2 \right)^{2n-2} - \nabla_2^{2n-2} \right\} U_N(\vec{r}) \psi_N(\vec{r}) \quad (7)$$

and

$$\epsilon_2(\vec{r}) = \sum_{n=2}^{\infty} \frac{1}{(n-1)!} \left(\frac{\beta^2}{4} \right)^{n-1} \left\{ \frac{1}{n} \left(\frac{\beta^2}{4} \right) \left(\nabla^2 - f_L(\vec{r}) \right) \right\} U_N(\vec{r}) \psi_N(\vec{r}). \quad (8)$$

By repeated use of Eq. (4), the over-all error of the approximation (1) at the point $\vec{r}$ can be expressed in terms of the coefficients of $\psi_N(\vec{r})$ and $\text{grad } \psi_N(\vec{r})$. The first term in the error $\epsilon_1(\vec{r}) + \epsilon_2(\vec{r})$ is

$$\frac{\beta^2}{4} \left\{ \nabla \psi_N \cdot \left[\nabla U_N + \frac{M\beta^2 U_N}{2\hbar^2} \nabla U_L \right] + \frac{1}{4} \psi_N \left[\nabla^2 U_N + \frac{M\beta^2 U_N}{\hbar^2} \nabla^2 U_L \right] \right\}.$$

Although the condition that the coefficients of ψ_N and $\text{grad } \psi_N$ vanish separately is not necessary for the validity of Eq. (1), it is sufficient and can therefore be used as a guide in the search for a convenient functional representation of $U_L(\vec{r})$. For example, if it is assumed that $U_N(\vec{r})$ and $U_L(\vec{r})$ have the same functional dependence on $\vec{r}$, then, for all points $\vec{r}$ such that $U_N(\vec{r})$ is less than zero, corresponding terms in $\epsilon_1(\vec{r})$ and $\epsilon_2(\vec{r})$ have opposite sign so that there is a partial cancellation of the errors introduced separately by approximations (3) and (6).

Figure 1 compares the differential cross section and polarization calculated by the nonlocal model of reference 1 (solid curve) with that of a local model using well depths derived from the parameters of the nonlocal potential (dashed curve). Both potentials were assumed to have the same functional dependence on $\vec{r}$. The agreement is not sufficient for most applications. It remains to determine whether meaningful improvement can be obtained by a slight relaxation of the condition of identical functional dependence of the two potentials. The figure also shows a curve (dotted) calculated on the basis of a local model but with different well-depth parameters than were used for the dashed curve. The agreement with the nonlocal-model calculations is distinctly better than it was for the dashed curve.

Several qualitative features of approximation (1) can be obtained from the expansions (7) and (8). Obviously $\epsilon_1(\vec{r}) + \epsilon_2(\vec{r})$ vanishes for those points $\vec{r}$ where U_N and U_L are independent of $\vec{r}$. However, there

is also a cancellation between corresponding terms in the over-all error for those points $\vec{r}$ where U_N and U_L vary rapidly with $\vec{r}$. The higher order terms in $\epsilon_1(\vec{r}) + \epsilon_2(\vec{r})$ are dependent on the energy E , and for increasing values of E the interval in r over which the error is appreciably different from zero increases. Finally, any correlation between the radial dependences of the real and imaginary parts of the optical potential U_N such that large values of one are correlated with large values of the derivative of the other can lead to large errors. There is some indication that this is the reason for the poor agreement in the example shown in the figure.

Fig. 1. The polarization and differential cross section as a function of $\theta_{c.m.}$ for neutrons scattered from Cd at 0.45 MeV. The solid curves are based on the nonlocal optical model of reference 1. The dashed curves are the result of a local-optical-model calculation using parameters based on the nonlocal values, and assuming the same functional dependence of the potentials on $\vec{r}$. The points marked x and the associated dotted curve are the results of a local-optical-model calculation using well-depth parameters modified somewhat from those corresponding to the dashed curve. For all three curves, therefore, the potentials are assumed to have the same functional dependence on $\vec{r}$ but with the well-depth parameters given by the following table.

	Real well depth (MeV)	Imag. well depth (MeV)	Real spin-orbit well (MeV)
Nonlocal (solid)	71	15	7.2
Local (dashed)	47.36	7.02	7.2
Local (dotted)	50	9	4

Also shown are the values of the total, shape-elastic, and compound-elastic cross sections (σ_t , $\sigma_{s.e.}$, and $\sigma_{c.e.}$, respectively) in barns, corresponding to the three curves. In these calculations inelastic scattering is assumed to be negligible.

I-31-2

Elastic Scattering of Protons

(51210-01)

L. L. Lee, Jr., and J. P. Schiffer

Reported by J. P. Schiffer

The elastic scattering of protons from two isotopes of Cu and two of Ni has been studied at bombarding energies of 7, 8, 9, 10, 11, and 12 MeV with the proton beam of the Argonne tandem Van de Graaff accelerator. Absolute cross sections were determined by comparing the scattering of 6-MeV alpha particles with the scattering of protons in the same scattering chamber. In this fashion, it was assured that the same portions of the targets were illuminated by the incident beam. It was found that the usual techniques of weighing the target foils could lead to errors as large as 25%, particularly for the Ni foils used in the measurements, because of nonuniformities. These are presumably due to the crystalline structure of Ni foils.

These cross sections are believed to be accurate to 10%. Additional measurements in progress at the present time will hopefully increase the accuracy of the absolute cross sections to 3%. The measured angular distributions compared to the angular distributions calculated from an optical model with parameters suggested by F. Perey of ORNL are shown in Fig. 2 for Cu^{63} and in Fig. 3 for Ni^{60} . These parameters

Fig. 2. Measured elastic-scattering cross sections from Ni^{60} expressed as the ratio to the Rutherford cross section. The line corresponds to the optical-model calculations with the parameters of Perey.

Fig. 3. Measured elastic-scattering cross sections from Cu^{63} expressed as the ratio to the Rutherford cross section. The line corresponds to the optical-model calculations with the parameters of Perey.

are defined by the relations

$$V = \frac{V_0}{1 + \exp[-(r - r_0)/a]},$$

$$V_0 = 53 - 0.55 E + \left[0.4 \frac{Z}{A^{1/3}} + 27 \frac{(N - Z)}{A} \right] \quad \text{in MeV,}$$

$$r_0 = 1.25^{1/3} F, \quad a = 0.65 F,$$

$$W = W_0 \exp \left\{ -[(r - r_0)/b]^2 \right\},$$

$$W_0 = 13.5 \text{ MeV,}$$

$$b = 1.0 F,$$

$$V_{s0} = 7.5 \text{ MeV.}$$

where the symbols have their usual meaning. The agreement is worse for the lower energies, particularly for Ni^{60} . Similar discrepancies were noted at energies below 11 MeV on Ni^{58} . The reason for this is apparent from the 90° excitation functions shown in Fig. 4. Since there are sharp fluctuations in the

Fig. 4. Excitation function for elastic scattering of protons. The target thickness was approximately 50 keV.

cross sections at the lower energies, it is obvious that the statistical assumption of the optical model (that of averaging over a large sample of states) is not satisfied and therefore one cannot expect good agreement. The presence of fluctuations implies that resonances occur in the different phase shifts and that these come in much more frequently than can possibly be predicted by the optical model. It is therefore meaningless to try to adjust the parameters in this region to give a better fit. The reason for the resonances occurring in some nuclei and not in others seems to be associated with the proximity of the neutron threshold. The results are consistent with the assumption that 1-2 MeV above the neutron threshold the levels become sufficiently broad and overlapping to allow the optical model to be applied. At lower energies the widths of the levels are too small, fluctuations occur as observed in the excitation function, and the optical model is not applicable.

Recently there has been some interest in interpreting fluctuations in cross sections in terms of "lifetimes of the compound

Fig. 5. Excitation function for elastic scattering of protons. The target thickness was approximately 5 keV.

nucleus." That such an analysis of the data shown in Fig. 4 can be misleading is shown in Fig. 5 where short sections of the excitation function on the two Ni isotopes and Cu^{63} are shown for the case of a thin, 5-keV target. It is clear that while Fig. 4 showed sharp fluctuations of the order of 50-100 keV where measurements were taken with 50-keV targets, the results in Fig. 5 show the sharper fluctuations of the order of 5 keV. It may well be that thinner targets would reveal even sharper fluctuations. It seems then that one should be ex-

ceedingly careful not to interpret fluctuations which are of the same magnitude as the target thickness as true fluctuations in the cross sections of the compound nucleus.

An attempt has been made to test Perey's prediction as to the dependence of the optical-model potential on neutron excess. It can be seen from Fig. 6 that this dependence is consistent with the experimental observations but that the experiment is also consistent with the difference obtained by assuming no explicit dependence on the neutron excess but allowing the radius to increase as $A^{1/3}$.

Fig. 6. Comparison of elastic scattering from Cu^{63} and Ni^{60} at 12 MeV. The line in the left figure corresponds to the difference predicted by Perey's parameters, the one in the right figure to the difference predicted if the radius were allowed to increase as $A^{1/3}$ with no change in potential depth.

II. MASS SPECTROMETRY

II-29-18 Gaseous Species in Equilibrium at High Temperatures
(51300-01)

J. Berkowitz and J. R. Marquart
Reported by J. Berkowitz

MASS SPECTROMETRIC STUDY OF THE MAGNESIUM HALIDES

The vapor composition of the four magnesium halides has been studied by mass-analyzing the vapor effusing from a Knudsen cell. Typical results for the saturated vapor at the indicated temperature are shown in Table I. In each instance, $\approx 1\%$ dimer and 0.01% trimer were

TABLE I. Partial pressures of magnesium
halide components at representative temperatures.

Partial pressure (atm.)	MgF ₂ (1429°K)	MgCl ₂ (919°K)	MgBr ₂ (798°K)	MgI ₂ (734°K)
MgX ₂	1.83×10^{-5}	3.8×10^{-5}	4.4×10^{-6}	5.18×10^{-6}
Mg ₂ X ₄	$\approx 3.77 \times 10^{-8}$	9.0×10^{-7}	8.8×10^{-8}	1.15×10^{-7}
Mg ₃ X ₆	$\approx 1 \times 10^{-10}$	2.1×10^{-8}	8.8×10^{-10}	2×10^{-9}

observed. This is only slightly smaller than the degree of polymerization observed in BeF₂¹ and BeCl₂,² although earlier vapor-density measurements indicated that the beryllium halides should be predominantly polymeric in the saturated vapor. Heats of sublimation have been measured by second-law and third-law methods. The results are shown in Table II. The discrepancy between the two methods varies from 2 kcal/mole for MgI₂ to 5 kcal/mole for MgF₂, with the third-law method giving consistently lower values. This could possibly be due to

¹ J. Berkowitz and W. A. Chupka, New York Acad. Sci. 79, 1073 (1960).

² L. N. Ryabchikov and G. F. Tikhinskii, Fiz. Metal. i. Metalloved. Akad. Nauk. S.S.S.R. 10, 635 (1960). Also, in unpublished experiments performed in 1959, J. Berkowitz, D. J. Meschi, and W. A. Chupka observed a yield of 1-2% dimer at 495°K and a BeCl₂ pressure of $\approx 10^{-5}$ atm.

TABLE II. Heats of sublimation of the magnesium halides.

	MgF ₂	MgCl ₂	MgBr ₂	MgI ₂
ΔH_T (slope, or 2nd-law method)	86	54.6	50.3	45.0
ΔH_T (absolute- pressure or 3rd-law method)	81.2	50.7	46.2	43.2
ΔH_0	85.6	53.9	50.9	47.4

the existence of electronic states that have not been considered in the entropy calculations. The experimental heats of ionization deduced from these data (Table III) are in fair agreement with calculations based on an ionic model.

TABLE III. Heats of ionization of the magnesium halides.

	Experimental (kcal/mole)	Calculated (kcal/mole)		
		Cubicciotti ^a	Cubicciotti ^b	Krasnov ^c
MgF ₂	604.5	591	560	564
MgCl ₂	543.6	533	511	532
MgBr ₂	519.5	513	490	...
MgI ₂	498.5	500	473	...

^aSee reference 3. Repulsive potential used was of the form Ar^{-n} .

^bSee reference 3. Repulsive potential used was of the form $Ae^{-r/\rho}$.

^cSee reference 4.

³ D. Cubicciotti, Stanford Res. Inst. Tech. Note No. 7, 1960.

⁴ K. S. Krasnov, Dokl. Akad. Nauk S.S.S.R. 128, 326 (1959).

The actual ions observed were Mg^+ , MgX^+ , MgX_2^+ , Mg_2X_3^+ , Mg_3X_5^+ , X^+ , X_2^+ . The ions Mg_2X_3^+ and Mg_3X_5^+ were attributed to dimer and trimer, respectively, while all of the others were considered to be due to the preponderant monomer. In the fragmentation pattern of the monomer, the ratio $\text{MgX}^+/\text{MgX}_2^+$ deduced from these data was less than unity for $\text{X} = \text{Cl}, \text{Br}, \text{and I}$, but was about 4 for $\text{X} = \text{F}$. This anomalously high fragmentation in the fluoride has also been observed with the alkali fluorides, and is presumably connected with the fact that the increase in internuclear distance when an initially neutral metal fluoride ion is ionized is greater than the increase for the other metal halides. This in turn can be rationalized as being due to the smaller electronic polarizability of the fluorine atom.

V. THEORETICAL PHYSICS, GENERAL

V-45-20.

Meson-Nucleon Interaction

(51151-01)

L. S. Liu^{*} and K. Tanaka

Regge introduced the idea of allowing the angular momentum ℓ , which appears as a parameter in the radial Schrödinger equation, to take not only continuous but even complex values, thus introducing the complex angular-momentum ℓ plane. The poles $\alpha(E)$ of the scattering amplitude $A(\cos \theta, E)$ in the ℓ plane can give rise to a family of bound states and resonances. The paths of the poles, which move continuously as functions of the energy E , are called Regge trajectories. We have studied the Regge trajectory of the pole $\alpha(t)$ as a function of the square of the momentum transfer t in field theory based on the sum over ladder diagrams in the crossed channel for scalar neutral particles of mass m . These diagrams are shown in Fig. 7.

Fig. 7. Ladder graphs for the scattering amplitude.

One assumes that a single Regge pole with complex α dominates the amplitude A at large values of s , the square of the total energy in the center-of-mass system, and obtains an equation for the trajectory of α . Its solution gives α as a function of t , the exchanged mass $\sqrt{a}$, and the coupling constant g^2 . The Regge pole that appears in the crossed channel t (the channel in which t is the square of the total energy) dominates the high-energy behavior in the s channel.

^{*} Northwestern University

A Regge trajectory, a plot of the real part of $\alpha(t)$ as a function of t is of the form shown in Fig. 8 for $a = 4m^2$ and when g^2 is

Fig. 8. Plot of the real part of the complex angular momentum as a function of t .

expected from a physical point of view,

normalized so that $\text{Re } \alpha(0) = 1$.

In the physical region of the s channel, $t \leq 0$.

For $t \leq 4m^2$, α is real and $d\alpha/dt > 0$, and $\text{Im } \alpha$ develops slowly from $t \geq 4m^2$. The various relations between g^2 , $\alpha(t)$, and a have been studied.

The features obtained are those

PUBLICATIONS SINCE THE LAST REPORT

PAPERS

ON THE FLUORESCENCE SPECTRUM AND DECAY TIME OF NAPHTHALENE

I. B. Berlman (RPY) and A. Weinreb (Project I-144)
Molecular Phys. 5, 313-319 (July 1962)

POLARIZATION AND DIFFERENTIAL CROSS SECTIONS IN n-d SCATTERING

A. J. Elwyn, R. O. Lane, and A. Langsdorf, Jr. . . (Project I-18)
Phys. Rev. 128, 779-789 (October 15, 1962)

SEARCH FOR AN EFFECT OF MAGNETIC FIELDS ON ELECTRO-MAGNETIC RADIATION

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PERSONNEL CHANGES IN THE ANL PHYSICS DIVISION
NEW MEMBERS OF THE DIVISION

Technician

Mr. Richard H. Davis joined the Physics Division on November 14,
1962 as a Research Technician (Junior) at the Van de
Graaff.