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Polarization in High-Energy Elastic Nucleon-Nucleus Scattering*

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The proposal to explain the comparatively large polarization observed in high-energy elastic proton-nucleus scattering by means of the spin-orbit interaction used in the shell model of the nucleus is examined. Three simple examples are considered in first Born approximation. The approximate magnitude and the approximate location of the maximum polarization for 340-Mev nucleons on carbon is in rough agreement with experiment. However, regions of negative polarization also seem to be predicted by the theory. Finally, an approximate method to calculate the polarization for high energy and small scattering angle is suggested.

I. INTRODUCTION

RECENT experiments[1–4] in the double scattering of high-energy protons by a nucleus indicate a considerably larger amount of spin polarization in the emerging proton beam than is observed in proton-proton scattering. These experiments also seem to indicate that the largest asymmetry occurs in the range of energy and scattering angle where the scattering can be expected to be mostly elastic.[3] In a previous communication,[5] the author suggested that these results might be explained by means of the spin-orbit interaction used in the shell model of the nucleus. This suggestion also has been advanced independently by other authors.[6,7] Some of the consequences will be elaborated here.

In the nuclear shell model,[8,9] it has been shown that the observed shell structure of nuclei can be explained by introducing a relatively strong spin-orbit interaction in the potential well model of the nucleus. We make the additional assumption that for elastic nucleon-nucleus scattering essentially the same spin-orbit force that acts on a bound nucleon also acts on the nucleon being elastically scattered by the nucleus. Since the strength of this interaction depends on the spin orientation of an incident nucleon relative to its orbital momentum, the scattered nucleon beam will be spin-polarized. For an unpolarized incident beam, the double scattering experiments referred to above which measure the right-left asymmetry of the second scattered nucleon beam are used to determine the amount of polarization.

We therefore suppose that the nucleus exerts a potential H' on the incident nucleon. This potential is of the form

$$H' = (1+i\epsilon)V_0(r) + V_1(r)\tfrac{1}{2}\boldsymbol{\sigma}\cdot\mathbf{r}\times(-i)\nabla, \qquad (1)$$

where the parameter ϵ is introduced to take account of nuclear absorption.[10] At 100 Mev, the value of ϵ is around $\frac{1}{2}$ and apparently slowly increases with higher energies.[11] The $V_0(r)$ is a central potential well and $V_1(r)$ describes the radial dependence of the spin-orbit interaction. The magnitude of $V_1(r)$ should be adjusted so that it gives the correct order of doublet splitting when calculated by bound state perturbation theory,[12] that is,

$$\Delta E = E_{l+\frac{1}{2}} - E_{l-\frac{1}{2}} = \langle \psi_{E,l}(r),\, V_1(r)\psi_{E,l}(r)\rangle (l+\tfrac{1}{2}), \qquad (2)$$

where $\psi_{E,l}(r)$ is the radial part of a bound state wave function for the unperturbed central nuclear potential well. According to Mayer,[8] ΔE should be roughly about -2 Mev for large l, say l about 5.

For an interaction given by (1), we assume that we can express the elastic scattering amplitude in the form[13]

$$f(\theta) = f_0(\theta) + \boldsymbol{\sigma}\cdot\mathbf{n} f_1(\vartheta), \qquad (3)$$

where $\mathbf{n}$ is a unit vector defined by

$$\mathbf{k}_0\times\mathbf{k} = \mathbf{n}k^2\sin\theta, \qquad (4)$$

where $\mathbf{k}_0$ and $\mathbf{k}$ are the initial and final nucleon propagation vectors. For an unpolarized incident beam, the polarization is then

$$P(\theta) = \frac{\frac{1}{2}\,\mathrm{Tr}[f^\dagger(\theta)\boldsymbol{\sigma}\cdot\mathbf{n}f(\theta)]}{\frac{1}{2}\,\mathrm{Tr}[f^\dagger(\theta)f(\theta)]} = \frac{f_0^*(\theta)f_1(\theta)+f_0(\theta)f_1^*(\theta)}{|f_0(\theta)|^2+|f_1(\theta)|^2}, \qquad (5)$$

where the trace is taken over the spin space. With this definition of $\mathbf{n}$, $P(\theta)>0$ corresponds to more nucleons with spin up being scattered to the left and $P(\theta)<0$ more with spin down being similarly scattered. We note that the differential cross section for single elastic scattering which is just the denominator of (5) depends on the sum of the absolute squares of the scattering

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† After September, 1954, at Washington University, St. Louis, Missouri.

[1] Oxley, Cartwright, and Rouvina, Phys. Rev. **93**, 806 (1954).
[2] Marshall, Marshall, and Carvalho, Phys. Rev. **93**, 1431 (1954).
[3] Chamberlain, Segrè, Tripp, Wiegand, and Ypsilantis, Phys. Rev. **93**, 1430 (1954).
[4] J. M. Dickson and D. C. Salter at Harwell, as reported in the Proceedings of the 1954 Rochester Conference on High Energy Physics (University of Rochester Press, Rochester, to be published).
[5] B. J. Malenka, Bull. Am. Phys. Soc. **29**, No. 4, 32 (1954).
[6] E. Fermi, Nuovo cimento (to be published).
[7] W. Heckrotte and J. V. Lepore, Phys. Rev. **94**, 500 (1954).
[8] M. G. Mayer, Phys. Rev. **78**, 16 (1950).
[9] Haxel, Jensen, and Suess, Ann. Physik **128**, 295 (1950).
[10] R. E. LeLevier and D. S. Saxon, Phys. Rev. **87**, 40 (1952).
[11] I. I. Shapiro and J. M. Teem (private communication).
[12] B. J. Malenka, Phys. Rev. **86**, 68 (1952).
[13] See for instance, J. Schwinger, Phys. Rev. **73**, 407 (1948) or L. Wolfenstein, Phys. Rev. **75**, 1664 (1949).

amplitudes, whereas $P(\theta)$ is sensitive to their relative phases and accordingly can have negative values.

From the shell model, we know that $V_0(r)$ is some modified square well extending a distance of the order of the nuclear radius. The depth of this well appears to be in the vicinity of 30 Mev, although current optical model calculations seem to indicate that for high-energy elastic nucleon-nucleus scattering, the effective well depth may be considerably reduced.[11] However, the radial dependence of $V_1(r)$ is not as well determined since several functions of r will yield approximately the right order of doublet splitting in (2) when the magnitude of $V_1(r)$ is suitably adjusted. The forms suggested so far for the radial dependence either assume an approximately uniform effect,[12] in which case $V_1(r)$ is itself some sort of modified square well, or represent an estimate the effect of a core acting on a nucleon,[14,15] in which case $V_1(r)$ approximately has the form $r^{-1}d\phi/dr$, where $\phi(r)$ characterizes the collective action of the core. The resulting form of the second category also includes a generalization of the relativistic Thomas effect which when unmodified gives a doublet splitting which is too small.[16] Here $V_1(r)$ would be proportional to $r^{-1}dV_0/dr$.

II. BORN APPROXIMATION

Since the range of energies used in the polarization experiments are relatively high and since the absorption parameter ϵ at these energies is larger than $\frac{1}{2}$, we can expect a first Born approximation calculation of the scattering amplitudes to give approximately quantitatively correct results for the polarization.†

For this approximation, we can use Eq. (5) directly. However, in many cases, it becomes convenient to make the rearrangement

$$f_0(\theta) = (1+i\epsilon) f_0'(\theta), \tag{6}$$

$$f_1(\theta) = i\tfrac{1}{2}\lambda\!\!\!^{-2}k^2 \sin\theta f_1'(\theta), \tag{7}$$

where $\lambda\!\!\!^{-}$ is the nucleon Compton wavelength, and where $f_0'(\theta)$ and $f_1'(\theta)$ are now real scattering amplitudes which are calculated in first Born approximation from the real potentials $V_0(r)$ and $\mathcal{V}_1(r) = -\lambda\!\!\!^{-\,-2}\int rV_1(r)dr$. The latter potential is obtained by partial integration in $f_1(\theta)$. The expression (5) for $P(\theta)$ then becomes

$$P_0(\theta) = \frac{\epsilon\lambda\!\!\!^{-2}k^2 \sin\theta f_0'(\theta) f_1'(\theta)}{(1+\epsilon^2)[f_0'(\theta)]^2 + \frac{1}{4}\lambda\!\!\!^{-4}k^4 \sin^2\theta [f_1'(\theta)]^2}. \tag{8}$$

In this form, it can be readily shown that $P_0(\theta)$ reaches its optimum value for the scattering angles, θ_0, which satisfy the equation

$$(1+\epsilon^2)[f_0'(\theta_0)]^2 = \tfrac{1}{4}\lambda\!\!\!^{-4}k^4 \sin^2\theta [f_1'(\theta_0)]^2. \tag{9}$$

The maximum and minimum values of $P_0(\theta)$ are then

$$P_0(\theta_0) = \pm\epsilon/(1+\epsilon^2)^{\frac{1}{2}}, \tag{10}$$

provided that θ_0 exists in the interval from 0 to π.

The effect of also including a shielded Coulomb interaction would be to add the Coulomb amplitude $f_c(\theta)$ to $f_0(\theta)$. Since this amplitude is real, it will only enter in the denominator of $P_0(\theta)$. It can be easily verified that except at very small angles, the contribution of the Coulomb term is comparatively quite small, so that neglecting it will not essentially alter any of the results presented in this paper.

As an elementary example of a uniform spin-orbit effect, we consider the potential defined by

$$\begin{aligned} &V_0(r) = -U_0, \quad V_1(r) = -U_1 = -\alpha_0 U_0, \quad r \le R, \\ &V_0(r) = 0, \qquad\;\; V_1(r) = 0, \qquad\qquad\qquad\quad r > R. \end{aligned} \tag{11}$$

The resulting expression for $P_0(\theta)$ is then

$$P_0^{(0)}(\theta) = \frac{\epsilon\alpha_0 kR \cos\frac{1}{2}\theta\, j_1(2kR\sin\frac{1}{2}\theta)\, j_2(2kR\sin\frac{1}{2}\theta)}{(1+\epsilon^2) j_1^2(2kR\sin\frac{1}{2}\theta) + \frac{1}{4}\alpha_0^2 k^2 R^2 \cos^2\frac{1}{2}\theta\, j_2^2(2kR\sin\frac{1}{2}\theta)}, \tag{12}$$

where j_1 and j_2 are regular spherical Bessel functions.

In Fig. 1, we have plotted $P_0^{(0)}(\theta)$ for carbon at 340 Mev using a nuclear radius $R = 3.2\times10^{-13}$ cm. The value of the absorption part of the well, ϵU_0, was taken to be about 16 Mev which is consistent with an optical model determination of this term from the sum of np and pp total cross sections. We also took U_0 to be about 27 Mev to facilitate comparison with reference 6. This gives $\epsilon \approx 0.6$ and from (10), a maximum polarization of approximately 51 percent. Finally from (2) and reference 12, we estimated U_1 to be about $\frac{1}{2}$ Mev and hence $\alpha_0 \approx 0.02$.

If we consider a generalization of the relativistic Thomas effect, in first Born approximation, we find that we can permit $V_0(r)$ to be any potential well where

$$V_1(r) = -\alpha_1 \lambda\!\!\!^{-2} \frac{1}{r}\frac{d}{dr} V_0(r). \tag{13}$$

Then $\mathcal{V}_1(r) = \alpha_1 V_0(r)$ and $f_1'(\theta) = \alpha_1 f_0'(\theta)$, so that

$$P_0^{(1)}(\theta) = \frac{\epsilon\alpha_1\lambda\!\!\!^{-2}k^2 \sin\theta}{1+\epsilon^2+\frac{1}{4}\alpha_1^2\lambda\!\!\!^{-4}k^4\sin^2\theta}, \tag{14}$$

[14] J. Keilson, Phys. Rev. 82, 759 (1951).

[15] J. Hughes and K. J. LeCouteur, Proc. Phys. Soc. (London) 63, 1219 (1951).

[16] B. H. Flowers in *Progress in Nuclear Physics*, Editor, O. R. Frisch (Academic Press, Inc., New York, 1952), Vol. 2, p. 271.

† We might note that at high energies the magnitude of $V_1(r)$ decreases and ϵ increases. However, as long as their product is smaller than the incident kinetic energy, the first Born approximation can be expected to be valid, and for large ϵ, contain the principal contribution to the polarization.

giving a polarization that is independent of the shape of $V_0(r)$. This expression for $P_0^{(1)}(\theta)$ agrees with that obtained by Fermi[6] for a square well. The $P_0^{(1)}(\theta)$ of (14) is plotted in Fig. 2 for an estimated $\alpha_1 \approx 15$.

Finally, we consider a simple example such as might be associated with the effect of a nuclear core interacting with a single nucleon. Here we again let

$$\begin{aligned} V_0(r) &= -U_0, \quad r \leq R \\ V_0(r) &= 0, \qquad r > R \end{aligned} \tag{15}$$

and take

$$V_1(r) = -\frac{\lambda\!\!\!^{-2}}{r}\frac{d}{dr}\mathcal{V}_1(r), \tag{16}$$

where

$$\begin{aligned} \mathcal{V}_1(r) &= -\alpha_2 U_0, & r \leq bR, \\ \mathcal{V}_1(r) &= -\alpha_2 U_0 (bR/r) e^{-\mu(r-bR)}, & r > bR. \end{aligned} \tag{17}$$

The resulting $P_0(\theta)$ is

$$P_0^{(2)} = \frac{\epsilon b^2 \alpha_2 \lambda\!\!\!^{-2} k^2 \sin\theta\, j_1(2kR \sin\tfrac{1}{2}\theta)\psi(k,bR,\theta)}{(1+\epsilon^2) j_1^2(2kR\sin\tfrac{1}{2}\theta) + \tfrac{1}{4} b^4 \alpha_2^2 \lambda\!\!\!^{-4} k^4 \sin^2\theta\, \psi^2(k,bR,\theta)}, \tag{18}$$

where

$$\begin{aligned} \psi(k,bR,\theta) = {} & j_1(2kbR \sin\tfrac{1}{2}\theta) \\ & + (\mu bR)^{-1}[1+(2k/\mu)^2 \sin^2\tfrac{1}{2}\theta]^{-1} \\ & \times [\sin(2kbR\sin\tfrac{1}{2}\theta) \\ & \quad + (2k/\mu)\sin\tfrac{1}{2}\theta \cos(2kbR \sin\tfrac{1}{2}\theta)]. \end{aligned} \tag{19}$$

The behavior of (18) is plotted in Fig. 3 for $\alpha_2 \approx 15$, $\mu^{-1} = 1.4 \times 10^{-13}$ cm and $b \approx 0.7$.

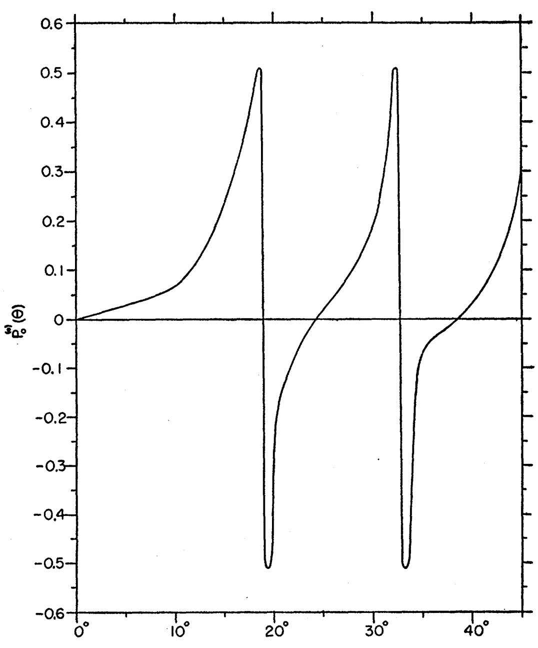

FIG. 1. The polarization as a function of scattering angle as plotted from Eq. (12).

Of the three examples considered, $P_0^{(1)}(\theta)$ and $P_0^{(2)}(\theta)$ appear to best fit the data of Chamberlain *et al.*[3] which indicate a maximum polarization at about 10°. The $P_0^{(0)}(\theta)$ has its first maximum at about 20°. The successive maxima in $P_0^{(0)}(\theta)$ and $P_0^{(2)}(\theta)$ would not be too readily detectable by present methods since they occur at the larger angles where the assumed unpolarized inelastic collisions become important and effectively reduce the observable asymmetry. A characteristic of this theory seems to be the presence of regions of negative polarization following the first maximum, the exception being the first Born approximation of the Thomas-like interaction where the $f_0'(\theta)$ and $f_1'(\theta)$ are proportional to each other, but presumably even here a more exact calculation would show regions of negative polarization.[17] The data referred to in this paper do not show any change in sign of polarization. However, since the experiments do not completely cover the regions where $P_0(\theta)$ can be negative, this possibility

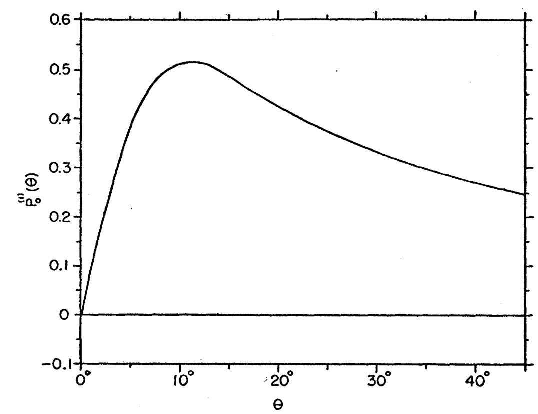

FIG. 2. The polarization as a function of scattering angle as plotted from Eq. (14).

cannot be ruled out. In general, increasing the incident nucleon energy tends somewhat to compress the graphs, towards the lower angles and decreasing the energy stretches them, so that at higher energies the regions of negative polarization may be in the more observable range of the lower scattering angles although the width of these regions may possibly become narrower.‡

The relative magnitudes α of the spin-orbit interaction were only estimated approximately. However, the condition for stationary α is just Eq. (9) for determining $P_0(\theta_0)$ so that small changes in α will leave the location

[17] A phase shift calculation shows that this is indeed the case. E. Fermi (private communication), and Snow, Sternheimer, and Yang, Phys. Rev. **94**, 1073 (1954).

‡ It is also possible that the experimental resolution is such that narrow regions of negative polarization cannot be readily detected. In this respect, we note that a more exact calculation which rounds off the potential wells involved also makes the regions of negative polarization considerably narrower as has been demonstrated by W. Heckrotte (private communication). For heavier nuclei, it may be possible to eliminate the regions of negative polarization altogether; R. M. Sternheimer (private communication).

of optimum polarization relatively unchanged. The magnitude of the maximum polarization is of the proper order, but the figure of 51 percent is not too significant and rather is an approximate lower limit on the maximum polarization. Actually, some of the current optical model calculations mentioned[11] above seem to indicate an $\epsilon \approx 1.2$ for carbon at 340 Mev. This would increase the optimum polarization for $P_0(\theta)$ to about 77 percent.

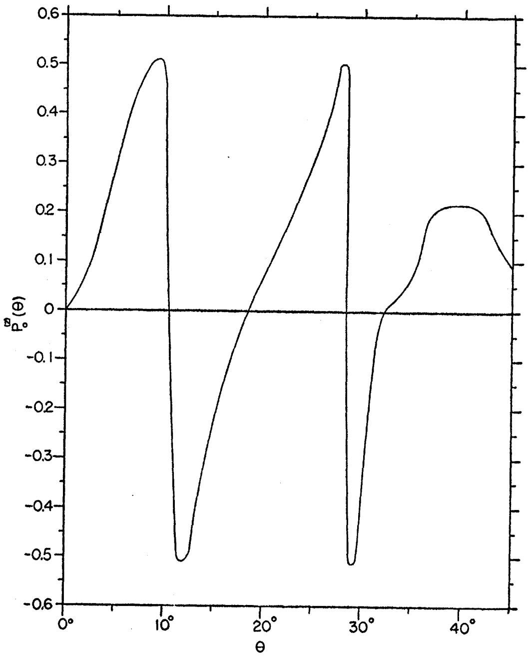

FIG. 3. The polarization as a function of scattering angle as plotted from Eq. (18).

III. SMALL-ANGLE APPROXIMATION

In an attempt to find an approximate treatment more accurate than the first Born approximation, yet less laborious than a phase shift calculation, we again notice that the high-energy elastic polarization experiments are carried out predominantly at relatively small scattering angles. A procedure which takes advantage of these properties for calculating the scattering amplitude has been suggested[18] and is essentially followed below.

The exact scattering amplitude may be written symbolically as

$$f(\theta)=-\frac{k}{2\pi\hbar v}\int \exp(-i\mathbf{k}\cdot\mathbf{r})\left(r\left|H'\frac{1}{1-GH'}\right|r'\right)\times\exp(i\mathbf{k}_0\cdot\mathbf{r}')(d\mathbf{r})(d\mathbf{r}'), \quad (20)$$

where G is the free nucleon Green's function. If we use the expansion

$$[1-GH']^{-1}=1+GH'+GH'GH'+\cdots \quad (21)$$

then $f(\theta)$ is expressed in terms of the Born series. Using cylindrical coordinates with the z axis parallel to $\mathbf{K}=\frac{1}{2}(\mathbf{k}_0+\mathbf{k})$ and evaluating each term of the Born series in the approximation of high energy and small scattering angle,[19] we find that the sum can be written in the form

$$f(\theta)=\frac{ik}{2\pi}\int_0^{2\pi}d\phi\int_0^{\infty}\rho d\rho\,\exp[-i(\mathbf{k}-\mathbf{k}_0)\cdot\boldsymbol{\rho}]\times\{1-\exp[2i\delta_0(\rho)+i\delta_1(\rho)\boldsymbol{\sigma}\cdot\boldsymbol{\rho}\times\mathbf{K}]\}, \quad (22)$$

where

$$\delta_0(\rho)=-\frac{1+i\epsilon}{\hbar v}\int_0^{\infty}V_0([\rho^2+z^2]^{\frac{1}{2}})dz,$$
$$\delta_1(\rho)=-\frac{1}{\hbar v}\int_0^{\infty}V_1([\rho^2+z^2]^{\frac{1}{2}})dz. \quad (23)$$

The expression for $f(\theta)$ can be regarded as a modification of the optical model to include the spin-orbit term provided that a suitable averaging over spin-space is understood. We can write Eq. (22) in the form of Eq. (3) by making use of the commutation properties of the σ and subsequently carrying out the angle integration.

[18] R. J. Glauber, Phys. Rev. 91, 459 (1953).
[19] See Appendix.

We find that

$$f_0(\theta)=ik\int_0^{\infty}\rho d\rho[1-e^{2i\delta_0}\cos(|\mathbf{K}|\rho\delta_1)]J_0(|\mathbf{k}-\mathbf{k}_0|\rho),$$
$$f_1(\theta)=ik\int_0^{\infty}\rho d\rho e^{2i\delta_0}\sin(|\mathbf{K}|\rho\delta_1)J_1(|\mathbf{k}-\mathbf{k}_0|\rho), \quad (24)$$

where $|\mathbf{K}|=\frac{1}{2}|\mathbf{k}_0+\mathbf{k}|=k\cos\frac{1}{2}\theta$ and $|\mathbf{k}-\mathbf{k}_0|=2k\sin\frac{1}{2}\theta$. If we keep only terms linear in δ_0 and δ_1, we obtain the first Born approximation scattering amplitudes such as were used to obtain the results of Sec. II. Presumably, using the scattering amplitudes of (24) instead of a first Born approximation in Eq. (5) would enable us to obtain a better approximation for $P(\theta)$, although as yet, no completely satisfactory estimate of the errors involved in this approximation has been carried out.

The author wishes to take this opportunity to thank K. Strauch for numerous conversations about the current polarization experiments and R. J. Glauber for several interesting discussions.

APPENDIX

The summation of the Born series of (21) is carried out by employing an approximate free particle Green's function[20] for high energy and small angle propagation. To do this, in the usual form of the Green's function

$$G(\mathbf{r},\mathbf{r}')=\frac{1}{(2\pi)^3}\int\frac{\exp[i\mathbf{p}\cdot(\mathbf{r}-\mathbf{r}')]}{p^2-(k^2+i\epsilon)}(d\mathbf{p}), \quad (A.1)$$

[20] This treatment follows the unpublished work of reference 18.

we make the substitution $\mathbf{p}=\mathbf{K}+\mathbf{q}$ where $\mathbf{K}=\frac{1}{2}(\mathbf{k}_0+\mathbf{k})$. Then

$$G(\mathbf{r},\mathbf{r}')=\exp[i\mathbf{K}\cdot(\mathbf{r}-\mathbf{r}')]\frac{1}{(2\pi)^3}\times\int\frac{\exp[i\mathbf{q}\cdot(\mathbf{r}-\mathbf{r}')]}{K^2+q^2+2\mathbf{K}\cdot\mathbf{q}-k^2-i\epsilon}(d\mathbf{q}). \quad (A.2)$$

At high energies, by assuming that

$$K^2-k^2=k^2(\cos^2\tfrac{1}{2}\theta-1)=-k^2\sin^2\tfrac{1}{2}\theta\approx -q^2 \quad (A.3)$$

and taking the z axis parallel to $\mathbf{K}$, we obtain

$$G(\mathbf{r},\mathbf{r}')\approx\exp[i\mathbf{K}\cdot(\mathbf{r}-\mathbf{r}')]\frac{1}{16\pi^3k}\times\int\frac{\exp[i\mathbf{q}\cdot(\mathbf{r}-\mathbf{r}')]}{q_z-i\epsilon}(d\mathbf{q}) \quad (A.4)$$

and when evaluated in cylindrical coordinates for positive $\epsilon\rightarrow 0$, this becomes

$$\begin{aligned} G(\mathbf{r},\mathbf{r}')&\approx\tfrac{1}{2}ik^{-1}\exp[i|\mathbf{K}|(z-z')]\delta(\boldsymbol{\rho}-\boldsymbol{\rho}'), && z>z', \\ G(\mathbf{r},\mathbf{r}')&\approx 0, && z<z'. \end{aligned} \quad (A.5)$$

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Ionization Chamber Measurements at 10 600 Feet of the Absorption of the *N* Component in Carbon and Hydrocarbon*

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A detector consisting of a lead-shielded ionization chamber with Geiger counters above and below the chamber is described. The charged *N* component of energy of several tens of Bev is detected.

The absorption of the charged *N* component in carbon is much smaller than that predicted from the absorption in air. This phenomenon may be explained by the fact that π mesons or other unstable particles which are produced in nuclear interactions can give rise to further nuclear interactions.

Approximate values for the interaction mean free path of *N* rays in carbon have been obtained. These values depend on the number of shielded counters below the ionization chamber which are struck. For zero counters discharged, the mean free path is 136 ± 12 g cm^{-2} and seems to decrease as the requirement on the number of shielded counters struck is increased. The interaction mean free paths in carbon are compared to previous results in lead.

The absorption of the *N* component in oil indicates a cross section of the hydrogen nucleus much smaller than that corresponding to the range of nuclear forces.

Events with large pulses from the ionization chamber or with multiple discharge of shielded counters below the chamber are more likely to be associated with the nucleonic component of extensive air showers.

I. INTRODUCTION

THE results reported here were obtained in a continuation of previous experiments[1–6] performed with a lead-shielded ionization chamber and an array of Geiger-Mueller counters above the lead shield. The time coincident pulses of the ionization chamber and of the Geiger counters recorded the arrival of penetrating ionizing particles which could produce ionization bursts below the lead shield. In a preceding paper[6] (referred to in what follows as I) we were able to separate the total burst rate recorded by the detector into two parts: one which arises from electromagnetic interactions of μ mesons, and another which arises from nuclear interactions of the so-called *N* component of cosmic rays.

The reader is referred to I for a discussion of the processes by which the bursts are produced. From the latitude effect observed with similar apparatus[5] and from correlation of the observed counting rate with the energy spectrum of the producing radiation, the mean energy of the detected radiation is determined to be about 10^{10} ev. This is in agreement with the energy transfer necessary to produce the smallest detected ionization burst (see I).

In the experiment described in I, we measured the absorption in air and in lead of the *N* component responsible for the detected bursts. We also obtained information about the collision mean free path for this radiation in lead.

In the present experiment, performed at 10 600 feet, we have extended the results of I to carbon and hydrocarbon absorbers. We have also investigated the absorption of the *N* component associated with extensive air showers.

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6 H. Bridge and R. Rediker, Phys. Rev. 88, 206 (1952).