# Inelastic corrections to diffraction scattering of high-energy particles by nuclei

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A general formalism is developed for description of the diffraction interaction of high-energy particles with nuclei. Corrections for inelastic screening to the inelastic cross section, the absorption cross section, and to quasielastic scattering in nuclei are obtained. A new method is proposed for analysis of data on diffraction dissociation in nuclei, which permits a lower bound to be found for the cross section for interaction of diffraction-produced systems with nucleons.

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## 1. INTRODUCTION

In the theory of the diffraction interaction of high-energy particles with nuclei, the situation is to a certain extent paradoxical. On the one hand the Glauber-Sitenko theory of multiple scattering,[2,3] which is reliably justified by the field-theory analysis of Gribov,[1] gives a reasonable description of the elastic and quasielastic scattering of hadrons by nuclei and also of the total cross sections and absorption cross sections over a wide range of energies. On the other hand, the rectilinear generalization of the theory of multiple scattering to diffraction-dissociation processes proposed by Kolbig and Margolis[4] and further developed by many authors (see the reviews of Refs. 5 and 6) has led to meaningless results. In the Kolbig-Margolis approximation the cross section for diffraction dissociation in a nucleus is sensitive to the cross section $\sigma_2^*$ for interaction of the diffraction-produced system with the nucleons of the nucleus. The values of $\sigma_2^*$ for $3\pi$, $5\pi$, or $N(n\pi)$ systems determined from the experimentally observed dependence of the diffraction-dissociation cross sections on the mass number of the nucleus $A$ turn out, as a rule, to be less than or of the order of $\sigma_{\pi N}$ and $\sigma_{NN}$,[5,6] and sometimes even turn out to be negative.[7,8]

The reason for this contradiction is clear in its general features. At high energies, diffraction properties are possessed not only by elastic scattering but also by processes of production of small masses as the result of exchange of vacuum quantum numbers in the $t$ channel—diffraction dissociation. The existence of diffraction dissociation means that in reality the diffraction scattering of hadrons is a multichannel process and that ordinary hadrons—which are eigenstates of the mass matrix—are not eigenstates of the diffraction scattering matrix $\hat{T}$. This formulation of the mechanism of diffraction dissociation goes back to the classical studies of Pomeranchuk and Feĭnberg,[9] Akhiezer and Sitenko,[10] Glauber,[11] and Good and Walker.[12]

In optical language the simple theory of multiple scattering corresponds to the gas approximation for the refractive index of nuclear matter:

$$n=1+2\pi\rho_A f_{el}/k^2. \tag{1}$$

Here $\rho_A$ is the density of nuclear matter, $k$ is the momentum of the particle (the system of units $\hbar=c=1$ is used), and $f_{el}$ is the amplitude for forward elastic scattering by a single nucleon. In the general case Eq. (1) involves instead of $f_{el}$ the matrix $\hat{f}$ of the amplitudes of all diffraction transitions at an angle 0°. The simple multiple scattering theory corresponds to approximation of the matrix $\hat{f}$ by the single number $f_{el}$, and the Kolbig-Margolis formula for the cross section in a nucleus corresponds to replacement of $\hat{f}$ by a $2\times 2$ matrix. These approximations are usually justified by the fact that for all dissociation processes $aN\to cN$ the amplitudes $f_{ac}$ are much smaller than the elastic scattering amplitudes $f_{aa}$ and formally perturbation theory in the nondiagonal amplitudes $f_{ac}$ seems to be applicable.

Corrections to the simple multiple scattering theory due to diffraction-dissociation processes were discussed for the first time by Gribov,[1] who calculated the correction to the total cross section for interaction with a deuteron. The similar correction $\Delta\sigma_{tot}^A$ to the total cross sections for interactions with heavy nuclei in perturbation theory in $f_{ac}$ was found by Karmanov and Kondratyuk.[13] Although $\Delta\sigma_{tot}^A/\sigma_{tot}^A \sim 5\cdot 10^{-2}$, a quantitative description of the total cross sections for $nA$ and $K_LA$ scattering is impossible without taking into account this correction.[14-17] At the same time it remains unclear why the absorption cross sections are described by the simple multiple scattering theory very well,[18,19] while in the case of diffraction dissociation the perturbation theory in which the Kolbig-Margolis formula was derived leads to completely meaningless results for $\sigma_2^*$.

In this situation it is important to understand why perturbation theory is applicable in some situations and poor in others, why it is necessary to describe diffraction dissociation outside the framework of perturbation theory, and what kind of connection there is between coherent and incoherent diffraction scattering outside the framework of perturbation theory. The analysis of this problem is the subject of the present work. This problem has been partially discussed already in a previous note by the author[20] regarding the limit of applicability of the Karmanov-Kondratyuk approximation for the total cross sections. General formulas for the amplitudes of coherent diffraction dissociation outside the framework of perturbation theory were obtained in

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the work of Good and Walker[12] and also by Czyz and Zielinski.[21] In the present work we formulate a general description of incoherent diffraction dissociation, quasielastic scattering of hadrons by nuclei, and the cross sections for absorption of hadrons by nuclei.

It will be shown that the corrections to the absorption cross section $\sigma^A_{abs}$ and to the differential cross section for quasielastic scattering, as well as to the total cross section, are expressed in terms of the complete differential cross section for all forward dissociation processes. The correction to the quasielastic scattering cross section is very large, while that to the absorption cross section is small. The latter explains why $\sigma^A_{abs}$ is satisfactorily reproduced by the simple multiple scattering theory. In the case of diffraction dissociation a multichannel treatment is necessary in principle. In the case of the simple three-channel problem it is shown how its approximation by the Kolbig-Margolis two-channel problem leads to paradoxes. A new procedure is proposed for analysis of experimental data on diffraction dissociation in nuclei, which permits a lower bound to be obtained on the true cross sections for interaction of diffraction-produced systems with nucleons.

The discussion is arranged as follows. In Sec. 2, which is introductory, we present a multichannel generalization of the multiple scattering theory of Glauber and Sitenko. In Sec. 3 we discuss elastic and quasielastic scattering and the inelastic corrections to the total cross section and to the absorption cross section. In Sec. 4 we formulate a new procedure for analysis of diffraction dissociation in nuclei.

The questions discussed here are comparatively simple, and we were stimulated to write this article only by the fact that up to this time numerous experimental data have been incorrectly analyzed. In a number of cases in which a simplified analysis gave a good quantitative description there was no understanding of the magnitude of the corrections for inelastic screening. The purpose of this work is to fill this gap. A number of remarks of a general nature have been placed in the Conclusion.

## 2. MULTICHANNEL GENERALIZATION OF MULTIPLE SCATTERING THEORY

Let $\hat{f}$ be the matrix of the amplitudes of diffraction transitions forward in a nucleon. The matrix $\hat{f}$ is diagonalized by the states $|\alpha\rangle$, which either are only absorbed or are elastically scattered (the eigenstates of diffraction scattering). For a fixed impact parameter $\mathbf{b}$ the matrix element of the $T$ matrix for the transition between the states $|a\rangle=\sum_\alpha a_\alpha|\alpha\rangle$ and $|c\rangle=\sum_\alpha c_\alpha|\alpha\rangle$ has the form (we are considering all diffraction amplitudes to be pure imaginary)[12,22,23]

$$\langle c|\operatorname{Im}\hat{T}|a\rangle=\sum_\alpha c_\alpha^* a_\alpha t_\alpha. \tag{2}$$

In particular, for elastic scattering

$$\langle a|\operatorname{Im}\hat{T}|a\rangle=\sum_\alpha |a_\alpha|^2 t_\alpha=\langle t_\alpha\rangle, \tag{3}$$

while the complete differential cross section for all diffraction dissociation processes[12,22,23] is

$$d\sigma_{DD}/d^2\mathbf{b}=\langle t^2\rangle-\langle t\rangle^2. \tag{4}$$

If we go over from the impact parameters to the momentum transfers $\mathbf{q}$, then[1] for $\mathbf{q}=0$ we have

$$\left.\frac{d^2\sigma}{d\mathbf{q}^2}\right|_{\mathbf{q}=0}=\frac{1}{16\pi}(\langle\sigma_\alpha^2\rangle-\langle\sigma_\alpha\rangle^2). \tag{5}$$

Here $\sigma_\alpha$ is the cross section for interaction of the eigenstates $|\alpha\rangle$ with the target.

No transitions occur between the eigenstates, so that their interaction with the nucleus is described by the simple multiple scattering theory. It is just for this reason that they are convenient for analysis of the interaction with the nucleus. The Glauber-Sitenko amplitude is well known; we shall recall its derivation, since the method of derivation will be needed by us below. In scattering by a nucleus the phase shifts of scattering by the different nucleons of the nucleus are assumed to be additive[2,3]:

$$\delta_A(\mathbf{b},\mathbf{s}_1,\ldots,\mathbf{s}_A)=\sum_{i=1}^{A}\delta_N(\mathbf{b}-\mathbf{s}_i). \tag{6}$$

Here $\mathbf{s}_i$ are the nucleon coordinates. The additivity of the phase shifts is valid in the eikonal approximation for potential scattering. Equation (6) was justified in field theory by Gribov.[1]

The element of the $T$ matrix corresponding to Eq. (6),

$$t_\alpha{}^A(\mathbf{b},\mathbf{s}_1,\ldots,\mathbf{s}_A)=1-\prod_{i=1}^{A}[1-t_\alpha{}^N(\mathbf{b}-\mathbf{s}_i)], \tag{7}$$

is an operator in the nucleon positions in the nucleus and it is necessary to calculate the matrix element between the initial and final wave functions of the nucleus $|in\rangle$ and $|f\rangle$. We shall assume that the nucleons in the nucleus are uncorrelated and shall neglect the motion of the center of mass of the nucleus. It is assumed that the energy is high and that longitudinal momentum transfers can be neglected. For the problems being considered by us, these approximations are unimportant; the corrections can be taken into account by standard methods (see Ref. 5).

We shall begin the analysis with the case of elastic scattering. Then

$$\langle in|t_\alpha{}^A(\mathbf{b},\mathbf{s}_1,\ldots,\mathbf{s}_A)|in\rangle=1-[1-\langle 1,in|t_\alpha{}^N(\mathbf{b}-\mathbf{s}_1)|1,in\rangle]^A. \tag{8}$$

The matrix element over the one-nucleon wave function $|1,in\rangle$ is

$$\langle 1,in|t_\alpha{}^N(\mathbf{b}-\mathbf{s}_1)|1,in\rangle=\frac{1}{A}\int dz_1\,d^2\mathbf{s}_1\rho_A(\mathbf{s}_1,z_1)t_\alpha{}^N(\mathbf{b}-\mathbf{s}_1)=\frac{\sigma_\alpha}{2A}T_\alpha(\mathbf{b}). \tag{9}$$

Here we have introduced the notation

$$T_\alpha(\mathbf{b})=\int\frac{d^2\mathbf{s}}{2\pi B_\alpha}T(\mathbf{s})\exp\left[-\frac{(\mathbf{b}-\mathbf{s})^2}{2B_\alpha}\right],\quad T(\mathbf{s})=\int dz\rho_A(\mathbf{s},z), \tag{10}$$

where $B_\alpha$ is the slope of the diffraction peak of elastic $\alpha N$ scattering: $d\sigma_{\alpha N}/d\mathbf{q}^2\sim\exp(-B_\alpha\mathbf{q}^2)$. In the ordinary situation $B_\alpha\ll R_A^2$, where $R_A$ is the radius of the nucleus, so that $T_\alpha(\mathbf{b})\approx T(\mathbf{b})$. We shall use this below to simplify the formulas.

Since $A\gg 1$, the power function in Eq. (8) can be replaced by an exponential, i.e.,

$$\langle in|t_\alpha{}^A(\mathbf{b},\mathbf{s}_1,\ldots,\mathbf{s}_A)|in\rangle=1-\exp[-{}^1/{}_2\sigma_\alpha T_\alpha(\mathbf{b})]. \tag{11}$$

Finally the exact amplitude for the coherent diffraction transition $aA \to cA$ has the form[12,22]

$$t_{ac}^A(\mathbf{b}) = \sum_\alpha c_\alpha^* a_\alpha \{1 - \exp[-{}^1/_2\sigma_\alpha T_\alpha(\mathbf{b})]\}. \quad (12)$$

A discussion of Eq. (12) and its comparison with the Kolbig-Margolis approximation can be found in Sec. 4.

Let us consider now incoherent scattering in which the nucleus is excited and summation is carried out over all final states of the nucleus. The amplitude of the transition $aA \to cA^*$ with excitation of the nucleus to the state $|f\rangle$ is

$$f_{ac}^{if}(\mathbf{q}) = \frac{ik}{2\pi} \int d^2\mathbf{b} e^{i\mathbf{q}\mathbf{b}} \sum_\alpha c_\alpha^* a_\alpha \langle f | t_\alpha^A(\mathbf{b}, \mathbf{s}_1, \ldots, \mathbf{s}_A) | in \rangle, \quad (13)$$

and the differential cross section is

$$\frac{d\sigma_{ac}^{if}}{d\mathbf{q}^2} = \frac{1}{4\pi} \int d^2\mathbf{b} d^2\mathbf{d} e^{i\mathbf{q}(\mathbf{b}-\mathbf{d})} \sum_\alpha c_\alpha^* a_\alpha$$
$$\times \sum_\beta c_\beta a_\beta^* \langle in | t_\alpha^A(\mathbf{b}, \mathbf{s}_1, \ldots, \mathbf{s}_A) | f \rangle \langle f | t_\beta^A(\mathbf{d}, \mathbf{s}_1, \ldots, \mathbf{s}_A)^* | in \rangle. \quad (14)$$

For the summation in Eq. (14) over all final states except the ground state, the transition to which is coherent and is given by the amplitude (12), we use the completeness relation[4]

$$\sum_{f \neq in} |f\rangle\langle f| = 1 - |in\rangle\langle in|.$$

This gives

$$\sum_{f \neq in} \langle in | t_\alpha^A(\mathbf{b}, \mathbf{s}_1, \ldots, \mathbf{s}_A) | f \rangle \langle f | t_\beta^A(\mathbf{d}, \mathbf{s}_1, \ldots, \mathbf{s}_A)^* | in \rangle$$
$$= \langle in | t_\alpha^A(\mathbf{b}, \mathbf{s}_1, \ldots, \mathbf{s}_A) t_\beta^A(\mathbf{d}, \mathbf{s}_1, \ldots, \mathbf{s}_A)^* | in \rangle$$
$$- \langle in | t_\alpha^A(\mathbf{b}, \mathbf{s}_1, \ldots, \mathbf{s}_A) | in \rangle \langle in | t_\beta^A(\mathbf{d}, \mathbf{s}_1, \ldots, \mathbf{s}_A)^* | in \rangle. \quad (15)$$

As above, the matrix elements which arise reduce to an average over one-nucleon states. What is new is the matrix element of the form

$$\langle 1, in | t_\alpha^N(\mathbf{b}-\mathbf{s}_1) t_\beta^N(\mathbf{d}-\mathbf{s}_1)^* | 1, in \rangle = \frac{1}{A} \int d^2\mathbf{s}_1 T(\mathbf{s}_1) t_\beta^N(\mathbf{d}-\mathbf{s}_1)^* t_\beta^N(\mathbf{b}-\mathbf{s}_1). \quad (16)$$

We introduce

$$\sigma_{\alpha\beta}(\Delta) = \int d^2\mathbf{s} t_\beta^N(\mathbf{s})^* t_\alpha^N(\mathbf{s}+\Delta). \quad (17)$$

Then Eq. (16) can be written as

$$\langle 1, in | t_\alpha^N(\mathbf{b}-\mathbf{s}_1) t_\beta^N(\mathbf{d}-\mathbf{s}_1)^* | 1, in \rangle = A^{-1} \sigma_{\alpha\beta}(\mathbf{b}-\mathbf{d}) T_{\alpha\beta}(\mathbf{b}-\Delta_{\alpha\beta}), \quad (18)$$
$$\Delta_{\alpha\beta} = (\mathbf{b}-\mathbf{d}) B_\beta / (B_\alpha + B_\beta),$$

where $T_{\alpha\beta}$ is calculated with the slope $B = B_{\alpha\beta} = B_\alpha B_\beta / (B_\alpha + B_\beta)$. In the ordinary situation both $\Delta_{\alpha\beta}$ and $B_{\alpha\beta}$ are small and $T_{\alpha\beta}(\mathbf{b} - \Delta_{\alpha\beta}) \approx T(\mathbf{b})$.

The final exact expression for the differential cross section of incoherent diffraction dissociation $aA \to cA^*$ has the form

$$\frac{d\sigma_{ac}^A(incoh)}{d\mathbf{q}^2} = \frac{1}{4\pi} \int d^2\mathbf{b} d^2\mathbf{d} e^{i\mathbf{q}(\mathbf{b}-\mathbf{d})}$$
$$\times \sum_{\alpha,\beta} a_\alpha a_\beta^* c_\beta c_\alpha^* \exp[-{}^1/_2\sigma_\alpha T_\alpha(\mathbf{b}) - {}^1/_2\sigma_\beta T_\beta(\mathbf{d})]$$
$$\times \{ [\sigma_{\alpha\beta}(\mathbf{b}-\mathbf{d}) T_{\alpha\beta}(\mathbf{b}-\Delta_{\alpha\beta})] - 1\}. \quad (19)$$

This formula is new and has not been derived previously.

## 3. ELASTIC AND QUASIELASTIC SCATTERING IN A NUCLEUS. TOTAL CROSS SECTION AND ABSORPTION CROSS SECTION

Here we shall discuss how the allowance for diffraction dissociation changes the description of elastic and quasielastic scattering by nuclei and also the description of the total cross section and the absorption cross section. For simplification of the formulas we shall neglect $B_\alpha$ in comparison with the radii of the nuclei. In the Glauber-Sitenko approximation

$$\sigma_{tot}^A = 2 \int d^2\mathbf{b} \{1 - \exp[-{}^1/_2\sigma_{tot}^N T(\mathbf{b})]\}, \quad (20)$$

$$\sigma_{abs}^A = \int d^2\mathbf{b} \{1 - \exp[-{}^1/_2\sigma_{tot}^N T(\mathbf{b})]\}^2. \quad (21)$$

The absorption cross section $\sigma_{abs}^A$ is the cross section for interaction with the nucleus with production of new particles. It does not include the cross section for elastic scattering with breakup of the nucleus $aA \to aA^*$, which is designated differently as quasielastic scattering. By integrating over the scattering angles in Eq. (19) it is easy to find that in the one-channel case ($\sigma_{abs}^N = \sigma^N = \sigma_{tot}^N - \sigma_{el}^N$)

$$\sigma_Q^A = \int d^2\mathbf{b} \{\exp[-\sigma_{abs}^N T(\mathbf{b})] - \exp[-\sigma_{tot}^N T(\mathbf{b})]\}, \quad (22)$$

$$\sigma_{abs}^A = \int d^2\mathbf{b} \{1 - \exp[-\sigma_{abs}^N T(\mathbf{b})]\}. \quad (23)$$

In the multichannel case Eq. (12) gives instead of Eq. (20)

$$\sigma_{tot}^A = 2 \int d^2\mathbf{b} \{1 - \langle \exp[-{}^1/_2\sigma_\alpha T(\mathbf{b})] \rangle\}. \quad (24)$$

For calculation of the inelastic correction to the Glauber-Sitenko formula we note that $\sigma_{tot}^N = \langle \sigma_\alpha \rangle$. The correction for inelastic screening is the difference between expressions (24) and (20):

$$\Delta\sigma_{tot}^A = 2 \int d^2\mathbf{b} \{exp[-{}^1/_2\langle\sigma_\alpha\rangle T(\mathbf{b})] - \langle \exp[-{}^1/_2\sigma_\alpha T(\mathbf{b})] \rangle\}$$
$$= -\int d^2\mathbf{b} T^2(\mathbf{b}) \exp[-{}^1/_2\langle\sigma_\alpha\rangle T(\mathbf{b})] \{{}^1/_4\langle\Delta\sigma_\alpha^2\rangle - {}^1/_{24}\langle\Delta\sigma_\alpha^3\rangle T(\mathbf{b}) + \ldots\}$$
$$\approx -4\pi \int d^2\mathbf{b} T^2(\mathbf{b}) \exp[-{}^1/_2\sigma_{tot}^N T(\mathbf{b})] (d\sigma_{DD}/d\mathbf{q}^2)_{q=0}. \quad (25)$$

In the transition to the last line in Eq. (25) we made use of Eq. (5).

The first nonvanishing term of the expansion in $\Delta\sigma_\alpha = \sigma_\alpha - \langle\sigma_\alpha\rangle$ in Eq. (25) coincides[20] with the Karmanov-Kondratyuk correction,[13] which was obtained in perturbation theory in the nondiagonal $f_{ac}$, on the assumption that all $f_{cc}$ are equal to each other: $\sigma_{cN} \equiv \sigma_{aN}$ (otherwise a closed expression for $\Delta\sigma_{tot}^A$ would not have been obtained). From this derivation it is clear why the Karmanov-Kondratyuk approximation is reasonable: allowance for the spread in $\sigma_{cN}$ would correspond to terms $\sim\Delta\sigma_\alpha^3$. While the value of $\langle\Delta\sigma_\alpha^2\rangle$ was determined beforehand by the experimental value $(d\sigma_{DD}/d\mathbf{q}^2)_{q=0}$ in a nucleon target, on the other hand $\langle\Delta\sigma_\alpha^3\rangle$ is actually a new parameter, knowledge of which would be interesting for construction of the parton wave functions of hadrons. Calculation of $\langle\Delta\sigma_\alpha^3\rangle$ for nucleons in the Miettinen-Pumplin model[23] gives

$$\langle(\Delta\sigma_\alpha/\sigma_{tot}^N)^3\rangle \approx -0.15 \langle(\Delta\sigma_\alpha/\sigma_{tot}^N)^2\rangle$$

and inclusion of the term $\langle\Delta\sigma_\alpha^3\rangle$ increases $\Delta\sigma_{tot}^A$ insignificantly: by 6% of $\Delta\sigma_{tot}^A$ for the Al nucleus and by 10% for the Pb nucleus.[2)] Experimentally $\Delta\sigma_{tot}^{nA}$ is actually

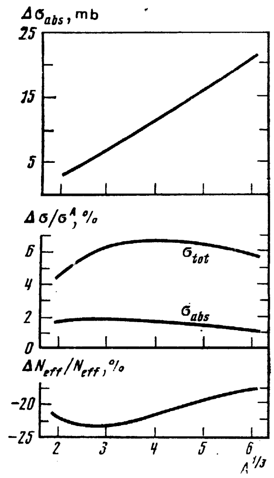


FIG. 1. Inelastic corrections to the total cross section, the absorption cross section, and the effective number of nucleons, as functions of the atomic number of the nucleus.

larger than given by the Karmanov-Kondratyuk formula, but the discrepancy between the experimental results of Refs. 16 and 24 on the $nA$ cross sections are of the same order as the discrepancy between any of the groups of data and the Karmanov-Kondratyuk approximation. It would be very interesting to have precision measurements of $\sigma_{tot}^{nA}$ and a direct analysis of data on the cross sections in terms of the moments $\langle\Delta\sigma_\alpha^n\rangle$. At an energy 200 GeV for all nuclei $\Delta\sigma_{tot}^{nA}\sim(0.05-0.07)\sigma_{tot}^{nA}$ (Fig. 1) and rises slowly with increase of the energy.

We shall calculate now the correction to the absorption cross section. We note at the beginning that in the multichannel case the cross section for elastic scattering by a nucleon is

$$\sigma_{el}{}^N=\sum_\alpha\sum_\beta |a_\alpha|^2|a_\beta|^2\sigma_{\alpha\beta}(0)=\langle\!\langle\sigma_{\alpha\beta}(0)\rangle\!\rangle \tag{26}$$

and that according to Eq. (19) the total cross section for quasielastic scattering is

$$\sigma_Q{}^A=\int d^2\mathbf{b}\langle\!\langle\exp[-{}^1/_2(\sigma_\alpha+\sigma_\beta)T(\mathbf{b})]\{\exp[\sigma_{\alpha\beta}(0)T(\mathbf{b})]-1\}\rangle\!\rangle. \tag{27}$$

Instead of Eq. (23) we then obtain the exact formulas

$$\sigma_{abs}^A=\int d^2\mathbf{b}(1-\langle\!\langle\exp\{-{}^1/_2[\sigma_\alpha+\sigma_\beta-2\sigma_{\alpha\beta}(0)]T(\mathbf{b})\}\rangle\!\rangle), \tag{28}$$

$$\Delta\sigma_{abs}^A=\int d^2\mathbf{b}\{\exp[-(\langle\sigma_\alpha\rangle-\langle\!\langle\sigma_{\alpha\beta}(0)\rangle\!\rangle)T(\mathbf{b})]-\langle\!\langle\exp\{-{}^1/_2[\sigma_\alpha+\sigma_\beta-2\sigma_{\alpha\beta}(0)]T(\mathbf{b})\}\rangle\!\rangle\}. \tag{29}$$

Let us consider again the first nonvanishing term of the expansion in $\Delta\sigma_{\alpha,\beta}$ in Eq. (29). Since

$$\sigma_{\alpha\beta}(0)=\sigma_\alpha\sigma_\beta/8\pi(B_\alpha+B_\beta),$$

the relation between $\sigma_\alpha$ and $B_\alpha$ is also important. We shall consider two extreme cases: global geometric scaling $B_\alpha/\sigma_\alpha$ = const, and the case of universal slopes $B_\alpha=B$ = const. In both cases $\Delta\sigma_{abs}^A$ has the form

$$\Delta\sigma_{abs}^A=-4\pi\int d^2bT^2(\mathbf{b})\exp[-\sigma_{abs}^NT(\mathbf{b})]\left.\frac{d\sigma_{DD}}{dq^2}\right|_{q=0}\left(1-\frac{\beta\sigma_{el}{}^N}{\sigma_{tot}^N}\right)^2, \tag{30}$$

where for geometric scaling $\beta=1$ and for universal slopes $\beta=2$. It is natural to expect that in the general case $1<\beta<2$.

Comparison of $KN$, $\pi N$, and $NN$ scattering shows that the slope of the diffraction peak increases with increase of the cross section, but substantially more slowly than in direct proportion to the cross section. A concrete model of eigenstates for $pp$ scattering was constructed by Miettinen and Pumplin.[23] In this model also the dependence of $B_\alpha$ on $\sigma_\alpha$ is very weak. Therefore we must expect that $\beta$ is close to 2.

The correction (30) for $nA$ interactions is given for $\beta=2$ in Fig. 1. For $\beta=1$ the value of the correction is higher by a factor 1.7. Also given are the ratios $\Delta\sigma^A/\sigma_A^A$ for $\sigma_{tot}^A$ and $\sigma_{abs}^A$. Numerically the value of $\Delta\sigma_{abs}^A$ is small, of the order of 2% of $\sigma_{abs}^A$. The suppression is due both to the factor in the large parentheses in Eq. (30) and to the fact that

$$\exp[-\sigma_{abs}^NT(\mathbf{b})]<\exp[-{}^1/_2\sigma_{tot}^NT(\mathbf{b})].$$

Experimentally $\sigma_{abs}^A$ is usually measured with a statistical accuracy of 0.5%, while the systematic errors are higher, of the order of (1–3)%.[18,19]

For completeness we give the correction to the inelastic scattering cross section $\sigma_{in}^A=\sigma_{tot}^A-\sigma_{el}^A$:

$$\Delta\sigma_{in}{}^A=-2\pi\int d^2\mathbf{b}\exp[-\sigma_{tot}^NT(\mathbf{b})]T^2(\mathbf{b})\left.\frac{d\sigma_{DD}}{dq^2}\right|_{q=0}. \tag{31}$$

It also is small and is numerically close to $\Delta\sigma_{abs}^A$ with $\beta=2$.

Inclusion of terms $\sim\Delta\sigma_\alpha^3$ in Eqs. (30) and (31) is not complicated, but their contributions are small and unimportant for the accuracies of conceivable experiments. The correction $\Delta\sigma_{in}^A$ has been discussed previously by Gaisser *et al.*[25] and by Nam *et al.*[26] in connection with the analysis of data on total cross sections measured in experiments with cosmic rays. It was correctly noted in Refs. 25 and 26 that $\Delta\sigma_{in}^A$ is small. In Ref. 25 $\Delta\sigma_{in}^A$ was found to be $\Delta\sigma_{in}^A(A)=\frac{1}{4}\Delta\sigma_{tot}^A(2A)$, which does not agree with the exact formula (31) either in its dependence on $A$ or in magnitude. In Ref. 26 a more general analysis was made, but an explicit Gaussian parametrization of the nuclear density was used, and the general equation (31) was not derived. In reality cosmic-ray experiments measure not $\sigma_{in}^A$, but $\sigma_{abs}^A$, and use of $\Delta\sigma_{in}^A$ in Refs. 25 and 26 instead of $\Delta\sigma_{abs}^A$ was incorrect. Shabel'skii[27] also has discussed the corrections to $\sigma_{in}^A$ and also found that $\Delta\sigma_{in}^A$ is small. However, he used the two-channel approximation, and the correction to quasielastic scattering and its dependence on the slopes of the diffraction peak were not discussed.

Let us consider now in more detail the quasielastic scattering $aA\to aA^*$. In Eq. (19) we expand the exponential in a series in $\sigma_{\alpha\beta}(\Delta)T(\mathbf{b})$. In the usual case this is a small parameter. The individual terms of the series, since it is found to be of constant sign, can arbitrarily be interpreted as the contributions of $n$-fold elastic scattering by the different nucleons of the nucleus. The important contribution is that from "single" scattering. In view of the inequality $B_{\alpha\beta}\ll R_A^2$, in the expression

$$\exp[-{}^1/_2\sigma_\alpha T(\mathbf{b})-{}^1/_2\sigma_\beta T(\mathbf{d})]$$

we can set $\mathbf{d}=\mathbf{b}$. Then the contribution of single scattering takes on a particularly nice form

$$\frac{d\sigma_Q{}^A}{d\mathbf{q}^2}=\frac{1}{4\pi}\int d^2\mathbf{b}\,d^2\Delta\,e^{i\mathbf{q}\Delta}T(\mathbf{b})\langle\!\langle\sigma_{\alpha\beta}(\Delta)\exp[-{}^1/_2(\sigma_\alpha+\sigma_\beta)T(\mathbf{b})]\rangle\!\rangle$$
$$=\int\frac{d^2\mathbf{b}}{k^2}T(\mathbf{b})\,|\langle f_\alpha(\mathbf{q})\exp[-{}^1/_2\sigma_\alpha T(\mathbf{b})]\rangle|^2. \qquad (32)$$

This expression permits a simple interpretation: the factors $f_\alpha(\mathbf{q})$ correspond to elastic scattering proper, and $\exp[-\frac{1}{2}\sigma_\alpha T(\mathbf{b})]$ describes the absorption of the wave before and after the quasielastic scattering. It is just the amplitude of the wave which is attenuated, as can be seen also from the fact that the exponential contains the total cross section, and not the absorption cross section as would be the case in a statistical attenuation of the intensity. This is a manifestation of the fact that the scale of longitudinal distances $\Delta z = E/\Delta m^2 \gg R_A$, so that the scattering by nucleons with different $z$ is coherent. However, interactions with different impact parameters are incoherent: the intensities of the scattered waves are summed over the impact parameters.

To estimate the inelastic correction to the differential cross section for quasielastic scattering it is sufficient to use Eq. (32). For $\mathbf{q} = 0$ we obtain

$$\left.\frac{d\sigma_Q{}^A}{d\mathbf{q}^2}\right|_{\mathbf{q}=0}\approx\left.\frac{d\sigma_{el}{}^N}{d\mathbf{q}^2}\right|_{\mathbf{q}=0}\int d^2\mathbf{b}T(\mathbf{b})\exp[-\sigma_{tot}^N T(\mathbf{b})]$$
$$\times\left\{1-\frac{d\sigma_{DD}}{d\sigma_{el}}\left[\sigma_{tot}^N T(\mathbf{b})-\frac{1}{4}[\sigma_{tot}^N T(\mathbf{b})]^2\right]\right\},$$
$$\frac{d\sigma_{DD}}{d\sigma_{el}}=\left[\frac{d\sigma_{DD}}{d\mathbf{q}^2}\Big/\frac{d\sigma_{el}{}^N}{d\mathbf{q}^2}\right]_{\mathbf{q}=0}. \qquad (33)$$

A specific feature of quasielastic scattering is that the relative size of the correction is not suppressed by small factors of the form $\exp[-\frac{1}{2}\sigma_{tot}^N\ T(\mathbf{b})]$ or $\exp[-\sigma_{abs}^N T(\mathbf{b})]$, which decrease the contribution of small impact parameters and accordingly of large $T(\mathbf{b})$. For all nuclei the correction decreases the quantity

$$N_{eff}=\left[\frac{d\sigma_Q{}^A}{d\mathbf{q}^2}\Big/\frac{d\sigma_{el}{}^N}{d\mathbf{q}^2}\right]_{\mathbf{q}=0}$$

by 20% (Fig. 1). In the limit of infinitely heavy nuclei the sign of the correction would change, but real nuclei are not sufficiently large and the term in Eq. (33) which is linear in $T(\mathbf{b})$ is dominant. Detailed calculations in terms of the model of Miettinen and Pumplin show that within the diffraction peak the relative size of the correction is a weak function of the scattering angle. This is a manifestation of the previously mentioned smallness of the fluctuations of the slopes of the diffraction peak in comparison with the fluctuations of the total cross sections.

For calculations of all of the cross sections being discussed it is necessary to know the hadronic size of the nuclei. As we have seen, the corrections for inelastic screening decrease all of the cross sections, so that if we analyze data on cross sections without inelastic screening, the nuclear radii are underestimated. Independent information on nuclear radii is given by the study of elastic scattering, in which the location of the first diffraction minimum should be especially sensitive to the radius of the nucleus.

The correction to $d\sigma_{el}^A/d\mathbf{q}^2$ for inelastic screening can easily be calculated by using Eq. (25): the integrand in

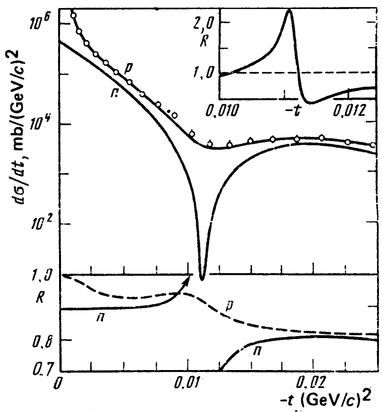


FIG. 2. Differential cross sections of protons and neutrons in carbon nuclei at energy 175 GeV, as functions of the momentum transfer. Also given is the ratio $R$ of the cross section with inclusion of the inelastic correction to the cross section without inclusion of the inelastic correction.

Eq. (25) is a correction to the partial wave. Results of the calculations for the lead nucleus are given in Figs. 2 and 3. We shall discuss first $nA$ scattering. Without taking into account the inelastic correction the first diffraction minimum occurs at $|t| = \mathbf{q}^2 = 0.0111$ (GeV/c)$^2$. Inclusion of the correction, with the same nuclear radius, shifts the minimum to the right by $\Delta t = 3\cdot10^{-4}$, i.e., $\Delta t/|t| = 3\cdot10^{-2}$. In Figs. 2 and 3 $R$ is the ratio of $d\sigma_{el}^A/dt$ with inclusion of the correction to $d\sigma_{el}^A/dt$ without inclusion of the correction. However, in description of the total cross section without inclusion of inelastic screening, if we decrease the nuclear radius, then

$$\Delta t/|t| = -\Delta R_A{}^2/R_A{}^2 \approx -\Delta\sigma/\sigma = 7\cdot10^{-2}.$$

This shows that accurate measurement of $\sigma_{tot}^A$ and $d\sigma_{el}^A/dt$ permits separation of elastic screening effects from changes of the nuclear radius.

It is interesting that the inelastic screening shifts the next diffraction minima in the direction of smaller $|t|$, which is evident from Figs. 2 and 3. However, if we take into account the uncertainties in the shape of the nuclear densities, it is hardly possible to predict re-

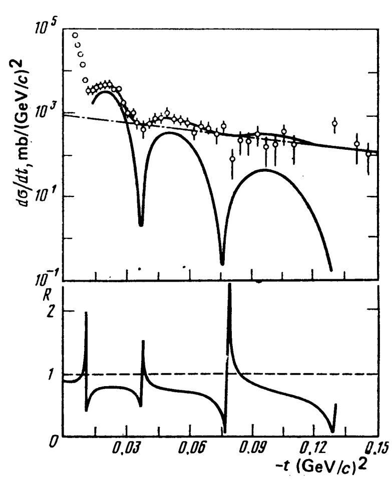


FIG. 3. The same as Fig. 2 but for higher values of momentum transfer.

liably the location of these minima.

In the case of scattering of protons by heavy nuclei it is important to take into account the Coulomb interaction, which is becoming strong. In potential theory the means of taking into account the Coulomb interaction are well known: it is necessary to add the Coulomb phase shift to the nuclear phase shift.[28] In a number of studies[29] this procedure has been applied to scattering of hadrons by nuclei, and proof of it in terms of field theory has been given by Kondratyuk and Kopeliovich.[30] In Figs. 2 and 3 we have given the results of our calculation of $d\sigma_{\text{cl}}^A/dt$ for $p$Pb scattering at 175 GeV. The theory and experiment[31] are in very good agreement. It must be emphasized that in that study Schiz *et al.*[31] analyzed the experimental data quite incorrectly: the interference of the Coulomb and strong amplitudes was neglected without justification. As a result the extrapolation of the differential cross section for nuclear scattering to 0° angle corresponded to a value of $\sigma_{\text{tot}}^A$ which was 1.5 times greater than that measured in direct experiments. Nuclear radii determined in this manner have no quantitative meaning.

In Figs. 2 and 3 we have given the sum of the differential cross sections for elastic and quasielastic scattering, since these two forms of scattering have not been resolved experimentally. The inclusion of the inelastic correction to quasielastic scattering improves the agreement of theory with experiment in the region of dominance of quasielastic scattering, but the accuracy of the existing data is still poor. As Fig. 2 shows, inclusion of the Coulomb interaction makes proton-nucleus elastic scattering comparatively insensitive to the inelastic correction. Therefore the experiments on elastic scattering of neutrons by nuclei being carried out at present by the group from the Institute of Theoretical and Experimental Physics and Freiburg at CERN[32] are very interesting.

## 4. WHAT IS MEASURED IN DIFFRACTION DISSOCIATION IN NUCLEI?

We shall show first that the cross sections for diffraction dissociation in the general case do not permit measurement of the cross section for interaction of diffraction-produced systems with nucleons. We recall that in the Kolbig-Margolis approximation for coherent diffraction dissociation[4] (here and subsequently $\sigma_1 = \sigma_{\text{tot}}^{aN}$, $\sigma_2{}^* = \sigma_{\text{tot}}^{cN}$)

$$t_{ac}{}^A(\mathbf{b}) = \frac{2f_{ac}{}^N(0)}{\sigma_2{}^* - \sigma_1}\left\{\exp\left[-\frac{1}{2}\sigma_1 T(\mathbf{b})\right] - \exp\left[-\frac{1}{2}\sigma_2{}^* T(\mathbf{b})\right]\right\}. \tag{34}$$

Let us consider the simple three-channel problem.[7] Let the physical states arbitrarily designated $|\pi\rangle$, $|3\pi\rangle$, and $|5\pi\rangle$ have the following expansions in terms of the eigenstates $|1\rangle$, $|2\rangle$, and $|3\rangle$:

$$|\pi\rangle = \frac{1}{2}|1\rangle + \frac{1}{\sqrt{2}}|2\rangle + \frac{1}{2}|3\rangle,$$
$$|3\pi\rangle = \frac{1}{2}|1\rangle - \frac{1}{\sqrt{2}}|2\rangle + \frac{1}{2}|3\rangle, \quad |5\pi\rangle = \frac{1}{\sqrt{2}}|1\rangle - \frac{1}{\sqrt{2}}|3\rangle \tag{35}$$

and let $\sigma_{1N} = \sigma_0$, $\sigma_{2N} = 2\sigma_0$, and $\sigma_{3N} = 3\sigma_0$, so that $\sigma_{\pi N} = \sigma_{(3\pi)N} = \sigma_{(5\pi)N} = 2\sigma_0$. Then according to Eq. (12) we obtain

$$t_{\pi\to3\pi}^A = {}^1/_4\{2\exp[-\sigma_0 T(\mathbf{b})] - \exp[-{}^1/_2\sigma_0 T(\mathbf{b})] - \exp[-{}^3/_2\sigma_0 T(\mathbf{b})]\},$$
$$t_{\pi\to5\pi}^A = 2^{-3/2}\{\exp[-{}^3/_2\sigma_0 T(\mathbf{b})] - \exp[-{}^1/_2\sigma_0 T(\mathbf{b})]\}, \tag{36}$$

which has nothing in common with the Kolbig-Margolis formula (34). Thus, in terms of the Kolbig-Margolis formula $t_{\pi\to5\pi}^A$ corresponds to the process in which a system with the cross section $\sigma_1 = 3\sigma_0 = (3/2)\sigma_{\pi N}$ is initially propagated in the nucleus and goes over into a system with an interaction cross section $\sigma_2{}^* = \sigma_0 = \frac{1}{2}\sigma_{(5\pi)N}$. However, an attempt to fit the amplitude $t_{\pi\to3\pi}^A$ to the Kolbig-Margolis formula would generally lead to $\sigma_2^* < 0$.[7] The reason for this is that $t_{\pi\to3\pi}^N = 0$ and the dissociation $\pi\to3\pi$ in the nucleus occurs as the result of the double transition $\pi\to5\pi\to3\pi$; here, although this process is formally of second order in the nondiagonal amplitude $f_{ac}$, in a heavy nucleus $t_{\pi\to3\pi}$ does not have a special smallness in comparison with $t_{\pi\to5\pi}$. This example shows clearly that perturbation theory is inapplicable to diffraction dissociation. In the case of the total cross sections or the absorption cross sections the situation was favorable in that one was calculating a correction to a term in the integrand which was already small, being proportional to $\exp[-\frac{1}{2}\sigma_{\text{tot}}^N T(\mathbf{b})]$ or $\exp[-\sigma_{\text{abs}} T(\mathbf{b})]$.

Let us now consider incoherent diffraction dissociation. In perturbation theory the correct expression for the differential cross section of incoherent diffraction dissociation was obtained by several workers.[33] It has a compact form only for universality of the slopes of the diffraction peak:

$$\frac{d\sigma_{ac}{}^A(incoh)}{d\mathbf{q}^2} = \frac{d\sigma_{ac}{}^N}{d\mathbf{q}^2}\int\frac{d^2\mathbf{b}T(\mathbf{b})}{(\sigma_2{}^* - \sigma_1)^2}\left\{\sigma_2{}^*\exp\left[-\frac{1}{2}\sigma_2{}^* T(\mathbf{b})\right] - \sigma_1\exp\left[-\frac{1}{2}\sigma_1 T(\mathbf{b})\right]\right\}^2. \tag{37}$$

For comparison with Eq. (37), we follow Eq. (32) and separate from the exact expression (19) for the differential cross section the contribution corresponding arbitrarily to single inelastic scattering:

$$\frac{d\sigma_{ac}{}^A(incoh)}{d\mathbf{q}^2} = \frac{1}{4\pi}\int d^2\mathbf{b}T(\mathbf{b})\left|\sum_\alpha c_\alpha{}^* a_\alpha f_\alpha(\mathbf{q})\exp\left[-\frac{1}{2}\sigma_\alpha T(\mathbf{b})\right]\right|^2. \tag{38}$$

In the case of the three-channel problem discussed above it is easy to see that in the general case it is not possible to approximate Eq. (38) by Eq. (37). In the case of strong coupling of the two channels (an example is the dissociation $\pi\to5\pi$) a formula is obtained which is similar to (37) but the cross sections which enter into it are neither $\sigma_{\pi N}$ nor $\sigma_{(5\pi)N}$.

What kind of information is provided by diffraction dissociation in nuclei? Let us introduce the moments

$$M_m = \sum_\alpha c_\alpha{}^* a_\alpha (\sigma_\alpha - \sigma_{tot}^{aN})^m. \tag{39}$$

Then the amplitude (12) and the integrand in (38) can be written in the form of an expansion in these moments:

$$\sum_{\alpha} c_\alpha^* a_\alpha \exp\left[-\frac{1}{2}\sigma_\alpha T(\mathbf{b})\right] = \exp\left[-\frac{1}{2}\sigma_{tot}^{aN} T(\mathbf{b})\right] \sum_{m=1}^{\infty} \frac{(-1)^m}{2^m m!} M_m T(\mathbf{b})^m,$$

$$\sum_{\alpha} c_\alpha^* a_\alpha \sigma_\alpha \exp\left[-\frac{1}{2}\sigma_\alpha T(\mathbf{b})\right] \tag{40}$$

$$= \exp\left[-\frac{1}{2}\sigma_{tot}^{aN} T(\mathbf{b})\right] \sum_{m=1}^{\infty} \frac{(-1)^{m-1}}{2^m m!} M_m T(\mathbf{b})^{m-1} (2m - \sigma_{tot}^{aN} T(\mathbf{b})).$$

The more accurately the differential cross sections for coherent and incoherent diffraction dissociation are measured, the larger the number of moments $M_m$ which can be determined from the dependence of the diffraction dissociation cross sections on the atomic number $A$ of the nucleus. The larger the number of moments $M_m$ which is known, the more severe are the restrictions imposed on the structure of the expansion of a diffraction-produced state in the system of eigenstates.

Traditionally the task of experiments on diffraction dissociation is considered to be the determination of $\sigma_2^*$. In the eigenstate formalism of scattering this is

$$\sigma_2^* = \sum_{\alpha} |c_\alpha|^2 \sigma_\alpha. \tag{41}$$

For a complete experiment—determination of all coefficients $c_\alpha$—it is necessary to know the systems of eigenstates and to have infinitely precise data on diffraction dissociation. In reality one can pose the question of estimates of upper and lower bounds on $\sigma_2^*$. For this purpose it is necessary to find the extrema of the expression (41) with several experimentally determined moments $M_m$, the orthogonality condition $M_0 = 0$, and the normalization $\sum_\alpha |c_\alpha|^2 = 1$ as constraints.

Let us consider what this procedure gives in the case of the diffraction dissociation

$$nA \to (p\pi^-)A, \tag{42}$$

using as the system of eigenstates the Miettinen-Pumplin function.[23] The cross section for interaction of the eigenstates with the nucleons $\sigma_\alpha$ and the coefficients of the expansion of a nucleon in these eigenstates $a_\alpha$ are given in the first two columns of the table. Note the "passive" state with $\sigma_\alpha = 0$. Its appearance is an inherent feature of quark-parton models.[34-36]

The reaction (42) has been studied by Mollet *et al.*[37] in the nonresonance mass region $1.35 \le m_{p\pi} \le 1.45$ (GeV/$c^2$) and at energies 100–300 GeV. Analysis by the Kolbig-Margolis method gave $\sigma_2^* = 27 \pm 3$ mb. Eight nuclear targets were used: Be, C, Al, Ti, Cu, Ag, Ta, and Pb. If these data are analyzed in terms of the moments $M_m$, one obtains $M_1 = 3.5$ mb, $M_2 = 40$ mb$^2$, and $|M_3| < 10^3$ mb$^3$. The error in determination of $M_2$ is of the order of 15 mb$^2$. The data of Ref. 37 are already insensitive to the higher moments.

Any state $|c\rangle$ normalized to unity which reproduces these two numbers $M_1$ and $M_2$ and in addition the orthogonality condition $M_0 = 0$ will describe at the same time also the dissociation cross section (42). Let us consider a state $|c\rangle$ such that $c_1 \approx -a_2 \approx -0.5$, $c_2 \approx a_1 \approx 0.44$, $c_{12} \approx 0.7$, and all remaining $c_i$ are negligibly small. Since $a_{12}$ is small, the condition of orthogonality is satisfied. Then

$$M_1 \approx a_2 a_1 (\sigma_2 - \sigma_1) + a_{12} c_{12} \sigma_{12} \approx 3.7\,\text{mb},$$

and $M_2 = -20$ mb$^2$, i.e., the dissociation cross section is reproduced in all nuclei. For the true cross section $\sigma_2^*$, however, we have from (41)

$$\sigma_2^* = |c_1|^2 \sigma_1 + |c_2|^2 \sigma_2 + |c_{12}|^2 \sigma_{12} = 65\,\text{mb}. \tag{43}$$

Thus, we have constructed an explicit model in which the state $|p\pi\rangle$ with an interaction cross section $\sigma_{(p\pi)N} = 65\,\text{mb} \approx \sigma_{pN} + \sigma_{\pi N}$ has the same diffraction-dissociation amplitude (42) as that given by the Kolbig-Margolis formula with the unphysical cross section $\sigma_2^* = 27$ mb. From the same example it is evident that in essence there is no upper bound on $\sigma_2^*$.

It is obvious that there are infinitely many sets of $c_\alpha$ and states $|c\rangle$ with amplitudes for dissociation in a nucleus which agree within experimental error but which have different $\sigma_2^*$. We can ask: What is the minimum allowed $\sigma_2^*$? If no additional restrictions are imposed on $c_\alpha$, then minimization of the expression (41) gives the solution I listed in the table, which corresponds to $\sigma_2^* = 16$ mb. It is unacceptable for the following reason.

In the quark model it is natural to assign passive states to the constituent quarks themselves.[35] A hadron is intrinsically passive if all the quarks are passive, i.e., for a nucleon of $N_q$ quarks the weight of a passive state $P_0$ is

$$P_0 \approx P_q^{N_q}, \tag{44}$$

where $P_q$ is the weight of the passive state for the constituent quark. For the initial neutron $N_q = 3$ and for the $p\pi$ system we have $N_q = 5$, so that the natural scale of $c_0$ is

$$c_0 \approx a_0^{5/3} = 0.11. \tag{45}$$

The state with the lowest nonzero cross section is naturally associated with the state in which all quarks are passive except one, i.e., for $P_q \ll 1$ it is natural to expect

$$|c_1|^2 \approx N_q P_q^{N_q - 1} (1 - P_q). \tag{46}$$

For the neutron proper, Eq. (46) would give $a_1 \approx 0.6$, and for the $p\pi$ system $c_1 \approx 0.3$. If we use (45) for $c_0$ and impose the restriction $|c_1| < 0.3$, then the minimum permissible cross section $\sigma_2^*$ is immediately increased to 41.6 mb (solution II in the table). For $c_0 = -(a_0)^{5/3}$ one also obtains a solution with $\sigma_2^* = 43.4$ mb (solution III in the table).

The example considered is purely illustrative. Both

TABLE I. Parameters of the expansion in eigenstates of the Miettinen-Pumplin model of the systems described in the text.

| Nucleon | | Solution I | | Solution II | | Solution III | |
|---|---|---|---|---|---|---|---|
| $\sigma_i$, mb | $a_i$ | $c_i$ | $c_i a_i$ | $c_i$ | $c_i a_i$ | $c_i$ | $c_i a_i$ |
| 0.00 | 0.2725 | 0.5755 | 0.1568 | 0.1146 | 0.0312 | −0.1146 | −0.0312 |
| 20.00 | 0.4394 | −0.7793 | −0.3398 | −0.2920 | −0.1283 | −0.2678 | −0.1177 |
| 35.00 | 0.5010 | 0.0062 | 0.0031 | −0.4486 | −0.2248 | −0.3503 | −0.1755 |
| 46.7 | 0.4664 | 0.1354 | 0.0631 | 0.8268 | 0.3857 | 0.8636 | 0.4028 |
| 56.0 | 0.3761 | 0.1518 | 0.0571 | −0.0614 | −0.0231 | −0.1415 | −0.0531 |
| 63.8 | 0.2712 | 0.1263 | 0.0342 | −0.0831 | −0.0226 | −0.1217 | −0.0330 |
| 70.3 | 0.1785 | 0.0902 | 0.0161 | −0.0641 | −0.0114 | −0.0847 | −0.0151 |
| 76.0 | 0.1088 | 0.0580 | 0.0063 | −0.0421 | −0.0046 | −0.0530 | −0.0058 |
| 81.0 | 0.0620 | 0.0343 | 0.0021 | −0.0250 | −0.0015 | −0.0307 | −0.0019 |
| 85.4 | 0.0333 | 0.0189 | 0.0006 | −0.0139 | −0.0005 | −0.0167 | −0.0006 |
| 89.4 | 0.0170 | 0.0099 | 0.0001 | −0.0072 | −0.0001 | −0.0086 | −0.0001 |
| 93.0 | 0.0083 | 0.0049 | 0.0000 | −0.0036 | 0.0000 | −0.0042 | 0.0000 |

in the analysis of Mollet *et al.*[37] by the Kolbig-Margolis method and in our model problem, spin effects were not taken into account. On inclusion of spin the eigenstates used by us are split into subsystems in spin and parity. It is necessary to determine the extrema of $\sigma_2^*$ individually for all spin states of the produced system. Nevertheless it is possible to draw a general conclusion: the anomolously low cross sections $\sigma_2^*$ which follow from the Kolbig-Margolis procedure are a result of the inadequacy of perturbation theory. It is necessary to give up the habit of analyzing results of diffraction-dissociation experiments by the Kolbig-Margolis method, since such $\sigma_2^*$ have no quantitative meaning and only lead one into error. The problem which must be solved is to determine the properties of the system of eigenfunctions of the scattering. The first useful step would be representation of the experimental data in the form of moments $M_m$, which could then be used to check quark-parton models for systems of eigenstates.

## CONCLUSION

We have found the corrections to the absorption cross sections and to quasielastic scattering and have obtained general formulas for description of diffraction dissociation in nuclei, which have been used to replace the Kolbig-Margolis formula (34) for coherent diffraction dissociation and the formula of Tarasov *et al.* (37) for incoherent diffraction dissociation. The problem of determination of $\sigma_2^*$ is not solved by the general formalism. In addition, direct determination of $\sigma_2^*$ turns out to be impossible. However, reduction of all diffraction dissociation in nuclei to determination of $\sigma_2^*$ is payment of unnecessary tribute to conservatism. The moments $M_m$ are a no less important quantity for the theory. Knowledge of them can make possible an improvement of the ideas regarding the quark-parton structure of hadrons. We have shown that even the lower bounds on $\sigma_2^*$ obtained with use of the moments $M_m$ clearly indicate that the true cross sections for interaction of multiparticle systems with nucleons are not small.

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[1)] In the work of Miettinen and Pumplin[23] a factor $1/16\pi$ was lost in this formula.

[2)] Here and below, all numerical estimates are carried out for scattering of 175-GeV nucleons.


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