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A MODEL FOR DIFFRACTIVE PRODUCTION OF PARTICLES IN HIGH-ENERGY COLLISIONS

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The general ideas of the Good and Walker diffraction dissociation mechanism are reviewed and a proposal of a specific realization of these ideas is presented.

The model, developed by Czyż, Kotański and myself, is based on the idea that the interacting particles should be considered as fluctuating objects. During the collision process these virtual fluctuations become real observable states. Thus, the model can be viewed as a generalization of the vector dominance model of ϱ photoproduction.

The model explains all essential features of diffractive production and it provides many important tests, particularly for the diffractive production from nuclei.

1. Good and Walker mechanism

The fact that the absorption of the incident light by the target implies (among other things) the elastic scattering of the wave is a well-known optical phenomenon. Such absorptive (diffractive) elastic scattering is characteristic to the interaction of any wave with the absorbing medium and was actually one of the basic observations proving the wave nature of the light.

The main characteristics of diffractive elastic scattering is that the scattering amplitude is dominantly imaginary and the cross-section depends only very weakly (if at all) on the primary energy. Since all the evidence we have about the elastic scattering of particles at high energies is consistent with these predictions, it is now commonly believed that diffractive scattering is a dominant mechanism also for elastic scattering of particles wave. In this picture the absorption is caused by numerous inelastic interactions appearing in the scattering process. They play an analogous role as the excitation of the absorber caused by the absorbed light (that is the heat production in the absorber). Good and Walker [1] pointed out that the absorption of the incident wave can lead not only to elastic but also to inelastic scattering of the incident wave. They proposed that it is therefore natural to

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expect a similar phenomenon in the scattering of high-energy particles. It was called diffractive dissociation.

The most characteristic feature of the inelastic scattering produced by such a diffractive interaction is that, similarly as elastic diffractive scattering, it should be essentially energy independent. The existence of such inelastic channels with energy-independent cross-sections is now a very well established property of high-energy interactions. It is therefore of considerable interest to find a proper description of the diffractive inelastic scattering.

The basic idea in the argument given by Good and Walker was presented on the example of absorption of the polarized light by an anisotropic absorber. Let the incident light moving in the z-direction be polarized in the direction $\vec{n} = (n_x, n_y)$ (orthogonal, of course, to the incident direction: z). Let this light be absorbed by an absorber. Assume furthermore that absorption depends on the direction of the polarization of the incident wave. One may imagine, *e. g.*, a small Nicol prismat, oriented in such a way that it stops all light polarized in the direction of the x axis and does not affect at all the light polarized in the direction of the y axis.

The wave function of the incident light can be written as a superposition of waves polarized in x and y directions:

$$\psi_n = n_x\psi_x + n_y\psi_y. \tag{1.1}$$

The component ψ_y is unchanged during the scattering. Consequently, there will be no elastically scattered wave polarized in y direction, and the scattered light is fully polarized in x direction. Thus, the scattered light has different polarization from the incident light, and we conclude that the absorption caused the inelastic scattering. To see how big is the cross-section for inelastic scattering, we can analyse polarization of the scattered wave with respect to the direction of polarization of the incident wave ($\boldsymbol{n}$) and the orthogonal direction ($\boldsymbol{n} \times z$). We obtain

$$\varphi_x = n_x\varphi_n + n_y\varphi_{n\times z}. \tag{1.2}$$

The ratio of inelastic to elastic cross-section is thus $(n_y/n_z)^2$.

The example we considered can be easily generalized for more general absorption properties of the target. It is instructive to write down the explicit formulae.

Assume that the target is characterized by two absorption coefficients η_x and η_y which describe the absorption of the light polarized (respectively) in x and y directions.

Furthermore, assume that no direct transition from x-polarized to y-polarized light is possible by the interaction. In other words, x and y directions are eigenvectors of the absorption matrix, and η_x and η_y are its eigenvalues.

Consider now the scattering of polarized light by this target.

The incident wave can be written as

$$\psi_n = n_x\psi_x + n_y\psi_y \tag{1.3}$$

where ψ_x and ψ_y represent the waves polarized in x and y directions. The outgoing wave (after absorption) is

$$\varphi = \eta_x n_x\psi_x + \eta_y n_y\psi_y. \tag{1.4}$$

Thus, the scattered wave is given by

$$\psi_n - \varphi = (1-\eta_x)n_x\psi_x + (1-\eta_y)n_y\psi_y = \\ = \lambda_{\text{el}}\psi_n + \lambda_{\text{inel}}\psi_{n\times z} \tag{1.5}$$

where λ_{el} and λ_{inel} are elastic and inelastic scattering amplitudes. It is easily seen that

$$\lambda_{\text{el}} = n_x^2(1-\eta_x) + n_y^2(1-\eta_y), \tag{1.6}$$

$$\lambda_{\text{inel}} = n_x n_y(\eta_y - \eta_x). \tag{1.7}$$

The formula (1.7) is very instructive. It shows that the absorption will always produce inelastic scattering, unless

(*i*) The incident light is polarized along the directions of an eigenvector of the absorption matrix ($n_x = 0$ or $n_y = 0$).

(*ii*) The absorption coefficients η_x and η_y are equal, that is the absorber is isotropic.

Before we extend the argument presented here to the scattering of elementary particles it may be worth while to emphasize that, although we write explicitly only the wave function of the incident light, we actually implicitly considered also the production of many other final states, responsible for the absorption. In this case these states are the various excitations of the target. They are orthogonal to the states (1.3) and (1.4), and their role is just to provide the absorption of the incident wave.

To see how the ideas presented above can be applied to scattering of particles at high energies, consider the scattering of a nucleon on any target. The initial state consists of nucleon and target. Assume the nucleon has very large momentum in the target rest frame (infinite momentum limit).

Let us consider at the same time other states, containing some extra particles (say pions), all of them also moving very fast (their momenta must be finite fractions of total momentum of the whole group in the target rest frame), and the unexcited target. We will denote such a set of states (including nucleon) by $|\tilde{\lambda}_i\rangle$, $i = 0, 1, 2 \ldots$ Note that the states with extra particles moving slowly in the rest frame of the target or in the CM system of the collision are excluded from the set $|\tilde{\lambda}_i\rangle$[1].

We are interested in the matrix elements

$$\langle\tilde{\lambda}_i|T|\tilde{\lambda}_j\rangle. \tag{1.8}$$

Good and Walker proposed to calculate these matrix elements (in close analogy to the optical example considered above), by assuming that the set of states $|\tilde{\lambda}_i\rangle$ can be

[1] The set of states $|\tilde{\lambda}_i\rangle$ $i = 0, 1, \ldots$ is analogous to the set of different polarization states of light ψ_n and $\psi_{n\times z}$ and the unexcited absorber in the example discussed above. All other states are analogous to the states with excited (heated) absorber. In the scattering of elementary particles the distinction between the two kinds of states is of course not so obvious as in the optical example of Good and Walker. It is hoped, however, that at very high energies the $|\tilde{\lambda}_i\rangle$ states which contain only a very fast group of particles will be better and better separated from other states which are responsible for absorption and which presumably will contain the products of "pionization", *i.e.* slow particles in the over-all CM system.

expanded into another set $|\lambda_i\rangle$ with the property that all states $|\lambda_i\rangle$ when scattered, are only absorbed.

Thus, we consider a set of states $|\lambda_i\rangle$, which are linear combinations of $|\tilde{\lambda}_i\rangle$

$$|\lambda_i\rangle = \sum_j c_{ij}|\tilde{\lambda}_j\rangle, \tag{1.9}$$

$$|\tilde{\lambda}_i\rangle = \sum_j d_{ij}|\lambda_j\rangle \tag{1.10}$$

and which satisfy the condition

$$T|\lambda_j\rangle = (1-\eta_j)|\lambda_j\rangle + \sum_k \gamma_{jk}|\mu_k\rangle \tag{1.11}$$

where all states $|\mu_k\rangle$ are orthogonal to the set $|\lambda_j\rangle$ and, consequently, also to the set $|\tilde{\lambda}_j\rangle$.

For instance in the single pion production in $N-N$ collision $|\tilde{\lambda}_0\rangle$ is a nucleon and $|\tilde{\lambda}_1\rangle$ is a π-nucleon state *etc...* Then $|\lambda^0\rangle$ and $|\lambda_1\rangle$ are linear combinations $|\tilde{\lambda}_0\rangle$ and $|\tilde{\lambda}_1\rangle$. The states $|\mu_k\rangle$ are all other states produced in $N-N$ interaction at high energies, *e. g.* the target excitations, the pionization *etc.* The condition (1.11) expresses the requirement that the states $|\lambda>$, when scattered, are subject only to elastic diffractive scattering (caused by absorption induced by production of the inelastic states $|\mu_i\rangle$) *i. e.* there are no transitions between different $|\lambda_i\rangle$ states. We call the states belonging to the $|\lambda\rangle$ set, the diffractive states.

Using (1.9), (1.10) and (1.11) we obtain

$$\langle\tilde{\lambda}_i|T|\tilde{\lambda}_j\rangle = (1-\eta_j)\delta_{ij} - \sum_k (\eta_k-\eta_j)d_{jk}c_{ki}. \tag{1.12}$$

This is equivalent to the Good and Walker formula.

As emphasized by Good and Walker, the inelastic scattering (which is contained in the second term of (1.12)) is different from zero only if the absorption coefficients for different $|\lambda_i\rangle$ states are different.

As seen from formula (1.12) in order to calculate the matrix elements $\langle\tilde{\lambda}_i|T|\tilde{\lambda}_j\rangle$, it is necessary to specify the absorption parameters η_j and the expansion coefficients c_{ij} or d_{ij} completely unknown at this stage. In fact, the formula (1.12) is so general that it does not seem to have any predictive power.

In order to obtain some possibilities of calculation of matrix elements (1.8) we have to make more specific assumptions, that is to propose a model for expansion coefficients and absorption parameters entering (1.12). An attempt in this direction will be described in the next Section.

2. A model for diffractive dissociation

As remarked in the previous Section, in order to make the Good and Walker idea a practical tool for calculation of the cross-sections, it is necessary to specify

(*i*) the absorption parameters η_i of the diffractive states $|\lambda_i\rangle$ on the target,

(*ii*) the expansion coefficients of the observed real states $|\tilde{\lambda}_i\rangle$ in terms of the diffractive states $|\lambda_i\rangle$.

In this Section I would like to describe a model specifying these parameters which was developed recently by Czyż, Kotański and the author [2].

The difficulty in determination of the absorptive parameters η_i is that the states $|\lambda_i\rangle$ are not the ones we observe in scattering experiments. Consequently η_i cannot in general be obtained from elastic scattering data. However, an important simplification can be obtained by exploiting the experimental evidence that the cross-sections for diffractive production processes are about one order of magnitude smaller than elastic ones. A possible way of including this information into the formalism of Good and Walker is to assume that the transformation from $|\tilde{\lambda}_i\rangle$ set to $|\lambda_i\rangle$ set is not far from unity:

$$c_{ij} = \delta_{ij} + \varepsilon_{ij} \tag{2.1}$$

where ε is small and ε^2 may be neglected. Then

$$d_{ij} = \delta_{ij} - \varepsilon_{ij} \tag{2.2}$$

and the formula (1.12) becomes (to first order in ε):

$$\langle\tilde{\lambda}_i|T|\tilde{\lambda}_j\rangle = (1-\eta_i)\,\delta_{ij} + (1-\eta_i)\,\varepsilon_{ij} - (1-\eta_j)\,\varepsilon_{ij}. \tag{2.3}$$

Thus, the elastic amplitude is

$$\langle\tilde{\lambda}_i|T|\tilde{\lambda}_i\rangle = 1-\eta_i. \tag{2.4}$$

The formula (2.4) shows that the absorption parameters η_i are determined by elastic scattering of real particles, (in this approximation). Obviously, this result greatly simplifies the interpretation and computation of absorption parameters.

The absorption of the initial diffractive state is then easily determined from elastic proton-proton scattering.

A difficulty arises, however, if one tries to determine the absorption of a multiparticle system (*e. g.* πN system). Then, it is necessary to provide a model for elastic scattering of such a system of real particles from the target. The simplest possibility, which we will exploit in this paper, is to assume that this elastic scattering is correctly described by the Glauber model.

This assumption is by no means obvious, but it seems reasonable to exploit its possibilities. The important point to notice at this stage is that, at least in principle, any assumption of this type can be verified by measurements of elastic scattering of the considered system of real particles off a given target.

The inelastic amplitude from formula (2.3) is the difference of the two terms:

$$\langle\tilde{\lambda}_i|T|\tilde{\lambda}_j\rangle = (1-\eta_i)\,\varepsilon_{ij} - (1-\eta_j)\,\varepsilon_{ij}. \tag{2.5}$$

This formula is illustrated in Figure 1, which shows that the amplitude for diffractive production of particles is proportional to the difference between the absorption of the produced particle and the absorption of the incident particle.

It should be remarked that the Glauber model approximation which we propose to use for calculation of the elastic scattering in the final state is not necessarily correct and

perhaps should be substituted by some better approximation. However, the great advantage of such a model is that, after it is introduced, the production matrix element can be expressed through the absorption parameters of real particles and the expansion coefficients ε_{ij}.

Thus, at this stage, only the expansion coefficients are left undetermined. To calculate them, we propose the following physical interpretation of the expansion (1.10).

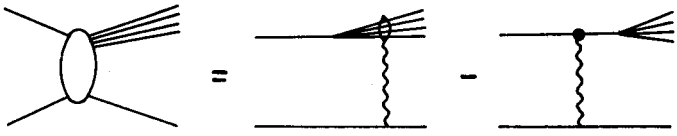

Fig. 1

The incident particle (which subsequently dissociates diffractively) is viewed as a fluctuating object with various fluctuations possible. For instance, a nucleon can fluctuate into π-nucleon, $\pi\pi$-nucleon system *etc.*... We interpret the expansion coefficients ε_{ij} as the probability amplitudes of such fluctuations. Such an interpretation implies an important fact that the properties of the expansion coefficients ε_{ij} are independent of the interaction with the target. There are two important consequences of this:

(a) All internal quantum numbers (except mass and spin) of the dissociated system should be the same as those of the incident particle.

(b) The expansion coefficients ε_{ij} are the same for scattering off any target.

The property (a) follows from the fact that the interaction Hamiltonian responsible for the fluctuation of the incident particle must obey the strong interaction selection rules.

Thus our model provides automatically the selection rule that the diffractive production of particles is possible only if there are no internal quantum numbers exchanged in the t-channel of the reaction.

The property (b) guarantees a very strong predictive power of the model. Namely, it relates the diffractive production on different targets. Thus, the comparison of diffractive dissociation on different targets (especially for various nuclear targets) is of crucial importance for our model, because it would check the "universality" of the expansion (1.10). Incidentally, let us note that the "universality" of ε_{ij} does not imply that the diffraction dissociation is identical for different targets because the absorption coefficients η_j are different.

It remains to give a prescription for calculating the expansion coefficient. As we have already mentioned, according to the experimental evidence the cross-sections for diffractive production processes are about order of magnitude smaller than the elastic scattering, hence we expect ε_{ij} to be small and we tentatively propose to compute them from a perturbation expansion. Obviously this should be considered as a rather preliminary suggestion. It is made just in order to make the simple calculation possible. We expect that the detailed comparison with the data will require probably fairly complicated ε_{ij}, and the important point is the universality property discussed above. To summarize, we have proposed a physical interpretation of the Good and Walker parameters which naturally incorporate

● the smallness of the diffractive production cross-section compared with the elastic one,

● the experimentally established selection rule of no internal quantum numbers exchange in the t-channel of the diffractive process.

Our interpretation has a further advantage in that it suggests a crude way of calculating (or measuring) the Good and Walker parameters.

Furthermore, which is perhaps the most important aspect, it gives a clear physical meaning to the Good and Walker formula and thus allows one to develop the intuition about the diffractive processes in the direction which has not been very much explored in recent years. It remains to hope that this may lead to a better undertanding of diffraction dissociation.

Finally, let me remark that our model is very analogous to the Cheng and Wu model of high energy electrodynamical processes [3]. Actually, the investigation described here was greatly influenced by their work. In particular, their model can also be considered as a particular realization of the general scheme of Good and Walker.

In the next Section I will proceed to develop a little more formal aspects of the model, in particular the formula for the transition matrix elements.

3. Explicit construction of the production matrix elements

Let us consider the coherent production process in which the particle a produces an object b on target T. The object b may be one particle or a group of particles $b_1 \dots b_n$. According to our model the production amplitude is a difference of two contributions (see Figure 1):

$$\mathcal{M}(\Delta, \boldsymbol{k}^{(i)}) = \mathcal{M}^{(b)}(\Delta, \boldsymbol{k}^{(i)}) - \mathcal{M}^{(a)}(\Delta, \boldsymbol{k}^{(i)}), \tag{3.1}$$

$$\mathcal{M}^{(a)}(\Delta, \boldsymbol{k}^{(i)}) = \gamma^{(a)}(\Delta) a_{\mathrm{fl}}(\boldsymbol{k}^{(i)}), \tag{3.2}$$

$$\mathcal{M}^{(b)}(\Delta, \boldsymbol{k}^{(i)}) = \int d\boldsymbol{p}^{(1)} \dots d\boldsymbol{p}^{(n)} a_{\mathrm{fl}}(\boldsymbol{p}^{(i)}) \Gamma^{(b)}(\Delta, \boldsymbol{p}^{(i)}, \boldsymbol{k}^{(i)}), \tag{3.3}$$

where $\Delta = \Sigma \boldsymbol{k}_\perp^{(i)}$ is the momentum transfer to the target, $\boldsymbol{k}^{(i)}$ $(i = 1, \dots, n)$ are the final momenta of the produced particles $b_1 \dots b_n$, a_{fl} is the probability amplitude for the fluctuation, $\gamma^{(a)}$ is the profile function of the particle a on the target T and $\Gamma^{(b)}$ is the profile operator of the set of particles $b_1 \dots b_n$ on the target T.

Since both $a_{\mathrm{fl}}(\boldsymbol{p}^{(1)}, \dots, \boldsymbol{p}^{(n)})$ and $\Gamma^{(b)}(\Delta, \boldsymbol{p}^{(1)} - \boldsymbol{k}^{(1)}, \dots, \boldsymbol{p}^{(n)} - \boldsymbol{k}^{(n)})$ contain the δ-functions which guarantee the conservation of the total 3-momentum of the system during fluctuation, as well as the longitudinal components of the momenta of all particles in the elastic scattering, the formula (3.3) can be integrated over longitudinal components of $\boldsymbol{p}^{(1)}, \dots, \boldsymbol{p}^{(n)}$ and be written in the form which contains only the integration over the transverse momenta

$$\mathcal{M}^{(b)}(\Delta, \boldsymbol{k}^{(i)}) = \int d^2 p_\perp^{(1)} \dots d^2 p_\perp^{(n)} a_{\mathrm{fl}}(\boldsymbol{p}^{(1)}, \dots, \boldsymbol{p}^{(n)}) \times$$

$$\times \gamma^{(b)}(\Delta, \boldsymbol{p}_\perp^{(1)} - \boldsymbol{k}_\perp^{(1)}, \dots, \boldsymbol{p}_\perp^{(n)} - \boldsymbol{k}_\perp^{(n)}). \tag{3.3}$$

As I have already stressed in the previous Section, the most important feature of the formulae (3.1)–(3.3′) is that the dependence on the target enters only the quantities $\gamma^{(a)}$ and $\gamma^{(b)}$, whereas the fluctuation probability amplitude depends only on the incident particle and is entirely independent of the target.

A serious complication arises from the fact that the elastic scattering of the final system $b = b_1, \ldots, b_n$ on the target is expressed in formula (3.3) by the scattering operator $\gamma^{(b)}(\Delta, \boldsymbol{p}_\perp^{(1)}, \ldots, \boldsymbol{p}_\perp^{(n)})$ and not by the elastic scattering amplitude of b which would read

$$\Phi(\Delta) = \langle b|\Gamma|b\rangle = \int d^2p_\perp^{(1)} \ldots d^2p_\perp^{(n)} a_{\mathrm{fl}}(\boldsymbol{p}^{(1)}, \ldots, \boldsymbol{p}^{(n)}) \times$$

$$\times \gamma^{(b)}(\Delta, \boldsymbol{p}_\perp^{(1)} - \boldsymbol{k}_\perp^{(n)}, \ldots, \boldsymbol{p}_\perp^{(n)} - \boldsymbol{k}_\perp^{(n)}) a_{\mathrm{fl}}^*(\boldsymbol{k}^{(1)}, \ldots, \boldsymbol{k}^{(n)}) d^2k_\perp^{(1)}, \ldots, d^2k_\perp^{(n)}). \tag{3.4}$$

Since the $a_{\mathrm{fl}}(\boldsymbol{p}^{(1)}, \ldots, \boldsymbol{p}^{(n)})$ enters the formula (3.3) linearly, the result of (3.3) seems much more sensitive to the details of the fluctuation probability amplitude $a_{\mathrm{fl}}(\boldsymbol{p}^{(1)}, \ldots, \boldsymbol{p}^{(n)})$ than the elastic amplitude (3.4). Therefore the interpretation of the quantities entering (3.3) is fairly complicated, except in the case when the produced object b consists of just one particle.

In such a case the integration in the formula (3.3) is trivial and we obtain

$$\mathcal{M}(\Delta, \boldsymbol{k}) = a_{\mathrm{fl}}(\boldsymbol{k}) \{\gamma^{(b)}(\Delta) - \gamma^{(a)}(\Delta)\}. \tag{3.5}$$

This formula is just the one we met already in the previous Section. The profile functions $\gamma^{(a)}$ and $\gamma^{(b)}$ are normalized as follows

$$\gamma(\Delta = 0) = \frac{\sigma^{\mathrm{tot}}}{2}, \tag{3.6}$$

$$\frac{d\sigma_{\mathrm{elastic}}}{d\Delta^2} = \frac{\pi}{(2\pi)^2} |\gamma(\Delta)|^2. \tag{3.7}$$

I would like to stress once more that the simple formula (3.5) is obtained only if the produced object consists of just one particle. Otherwise both calculations and interpretation are much more complicated. From the formulae (3.5) and (3.6) we see explicitly that if the total cross-sections of particles a and b on a given target are the same, the forward production cross-section of the b particle should vanish. Thus, the immediate consequence of our model and the experiments of the coherent production of unstable systems on nuclei is that the coherently produced objects like A_1, Q and B^* bumps cannot be interpreted as elementary. Indeed, if *e. g.* A_1 is interpreted as a single body without internal degrees of freedom, the formula (3.5) should be valid. Since the experiments on nuclei indicate that $\sigma_{\pi N} = \sigma_{A_1 N}$ [4], we obtain the result that the A_1 bump should not be produced in the πN interactions. This conclusion clearly contradicts the data and the only way out is to reject the formula (3.5), that is to admit that A_1 is not an elementary object. As I already pointed out, the formula (3.5) is valid for any target, provided the object b is elementary.

An interesting insight into the physical meaning of our approach can be obtained

by applying (3.5) to the production off nuclei. In this case the profile functions $\gamma^{(a)}$ and $\gamma^{(b)}$ are given (in the optical approximation) by

$$\gamma^{(a,b)} = \int d^2b[1-e^{-\frac{\sigma_{a,b}}{2}T(b)}] \tag{3.8}$$

where σ_a snd σ_b are the total aN and bN cross-sections and

$$T(b) = A\int_{-\infty}^{+\infty} \varrho(b, z)dz. \tag{3.9}$$

Here A is the mass number and $\varrho(\boldsymbol{r})$ is the density of the nucleus in question. Using (3.8) and (3.5) we obtain

$$\mathcal{M}_A(\Delta) = a_{\mathrm{fl}}(\boldsymbol{k})\int d^2b\{e^{-\frac{\sigma_a}{2}T(b)}-e^{-\frac{\sigma_b}{2}T(b)}\} \tag{3.10}$$

but $a_{\mathrm{fl}}(\boldsymbol{k})$ can be determined from the experiments with free nucleons:

$$a_{\mathrm{fl}}(\boldsymbol{k}) = \mathcal{M}_N(\Delta = 0, \boldsymbol{k})\,[\gamma_N^{(b)}(\Delta = 0)-\gamma_N^{(a)}(\Delta = 0)]^{-1} \tag{3.11}$$

where $\mathcal{M}_N(\Delta = 0)$ is the forward production amplitude off nucleon.

Using (3.6) we obtain

$$a_{\mathrm{fl}}(\boldsymbol{k}) = \mathcal{M}_N(\Delta = 0)\frac{2}{\sigma_b-\sigma_a}. \tag{3.12}$$

Substituting (3.11) into (3.10) we get finally

$$\mathcal{M}_N(\Delta) = \mathcal{M}_N(\Delta = 0)\frac{2}{\sigma_b-\sigma_a}\int d^2b\{e^{-\frac{\sigma_a}{2}T(b)}-e^{-\frac{\sigma_b}{2}T(b)}\}. \tag{3.13}$$

The interesting point about the formula (3.13) is that this is just standard optical model formula [5] used for analysing the experiments of production from nuclei. Thus, in this particular case our model gives results analogous to the standard approach. This I find a rather pleasant feature because it indicates that the assumptions we made seem fairly reasonable.

To conclude, we proposed a rather well-defined scheme of calculating the diffractive production processes. It incorporates automatically the gross-features of the data. It remains to be seen, however, whether it will survive the detailed comparison with experiment.

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