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HIGH ENERGY PHOTOPRODUCTION OF PSEUDOSCALAR MESONS

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Talk given at the
International School of Elementary Particle Physics
Basko-Polje, Yugoslavia
September 1971

Ref.TH.1433-CERN
30 November 1971

INTRODUCTION

Low energy photoproduction may be studied in an almost model-independent way by means of phase shifts, but high energy photoproduction has to be studied in terms of models simply because phase shift analysis becomes impossible at sufficiently high energies.

Since good data first became available a large number of models has been proposed to "explain" the data. Unfortunately, none of these distinguish themselves as appearing to be definitely more correct than the others, although it seems that the models that have done best so far are of the Regge pole plus Regge cut types. The cuts being either absorptive or more phenomenological. However, also these models have met difficulties¹⁾ and should not be taken too seriously.

Recognizing these facts the following discussion will be to a large extent qualitative, and we will try to extract consequences of the models that can be understood without calculations and that are common to as many models as possible. Below we list some of the models that have been tested on photoproduction. Some of them exist in several versions and some are incomplete in the sense that they try to describe only a limited amount of data. Together with the list are given some references, but this list is not claimed to be complete.

- 1) Strong absorption²⁾ (Michigan)
- 2) Mixed model³⁾
- 3) Weak absorption^{3),4)} (Argonne)
- 4) Other absorptive models⁵⁾⁻¹¹⁾
- 5) Multiple scattering¹²⁾
- 6) Pseudo-model¹³⁾
- 7) Phenomenological cuts¹⁴⁾⁻¹⁷⁾
- 8) Complex poles¹⁸⁾
- 9) Conspiracy models¹⁹⁾⁻²²⁾
- 10) Electric Born term²³⁾
- 11) Veneziano model²⁴⁾⁻²⁸⁾
- 12) Fixed poles²⁹⁾
- 13) Vector dominance model

The order in which the models appear does not imply any judgement on their relative correctness.

Most of what will be said in the following sections is fairly easily adapted to models 1)-6). However, model 3) in its original form is unable to fit the data ³⁾.

Complex poles will be neglected essentially because of their lack of predictive power. The conspiracy models and the electric Born term models seem unlikely to be correct for reasons to be given later.

There have also been some attempts to use the Veneziano model. In principle, it probably makes more sense to study the Veneziano four-point function in πN photoproduction than in πN scattering because of the absence of the Pomeron. The problem of introducing spin in the Veneziano model still seems hard and is possibly connected to the fact that the models, that have been tried, suffer from having parity doublets and a pion conspirator. Satellite terms are needed to connect the forward and backward directions ²⁸⁾. This is also considered undesirable.

The linearity of the unitarity relations for photoproduction implies that fixed poles are not forbidden. The fact that (at least for charged photoproduction) $d\sigma/dt \sim s^{-2}$ has lent support to the assumption that fixed poles rather than moving (Regge) poles are dominant. Essentially, the arguments that will be given later against the electric Born term model also apply to fixed poles.

The vector dominance model will hardly be mentioned because it is the topic of another lecture at the school.

GENERAL CONSIDERATIONS

The basic ingredients in all models trying to describe two-body or quasi two-body processes is the assumed dominance of t channel exchanges ^{*)}. Therefore, the first step in any conventional analysis of two-body scattering processes is to establish the possible exchanges.

^{*)} Of course, the duality picture tells us that we gain extra information if s channel resonances are included in the analysis, but at the moment we will ignore this.

A list of the processes so far studied experimentally is given in Table 1. All processes listed (above) below the double line proceed via (non) strange exchanges.

As is seen from the Table the most extensive experimental studies have been done on photoproduction of pions off nucleons. The same statement is no doubt true for the theoretical efforts. Most of the features of these processes cannot teach us very much without some labour on the spinology of the problem. It is, therefore, interesting to note that one can get essential pieces of information from the data on $\gamma N \rightarrow \pi \Delta$ and $\gamma N \rightarrow K \Sigma$ without elaborate models.

If we believe in any type of absorbed Regge model, it implies that each amplitude can be written symbolically as $R + RP$, where R represents the contribution from a pure Regge pole and RP represents the Regge-Pomeron cut induced by the absorption. Allowing the Pomeron to be exchanged in the initial and final states does not change the quantum numbers of the t channel exchanges with the exception of parity as we shall see later.

In first order perturbation theory we have to consider simple one-particle exchange diagrams. These will give rise to factors $\langle \pi | j_\mu(0) | B \rangle$ where B represents the exchanged particle. The electromagnetic current may be written as a sum of an isovector and an isoscalar current ^{*)}

$$j_\mu(x) = j_\mu^{(s)}(x) + j_\mu^{(v)}(x) \quad (1)$$

Under G parity

$$j_\mu^{(s)} \rightarrow -j_\mu^{(s)}$$

$$j_\mu^{(v)} \rightarrow j_\mu^{(v)}$$

(2)

^{*)} Here we have ignored the possibility that the current may also have an isotensor part. This point has been discussed by Professor Rollnik in his lectures.

and from this follows that $I^G = 0^+, 2^+$ are not allowed as quantum numbers for the exchanged particles.

If we look at the simple one-particle exchange diagram for $\gamma N \rightarrow \pi \Delta$ we see that only $I=1,2$ may be exchanged, but no particle with $I=2$ is known, so, if we have an R+RP model $I=1$. Looking at the Clebsch-Gordan coefficients, we find

$$\begin{aligned} \frac{d\sigma}{dt} (\gamma p \rightarrow \pi^- \Delta^{++}) &\sim 3 \sum_{i=1}^8 |M_i^{(s)} - M_i^{(v)}|^2 \\ \frac{d\sigma}{dt} (\gamma n \rightarrow \pi^+ \Delta^-) &\sim 3 \sum_{i=1}^8 |M_i^{(s)} + M_i^{(v)}|^2 \\ \frac{d\sigma}{dt} (\gamma n \rightarrow \pi^- \Delta^+) &\sim \sum_{i=1}^8 |M_i^{(s)} - M_i^{(v)}|^2 \\ \frac{d\sigma}{dt} (\gamma p \rightarrow \pi^+ \Delta^0) &\sim \sum_{i=1}^8 |M_i^{(s)} + M_i^{(v)}|^2 \end{aligned}$$

(3)

where M_i are the eight-independent t channel amplitudes and the indices s and v denotes photon isospin.

This means that with the assumptions of t channel dominance and no $I=2$ exchange it follows that

$$\frac{d\sigma}{dt} (\gamma p \rightarrow \pi^- \Delta^{++}) = 3 \frac{d\sigma}{dt} (\gamma n \rightarrow \pi^- \Delta^+) \quad (4)$$

and

$$\frac{d\sigma}{dt} (\gamma n \rightarrow \pi^+ \Delta^-) = 3 \frac{d\sigma}{dt} (\gamma p \rightarrow \pi^+ \Delta^0) \quad (5)$$

Disregarding deuteron effects

$$R_{\pi^+} \equiv \frac{d\sigma}{dx} (\gamma d \rightarrow \pi^+ \Delta N) / \frac{d\sigma}{dx} (\gamma p \rightarrow \pi^+ \Delta^0) = 4 \quad (6)$$

$$R_{\pi^-} \equiv \frac{d\sigma}{dx} (\gamma d \rightarrow \pi^- \Delta N) / \frac{d\sigma}{dx} (\gamma p \rightarrow \pi^- \Delta^{++}) = 4/3 \quad (7)$$

Correcting for deuteron effects brings down the ratio R_{π^+} a bit but not sufficiently to be in agreement with the experiment ³⁰⁾. See Fig. 1.

Requiring the absence of $I = \frac{3}{2}$ exchanges in the $\gamma N \rightarrow K \Sigma$ processes gives

$$\frac{d\sigma}{dt} (\gamma d \rightarrow K^+ \Sigma N) / \frac{d\sigma}{dt} (\gamma p \rightarrow K^+ \Sigma^0) = 3 \quad (8)$$

which again is a bit bigger than the experimental values ³¹⁾.

Doing the analysis with the "exotic" $I = 2$ and $I = \frac{3}{2}$ exchanges included, it can be seen that a contribution of about 10-20% of exotic exchanges in the amplitudes is sufficient to account for the data. The principal explanations for the observed discrepancies are:

1. Regge-Regge cuts;
2. s channel effects;
3. exotic particle exchanges;
4. tensor current;
5. deuteron effects;
6. background effects.

The first three explanations accept that there is an effect, while the two last ones cast some doubt on the results. The iso-tensor current explanation is impossible because the Clebsch-Gordan coefficients are so that it cannot be distinguished from the iso-scalar part. If any one of the first three explanations is right it is a bit disturbing to note that sum rules following from the requirement of absence of exotics have recently been tested ³²⁾ and found always to be satisfied within errors for purely hadronic processes above 4-5 GeV. If one believes in the vector dominance model there should be no reason why the sum rules should not be satisfied also in photoproduction processes.

The Glauber theory is normally used to extract neutron target data from deuteron data. The theory has been tested in a large number of cases and seems to be true to a quite good accuracy. This makes also explanation 5. unlikely.

The possibility that the experimental results (at least for $\gamma N \rightarrow \pi \Delta$) may not be entirely above question is mainly due to the fact that the cross-sections have to be extracted from data with background that may not be the same in all charge modes.

It would be very useful if the experiments were done at several energies. If RR cuts are in operation they would probably show up in the energy dependence. At this point it should be mentioned that a backward peak has been observed in $K^-p \rightarrow K^-p$ scattering³³⁾. This process can proceed only by "exotic" exchanges in the u channel. A theoretical calculation using RP cuts³⁴⁾ is in rough agreement with the data, but it is not clear if the predicted energy dependence will be correct.

Before we proceed to look in greater detail at the other processes, we should bear in mind the lessons we have learned from the preceding, namely that there may be contributions of the order of 10-20% in the amplitudes that may not be explained by any R+RP model.

ABSORPTION TYPE MODELS

Concepts like Regge-Pomeron cuts or Regge-Regge cuts are very important in our understanding (or lack of understanding) of two-body or quasi two-body processes, so a short review of some basic ideas may be useful.

Qualitatively speaking, the procedure is to Reggeize in the t channel and then cross to the s channel where the absorption is performed. For the time being let us assume that the Reggeization and crossing have been done. The s channel Regge pole amplitudes may then be written as $M^R_{\lambda' \mu'; \lambda \mu}$ where λ, μ and λ', μ' are the helicities of the incoming and outgoing particles respectively. For later use we make the following definitions:

$$n \equiv |(\lambda' - \mu') - (\lambda - \mu)| \quad (9)$$

$$x \equiv |\lambda - \lambda'| + |\mu - \mu'| - n \quad (10)$$

In order to get a feeling of the significance of these variables let us look at the partial wave expansion of the amplitudes

$$M_{\lambda'\mu'; \lambda\mu} = \sum_J (2J+1) a_{\lambda'\mu'; \lambda\mu}^J d_{\lambda'-\mu'; \lambda-\mu}^J(\theta) \quad (11)$$

The d^J functions contain factors $\sin^n(\theta/2)$ and, therefore, $M(\theta=0)=0$ if $n \neq 0$. Let us also neglect the problems of unequal mass kinematics so that $t=0$ corresponds to the forward direction. Then, if the Regge poles factorize in the s channel (which we shall make plausible later) we see that there is a factor $(\sqrt{-t})^{|\mu - \mu'|}$ associated with the upper vertex of a simple one-Reggeon-exchange diagram, and a similar factor $(\sqrt{-t})^{|\lambda - \lambda'|}$ associated with the lower vertex which give altogether a factor $(\sqrt{-t})^{n+x}$. Consequently, since $x \geq 0$ a Regge pole amplitude will be non-zero in the forward direction only when $n=0$ and $x=0$, while kinematics alone demands only $n=0$. Already from this one sees immediately that all t channel pole contributions have to vanish at $t=0$ in pseudoscalar meson photoproduction simply because

$$n+x = |\lambda_\gamma - \lambda_\pi| + |\lambda_{N_i} - \lambda_{N_f}| = 1 + |\lambda_{N_i} - \lambda_{N_f}| > 0$$

Let us see what happens if we include absorption. According to the Sopkovich prescription³⁵⁾, each partial wave amplitude of the Regge pole amplitude M_ℓ^R is transformed in the following way

$$M_\ell^{abs} = (S_\ell^el)^{1/2} M_\ell^R (S_\ell^el)^{1/2} \quad (12)$$

S_ℓ^{el} is the S matrix element for elastic scattering in the state with angular momentum ℓ . We have not distinguished between initial and final state elastic scattering and all helicity indices have been suppressed.

Putting

$$S^{\ell} = 1 + i\rho M^{\ell} \quad (13)$$

where ρ is a simple kinematical factor we get

$$M_l^{abs} = M_l^R + i M_l^R M_l^{\ell} \rho \quad (14)$$

The last term now represents the partial waves of the RP cut. Since the elastic scattering is predominantly imaginary and positive, it follows from Eq. (14) that the interference between the cut and the pole terms will be essentially destructive.

It is customary to make the impact parameter representation which means

$$\begin{aligned} l &\rightarrow qr - \frac{1}{2} \\ \sum_{l=0}^{\infty} &\rightarrow \int_0^{\infty} q dr \end{aligned} \quad (15)$$

and

$$d_{\lambda\mu}^J(\theta) \approx f_n(r\sqrt{-t}) \quad (16)$$

Parametrizing

$$M^{\ell} \sim q^2 e^{\Lambda t} \quad (17)$$

one gets

$$M_l^{\ell} \sim \frac{1}{A} e^{-\frac{l^2}{4q^2\Lambda}} \quad (18)$$

which is a Gaussian in l . This means that the lowest partial waves are more absorbed than the higher partial waves.

Following Ref. 2) we will be a bit more precise and assume

$$M_{\lambda'\mu';\lambda\mu}^{\mathcal{L}} = 4 q^2 \sigma_{tot} (i + \rho) e^{A t/2} \delta_{\lambda\lambda'} \delta_{\mu\mu'} \quad (19)$$

where $\rho = \text{Re } M(0) / \text{Im } M(0)$. Using the impact parameter representation and doing a somewhat lengthy calculation one ends up with a formula for the cut in the form of a convolution integral

$$M_{\lambda'\mu';\lambda\mu}^{\mathcal{C}} = - \frac{\sigma_{tot}}{4\pi} (1 - i\rho) e^{A t/2} \times \int_{-\infty}^0 \frac{dt'}{2} e^{A t'/2} I_n(A \sqrt{t t'}) M_{\lambda'\mu';\lambda\mu}^{\mathcal{R}}(t') \quad (20)$$

where

$$I_n(z) = (-i)^n f_n(i z) \quad (21)$$

is real for real z .

Let us define

$$F_n(z) \equiv I_n(z) / z^n \quad (22)$$

then $F_n(0)$ is finite and different from 0. From this it is easy to see some of the most characteristic features of the absorptive Regge cut models. Remember that at some fixed energy we may write

$$M^{\mathcal{R}}(t) = (-t)^{\frac{1}{2}(n+x)} f(t) \quad (23)$$

where $f(t)$ contains all residue functions, choosing mechanism factors, signature factors and propagators. Inserting in Eq. (20) the result is

$$M_n^c \sim e^{At/2} (\sqrt{-t})^n \int_{-\infty}^0 \frac{dt'}{2} (-t')^{n+x/2} e^{\lambda t'/2} F_n(\lambda \sqrt{t t'}) \quad (24)$$

The factor $(\sqrt{-t})^n$ in front of the integral is exactly the required kinematical factor. $F_n(z)$ is a slowly varying function with z . Since $e^{At'/2} f(t')$ decreases very rapidly with increasing $-t'$ it is clear that the essential part of the integral will come from the small $-t'$ region and, therefore, depend very crucially on the value of $n+x/2$. Clearly, the effect of absorption will be big in $n=0$ amplitudes, less in $n=1$ amplitudes and small in $n=2$ amplitudes. It is also obvious that the effect will be especially dramatic in amplitudes of the type $(n,x)=(0,x>0)$ which will be non-vanishing at $t=0$. (Remember that all pole contributions to such amplitudes have to vanish at $t=0$.)

The Michigan group ²⁾ argues that in the formula for M^c one has not taken into account the fact that there may be inelastic intermediate states induced by the same exchanges. Because of the energy phase relation such states will add coherently. For this reason the replacement $M^c \rightarrow \lambda M^c$ is made with $\lambda \geq 1$ real to take care of the inelastic intermediate states.

When λ gets so big that the lowest partial waves change sign, one characterizes the situation by saying that the partial waves have been over-absorbed. If there is no over-absorption in an $(n,x)=(0,x \geq 0)$ amplitude there will be, in general, two dips in the modulus of the amplitude simply because the cut will dominate at very small $-t$ and at large $-t$ and the pole will dominate in between. When the over-absorption is made sufficiently large the two dips collapse to one. This is, in fact, what happens in the Michigan model ²⁾ for the π exchange in charged π photoproduction from nucleons. The large over-absorption is "necessary" to fit the data at $t=0$ but has the secondary effect of giving a cut that dominates at large t .

In spite of the relative successes of the absorptive type models and their intuitive appeal, it should be emphasized that there is no fundamental theory behind them, and consequently, the absorption may be performed in other equally well (or badly) justified ways.

KINEMATICS FOR $\gamma N \rightarrow 0 \frac{1}{2}^+$

For our analysis it is convenient to construct the t channel parity conserving helicity amplitudes ³⁶⁾. For simplicity we shall write

$$M_{\lambda_\gamma - \lambda_\pi; \lambda_{\bar{N}} - \lambda_N}^{(t) \pm} \equiv M_{\lambda_\gamma \lambda_\pi; \lambda_{\bar{N}} \lambda_N}^{(t) \pm} \quad (25)$$

where M^+ and M^- represent parity conserving helicity amplitudes with contributions from natural $[\bar{P} = (-1)^J]$ and unnatural $[\bar{P} = (-1)^{J'}]$ parity states, respectively.

By analyzing the quantum numbers of the $\bar{N}N$ system one easily finds to which amplitudes the various poles contribute. A listing is given in Table 2. Doing the crossing ³⁷⁾ and keeping only terms of leading order in s we get the following results:

$$M_{0-\frac{1}{2}; 1\frac{1}{2}}^{(s)} \sim M_{10}^{(t)-} + M_{11}^{(t)+} \quad (n=0)$$

$$M_{0\frac{1}{2}; 1-\frac{1}{2}}^{(s)} \sim M_{10}^{(t)-} - M_{11}^{(t)+} \quad (n=2)$$

$$M_{0\frac{1}{2}; 1\frac{1}{2}}^{(s)} \sim M_{10}^{(t)+} + M_{11}^{(t)-} \quad (n=1)$$

$$M_{0-\frac{1}{2}; 1-\frac{1}{2}}^{(s)} \sim M_{10}^{(t)+} - M_{11}^{(t)-} \quad (n=1)$$

(26)

Here some kinematic factors irrelevant for our purposes have been neglected.

The t channel parity conserving helicity amplitudes are supposed to factorize and from Eq. (26) it is clear that also the pole contributions to the s channel amplitudes will factorize.

From what we know about kinematics and absorptive cuts we see that in general

$$M_{0-1/2; 11/2}(t=0) \neq 0 \quad (n, x) = (0, 2) \quad (27)$$

but

$$M_{01/2; 1-1/2}(t=0) = 0 \quad (n, x) = (2, 0) \quad (28)$$

Translated into t channel language Eq. (28) reads

$$M_{10}^-(t=0) = M_{11}^+(t=0) \quad (29)$$

which is often referred to as the conspiracy condition because it implies that at $t=0$ the natural and unnatural parity contributions are not independent of each others. The absorptive cuts automatically give non-trivial solutions to Eq. (29) and they are, therefore, said to be self-conspiring.

It is easily seen from Eq. (26) that the $\pm$ nature of the amplitudes M_{10}^+ and M_{11}^- are preserved after the absorption, but this is not the case for the amplitudes M_{10}^- and M_{11}^+ .

Let us establish that the asymmetry measured with a plane-polarized beam can distinguish between natural and unnatural parity exchanges. The polarization vectors for photons of helicities ± 1 are

$$\epsilon_{\pm 1} = \frac{1}{\sqrt{2}} (\mp 1, -i, 0) \quad (30)$$

and those representing plane-polarized photons

$$\epsilon^{\perp} = \frac{i}{\sqrt{2}} (\epsilon_{-1} + \epsilon_{+1})$$

and

$$\epsilon^{\parallel} = \frac{1}{\sqrt{2}} (\epsilon_{-1} - \epsilon_{+1}) \quad (31)$$

Defining the asymmetry as

$$\Sigma = \frac{d\sigma_{\perp} - d\sigma_{\parallel}}{d\sigma_{\perp} + d\sigma_{\parallel}} \quad (32)$$

and using Eq. (26) and the parity relations

$$M_{\lambda_{\pi}\lambda_N;\lambda_Y\lambda_N}^{(s)} = -(-1)^{\lambda_Y - \lambda_N - \lambda_{\pi} + \lambda_N} M_{-\lambda_{\pi} - \lambda_N; -\lambda_Y - \lambda_N}^{(s)} \quad (33)$$

one obtains ³⁸⁾

$$\Sigma = \frac{|M_{10}^+|^2 + |M_{11}^+|^2 - |M_{10}^-|^2 - |M_{11}^-|^2}{|M_{10}^+|^2 + |M_{11}^+|^2 + |M_{10}^-|^2 + |M_{11}^-|^2} \quad (34)$$

In a very similar manner one may find the asymmetry with a polarized target, A, and the polarization, P, of the recoiling nucleon.

$$A = \frac{2J_m [M_{10}^+ M_{11}^{+*} + M_{10}^- M_{11}^{-*}]}{|M_{10}^+|^2 + |M_{11}^+|^2 + |M_{10}^-|^2 + |M_{11}^-|^2} \quad (35)$$

$$P = \frac{2J_m [M_{10}^+ M_{11}^{+*} - M_{10}^- M_{11}^{-*}]}{|M_{10}^+|^2 + |M_{11}^+|^2 + |M_{10}^-|^2 + |M_{11}^-|^2} \quad (36)$$

A glance at Table 2 tells us that only A_1 exchange may contribute to M_{11}^- for charged π photoproduction, and for π^0 photoproduction there is no candidate to contribute to M_{11}^- . We also remember that if the elastic scattering is diagonal in the helicities, then also the absorption will not generate contributions to M_{11}^- . This means that in a quite general way A_{π^0} is equal to P_{π^0} . This may also be approximately true in charged photoproduction because the A_1 and its associated Pomeron cut will not interfere with other poles or their Pomeron cuts. Consequently, the A_1 exchange will constitute some sort of common background to π^- and π^+ photoproduction which certainly has to be small because of the smallness of the π^-/π^+ ratio.

Especially, if one finds $A_{\pi^0} \neq P_{\pi^0}$ one would probably have to ascribe the difference to RR cuts, or an off-diagonal Pomeron contribution to a RP cut. However, the difference between A_{π^0} and P_{π^0} cannot be very big because A, P and Σ are not independent of each other. To see this we note that

$$\begin{aligned} A + P &\sim 2Jm M_{10}^+ M_{11}^{+*} \\ A - P &\sim 2Jm M_{10}^- M_{11}^{-*} \\ 1 + \Sigma &\sim |M_{10}^+|^2 + |M_{11}^+|^2 \\ 1 - \Sigma &\sim |M_{10}^-|^2 + |M_{11}^-|^2 \end{aligned} \quad (37)$$

Since

$$\begin{aligned} |M_{10}^+ \pm i M_{11}^+|^2 &\geq 0 \\ |M_{10}^- \pm i M_{11}^-|^2 &\geq 0 \end{aligned} \quad (38)$$

it follows that

$$\begin{aligned} |A - P| &\leq 1 - \Sigma \\ |A + P| &\leq 1 + \Sigma \end{aligned} \quad (39)$$

The data ^{39),40)} for Σ for $\gamma p \rightarrow \pi^0 p$ (Fig. 2) show that for small $|t|$ Σ_{π^0} is not far from 1 and thus $A \approx P$ no matter how big RR contributions may be. Evidently, the best place to look for a difference between A and P is in the bottom of the dip at about $t = -0.5 \text{ GeV}^2$.

In $\gamma p \rightarrow \pi^+ n$ both ⁴¹⁾⁻⁴³⁾ Σ and ⁴⁴⁾ A have been measured and assuming $A = P$, we can make a comparison between $|A|$ and $\frac{1}{2}(1 + \Sigma)$. The comparison has been given in Fig. 3.

It is interesting to note that for $t \approx -0.1 \text{ GeV}^2$ the bound is almost saturated which can only come about if the M_{10}^+ and M_{11}^+ amplitudes are orthogonal to each other.

ISOSPIN FOR π N PHOTOPRODUCTION

The isospin decomposition which is most suitable when one considers t channel exchanges is given in Ref. 45) and quoted here for completeness

$$M(\gamma p \rightarrow \pi^+ n) = \sqrt{2} M^{(-)} + \sqrt{2} M^{(0)}$$

$$M(\gamma n \rightarrow \pi^- p) = -\sqrt{2} M^{(-)} + \sqrt{2} M^{(0)}$$

$$M(\gamma p \rightarrow \pi^0 p) = M^{(0)} + M^{(+)}$$

$$M(\gamma n \rightarrow \pi^0 n) = -M^{(0)} + M^{(+)}$$

(40)

CONSTRUCTION OF MODELS

Photoproduction of π^0

By now we have established the general formalism and may start looking at models. Probably the simplest case is π^0 photoproduction, so let us try to build a model from scratch and as far as possible let the experiments tell us what ingredients we need.

As is seen from Table 2, the well-established particles that may be exchanged are all spin one particles. The simple Born term with elementary particles exchanged will, therefore, give a differential cross-section independent of energy. Experimentally, there is a fairly rapid decrease with energy³⁹⁾, (Fig. 4) and hence the necessity to Reggeize.

The general shape of the differential cross-section ³⁹⁾, Fig. 5, gives a strong hint that the process is dominated by ω exchange. The ω exchange is supposed to contribute dominantly to amplitudes that do not change the nucleon helicity, i.e., $n=1$ amplitudes. From pure kinematics it then follows that there will be a forward dip, and, furthermore, this dip will remain after absorption. Also the dip at $t = -0.5 \text{ GeV}^2$ agrees nicely with the Regge picture which predicts a dip at the nonsense wrong signature point $\alpha_\omega(t) = 0$ due to the vanishing of the signature factor at this point. In the Michigan model ²⁾ the dip comes about as the result of interference between the pole and the cut term, but the arguments leading to ω exchange dominance are independent of a particular dip mechanism.

Our model will predict a nice clean dip in the forward direction. Experimentally ⁴⁶⁾, however, the differential cross-section rises very steeply as one approaches the extreme forward direction. This is due to the exchange of a γ ; the Primakoff effect. Because of the factor $1/t$ in the photon propagator this term becomes dominant at sufficiently small $|t|$ in spite of the smallness of the electromagnetic coupling of the γ to the proton. Notice that $t=0$ is just outside the physical region. Given the π^0 lifetime there are no unknowns in the calculation of the Primakoff effect. Since the virtual γ couples to the electric charge essentially in the nucleon helicity non-flip amplitudes it follows that $n=1$ for the amplitudes where the γ exchange is important. This again means that also the Primakoff effect has to vanish in the exact forward direction. The photon has spin one and mass zero. Therefore, the one-photon exchange has $\alpha = 1$ in Regge language. Nevertheless, the Primakoff spike will continue to grow with energy, a fact that can easily be checked from the parametrization, in Eq. (41), and the fact that $t_0 \approx -m_\mu^4/s^2$. Because the width of the Primakoff spike is proportional to t_0 it is seen to shrink very rapidly, with energy. However, the integral under the peak grows like $\ln s$. In π^0 photoproduction this will probably be impossible to detect because of the narrowness and rapid shrinkage of the spike. However, in photoproduction of η the Primakoff peak should be detectable up to much higher energies because of the big mass of the η .

Though it may be difficult to observe the Primakoff peak directly at high energies one might make some interesting observations of the γ exchange through its interference with the hadronic part of the exchanges which is dominant a bit away from the forward direction. This interference term is proportional to

$$\frac{-t+t_0}{t} \propto R_2 \omega \approx -\propto R_2 \omega$$

because of the smallness of t^0 . Therefore, the interference term will behave more or less as a simple decreasing (with t) background, but which is more slowly falling with energy than the purely hadronic exchange contributions. Already at 5-6 GeV/c the interference term may constitute about 15% of the differential cross-section near the maximum at $t = -0.1$. From this it is quite clear that the Primakoff effect should not be left out in realistic calculations. Clearly, if the Primakoff effect is included in the analysis it makes the effective α of the hadronic exchanges smaller if the interference is constructive and bigger if the interference is destructive. A recent and accurate determination³⁹⁾ of the total effective α shows that it is embarrassingly low (Fig. 4), a possible indication of the interference being destructive. On the other hand, the assumption of destructive interference may create difficulties for the $-t$ dependence, but to the author's knowledge no analysis has yet been carried out with such assumptions.

Since our model so far only has contributions from natural parity exchanges and only to the $n=1$ amplitudes we immediately predict that $\Sigma = 1$ even if absorption is used. Since Σ_{π^0} has been measured^{39),40)} (Fig. 2) we know that this is not true, so we need something in the two other amplitudes. The candidates are either ρ or B or both. This type of contribution must also be there because of the known difference between π^+ and π^- photoproduction off nucleons. The B would make $\Sigma \neq 1$ simply by being an unnatural parity object and the ρ exchange would make $\Sigma \neq 1$ through its ρP cut which has both natural and unnatural parity parts. If the B trajectory has anything close to the universal slope one should be able to tell which mechanism is in operation through the energy dependence. If B exchange dominated there would be a fairly rapid approach to $\Sigma = 1$, while if ρP is dominant one would expect a very slow energy dependence. Another point is that if B exchange is dominant then the polarization will be roughly proportional to the BP cut which again should show up in the energy dependence of the polarized target asymmetry.

Obviously our model by now will predict a ratio

$$R_{\pi^0} \equiv d\sigma(\gamma n \rightarrow \pi^0 n) / d\sigma(\gamma p \rightarrow \pi^0 p) = 1$$

apart from the very forward direction where the absence of the Primakoff peak in $\gamma n \rightarrow \pi^0 n$, makes the ratio smaller. R_{π^0} has been measured ⁴⁷⁾ and found to be less than one. This is not surprising since both the magnetic and the electric coupling of the ω and ρ mesons to the $\bar{N}N$ are not zero at the poles. Therefore, we have to introduce a small contribution of the ω to the $n=0$ and $n=2$ amplitudes, and also a contribution of the ρ to the $n=1$ amplitudes.

We therefore arrive at the following Regge pole parametrization for photoproduction of neutral pions from protons.

$$M_{0-1/2; 1/2}^R = t (\gamma_{11}^\omega \omega + \gamma_{11}^\rho \rho + \gamma_{11}^B B)$$

$$M_{0 1/2; 1-1/2}^R = t (-\gamma_{11}^\omega \omega - \gamma_{11}^\rho \rho + \gamma_{11}^B B)$$

$$M_{0 1/2; 1 1/2}^R = \sqrt{-t+t_0} (\gamma_{10}^\omega \omega + \gamma_{10}^\rho \rho + \frac{2 \gamma_{\pi^0} s}{-t})$$

$$M_{0-1/2; 1-1/2}^R = M_{0 1/2; 1 1/2}^R$$

(41)

Here

$$t_0 \approx -m^2 \mu^4 / s^2$$

$$\omega = \Gamma(1-\alpha_\omega) (s/s_0)^{\alpha_\omega} \zeta(\alpha_\omega)$$

$$\rho = \omega(\alpha_\omega \rightarrow \alpha_\rho)$$

$$B = \Gamma(-\alpha_B) (s/s_0)^{\alpha_B} \zeta(\alpha_B)$$

$$\zeta(\alpha) = \begin{cases} (1 - e^{-i\pi\alpha}) \\ i e^{-i\frac{\pi}{2}\alpha} \end{cases}$$

(Argonne model)

(Michigan model)

and $\gamma\pi^0$ is simply related to the π^0 lifetime. The rest of the γ 's are residue functions containing possible α and t dependent factors. The factor $\Gamma(-\alpha_B)$ occurs when the B is taken to be exchange degenerate with the π in which case the B does not have to choose nonsense at $\alpha_B=0$.

The signature factor in the Michigan model does not vanish at $\alpha_B=0$, so to avoid a pole there there has to be a change of the factor $\Gamma(-\alpha_B)$. For instance, $\Gamma(-\alpha_B) \rightarrow \Gamma(\frac{1}{2}(1-\alpha_B))$. The final step to put the model to work is then to perform the absorption according to some prescription, for instance Eq. (20), or to introduce cuts in some more phenomenological way.

The presence of the ρ (or B), in the $n=0$ and $n=2$ amplitudes is seen dramatically in the photoproduction of charged pions where ρ quantum number (or again possibly B quantum number) exchange is responsible for the smallness of the π^-/π^+ ratio.

As no particular structure is seen in this ratio around $t \sim -0.5 \text{ GeV}^2$ one may conclude that the effective ρ (or B) contribution has no structure in the $n=0$ plus $n=2$ amplitudes. Since the ρP cut from these amplitudes has an unnatural parity part, one expects $d\sigma''/dt$ in π^0 photoproduction to be structureless around $t \sim -0.5 \text{ GeV}^2$. This is indeed the case as seen from Fig. 6. From what we have learned the dip structure in $\pi^0 p$ photoproduction arises from the $n=1$ amplitude which has only natural parity exchange. Since the data ⁴⁷⁾ for R_{π^0} indicate a weak minimum around $t \sim -0.5 \text{ GeV}^2$ this is most easily interpreted as an enhancement of $n=1$ amplitudes relative to $n=0$ and $n=2$ amplitudes in going from $\pi^0 p$ to $\pi^0 n$ photoproduction. This immediately tells us that

$$\Sigma \pi^0 n > \Sigma \pi^0 p \quad (42)$$

By similar arguments one finds

$$|A_{\pi^0 n}| < |A_{\pi^0 p}| \quad (43)$$

Again these conclusions are not necessarily correct for very small t because of the absence of the Primakoff effect in the $\pi^0 n$ case so that $\Sigma \pi_{on}$ will go nicely to zero in the forward direction while $\Sigma \pi_{op} \rightarrow 1$ for the smallest angles attainable because of finite resolutions. Preliminary data ⁴⁸⁾ for $A \pi_{op}$ show that it is quite large and negative, in general agreement with most models.

η photoproduction

Because η is an isoscalar and π^0 an isovector, the role of the isoscalar and isovector parts of the current are reversed in the cases of η and π^0 photoproduction.

Using SU_3 and the vector dominance model we have for the upper vertex ⁴⁹⁾

$$g_{\gamma\eta\omega} / g_{\gamma\eta\rho} = 1/3 \quad (44)$$

while for π^0 photoproduction

$$g_{\gamma\pi\omega} / g_{\gamma\pi\rho} = 3 \quad (45)$$

so that in η photoproduction the importance of ρ exchange is enhanced a factor 9 relative to π^0 photoproduction. On the general, semi-experimental grounds given previously, we then expect the η photoproduction to have little or no structure around $t = -0.5 \text{ GeV}^2$. This is indeed the case experimentally ^{50),51)}. In actual models this lack of structure comes about in a natural way in the Michigan model, while the NWS zero from weak cut models can be filled by the B meson exchange, but the energy dependence of the two types of models is quite different because the B trajectory is low lying. The DESY and SLAC data ⁵⁰⁾ (Fig. 8) are in best agreement with no B exchange, while the Cornell data ⁵¹⁾ agree with the B exchange model. A detailed discussion of these points can be found in Refs. 3) and 4).

Maybe it is not surprising that the ρ behaves like an irregular object in view of the very big difficulties that simple models have had in fitting πN CEX data. In fact, in a recent paper ¹⁾ the amplitudes have been reconstructed from existing data, and the only model consistent with these amplitudes is a $\rho + \rho'$ model ⁵²⁾.

Because the ρ contribution introduces a lot of unnatural parity via the absorption, we expect Σ_γ to be substantially smaller than

$$\Sigma_{\pi^0}.$$

$\pi^\pm$ photoproduction

Typical data for these processes are shown in Figs. 9-13. The analysis is made difficult by the many different ways of explaining the forward spike ²³⁾. The explanation chosen unfortunately has consequences also at larger angles.

The principal explanations are: destructive interference of the pole with a background provided by

1. cuts;
2. conspiring complex poles;
3. direct channel terms;
4. unknown background; or
5. the π trajectory is supposed to conspire with some other trajectory.

The last explanation can probably be ruled out because factorization then would demand ⁵³⁾ that there is a forward dip in

$$\pi^+p \rightarrow \rho^0 \Delta^{++}$$

whereas a peak has been observed ⁵⁴⁾. Using the explanation of the direct channel effects as the background, one immediately admits defeat in $np \rightarrow pn$ which has a spike ⁵⁵⁾ similar to the one seen in charged pion photoproduction. In that case there is no direct channel resonances to help ^{*)}. Anyhow, it is necessary to multiply the Born term with a quite strongly t dependent form factor ⁵⁶⁾ in order that the cross-section shall not become completely unreasonable at $-t > \mu^2$. The meaningfulness of such a procedure may, however, be questionable.

The complex poles clearly can fit the data, but the physical interpretation of such fits is difficult, so we shall concentrate on the cut explanation.

*) Of course, there is the π^0 in the u channel which might play the same role as the nucleon in photoproduction, but many good arguments can be given against such an explanation.

Let us start by considering only π exchange and call the pion contribution to the $n=0$ amplitude

$$M_{\pi}(t) \sim \frac{t}{t-\mu^2} \quad (46)$$

so that $M_{\pi}(0)=0$. Approximating $d_{\frac{1}{2},\frac{1}{2}}^I(\theta)$ by $P^{\ell}(\cos\theta)$ we obviously must have

$$\sum_{\ell} (2\ell+1) a_{\pi}^{\ell}(s) = 0 \quad (47)$$

The highest partial waves will be negative and depend very little on what form factors one uses for the π exchange, while the lowest partial waves will be positive and strongly dependent on form factors, but in such a way that

$$\sum_{\text{positive } a_{\ell}} (2\ell+1) a_{\pi}^{\ell}$$

does not depend much on the form factor. A figure displaying $(2\ell+1)a^{\ell}$ as a function of ℓ for Reggeized one-pion exchange in np CEX is shown in Fig. 14. What the absorption model does is to subtract something from the lowest partial waves and as a result one gets a non-vanishing cross-section at $t=0$. However, if $\lambda_{\pi}=1$, there is still too much left of the positive part of the partial waves, i.e., the cross-section is too small. Therefore, the Michigan model increases λ_{π} until the forward differential cross-section is correct. As is seen from Fig. 14, the lowest partial waves are now negative and not very small and the sum of the lowest 10 or so partial waves is roughly zero. Clearly, any model that wants to fit the data must have a sum of the lowest partial waves that is small. But, if this is achieved by a cancellation like in the Michigan model, the result will be that the amplitude becomes big again at larger t . In pole cut language it means that the pole will cancel the cut at some small $-t$, typically $-t \approx \mu^2$ but the cut contribution will be bigger than the pole contribution at larger $|t|$, and, finally become so big that it dominates the theoretical cross-section which is bigger than the experimental one. This phenomenon may be taken as a possible indication that the cut extends further out in partial waves than in present absorption type models.

If one assumes that it is the electric Born term ²³⁾ that interferes with the pure one-pion exchange, one effectively subtracts out a very large s wave which then is sufficient to achieve the necessary cancellation of the rest of the low partial waves. However, the low partial waves are now very big and if no form factor is used one will again get a much too large cross-section at larger angles.

It seems, therefore, that a reasonable model should have all low partial waves small not to be in conflict with the large angle data, and the absorption should extend to higher partial waves to be in agreement with the small angle data ^{*) 57)}.

Clearly, the question of how the π P cut behaves at larger t is very important because an incorrect π P cut will distort the other contributions, especially the A_2 contributions to the n=0 and n=2 amplitudes that can hardly be distinguished from the π P cut except through their energy dependence and corresponding phases. Therefore, one should not consider present day fits as anything final.

All models that want to fit the data have to have a very small n=0 amplitude at $|t| > \mu^2$. At larger angles the double flip amplitude must be dominant at least in π^+ photoproduction because of the large π exchange contribution to this amplitude which is not much changed by absorption. Around $t = -0.25$ the ratio $R = d\sigma(\gamma n \rightarrow \pi^- p) / d\sigma(\gamma p \rightarrow \pi^+ n)$ is small and, therefore, there has to be a strong interference between the $\rho + B$ types of exchanges and the $\pi + A_2$ types of exchanges in the n=2 amplitude. Since the absorption is unimportant in this amplitude one may think in terms of pure Regge poles. Taking ρ, A_2 to be strongly exchange degenerate the phase of the $\rho + A_2$ contribution may be $e^{-i\pi\alpha}$ in the π^+ case and 1 in the π^- case ^{**) 58)}. This makes possible the strongly destructive interference with the π exchange contribution that we need to explain R. At the nonsense wrong signature point of the ρ

*) After this talk was given we have received a preprint ¹¹⁾ by Worden who has successfully used an approach like the one suggested here.

**) A slightly different result is obtained assuming the absence of exotics in $K^* N \rightarrow KN$ and using SU_3 and the vector dominance model.

obviously the $n=2$ amplitudes for π^+ and π^- production are roughly the same. This means that also the $n=0$ plays an important role in explaining R but the interferences between the various poles and their cuts are now quite complicated. The cuts in the $n=0$ amplitude are proportional to the corresponding pole strengths, and since the natural parity contributions have opposite signs in the $n=0$ and $n=2$ amplitudes, while the unnatural parity contributions have the same sign, we predict $R > 1$ for $t=0$ and $R=1$ at about $t = -\mu^2$. Experimentally, this seems to be well satisfied, (Fig. 10).

In the Michigan model the ρ contribution does not vanish at $\alpha_\rho = 0$ and, therefore, the link between the ρ and A_2 contribution becomes weaker than in a strong exchange degeneracy model. Also in the Michigan model the details of the various interferences are quite complicated. The interference between the π exchange and the πP cut almost extinguishes the unnatural parity contributions close to $t = -\mu^2$. Accordingly, we expect $\Sigma \pi_+ \approx 1$ at this point which is indeed observed ⁴³⁾.

The asymmetry with a polarized target has been measured at 5 GeV/c and 16 GeV/c and is large and negative ⁴⁴⁾. Obviously, therefore, the $n=1$ amplitudes are non-zero. Candidates to contribute are the $\rho + A_2$ which we expect to contribute since both πN CEX and $\pi N \rightarrow \eta N$ do not vanish in the forward direction.

It is easy to see from Eqs. (26) and (35) that the polarization gets its main contributions from the interference between πP and the $\rho + A_2$ contributions to the $n=1$ amplitudes. At small angles these amplitudes will be almost orthogonal to each other and this permits the polarization to become quite big. However, the phase of the $(\rho + A_2)$ contribution will rotate faster with t (how much faster depends on details of the models) than the πP cut and therefore one expects the polarization to become small again. The most naïve guess is that it would be zero at an angle somewhat larger than that which corresponds to the NWS zero of the ρ contribution. Experimentally it is still sizeable at $\alpha_\rho(t) = 0$.

The dip mechanism in the Michigan model and also Harari's model ⁵⁸⁾ may be said to be kinematical or geometrical since the absence or presence of a dip is supposed to depend only on the spin structure of the amplitudes. The Argonne type of models may be said to have dynamical dips. The presence of dips depending only on the particles exchanged and not on the spin states. Harari ⁵⁹⁾ has recently proposed a test of the two types

of dip mechanism by comparing π^0 photoproduction to π^0 electroproduction. If the dip mechanism is of geometrical nature, the dip should go away with increasing $|q^2|$ for the virtual photon.

Since good phase shift analysis exists for low energy photoproduction, one may use finite energy sum rules (FESR) or continuous moment sum rules (CMSR) to obtain predictions for high energies ^{13),60)}. This approach is a check on the duality idea, and secondly it may simplify the fitting procedure only to fit the FESR and then predict the high energy results, which, if in agreement with data, would give us confidence in the high energy model we were testing.

Such a test has been carried out for the very forward direction for π^+ photoproduction ¹³⁾. The result was that, at present, the FESR could not distinguish between a conspiracy model and an absorptive cut model. However, one may expect the FESR to become a more sensitive test as more accurate phase shifts become available ^{*)}.

Photoproduction of K mesons

The processes best studied are

$$\gamma p \rightarrow \begin{cases} K^+ \Lambda \\ K^+ \Sigma^0 \end{cases}$$

The particles that may be exchanged are K , K^* and K^{**} analogous to π , ρ and A_2 in charged pion photoproduction. Naïvely, one might then expect a shape similar to that which is seen in charged photoproduction with the proper corrections for the π and K mass differences. Experiments ⁶¹⁾ (Fig. 15), however, show a quite different behaviour with a dip in the forward direction, a maximum at about $t = -0.2 \text{ GeV}^2$ and then a smooth fall off with increasing angle.

A very simple but nice analysis has been performed in Ref. 62). K exchange is neglected, and since a dip is observed this indicates $K^* - K^{**}$ dominance in the $n=1$ amplitudes. From the Λ / Σ^0 ratio at $|t| > 0.2 \text{ GeV}^2$ F/D ratios for the baryon non-flip amplitudes are obtained in good agreement with other data. The Λ / Σ^0 ratio at $t=0$ is determined by the $n=0$ amplitude and is supposedly dominated by

*) Several new FESR tests have been carried out in Ref. 11).

$P(K^* + K^{**})$ cuts. Since the strength of the cuts are proportional to the strength of the poles one can also deduce the F/D ratio for the baryon flip amplitude; correct to the extent one can neglect K exchange. The F/D ratio thus obtained is in agreement with other data.

Measurements with a polarized target ¹⁴⁾ has been done at SLAC without distinguishing between the Λ and the Σ^0 . The result is a quite large and negative polarization. This looks a bit embarrassing to the $K^* + K^{**}$ model that one may expect to give only a small polarization. However, in other processes such a model can give large polarizations ⁶³⁾. It might be that the K exchange plays a role in getting the polarization correct, but no actual attempt at fitting the polarization is known.

For a more detailed discussion, see Ref. 64).

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	$d\sigma/dt$	Pol. beam	Pol. target
$\gamma p \rightarrow \pi^+ n$ $\gamma n \rightarrow \pi^- p$ $\gamma p \rightarrow \pi^0 p$ $\gamma n \rightarrow \pi^0 n$	x x x x	x x x 	x x
$\gamma p \rightarrow \gamma p$ $\gamma n \rightarrow \gamma n$	x 		
$\gamma p \rightarrow \pi^+ \Delta^0$ $\gamma n \rightarrow \pi^+ \Delta^-$ $\gamma p \rightarrow \pi^- \Delta^{++}$ $\gamma n \rightarrow \pi^- \Delta^+$ $\gamma p \rightarrow \pi^0 \Delta^+$ $\gamma n \rightarrow \pi^0 \Delta^0$	x x x x 	x x 	
$\gamma p \rightarrow K^+ \Lambda$ $\gamma n \rightarrow K^0 \Lambda$	x 		x
$\gamma p \rightarrow K^+ \Sigma^0$ $\gamma n \rightarrow K^+ \Sigma^-$ $\gamma p \rightarrow K^0 \Sigma^+$ $\gamma n \rightarrow K^0 \Sigma^0$	x x x 		x
$\gamma p \rightarrow K^+ Y^0$ $\gamma n \rightarrow K^+ Y^-$ $\gamma p \rightarrow K^0 Y^+$ $\gamma n \rightarrow K^0 Y^0$	x x 		

TABLE 1 Table of pseudoscalar photoproduction processes.

The crosses indicate that the experiment has been done. The line between the measured target asymmetry for $\gamma p \rightarrow K^+ \Sigma^0$ and $p \rightarrow K^+ \Lambda$ indicates that Λ and Σ^0 were not distinguished from another.

Isospin index $(\lambda\mu)^\pm$	(-)	(0)	(+)
$(10)^+$	A_2	ρ	ω
$(10)^-$	π	B	—
$(11)^+$	A_2	ρ	ω
$(11)^-$	A_1	—	—

Table 2 Pole contributions to the
various t channel parity
conserving helicity amplitudes.

FIGURE CAPTIONS

- Figure 1 : The ratios $(\gamma d \rightarrow \pi^+ \Delta N)/(\gamma p \rightarrow \pi^+ \Delta^0)$ (open circles) and $(\gamma d \rightarrow \pi^- \Delta N)/(\gamma p \rightarrow \pi^- \Delta^{++})$ (closed circles) versus $(-t)^{\frac{1}{2}}$ at 16 GeV/c. From Ref. 30).
- Figure 2 : Σ_{π^0} at 3, 4 and 10 GeV/c. From Ref. 39).
- Figure 3 : Comparison of $|A_{\pi^+}|$ (open circles) at 16 GeV/c, [Ref. 44] to $\frac{1}{2}(1 + \Sigma_{\pi^+})$ (closed circles) at 12 GeV/c, [Ref. 42].
- Figure 4 : α_{eff} for $\gamma p \rightarrow \pi^0 p$. From Ref. 39).
- Figure 5 : $d\sigma/dt$ for $\gamma p \rightarrow \pi^0 p$. From Ref. 39).
- Figure 6 : $d\sigma^\perp/dt$ (closed circles) and $d\sigma''/dt$ (open circles) at 6 GeV/c. From Ref. 39).
- Figure 7 : Ratio $(\gamma n \rightarrow \pi^0 n)/(\gamma p \rightarrow \pi^0 p)$. From Ref. 47).
- Figure 8 : $d\sigma/dt$ for $\gamma p \rightarrow \eta p$. From Ref. 3). Data from Ref. 50).
- Figure 9 : $d\sigma/dt$ for $\gamma p \rightarrow \pi^+ n$. From Ref. 23).
- Figure 10 : Ratio $(\gamma n \rightarrow \pi^- p)/(\gamma p \rightarrow \pi^+ n)$. From Ref. 23).
- Figure 11 : Σ_{π^+} from Ref. 42).
- Figure 12 : A_{π^+} from Ref. 44).
- Figure 13 : Σ_{π^-} from Ref. 65).
- Figure 14 : Partial wave amplitudes times $(2J+1)$ for the real part of the unabsorbed Reggeized π non-flip and double-flip amplitudes compared to the same for the Michigan model for the process $np \rightarrow pn$. From Ref. 57).
- Figure 15 : $k^2 d\sigma/dt$ versus $\sqrt{-t}$ for $\gamma p \rightarrow K^+ \Sigma^0$ and $p \rightarrow K^+ \Lambda$. From Ref. 62). Data from Ref. 61).

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11

Fig. 12

Fig. 13

Fig. 14

Fig. 15