CM-P00057199

THEORY OF STRONG INTERACTIONS OF ELEMENTARY PARTICLES
IN THE GEV REGION

L. Van Hove

(Invited talk, International Symposium on Electron and
Photon Interactions at High Energies,
Hamburg, June 8 - 12, 1965)

This lecture attempts to review the main aspects of strong interaction processes in the high energy region where cross-sections show a smooth energy dependence suggestive of an asymptotic behaviour. Due to a steady increase of accurate and reliable experimental data the general properties of this asymptotic region manifest themselves more and more clearly. In absence of a true theory, the theoretical analysis is essentially of a phenomenological character. Either it correlates data on the basis of a few general principles and assumptions, or it fits them by means of tentative theoretical models. Despite its modest scope, this type of analysis appears to be quite successful in developing gradually a general picture of strong interactions at high energies.

I. ELASTIC SCATTERING AT SMALL MOMENTUM TRANSFER

We discuss first elastic scattering processes $A+B \rightarrow A+B$ (A, B are hadrons) in the region of c.m. momentum transfers $\Delta = \sqrt{-t}$ defined by

$$0 < -t \lesssim 1(\text{GeV}/c)^2. \quad (\text{I.1})$$

Good data are available for a number of reactions, and the general features appear to be the same for all. For laboratory momentum $p_L \gtrsim 6 \text{ GeV}/c$ a good description is obtained by means of an asymptotic series expansion of the scattering amplitude containing two or three terms. We give it the general form

$$T(s,t) \underset{s \rightarrow +\infty}{\sim} \sum_{j=1,2,\dots} c_j(t) s_t^{\alpha_j(t)} (\ln s_t)^{\beta_j(t)} \quad (\text{I.2})$$

2.

s is the square of the c.m. energy, and for later convenience the expansion is taken to be in decreasing powers of the quantity

$$s_t = s + \frac{t}{2} - m_A^2 - m_B^2. \quad (I.3)$$

$\alpha_j(t)$, $\beta_j(t)$ are real exponents, $c_j(t)$ can be complex. One has $\alpha_1(t) > \alpha_2(t) > \dots$.

(I.2) is supposed to represent the spin-independent part of the elastic scattering amplitude. Spin-dependent terms of the amplitude can be expanded in the same way, but they are generally expected to grow less rapidly with s than the spin-independent term ¹⁾.

One could add in each term of (I.2) factors having a variation slower than logarithmic, for example $(\ln \ln s_t)^{\delta_j(t)}$. The data available, however, are much too limited to decide on the presence or absence of such factors. In fact it is even impossible to establish or reject the presence of the factors $(\ln s_t)^{\beta_j(t)}$ in (I.2). The analysis is therefore usually carried out by putting $\beta_j(t) = 0$.

The expansion (I.2) is general enough to accommodate various types of singularities of the scattering amplitude in the complex angular momentum plane of the t channel. Branch points give $\beta_j(t) \neq 0$, poles $\beta_j(t) = 0$. The poles can be moving (Regge poles) or fixed depending on whether $\alpha_j(t)$ depends on t or not. As mentioned above experiment does not allow to decide between the occurrence of poles or branch points. It is now generally recognized that both types of singularities are theoretically possible. In fact, branch points are expected to be present ²⁾.

Using analyticity of T in s and crossing symmetry ³⁾, the expansion (I.2) can be related to the corresponding asymptotic expansion for the crossed reaction $\bar{A}+B \rightarrow \bar{A}+B$

$$\bar{T}(ut) \underset{u \rightarrow +\infty}{\simeq} \sum_{j=1,2,\dots} c_j^*(t) e^{-i\pi \alpha_j(t)} u_t^{\alpha_j(t)} (\ln u_t)^{\beta_j(t)} \quad (I.4)$$

where u is the square of the c.m. energy and

$$u_t = u + \frac{t}{2} - m_A^2 - m_B^2. \quad (I.5)$$

The coefficients $c_j(t)$, $\alpha_j(t)$, $\beta_j(t)$ are the same as in (I.2). The variables s_t , u_t in (I.2) and (I.4) have been so chosen as to give this very simple form to the high energy consequences of analyticity and crossing symmetry ⁴⁾.

All data on direct and crossed reactions are consistent with (I.2) and (I.4). It is in fact very useful to use these equations simultaneously in order to determine the parameters $c_j(t)$, $\alpha_j(t)$ as accurately as possible. We recall that for the time being $\beta_j(t)$ cannot be determined from the data and is usually put equal to zero, so that (I.2) and (I.4) reduce to

$$T(st) \underset{s \rightarrow +\infty}{\simeq} \sum_{j=1,2,\dots} c_j(t) s_t^{\alpha_j(t)} \quad (I.6)$$

$$\bar{T}(ut) \underset{u \rightarrow +\infty}{\simeq} \sum_{j=1,2,\dots} c_j^*(t) e^{-i\pi \alpha_j(t)} u_t^{\alpha_j(t)}. \quad (I.7)$$

1. The forward scattering amplitude

Our knowledge of the forward scattering amplitude $T(s0)$ stems from the data on total cross-sections σ_T , related to $T(s0)$ by the optical theorem

$$\sigma_T = (8\pi^2/ks^{\frac{1}{2}}) \operatorname{Im} T(s0), \quad k = \text{c.m. momentum} \quad (I.8)$$

and from the recently acquired data on the ratio

$$X = \frac{\operatorname{Re} T(s_0)}{\operatorname{Im} T(s_0)} . \quad (\text{I.9})$$

The by now familiar total cross-section curves are reproduced in Fig. 1⁵⁾. The ratio X is obtained from elastic scattering measurements at very small $|t|$ where the nuclear and Coulomb amplitudes interfere. One assumes the scattering to be spin-independent and one assumes $T(st)$ to be constant between $t = 0$ and the small $|t|$ where the interference is measured. The results available to date are given in Figs. 2 and 3^{6),7)}. There appears to be compatibility with dispersion relations for pp scattering. The situation is less clear for $\bar{p}p$ where the data for forward charge exchange scattering⁸⁾ are also to be taken into account.

We discuss the general behaviour of σ_T and X in terms of (I.6-7). The total cross-section curves indicate with good accuracy

$$\alpha_1(0) = 1. \quad (\text{I.10})$$

Keeping only the leading term with $j = 1$, (I.6) and (I.10) then give

$$\operatorname{Im} T(s_0) \simeq \operatorname{Im} \bar{T}(s_0), \quad \operatorname{Re} T(s_0) \simeq -\operatorname{Re} \bar{T}(s_0) \quad (\text{I.11})$$

for $s \rightarrow +\infty$. The first relation is the Pomeranchuk theorem $\sigma_T \simeq \bar{\sigma}_T$ for total cross-sections⁹⁾. Its asymptotic validity is compatible with the data, although up to 30 GeV there are considerable differences

$\sigma_T - \bar{\sigma}_T$ which, for $p_L \gtrsim 6$ GeV/c, can be well described by the second term $j = 2$ of (I.6) and (I.10), $\alpha_2(0)$ being of order 0.5.

The second relation (I.11) gives $X \simeq -\bar{X}$ for $s \rightarrow +\infty$, where $\bar{X}$ is the analogue of (I.9) for the crossed reaction. Experiment shows X and $\bar{X}$ to be negative where measured, so that (I.11) suggests

$$X \rightarrow 0 \quad \text{and} \quad \bar{X} \rightarrow 0 \quad \text{for} \quad s \rightarrow \infty. \quad (\text{I.12})$$

This would mean that the coefficient $c_1(0)$ in (I.6) is purely imaginary

$$c_1(0) = i |c_1(0)|. \quad (\text{I.13})$$

The experimental results are entirely compatible with this expectation which implies that the measured values of X are due to the $j = 2$ term of the expansions (I.6-7), with $\alpha_2(0) \simeq 0.5$ as before. It should be noted, however, that there would be no theoretical objection against X becoming positive for some reactions in the region $p_L > 30 \text{ GeV}/c$, if σ_T would show a minimum in this region.¹⁰⁾

A detailed analysis of σ_T and X for pp and $\bar{p}p$ has been carried out by A. Biaľas and E. Biaľas on the basis of expansions of type (I.6-7)¹¹⁾. It reveals that the high energy limit of (I.9) for pp is compatible with 0 and is very likely to be contained between -0.2 and 0. Furthermore it shows that measurement of (I.9) for $\bar{p}p$ at laboratory momenta above 5 GeV/c is by far the best way of reaching a better determination of the expansion parameters $c_j(0)$, $\alpha_j(0)$ for the proton-proton system.

2. The shape of the diffraction peak

All diffraction peaks measured until now have in the momentum transfer interval (I.1) a common shape given by

$$\frac{d\sigma}{dt} = \left(\frac{d\sigma}{dt} \right)_{t=0} e^{at+bt^2}. \quad (\text{I.14})$$

The coefficient a , although depending somewhat on the reaction and the energy, is always of order $10(\text{GeV}/c)^{-2}$. The dimensionless quantity b/a^2 is not well determined, but it is always positive and has an order of magnitude

$$\frac{b}{a^2} \simeq 0.01 \text{ to } 0.04. \quad (\text{I.15})$$

From (I.6-7) taken at energies sufficiently high for the leading term $j = 1$ to be the only important one, one has $|T| \simeq |\bar{T}|^2$ and $\frac{d\sigma}{dt}$ therefore becomes the same for the direct and the crossed reaction :

$$\frac{d\sigma}{dt} \simeq \frac{d\bar{\sigma}}{dt} \quad (\text{I.16})$$

where both sides are taken at the same energy, $s = u \rightarrow +\infty$. (I.16) is the generalized Pomeranchuk theorem ¹²⁾. It implies

$$a \simeq \bar{a}, \quad b \simeq \bar{b} \quad (\text{I.17})$$

at very high energy (the bar always refers to the crossed reaction).

(I.17) is in good agreement with experiment for π^+p scattering for $p_L > 6 \text{ GeV}/c$. Both a and b appear to be constant with energy (this result is only significant for a , the precision on b being still poor). For pp and $\bar{p}p$ scattering, there is a pronounced energy variation of a (referring to pp), which increases, as well as of $\bar{a}$ (referring to $\bar{p}p$) which decreases. As shown by Czyzewski et al. ¹³⁾, from a compilation of all data, a and $\bar{a}$ appear to tend to a common limit $\simeq 9(\text{GeV}/c)^{-2}$, about the same value as obtained for π^+p . This is clearly indicated by Fig. 4, taken from Ref. ¹³⁾, where our quantity a is denoted by A .

Therefore, in all cases mentioned, we conclude that the leading term of the expansions (I.6), (I.7) can be given the exponent

$$\alpha_1(t) = 1 \quad \text{for} \quad 0 < -t < 1(\text{GeV}/c)^2. \quad (\text{I.18})$$

From general considerations ¹⁴⁾ one is then led to expect

$$c_1(t) = i |c_1(t)| = i e^{\frac{1}{2}(at+bt^2)} \quad (\text{I.19})$$

for all t values in the same interval, thereby extending (I.13). The results (I.18-19) for the leading term of (I.6-7) correspond to the familiar diffraction picture characterized by constant cross-sections and purely imaginary amplitude. Deviations thereof will be produced by the higher terms of the expansion; the data roughly suggest $\alpha_2(t) \simeq 0.5$ for the t interval considered.

Although our conclusion is that the diffraction picture provides by far the most natural interpretation of the experimental facts, the latter are also compatible with a Regge pole model using several Regge trajectories with positive slope. When several trajectories contribute, the shrinkage effects resulting from their positive slope can be cancelled over a finite energy interval by expansion (antishrinkage) effects produced through appropriate compensations between the various trajectories. For example, a Regge pole fitting of this sort has been made recently by Phillips and Rarita for πN , KN and $\bar{K}N$ scattering, using three trajectories for πN and five for KN , $\bar{K}N$ ¹⁵⁾. The final decision between diffraction picture and Regge pole model will probably require data at much higher energies than now available.

II. DOMINANT INELASTIC COLLISIONS

Our knowledge of the most common inelastic collisions, i.e., those collisions which make up most of the inelastic cross-section, is still very limited, both theoretically and experimentally. We limit our discussion to two aspects of the theoretical work devoted recently to this problem.

1. Peripheral, multiperipheral and centropерipheral models

Because of the striking property of the dominant inelastic collisions to have low transverse momenta for all outgoing particles, one has tried for many years to account for them in terms of one meson exchange diagrams. This so-called peripheral model, originally based on diagram (a) of Fig. 5, was soon realized to be inconsistent with itself because, for increasing energy of the collision AB, at least one of the collisions πA , πB in diagram (a) has an increasing energy and should itself become peripheral. Self-consistency is re-established by considering the sum of all diagrams of type (b) with 1,2,3 vertices, and specifying some restrictive property for the vertex which prevents it from becoming itself a peripheral collision. Two ways of doing this have been proposed.

The first procedure is to say that the energy of each vertex in its own rest frame remains below a certain maximum. This gives the multiperipheral model of Amati, Stanghellini and Fubini¹⁶⁾. The resulting shadow scattering at high energy is similar to the elastic scattering produced by a single Regge trajectory with positive slope and is therefore not compatible with the diffraction picture.

In the second procedure, recently proposed by Dremin et al.¹⁷⁾, each vertex is regarded to represent a central collision. To this effect a specific algebraic form is assumed for the shadow scattering amplitude

resulting from a single vertex, i.e., from central collisions, giving it the form $sf(t)$ when the two incident particles are on their mass shell. For the shadow scattering of the total amplitude, this gives the diffraction picture at very large energy, with additional contributions of peripheral nature in the second and higher terms of the expansion (I.2). The singularities in the angular momentum plane in the t channel are a fixed pole, a moving pole and a moving branch point. It is now recognized that the theoretical objections raised against such singularities ¹⁸⁾ are not valid for $t < 0$.

The latter model, which might be called centrop peripheral, is undoubtedly more adequate than the multiperipheral one for the energy region up to 30 GeV. It allows for a variable proportion of central and peripheral collisions. Thus πp and $\bar{p}p$ would be mainly central, pp being mainly peripheral.

2. Shape of the diffraction peak in relation to inelastic processes

As mentioned above, all diffraction peaks measured up to now have in the interval

$$0 < -t \lesssim 1(\text{GeV}/c)^2 \quad \text{or} \quad 0 < -at \lesssim 10 \quad (\text{II.1})$$

the same shape given by (I.14-15). This shape is best discussed in terms of the partial wave amplitude

$$\eta_\ell = 1 - e^{2i\delta_\ell} \quad (\text{II.2})$$

where δ_ℓ is the complex phase shift. The function η_ℓ is real when the diffraction picture holds (purely imaginary scattering amplitude T). We first concentrate our discussion on this case. The effect of a non-vanishing ratio $\text{Re } T / \text{Im } T$ will be commented upon later.

The most conventional assumption concerning η_ℓ , which is given by

$$\begin{cases} \eta = \eta_0 & \text{for } 0 \leq \ell \leq L \\ \eta_\ell = 0 & \text{for } \ell > L \end{cases} \quad (\text{II.3})$$

and corresponds to absorption by a black sphere, is in clear disagreement with the facts. It gives

$$\frac{b}{a^2} \simeq -0.08 \quad (\text{II.4})$$

and produces a diffraction minimum for $-at \simeq 3.7$. This disagreement is not suppressed by a slight rounding off of the edge at $\ell = L$.

Much better agreement is obtained with a Gaussian distribution

$$\eta_\ell = \eta_0 \exp(-\ell^2/L^2) \quad (\text{II.5})$$

which gives a vanishing b . The ratio $\sigma_{\text{el}}/\sigma_T$ (σ_{el} = total elastic cross-section) ranges from 0 to 0.25 as η_0 grows from 0 to 1¹⁹). The experimental values of $\sigma_{\text{el}}/\sigma_T$ being of order 0.2 (see below), one concludes that η_0 is a little smaller than 1.

A further improvement is obtained if η_ℓ is taken to be the smaller root of the equation

$$\eta_\ell - \frac{1}{2} \eta_\ell^2 = f_0 \exp(-\ell^2/L^2). \quad (\text{II.6})$$

This gives a positive b/a^2 ranging from 0 to 0.02 as f_0 grows from 0 to the maximum value $\frac{1}{2}$ allowed by unitarity. As to the ratio $\sigma_{\text{el}}/\sigma_T$, it increases from 0 to a maximum of 0.185. Equation (II.6) is a special case of the unitarity condition for elastic scattering which can be put in the form

$$\operatorname{Re} \gamma_\ell - \frac{1}{2} |\gamma_\ell|^2 = f_\ell \quad (\text{II.7})$$

The left-hand side is the contribution of the elastic channel to the unitarity condition, the right-hand side being the total contribution of all inelastic channels. The precise definition of f_ℓ is given by the equation

$$\langle \psi(\vec{k}_A' \vec{k}_B') | \psi(\vec{k}_A \vec{k}_B) \rangle = \frac{\delta(P'-P) \sqrt{s}}{2\pi k \epsilon_A \epsilon_B} \sum_{\ell=0}^{\infty} (2\ell+1) f_\ell P_\ell(\cos\theta) \quad (\text{II.8})$$

$\psi(\vec{k}_A \vec{k}_B)$ = inelastic final state produced by initial state $|\vec{k}_A \vec{k}_B\rangle$
 where $\vec{k}_A, \vec{k}_B$ are the momenta of the incident particles
 A, B ²⁰⁾.

k = length of momentum of incident particles in c.m. system.

ϵ_A, ϵ_B = energies of incident particles in c.m. system.

θ = angle between $\vec{k}_A$ and $\vec{k}_A'$ in c.m. system.

$\delta(P'-P)$ = four-dimensional delta function for total momentum-energy conservation.

The Gaussian ℓ dependence adopted for f_ℓ in (II.6) can be shown to hold true when the inelastic final state is mainly composed of many-particle states with small transverse momenta and with weak correlations between the particles. This model therefore proposes an explanation of the diffraction peak shape in terms of general properties of the dominant inelastic collisions ^{21), 22)}.

As shown by Cottingham and Peierls ²³⁾, the model just discussed gives an excellent description of πp diffraction scattering at $p_L \gtrsim 15$ GeV/c where the real part of the scattering amplitude is sufficiently small to neglect its contribution to (II.7) ²⁴⁾. The parameter f_0 is fixed by the ratio σ_{el}/σ_T which is 0.16, somewhat below the allowed maximum 0.185.

For pp collisions the ratio σ_{el}/σ_T is larger than 0.185 (its value is 0.276 ± 0.008 at $p_L = 12.8$ GeV/c and 0.244 ± 0.12 at 19.6 GeV/c). Cottingham and Peierls attributed this to a non-vanishing real part of the scattering amplitude T , the existence of which was also predicted by dispersion relations²⁵⁾ and was confirmed experimentally (see Fig. 3). By giving to this real part a short range they were able to account simultaneously for the observed pp elastic scattering at small and at large $|t|$ values²³⁾. If on the other hand one would assume the real part of T to have the same t dependence as the imaginary part, the partial wave amplitude would verify

$$\eta_\ell = |\eta_\ell| e^{i\alpha}$$

with α independent of ℓ , and (II.6) would become

$$|\eta_\ell| \cos \alpha - \frac{1}{2} |\eta_\ell|^2 = f_0 \exp(-\ell^2/L^2)$$

which is equivalent to

$$|\eta_\ell| \cos^{-1} \alpha - \frac{1}{2} |\eta_\ell|^2 \cos^{-2} \alpha = f_0 \cos^{-2} \alpha \exp(-\ell^2/L^2).$$

The shape of the diffraction peak would then be the same as follows from (II.6) with real η_ℓ , but with f_0 replaced by $f_0 \cos^{-2} \alpha$; this gives a larger value of b/a^2 . Also the ratio σ_{el}/σ_T is becoming larger.

Since we expect the diffraction picture to become valid at high energy, the applicability of the model (II.6) to pp scattering requires that the ratio σ_{el}/σ_T decrease below 0.185 at very high energy. It is interesting in this connection that, using a simultaneous extrapolation of pp and $\bar{p}p$ data, Czyzewski et al. find for the asymptotic limit of this ratio the value 0.178 ± 0.006 ¹³⁾.

III. INELASTIC TWO-BODY COLLISIONS

Among the inelastic collisions, considerable interest is attracted by two-body collisions $A+B \rightarrow A'+B'$ where particles A' and/or B' are different from A and B . In fact A', B' can be unstable resonances of definite quantum numbers. The cross-section for such processes is surprisingly high at low and intermediate energies ($p_L \lesssim 5 \text{ GeV/c}$). For example in K^+p collisions at $p_L = 3 \text{ GeV/c}$, about 40% of the $K^0 \pi^+ p$ final states are in the two-body channel $K^{*+}p$ and 40% in the two-body channel $K^0 N^{*++}$ whereas only 20% are genuine three-body states²⁶⁾. It would be desirable to find a general, qualitative explanation for this remarkable fact. As the energy increases, it appears to be a general rule that the cross-section for any individual inelastic two-body process $A+B \rightarrow A'+B'$ decreases rapidly, at least when the internal quantum numbers of A' and/or B' differ from those of A and B . Typical examples are a cross-section decreasing like s^{-1} for $\pi^- p \rightarrow \pi^0 n$ and one decreasing like s^{-2} for $\pi p \rightarrow \rho p$. The number of two-body channels increases with s , however, and it is not known whether the total cross-section into all two-body channels decreases or becomes constant. Experimental indications on this point would be very helpful for our general understanding of high energy collisions. Thus, the success of the model discussed in Section II.2 and based on Eq. (II.6) is readily understood if at high energy most inelastic collisions are of many-body type. Whether the Gaussian ℓ dependence of f_ℓ can also be derived on general grounds when an appreciable fraction of the inelastic collisions leads to two-body states is an unsolved question.

1. One meson exchange models for inelastic two-body collisions

The model most commonly used to describe inelastic two-body collisions is the one meson exchange model (OME) with various improvements, mainly form factors²⁷⁾ and absorption corrections²⁸⁾. Very extensive work has been devoted to these models. We shall limit ourselves to a few brief comments because excellent reviews have been given recently by Jackson²⁹⁾ and Jackson et al.³⁰⁾.

For processes dominated by one pion exchange the OME model, either with form factors or with absorption, has reached an appreciable degree of success. A favourable example is provided by the reaction $\pi p \rightarrow \rho p$ where both models give a very successful description of the differential cross-section as function of p_L and t in the intervals $2 < p_L < 8 \text{ GeV/c}$, $0 < -t < 0.5(\text{GeV/c})^2$. The absorption model has the advantage of predicting decay correlations and a non-vanishing difference between $\pi^+ p \rightarrow \rho^+ p$ and $\pi^- p \rightarrow \rho^- p$, quantities which are absent in the OME with or without form factors but are present in the data. The success of the one pion exchange models for nucleon-nucleon and antinucleon-nucleon collisions with isobar excitation is more limited and seems to be restricted to smaller energies ³¹⁾.

More fundamental difficulties are encountered by the OME models in those cases, like $Kp \rightarrow K^*p$, where vector-meson exchange is found to be required by the experimental data. Irrespective of whether form factors or absorption corrections are introduced, the model gives a cross-section independent of energy for $s \rightarrow +\infty$, in clear contradiction with the facts. The shape of the angular distribution, on the other hand, can be accounted for with more success, especially when absorption and form factor effects are included simultaneously. An example is given by charge exchange scattering, as discussed below.

A very interesting modification of the OME model, introducing the OME Born amplitudes in the K matrix rather than in the S matrix, has been recently studied by Dietz and Pilkuhn ³²⁾. Although the method is quite different, an analysis of K^+p reactions shows that the main conclusions of good agreement with intermediate energy data and of difficulties with the energy dependence at large energy remain essentially valid. The model also allows to calculate elastic scattering and finds a satisfactory t dependence but too large an absolute value of the differential cross-section.

2. Charge exchange scattering

Among inelastic two-body collisions, charge exchange scattering processes like $\pi^-p \rightarrow \pi^0n$ and $K^-p \rightarrow K^0n$ are of particular interest because of their close relation to elastic scattering and also because accurate experimental information is becoming available. We shall discuss only meson-nucleon charge exchange which is appreciably simpler than the proton-neutron case. Both K^-p and π^-p charge exchange data are available. They are characterized by the same general shape, as shown on Fig. 6³³⁾ and Fig. 7³⁴⁾. This shape shows a remarkable difference from elastic scattering, in the form of a small maximum at very low momentum transfer. The most natural explanation is in terms of a spin-flip amplitude of the same order of magnitude as the non-spin-flip amplitude (both being very small compared to the elastic amplitude which is believed to be mainly of non-spin-flip type). Beyond the small maximum the cross-section drops exponentially in $|t|$, with an exponent which is about the same as the exponent a in the elastic cross-section (I.14). At still larger $|t|$ there is indication of a second maximum decreasing much faster with energy than the first one.

Figure 6 illustrates, following H. Högaasen and J. Högaasen³⁵⁾, how vector meson exchange corrected for absorption can reproduce the general shape of the charge exchange data but neither the width nor the absolute magnitude. While these discrepancies can be reduced at moderate energies by the introduction of form factors, the constancy of the calculated cross-section with energy is certainly in strong disagreement with the facts.

The energy dependence of the charge exchange cross-section is known for the π^-p case⁸⁾ and is illustrated in Fig. 7 where the ordinate is

$$p_L \frac{d\sigma}{dt} \propto s \frac{d\sigma}{dt}$$

It raises the very interesting question of whether it is better interpreted by means of one single term or of two terms in an asymptotic expansion (I.6) for the charge exchange amplitude. Let us use first one single term

$$T_{\text{ch.ex.}} \simeq c_1(t) s_t^{\alpha_1(t)} \quad (\text{III.1})$$

The data give $\alpha_1(0) \simeq 0.5$ with good precision and suggest roughly $\alpha_1(-0.5) \simeq 0$, $\alpha_1(-1) \simeq -0.5$, t being measured in $(\text{GeV}/c)^2$. If the decrease of $\alpha_1(t)$ for $-t$ going from 0 to 1 is gradual, the single-term analysis (III.1) is quite satisfactory and the exponent $\alpha_1(t)$ has a large positive slope, of order $\alpha_1'(0) \simeq 1$. If on the other hand, as the data suggest, $\alpha_1(t)$ changes slowly for $0 < -t < 0.5$, then drops rapidly for $-t$ around 0.5 to vary again more slowly for $-t > 0.5$, it would be more natural to use instead of (III.1) a two-term expansion

$$T_{\text{ch.ex.}} \simeq c_1(t) s_t^{\alpha_1(t)} + c_2(t) s_t^{\alpha_2(t)} \quad (\text{III.2})$$

with $\alpha_1(t)$ and $\alpha_2(t)$ changing slowly with t and having orders $\alpha_1 \simeq 0.5$ and $\alpha_2 \simeq -0.5$ for the whole t range. $c_1(t)$ should be more peaked than $c_2(t)$; then, at any energy s , there would be a value $t_0(s)$ such that the first term of (III.2) would dominate for $t_0 < t < 0$ and the second for $t < t_0$.

A decision between the two possibilities just discussed seems to be difficult with the statistics now available. It will probably have to await further experiments. One point is already clear, however, namely that $\alpha_1(0) \simeq 0.5$ in all cases. This explains the failure of the vector meson exchange model, which predicts $\alpha_1(0) = 1$, and is a clear and unambiguous success for the Regge pole model. Indeed, among

all trajectories commonly considered in the Regge pole model, only the ρ trajectory has the quantum numbers involved in π^-p charge exchange, and the model predicts a value of $\alpha_\rho(0)$ which is close to 0.5.

Isospin invariance implies that the crossed amplitude $\bar{T}_{\text{ch.ex.}}$ of the π^-p charge exchange amplitude $T_{\text{ch.ex.}}$ verifies

$$\bar{T}_{\text{ch.ex.}}(st) = -T_{\text{ch.ex.}}(st).$$

Comparing (I.6) and (I.7) we can then derive the phases of the coefficients $c_j(t)$ modulo π

$$c_j(t) = \pm |c_j(t)| \exp\left[-\frac{i\pi}{2} (\alpha_j(t)+1)\right]. \quad (\text{III.3})$$

This has interesting consequences for the polarization of the recoil neutron ³⁶⁾. Assume (III.1) to hold with good accuracy both for the spin-flip and the non-spin-flip amplitudes. From (III.3) these amplitudes will have the same phase and zero polarization will result. In the higher approximation (III.2) polarization will occur through interference between the two terms of the expansion and its magnitude will decrease as $s_t^{\alpha_2(t)-\alpha_1(t)}$. Among the two possibilities mentioned above to account for the energy dependence of $d\sigma/dt$, the second suggested $\alpha_1 \sim 0.5$, $\alpha_2 \sim -0.5$. If it turns out to be valid, the polarization will decrease roughly as s^{-1} .

IV. LARGE MOMENTUM TRANSFERS - HIGHER SYMMETRIES

We end by a few comments on some fundamental problems which begin to be tackled experimentally and theoretically from various sides, and which may well turn out to become in the future the most important items in the study of high-energy dynamics.

1. Large angle scattering

As discussed above, elastic scattering at small momentum transfer $[0 < -t < 1(\text{GeV}/c)^2]$ appears to have the same properties for all hadrons and is probably just shadow scattering, thereby reflecting only very little of the more specific features of high energy dynamics. One may hope that at large momentum transfers elastic scattering becomes more sensitive to the underlying dynamics. We are far from knowing whether this will be so, but a hopeful indication is given by the fact that pp and πp elastic scattering become different from each other at large angle. This is seen from Fig. 8, which gives the pp data at all energies available ³⁷⁾, and Fig. 9 giving the πp results at $p_L = 8 \text{ GeV}/c$ ³⁸⁾. Only one feature appears at present to be common to these results, namely the exponential drop of the differential cross-section at $\theta_{\text{c.m.}} \sim 90^\circ$ with energy, roughly proportional to $\exp(-3.2 \sqrt{s})$ in agreement with the predictions of the statistical model ³⁹⁾. It will be extremely interesting to study experimentally further cases of two-body elastic and inelastic processes at large momentum transfer. Only in this way will one be able to disentangle the many lines of interpretation which are possible at present. In addition to the statistical model, one has mentioned as possibilities a minimal interaction allowed by analyticity ⁴⁰⁾ and universal form factors ⁴¹⁾, and there is little doubt that still other concepts could be introduced to analyze the same phenomena. Only experiment can decide about their usefulness. We note also that, as shown by Cottingham and Peierls ²³⁾, the observed large angle pp scattering can be explained by giving a short range to the real part of the scattering amplitude, an explanation which does not apply to πp and may account for the very different shapes of the angular distributions.

Figure 10 shows πp large angle scattering plotted in $\cos \theta_{\text{c.m.}}$, up to the backward region which is reproduced in enlarged scale ³⁸⁾. It reveals a pronounced backward peak for the $\pi^+ p$ system, with a surprisingly narrow width of the order of half the width of the forward diffraction peak. It is natural to associate this peak with

nucleon exchange, and it is clear that a reggeized nucleon exchange model will have enough free parameters and functions to fit the data. It would be nice, however, to correlate this phenomenon with others in order to arrive at non-trivial conclusions. It should be recalled in this connection that the crossed reaction is $\bar{p}p \rightarrow \pi \pi$ ⁴²⁾ so that a common analysis with help of expansions of type (I.6-7) can be carried out.

2. Unitary symmetry

SU(3) and SU(6) symmetry has been very successful in classifying particles and in predicting properties of three-particle vertex functions like those involved in electromagnetic and weak interactions. The success is more limited for scattering and particle production processes, where it appears empirically that symmetry breaking effects considerably obscure the situation. A detailed discussion of this question has recently been given by Harari ⁴³⁾. At high energies, most elastic and inelastic collisions still involve small momentum transfers, so that the effect of symmetry breaking has no reason to become less important. Following Harari, we illustrate this fact by the following considerations.

Using U spin conservation one can demonstrate ⁴⁴⁾ that exact SU(3) symmetry implies

$$|T(K^-p \rightarrow K^-p) - T(\pi^-p \rightarrow \pi^-p)| = |T(K^-p \rightarrow \pi^- \Sigma^+)| \quad (\text{IV.1})$$

where all scattering amplitudes are taken at the same s and t . For $s \rightarrow +\infty$ at fixed t , the ratio

$$T(K^-p \rightarrow \pi^- \Sigma^+)/T(K^-p \rightarrow K^-p) \quad (\text{IV.2})$$

tends rapidly to zero because it involves the strangeness exchange reaction $K^-p \rightarrow \pi^- \Sigma^+$. This implies

$$\left[\frac{T(K^- p \rightarrow K^- p)}{T(\pi^- p \rightarrow \pi^- p)} \right] \rightarrow 1. \quad (\text{IV.3})$$

Taking this relation for $t = 0$ and applying the optical theorem, one concludes that it contradicts the fact that $K^- p$ and $\pi^- p$ have different total cross-sections at high energy. In other words, for s sufficiently large to neglect the ratio (IV.2), any difference between the left-hand side of (IV.3) and unity is due to symmetry breaking.

This has an important consequence for the relations proposed by Johnson and Treiman⁴⁵⁾ to test $SU(6)$ symmetry

$$\frac{1}{2} [f(K^+ p) - f(K^- p)] = f(\pi^+ p) - f(\pi^- p) = f(K^0 p) - f(\bar{K}^0 p) \quad (\text{IV.4})$$

where f denotes the forward elastic scattering amplitude. As soon as the energy is high enough for strangeness exchange processes to become negligible, all quantities

$$\frac{f(K^- p)}{f(\pi^- p)}, \quad \frac{f(K^+ p)}{f(\pi^+ p)}, \quad \frac{f(K^+ p)}{f(K^- p)}$$

differ from 1 only through $SU(3)$ symmetry breaking effects⁴³⁾. Under these circumstances, any equality involving non-vanishing experimental values of the first and second differences in (IV.4) is a relation between $SU(3)$ violating terms and is no test of $SU(3)$ nor of $SU(6)$. Consequently, the Johnson-Treiman relation can only be used at low energy where strangeness exchange is appreciable and where one might hope that $SU(3)$ symmetric contributions to (IV.4) would no longer be overshadowed by the symmetry breaking effects. Whether this hope is justified is unknown.

We remain with the central question whether unitary symmetry will manifest itself more clearly in processes where not only energy but also all relevant momentum transfers are large compared to the baryon masses. Experimentally, one should for example check the triangular inequalities between cross-sections following from (IV.1), and many other similar tests of $SU(3)$, in two-body processes at high energy and large angle. Theoretically one should extract from future experiments a more definite picture of the dynamical behaviour of large momentum transfer phenomena, and guess how it accommodates SU_3 symmetry. The undeniable progress made in the last years in the phenomenological study of high energy phenomena at low momentum transfers gives us good hope that similar advances will be possible in the more fundamental but less accessible region where all kinematical variables are large ⁴⁶⁾.

REFERENCES AND FOOTNOTES

- 1) Experimental confirmation of this conjecture would be highly desirable. For πp collisions it would require polarization and spin correlation measurements by means of polarized proton targets. For pp collisions such measurements should be done with incident and target protons being both polarized.
- 2) S. Mandelstam, *Nuovo Cimento* 30, 1127, 1148 (1963).
- 3) One is now very close to being able to derive the necessary properties from axiomatic field theory (V. Glaser, private communication).
- 4) For a proof, see L. Van Hove, Lectures at the Seminar on High Energy Physics and Elementary Particles, International Centre for Theoretical Physics of I.A.E.A., Trieste, May-June 1965, to be published.
- 5) S.J. Lindenbaum, Rapporteur's talk, International Conference on High Energy Physics, Dubna, August (1964).
- 6) K.J. Foley, R.S. Gilmore, R.S. Jones, S.J. Lindenbaum, W.A. Love, S. Ozaki, E.H. Willen, R. Yamada and L.C.L. Yuan, *Phys.Rev. Letters* 14, 862 (1965).
- 7) G. Bellettini, G. Cocconi, A.N. Diddens, E. Lillethun, J. Pahl, J.P. Scanlon, J. Walters, A.M. Wetherell and P. Zanella, *Phys.Letters* 14, 164 (1965).
- 8) I. Mannelli, A. Bigi, R. Carrara, M. Wahlig and L. Sodickson, *Phys.Rev.Letters* 14, 408 (1965) ;
A.V. Stirling, P. Sonderegger, J. Kirz, P. Falk-Vairant, O. Guisan, C. Bruneton, P. Borgeaud, M. Yvert, J.P. Guillard, C. Caverzasio and B. Amblard, *Phys.Rev.Letters* 14, 763 (1965).
- 9) I.Ia. Pomeranchuk, *Zhur.Eksp.i Teoret.Fiz.* 34, 725 (1958) ;
(English translation : *Soviet Phys. - JETP* 7, 499 (1958)).

- 10) The author thanks Professors V.S. Barashenkov and G. Höhler for a discussion of this point. Qualitative arguments tending to explain the negative sign of X , as given for example in H.E. Conzett, Phys.Letters 16, 189 (1965), are not compelling because the sign of X can always be modified by a real optical potential. In fact, as shown by A. Biaľas and E. Biaľas, [CERN preprint TH. 504 and Nuovo Cimento 37, 1686 (1965)], the negative sign of X is related to the fact that σ_T and $\overline{\sigma}_T$ decrease toward their asymptotic limit ; an increase would make X positive.
- 11) A. Biaľas and E. Biaľas, as quoted in Ref. ¹⁰).
- 12) L. Van Hove, Phys.Letters 5, 252 (1963) ;
A.A. Logunov, Nguyen Van Hieu, I.T. Todorov and O.A. Khrustalev, Phys.Letters 7, 69 (1963).
- 13) O. Czyzewski, B. Escoubès, Y. Goldschmidt-Clermont, M. Guinea-Moorhead, D.R.O. Morrison and S. de Unamuno-Escoubès, Phys. Letters 15, 188 (1965).
- 14) L. Van Hove, Phys.Letters 7, 76 (1963).
- 15) R.J.N. Phillips and W. Rarita, Lawrence Radiation Laboratory preprint UCRL 16033.
- 16) D. Amati, A. Stanghellini and S. Fubini, Nuovo Cimento 26, 896 (1962).
- 17) I.M. Dremin, I.I. Roysen, E.L. Feinberg, D.S. Chernavski and R. White, Proceedings, 1964 International Conference on High Energy Physics at Dubna, to be published. See also E.L. Feinberg's review at the same conference.
- 18) V.N. Gribov, Nuclear Phys. 22, (1961) ;
R. Oehme, Phys.Rev.Letters 9, 358 (1962).

- 19) In contrast herewith, the black sphere distribution (II.3) gives a ratio σ_{el}/σ_T ranging from 0 to 0.5 as η_0 increases from 0 to 1.
- 20) The normalization is such that $\langle \vec{k}'_A \vec{k}'_B | \vec{k}_A \vec{k}_B \rangle = \int \delta(\vec{k}'_A - \vec{k}_A) \delta(\vec{k}'_B - \vec{k}_B)$.
- 21) L. Van Hove, *Nuovo Cimento* 28, 798 (1963) and *Revs. Modern Phys.* 36, 655 (1964). A preliminary discussion of this method, and of the relation between inelastic collisions and diffraction peak width, was given in "Physikertagung Stuttgart, 1962", (Physik. Verlag, Mosbach/Baden, p. 84, 1963).
- 22) A concrete, bremsstrahlung-type model of weakly correlated inelastic final states has recently been discussed for proton-proton collisions by A. Białas and Th. Ruijgrok, CERN preprint TH. 535.
- 23) W.N. Cottingham and R.F. Peierls, *Phys. Rev.* 137, B147 (1965).
- 24) This contribution is of order $(\text{Re } T / \text{Im } T)^2$, which is a few percent for very small $|t|$ and can be safely assumed to remain small for larger $|t|$.
- 25) P. Söding, *Phys. Letters* 8, 285 (1964).
- 26) G.R. Lynch, M. Ferro-Luzzi, R. George, Y. Goldschmidt-Clermont, V.P. Henri, B. Jongejans, D.W.G. Leith, F. Müller and J.M. Perreau, *Phys. Letters* 9, 359 (1964) and *Phys. Rev.*, to be published.
- 27) E. Ferrari and F. Selleri, *Suppl. Nuovo Cimento* 24, 453 (1962).
- 28) N.J. Sopkovich, *Nuovo Cimento* 26, 186 (1962).
- 29) J.D. Jackson, *Revs. Modern Phys.* (to be published).
- 30) J.D. Jackson, J.T. Donohue, K. Gottfried, R. Keyser and B.E.Y. Svensson, *Phys. Rev.* (to be published).
- 31) E. Ferrari, S. Gennarini and P. Lariccia, CERN preprint TH. 521 (to be published in *Nuovo Cimento*) ;
B.E.Y. Svensson, CERN preprint TH. 549 (to be published in *Nuovo Cimento*).

- 32) K. Dietz and H. Pilkuhn, *Nuovo Cimento* 37, 1561 (1965) and CERN preprint TH. 563.
- 33) P. Astbury, G. Finocchiaro, A. Michelini, C. Verkerk, C.H. West, W. Beusch, B. Gobbi, M. Pepin, M.A. Pouchon and E. Polgar, *Phys.Letters* 16, 328 (1965).
- 34) A.V. Stirling et al., Ref. 8).
- 35) H. Högaasen and J. Högaasen, CERN preprint TH. 554, to be published in *Nuovo Cimento*.
- 36) The author acknowledges useful discussions with Professor G. Höhler and a very enlightening correspondence with Professor G.F. Chew on this question.
- 37) G. Cocconi, V.T. Cocconi, A.D. Krisch, J. Orear, R. Rubinstein, D.B. Scarl, B.T. Ulrich, W.F. Baker, E.W. Jenkins and A.L. Read, *Phys.Rev.* 138, B165 (1965).
- 38) J. Orear, R. Rubinstein, D.B. Scarl, D.H. White, A.D. Krisch, W.R. Frisken, A.L. Read and H. Ruderman, (private communication from J. Orear).
- 39) G. Fast and R. Hagedorn, *Nuovo Cimento* 27, 208 (1963) ;
 G. Fast, R. Hagedorn and L.W. Jones, *Nuovo Cimento* 27, 856 (1963) ;
 A. Bialas and V.F. Weisskopf, *Nuovo Cimento* 35, 1211 (1965) ;
 R. Hagedorn, *Nuovo Cimento* 35, 216 (1965) ; s is in $(\text{GeV}/c)^2$.
- 40) F. Cerulus and A. Martin, *Phys.Letters* 8, 80 (1964) ;
 T. Kinoshita, *Phys.Rev.Letters* 12, 256 (1964) ;
 A. Martin, *Nuovo Cimento* 37, 671 (1965).
- 41) T.T. Wu and C.N. Yang, *Phys.Rev.* 137, B708 (1965).
 A modified and somewhat more intuitive picture of the Wu-Yang form factors can be obtained if one thinks of the baryons as bound states of three quarks and of the mesons as quark-antiquark bound states.

- 42) A. Białas and O. Czyzewski, Phys.Letters 13, 337 (1964).
In Eq. (2) of this paper, a factor 2 is missing in the right-hand side (private communication from Dr. A. Białas).
- 43) H. Harari, Lectures at the Seminar on High Energy Physics and Elementary Particles, International Centre for Theoretical Physics of I.A.E.A., Trieste, May-June 1965, to be published.
- 44) H. Harari and H.J. Lipkin, Phys.Rev.Letters 13, 208 (1964).
- 45) K. Johnson and S.B. Treiman, Phys.Rev.Letters 14, 189 (1965).
- 46) The success of $SU(3)$ and $SU(6)$ symmetry suggests that baryons might be bound states of three triplets, quarks for example, mesons being quark-antiquark bound states. This may give a simple explanation for the large proportion of inelastic two-body collisions at relatively high energies (see Section III), the reason being that in these collisions the incoming composite systems are simply excited, without creation of further quark-antiquark pairs. A discussion of the suppression of particle production in the quark model is given in Z. Koba, Phys. Letters 16, 326 (1965).

FIGURE CAPTIONS

- Figure 1 : Total cross-sections for nucleon-nucleon and meson-nucleon scattering, reproduced from Ref. 5).
- Figure 2 : The ratio X of real to imaginary parts of the forward scattering amplitude for pion-proton scattering, following Ref. 6).
- Figure 3 : The ratio X for proton-proton scattering from various experiments, as given in Ref. 7).
- Figure 4 : The width of the diffraction peak for pp and $\bar{p}p$ as a function of laboratory momentum, following Ref. 13).
- Figure 5 : The diagrams characterizing the peripheral model in its original (a) and extended forms (b).
- Figure 6 : $K^-p \rightarrow \bar{K}^0 n$ differential cross-section at $p_L = 9.50$ GeV/c, following Ref. 33).
- Figure 7 : $\pi^-p \rightarrow \pi^0 n$ differential cross-section at four incident energies, following Stirling et al., Ref. 8).
- Figure 8 : pp large angle elastic scattering, following Ref. 37).
- Figure 9 : πp large angle elastic scattering at $p_L = 8$ and 8.5 GeV/c, see Ref. 38).
- Figure 10 : πp large angle elastic scattering at $p_L = 8$ and 8.5 GeV/c, with special reference to backward region, see Ref. 38).

FIG.1

FIG. 2

FIG 3

OPTICAL MODEL RADIUS FOR PROTONS AND ANTIPROTONS SCATTERING

$$\left(\frac{d\sigma}{dt}\right)/\left(\frac{d\sigma}{dt}\right)_0 = e^{At} = e^{-\left(\frac{R}{2}\right)^2 \cdot t}$$

* p $\bar{p}$
○ p p

FIG. 4

FIG. 5

FIG. 6

FIG. 7

FIG. 8

FIG.9

FIG.10