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Dual Quarks and Parton Quarks

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DUAL QUARKS AND PARTON QUARKS

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Abstract : The constraints of duality are imposed on the structure functions of total as well as one particle inclusive deep inelastic electron-nucleon scattering.

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INTRODUCTION

By the dual quarks we mean the quarks which appear in duality diagrams and by parton quarks we mean partons that are identified with quarks. As is known, these two types of quarks play a very different role in high energy physics. The former was introduced in order to describe the idea of duality, especially in its algebraic sense like absence of exotics in the appropriate channels. The interactions among them are considered to be coherent. The latter on the other hand was introduced as a point like substructure of hadrons in order to explain the scaling in deep inelastic scattering[1)] and the individual current-quark interactions are incoherent. There is another framework called the light cone approach[2)], which turned out to be equivalent[3)] (at least formally) to the usual quark-parton model as far as the total deep inelastic scattering is concerned. However, it cannot describe the single particle inclusive processes in the current fragmentation region, while the parton model was extended quite naturally to describe the latter[4)].

Taking this into account we shall devide our arguments into two parts according to the processes we discuss, namely

(a) total deep inelastic scattering and
(b) single particle inclusive reactions in the current fragmentation region.

As was mentioned above, for the process (a), the quark-parton model has a sounder footing of the light cone approach, while for (b) there is no such basis.

Duality leads us to expect two components in the deep inelastic amplitude (the imaginary part of the forward current-hadron scattering amplitude). One component is restricted to nonexotic representations in both s- and t-channels and the other is associated with diffractive contribution which is dual to the pomeron exchange. In the parton model this corresponds to the decomposition of the quark distribution within a hadron into a valence (non exotic in both s- and t-channels) and a sea (SU(3)-singlet in t-channel) components[5)].

It is possible, however, to reproduce the same results without dealing with quarks explicitly[6)], at least in the case of the total deep inelastic process, which we discuss in section 2 after a brief description of the kinematics. We show that the non-diffractive component of the structure functions are described by only two independent functions, which are nothing but the Fourier transforms of D- and F-coupled nucleon matrix elements of the bilocals. The behaviour of the proton and neutron structure functions near the threshold $x \simeq 1$, determines the behaviour of the D/F ratio there, which in turn leads to a particular structure of s-channel contributions[7]. We find complete absence of contributions from baryon resonances belonging to the abnormal parity trajectory. This implies that $SU(6)_W$-symmetry, where the nucleon (member of normal parity trajectory) belongs to 56-plet together with the Δ-resonance (member of abnormal parity trajectory), cannot be valid in this region and we are lead either to break the symmetry[7,8] itself or to introduce a configuration mixing mechanism if we wish to maintain $SU(6)_W$-symmetry[9]. Similar arguments are applied to deep inelastic scattering from polarized nucleons.

In section 3 we generalize the two component duality idea to the single particle inclusive distributions in deep inelastic scattering[10]. Compared to the unconstrained quark-parton model the number of independent distribution functions is reduced* and we can predict for example the behaviour of the π^+/π^- ratio for both proton and neutron targets. Then we study in our model the problem of the approach to Feynman scaling and the related problem of exoticity criteria[11]. We find that the early approach to Feynman scaling is guaranteed if ab or $ab\bar{c}$ is exotic. This leads us to strong exchange degeneracy relations among inclusive vertex functions. The Regge-Mueller analysis of the inclusive electroproduction of pions with nucleon targets in the current fragmentation region has been performed recently[12] and the authors have found that such relations are in good agreement with experiment.

We conclude in section 4 with a few remarks and discussions.

TOTAL DEEP INELASTIC SCATTERING

- structure of s-channel resonance contributions -

In this section we consider total deep inelastic electron-nuclon scattering ($J_V(q) + N(p) \rightarrow$ anything).

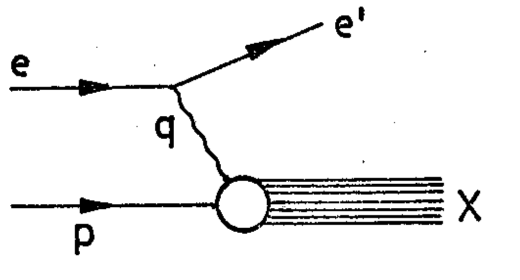

Fig.1

The virtual current carries four-momentum q; p is the momentum of the nucleon target. This cross section can be viewed as the imaginary part of the forward compton scattering amplitude.

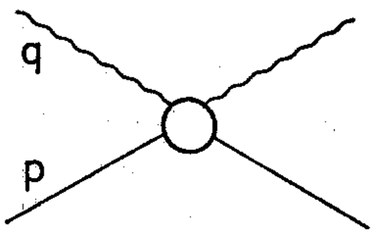

Fig.2

In the Bjorken limit, $\nu = pq$, $-q^2 \to \infty$ with the ratio $x = -\frac{q^2}{2\nu}$ finite, the structure functions scale i. e. become functions of only x.

Let us shortly recapitulate the basic ideas of the dual quark-parton model and the light cone version of it. In the parton model[1] the diagram corresponding to fig. 2 will be

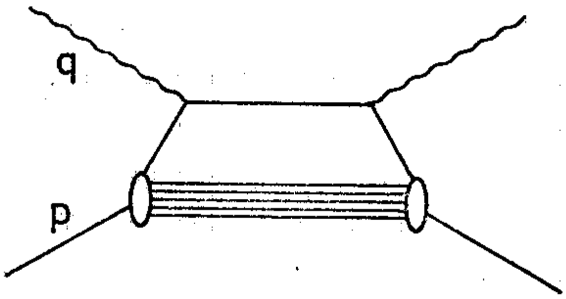

Fig.3

where the nucleon is considered to consist of many point like constituents called partons and the scaling behaviour of the structure functions is explained in terms of incoherent current-parton interactions. The structure function say of proton is expressed in this model in terms of parton distribution functions as

$$F_1^{ep}(x) = \sum_i Q_i^2 \, u_i(x), \tag{1}$$

where i runs over the type of partons, which we identify with quarks and antiquarks. Q_i is the charge of parton of type i and $u_i(x)$ is its longitudinal momentum distribution within a proton.

Two component duality leads to decomposition[5] of the six distribution functions of quarks and antiquarks into those of valence and sea quarks:

$$\begin{aligned} u_p &= v_p + s, \quad u_n = v_n + s \\ u_{\bar{p}} &= u_{\bar{n}} = u_\lambda = u_{\bar{\lambda}} = s, \end{aligned} \tag{2}$$

corresponding to the following two duality diagrams.

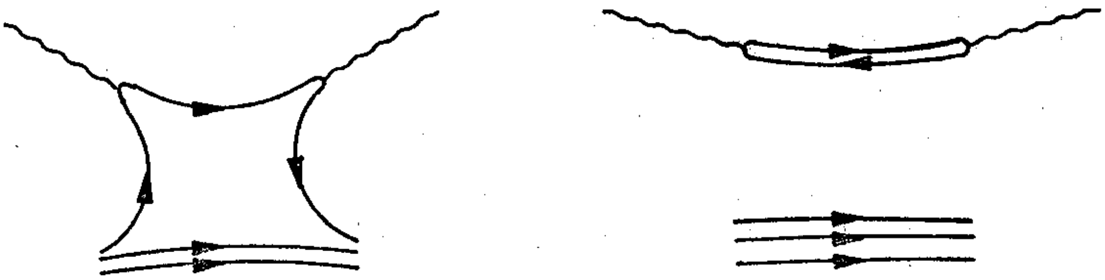

Fig. 4 Two duality components for the forward compton scattering

The structure functions of proton and neutron are expressed now as

$$F_1^{ep}(x) = \frac{4}{9} v_p + \frac{1}{9} v_n + \frac{4}{3} s$$
$$F_1^{en}(x) = \frac{1}{9} v_p + \frac{4}{9} v_n + \frac{4}{3} s \tag{3}$$

In the light cone approach[2] the structure functions are related to the matrix elements of bilocals:

$$\langle \alpha, p | V_\sigma^c(z,0) | \beta, p \rangle = 2 p_\sigma \Big[i\, f_{\alpha c \beta}\, g_F(pz) + d_{\alpha c \beta}(1 - \delta_{co})\, g_D(pz) + d_{\alpha o \beta}\, \delta_{co}\, g_S(pz) \Big]. \tag{4}$$

where

$$V_\sigma^c(x,y) = \bar{q}(x) \gamma_\sigma \frac{\lambda^c}{2} q(y).$$

Here g_F, g_D are F- and D-coupled matrix elements of bilocals. The structure functions are expressed in terms of the following six functions.

$$A_\pm(x) = \pm \frac{1}{4\pi} \int_{-\infty}^{\infty} e^{\pm ix(pz)}\, g_A(pz)\, d(pz), \tag{5}$$

where A stands for F, D or S. These functions correspond to the six quark-parton distribution functions in the parton model (A_+ to quark and A_- to antiquark distributions, see Ref. 6 for details).

Two component duality requires the following relations among these functions[6]:

$$\hat{S}_+ = \frac{1}{2}(3 F_+ - D_+), \quad F_- = D_- = S_- = 0, \tag{6}$$

where S_+ is the non-pomeron part in the decomposition

$$S_+ = \hat{S}_+ + S_+^P. \tag{7}$$

From now on we drop the suffices + and write simply F, D etc. The nucleon structure functions are expressed now as:

$$F_1^{ep}(x) = 2\,F(x) - \frac{2}{9}\,D(x) + \frac{8}{9}\,S^P(x),$$
$$F_1^{en}(x) = \frac{4}{3}\,F(x) - \frac{8}{9}\,D(x) + \frac{8}{9}\,S^P(x). \quad (8)$$

It is interesting to note here that one of the well known Nachtman's inequalities[13],

$$\frac{1}{4} \leq \frac{F^{en}}{F^{ep}} \leq 4, \quad (9)$$

which is obtained from the positivity of the imaginary part of the forward current hadron amplitude in the general quark-parton model or the light cone approach, is replaced now by[6]

$$\frac{1}{4} \leq \frac{F^{en}}{F^{ep}} \leq \frac{3}{2}, \quad (10)$$

after two component duality constraints are imposed. The D/F ratio is correspondingly restricted as follows:

$$1 \geq \frac{D}{F} \geq -3. \quad (11)$$

SU(6)-symmetry leads to the simple valence quark-parton model where $v_n = \frac{1}{2}\,v_p$ or to pure F-type coupling of bilocals to nucleons and the corresponding inequality is

$$\frac{2}{3} \leq \frac{F^{en}}{F^{ep}} \leq 1, \quad (12)$$

which is definitely violated by the experimental results, especially near the threshold $x \simeq 1$, where the ratio seems to approach the lower limit $\frac{1}{4}$ of Eq's (9) or (10).

Now we turn to a more direct duality argument and examine the behaviour of baryon-resonance contributions in the direct channel. For pseudo-scalar meson-baryon scattering the problem of what pattern of s-channel resonances would reproduce the t-channel (Regge) description of the same process, has been exploited some time ago[14]. Two simple solutions were found and the pattern of the baryon resonances dominating in the s-channel was given by the following two exchange degenerate sets (see Fig. 13)

$$8_\alpha(J^P = \frac{1}{2}^+, \frac{5}{2}^+ \ldots) - (8+1)_\gamma\,(J^P = \frac{3}{2}^-, \frac{7}{2}^- \ldots) \quad (13)$$

normal parity series,

$$(10 + 8)_\delta \ (J^P = \frac{3^+}{2}, \frac{7^+}{2} \ldots) - 8_\beta(J^P = \frac{5^-}{2}, \frac{9^-}{2} \ldots) \tag{13}$$

abnormal parity series,

with definite relations among the couplings of resonances of each set to M-B system. An interesting property of the solutions is that these two sets satisfy the constraints of duality separately and the relative importance of them is controlled by the D/F ratio of couplings to baryons of the nonet of Reggeons exchanged in the t-channel.

In the deep inelastic case, we know from experiment that the D/F ratio in the t-channel changes as a function of x, which implies that the pattern of s-channel resonances should be x-dependent[7]. Especially, we find a complete absence of contributions from baryons belonging to abnormal parity series as $x \to 1$.

To see this let us use the language of the quark-parton model constrained by duality (Eq. (2)). We rewrite this expression in terms of slightly different distribution functions v_o and v_1 defined as:

$v_o(x)$ - distribution of a quark, which interacts with the electromagnetic field, when the remainder (two spectator quark system with possible gluons) is in $I = 0$ state.

$v_1(x)$ - the same as above when the remainder is in an $I = 1$ state.

In terms of these functions we have

$$v_p = v_o + \frac{1}{3} v_1, \quad v_n = \frac{2}{3} v_1, \tag{14}$$

and Eq. (3) becomes

$$\begin{aligned} F_1^{ep}(x) &= \frac{4}{9} v_o + \frac{2}{9} v_1 + \frac{4}{3} s, \\ F_1^{en}(x) &= \frac{1}{9} v_o + \frac{1}{3} v_1 + \frac{4}{3} s. \end{aligned} \tag{15}$$

It is now clear that $v_1(x) \to 0$ (and also $s(x) \to 0$) as $x \to 1$, if we assume that the lower limit of Eq. (10) is actually approached in this region. This in turn implies that as $x \to 1$ the spectator quarks are mostly in $I = 0$ (antisymmetric with respect to SU(3)-indices) state, which forbids the decuplet contribution in the s-channel, and with this the contribution from abnormal parity series.

This shows that in this particular region of $x \to 1$ the nucleon does not want to transit to the other series of abnormal parity and prefers to remain in the same series, while in the hadronic processes the transition to both series occurs almost with the same probability. This can also be seen from the fact that the transition form factor of P_{33} ($\Delta(1236)$, $J^P = \frac{3^+}{2}$) which belongs to the abnormal parity series falls faster[15] than those of the normal parity resonances, i. e. elastic form factor, the transition form factors of D_{13} ($N(1520)$, $J^P = \frac{3^-}{2}$) and F_{15} ($N(1688)$, $J^P = \frac{5^+}{2}$). The form factor of D_{15} ($N(1670)$, $J^P = \frac{5^-}{2}$) should also fall rapidly, but the latter does not couple to the photon-nucleon system so strongly.

Now it is clear that $SU(6)_W$ is not a good symmetry to be used in deep inelastic region, because it puts on the same footing the resonances belonging to normal and abnormal parity series - nucleon and Δ belong to the same 56-representation. Thus we should either introduce the $SU(6)_W$ breaking or a configuration mixing where the nucleon is in a mixture of 56- and 70-representation. Similar representation mixing is naturally expected when we think of the nucleon as a 56-plet in the constituent quark basis and transform it to the current quark basis[16]. However, this difficulty of $SU(6)_W$ is due to the pure F-type coupling property of the vector bilocals to the nucleons. This comes from the fact that the three valence quarks in the nucleon are treated symmetrically in $SU(6)_W$ and this property does not change even after the current-constituent transformation is applied (to each quark indices of the nucleon). Therefore we should introduce a configuration mixing in the constituent quark basis[9].

Finally we would like to apply our arguments to the deep inelastic processes with polarized targets[17]. In the light cone approach the structure functions will be related to the matrix element of the axial vector bilocals

$$A^c_\sigma(x,y) = \bar{q}(x)\,\gamma_\sigma\,\gamma_5\,\frac{\lambda^c}{2}\,q(y). \tag{16}$$

In the quark-parton model we decompose the quark distribution functions as

$$u_i = u_i{\uparrow} + u_i{\downarrow} \tag{17}$$

and introduce the distribution functions for the "axial quark number"

$$\bar{u}_i = u_i{\uparrow} - u_i{\downarrow}. \tag{18}$$

These distribution functions correspond to the matrix elements of the axial vector bilocals and noting that $A^e_\sigma(x,x)$ is nothing but the axial vector current we have

$$-\,g_A = \int_o^1 (\bar{v}_p - \bar{v}_n)\,dx. \tag{19}$$

Let us now define the polarization asymmetry, which is measured in these processes[1].

$$A(x) = \frac{d\sigma(\uparrow\downarrow) - d\sigma(\uparrow\uparrow)}{d\sigma(\uparrow\downarrow) + d\sigma(\uparrow\uparrow)}, \tag{20}$$

where $d\sigma(\uparrow\downarrow)$ $(d\,\sigma(\uparrow\uparrow))$ denotes the differential cross sections for scattering in the anti-parallel (parallel) spin configuration and in the quark-parton model the asymmetry can be expressed as

$$A(x) = \frac{\sum \overline{u}_i\ Q_i^2}{\sum u_i\ Q_i^2}. \tag{21}$$

$SU(6)_w$ without configuration mixing tells us that

$$v_o = v_1\ (\equiv v),\quad \overline{v}_o = v \text{ and } \overline{v}_1 = -\frac{1}{3}v. \tag{22}$$

which reproduces the standard results[1]:

$$-g_A = \frac{5}{3},\quad A_p = \frac{5}{9} \text{ and } A_n = 0. \tag{23}$$

However, we know from the previous arguments that $SU(6)_w$ results have no chance to be correct in the region of $x \simeq 1$. Duality on the other hand does not provide us with any information about $\overline{u}_i$'s.

We can try, however, a simple and intuitive assumption[7,8] that the system of spectators tend to be in $J = 0$ state as well as $I = 0$ state as we found before in the region of $x \simeq 1$. This gives immediately

$$A_p,\ A_n \to 1 \quad \text{as} \quad x \to 1. \tag{24}$$

We have seen in this section how duality restricts the form of the structure functions and how it helps us to build a physical picture of deep inelastic processes. We have seen also that $SU(6)_w$, which was approximately correct for the hadronic processes, fails to describe deep inelastic scattering in the region of $x \simeq 1$.

ONE PARTICLE INCLUSIVE SCATTERING

\- strong exchange degeneracy relations among inclusive vertex functions -

In this section we consider the process $J_V(q) + N(p) \to h + \text{anything}$, where the detected hadron carries four momentum h.

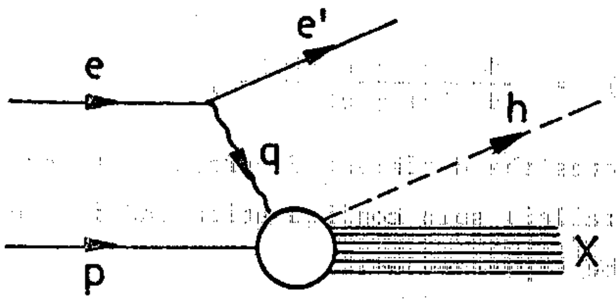

Fig.5 Single particle inclusive electron nucleon scattering

The current fragmentation region is defined[4)] as

$$-q^2 \to \infty, \quad \nu = pq \to \infty, \quad x = -\frac{q^2}{2\nu} \text{ finite,}$$

$$hq \to -\infty, \quad z = \frac{h \cdot p}{\nu} = \frac{2 \cdot hq}{q^2} \text{ finite, } h_T \text{ finite.} \tag{25}$$

It has been proposed by Feynman[4)] that this process occurs in two steps: first the photon is absorbed by a parton whose momentum is a fraction x of the nucleon momentum and second this parton fragments into hadrons which are current fragments and the rest of the partons materialize in the form of hadrons which are target fragments. As a result we have the following picture for this process:

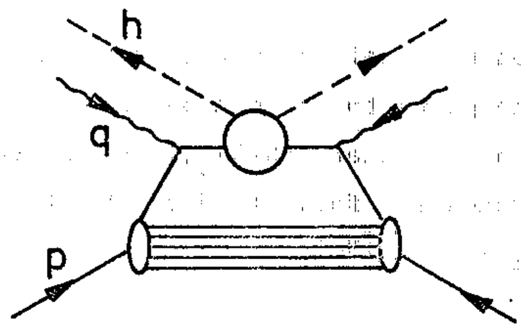

Fig.6 Forward $\gamma N\bar{h} \to \gamma N\bar{h}$ amplitude in the parton model

and the structure function for the inclusive electroproduction of a hadron h from say a proton in the current fragmentation region can be written as

$$L^{ep}_{1,h}(x,z,h_T) = \sum_i Q_i^2 \, u_i(x) \, D_i^h(z,h_T), \tag{26}$$

which is the analogue of $F_1(x)$ in the total deep inelastic process, $D_i^h(z,h_T)$ is the function describing the fragmentation of this quark i into h + anything and h carries a longitudinal momentum which is a fraction z of the quark momentum.

Two component duality leads to decomposition of the distribution functions $u_i(x)$ as is given in Eq. (2) and we should decompose also the fragmentation functions[10] $D_i^h(z,h_T)$ into a sum of two terms

$$D_i^h = V_i^h + S^h, \tag{27}$$

where V_i^h is that part of D_i^h in which the fragmenting quark can be found in the hadron h as a valence quark and S^h is that part in which the quark i ends up anywhere else in the final state. This decomposition corresponds to the following duality diagrams:

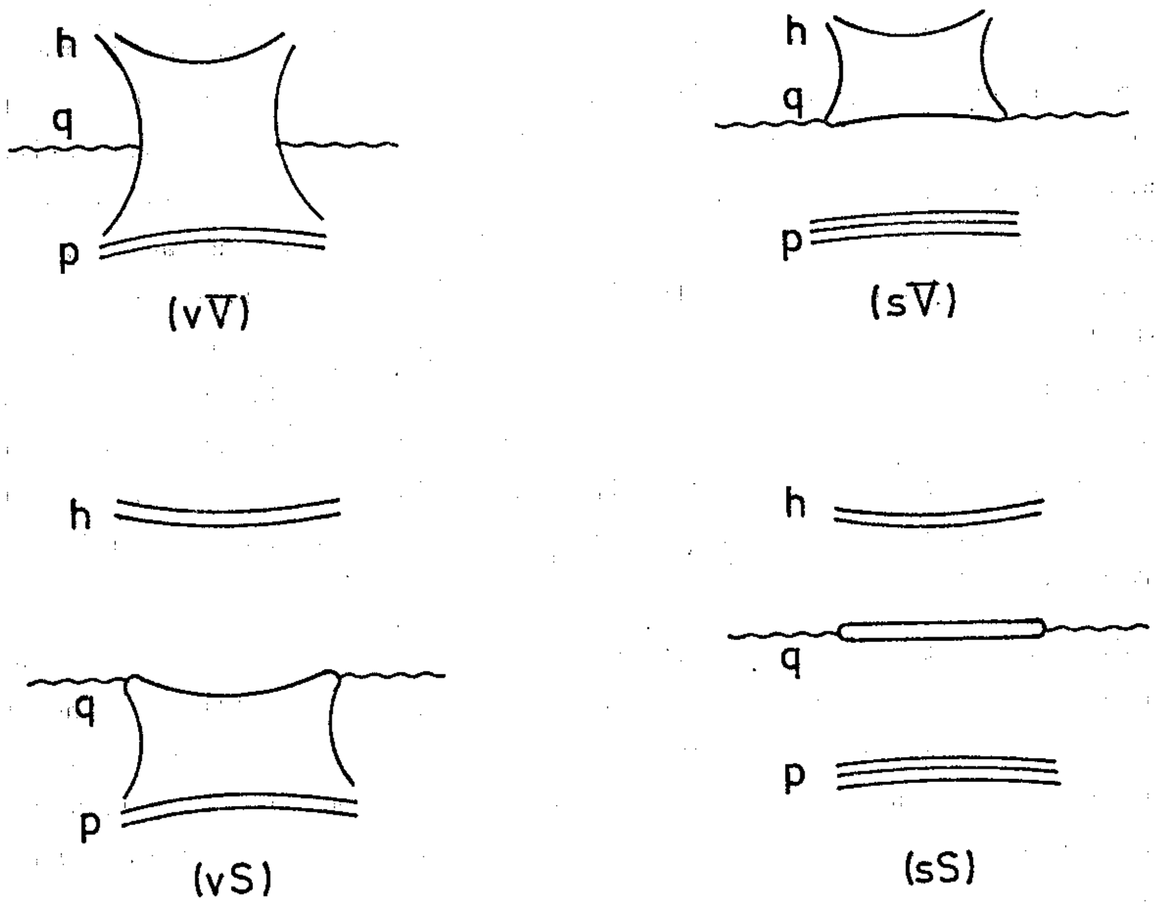

Fig 7. Four duality components for the forward three body amplitude

In the case of pion emission we have

$$D_p^{\pi^+} = V^\pi + S^\pi ,$$
$$D_n^{\pi^+} = D_S^\pi = S^\pi . \quad (28)$$

The other D_i functions are related to those of Eq. (28) by isospin and charge conjugation invariance. Using these relations we can express the structure functions of nucleons as follows:

$$L^{ep}_{1,\pi^+}(x,z,h_T) = (\tfrac{4}{9} v_p + \tfrac{5}{9} s)\, V^\pi + (\tfrac{4}{9} v_p + \tfrac{1}{9} v_n + \tfrac{4}{3} s)\, S^\pi$$
$$L^{ep}_{1,\pi^-}(x,z,h_T) = (\tfrac{1}{9} v_n + \tfrac{5}{9} s)\, V^\pi + (\tfrac{4}{9} v_p + \tfrac{1}{9} v_n + \tfrac{4}{3} s)\, S^\pi$$
$$L^{en}_{1,\pi^+}(x,z,h_T) = (\tfrac{4}{9} v_n + \tfrac{5}{9} s)\, V^\pi + (\tfrac{4}{9} v_n + \tfrac{1}{9} v_p + \tfrac{4}{3} s)\, S^\pi \quad (29)$$
$$L^{en}_{1,\pi^-}(x,z,h_T) = (\tfrac{1}{9} v_p + \tfrac{5}{9} s)\, V^\pi + (\tfrac{4}{9} v_n + \tfrac{1}{9} v_p + \tfrac{4}{3} s)\, S^\pi$$

From these expressions we immediately see that the π^+/π^- ratio for the proton is always larger than unity and increases as $x \to 1$, which is compatible with what is observed experimentally[18] (see Fig. 8). An interesting argument can be given for $(\pi^+/\pi^-)_n$. Namely we observe that this ratio crosses unity at a particular value of x, independent of z and h_T. This value can be obtained from total deep inelastic processes, i. e. we should find the point x where $4\, v_n = v_p$ or $D/F = 1/2$ which corresponds to the point where $F^{en}(x)/F^{ep}(x) = 8/17$ when the contribution from sea is neglected. We find from Fig. 10 $x \simeq 0.6 \sim 0.7$. If we take into account the effect from sea component we have only an inequality $x \le 0.6 \sim 0.7$. On the other hand experimental data on $(\pi^+/\pi^-)_n$ seem to cross unity around $x \sim 0.2$ (see Fig. 9).

The rest of this section will be devoted to the problem of the approach to Feynman scaling encountered in the inclusive reactions. Although the problem is still unsettled in the central region, the criterion, $ab\bar{c}$ exotic, for early scaling in the fragmentation region of $ab \to c\, X$ seems to work quite successfully. (In our case $a = \gamma_v$, b = nucleon, c = produced hadron.)

In the Regge-Mueller approach[20], the invariant cross section for the inclusive process $\gamma_v + N \to c$ + anything in the photon fragmentation region can be written as

$$f(\gamma \overset{N}{\rightarrow} c) \;=\; P + \Big[\beta_N^f\, F_f(\gamma \rightarrow c) + \beta_N^\rho\, F_\rho(\gamma \rightarrow c) + \beta_N^\omega\, F_\omega(\gamma \rightarrow c) + + \beta_N^{A_2}\, F_{A_2}(\gamma \rightarrow c) \Big]\, \nu^{-\frac{1}{2}} \qquad (30)$$

where P represents the Pomeron contribution and β_N^i and $F_i(\gamma \rightarrow c)$ are the Regge couplings $(i = f, \rho, \omega, A_2)$ to the nucleon and the fragmentation verteces respectively. Feynman scaling implies that P and the $F_i(\gamma \rightarrow c)$ are functions of q^2, x_F and h_T, where x_F is the usual Feynman scaling variable.

Two component duality extended to the forward elastic three-body amplitude $ab\bar{c} \rightarrow ab\bar{c}$ says that the following four duality diagrams contribute to the energy dependent part of the cross section for the fragmentation process $b \overset{a}{\rightarrow} c$.

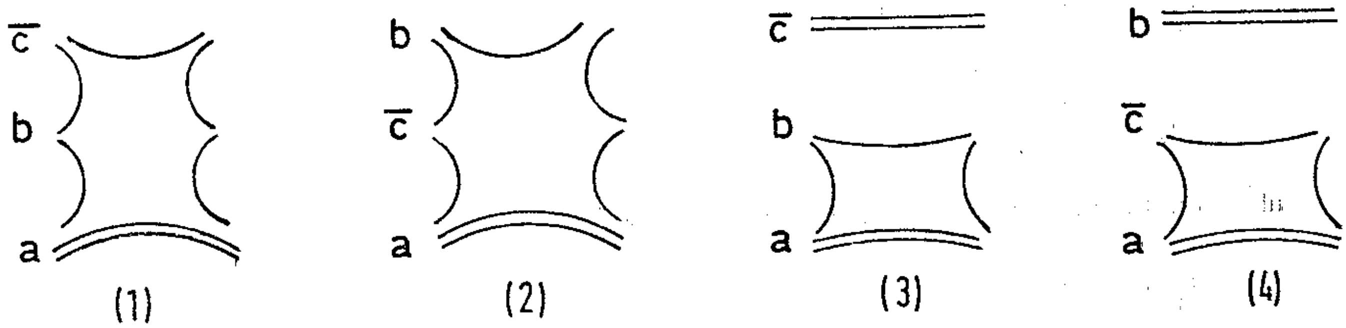

Fig.11 Four duality components contributing to the energy dependent part of the fragmentation process $b \xrightarrow{a} c$

(1) and (3) correspond to νV and νS terms in the quark-parton model, while there is no corresponding term to (2) and (4).

$ab\bar{c}$ exoticity criterion for rapid scaling, which works quite satisfactory in the fragmentation regions, implies that the diagrams (3) and (4) are small compared to (1) and (2). If we now assume according to the dual quark-parton model[10] that the contribution from (2), which is completely absent there, can be neglected, we are left with only one diagram (1) and the criterion for rapid scaling now becomes: "$ab\bar{c}$ or ab exotic". This new criterion obviously gives more stringent exchange degeneracy relations among the inclusive vertex functions, which we are going to study for the case of $\gamma_V \overset{N}{\rightarrow} \pi$.

As is known, isospin relates the nucleon vertices.

$$\beta^f_p = \beta^f_n, \quad \beta^\rho_p = -\beta^\rho_n, \quad \beta^\omega_p = \beta^\omega_n, \quad \beta^{A_2}_p = -\beta^{A_2}_n. \tag{31}$$

Exchange degeneracy relations from the two-body amplitudes gives

$$\beta^f_N = \beta^\omega_N, \quad \beta^\rho_N = \beta^{A_2}_N, \tag{32}$$

and these two are related to D- and F-couplings of the exchange degenerate nonets of Reggeons to nucleons as follows:

$$\frac{\beta^f_N}{\beta^{A_2}_N} = \left[\frac{3\,F - D}{F + D}\right]_{\text{Reggeon}} . \tag{33}$$

For inclusive vertex functions we have similarly

$$\begin{aligned} F_f\,(\gamma \to \pi^+) &= F_f\,(\gamma \to \pi^-), & F_\rho(\gamma \to \pi^+) &= -F_\rho(\gamma \to \pi^-), \\ F_{A_2}(\gamma \to \pi^+) &= F_{A_2}(\gamma \to \pi^-), & F_\omega(\gamma \to \pi^+) &= -F_\omega(\gamma \to \pi^-). \end{aligned} \tag{34}$$

$ab\bar{c}$ exoticity criterion gives further

$$F_f\,(\gamma \to \pi^-) = F_\rho\,(\gamma \to \pi^-), \quad F_\omega(\gamma \to \pi^-) = F_{A_2}(\gamma \to \pi^-) . \tag{35}$$

On the other hand we know that the energy dependent part of the cross section is described by only one diagram (1) of Fig. 11 and this allows us to express the coupling of the inclusive vertices as:

$$F_V(B \to C)\; T_r\left[(B\bar{C}C\bar{B} - \bar{B}C\bar{C}B)\; V\right] + F_T(B \to C)\; T_r\left[(B\bar{C}C\bar{B} + \bar{B}C\bar{C}B)\; T\right], \tag{36}$$

where in our care B is the virtual photon, C is the pion and V and T are the Reggeons exchanged in the t-channel. B, C, V and T are 3 x 3 matrices of SU(3). The usual exchange degeneracy of Eq. (35) implies $F_V = F_T$. Therefore Eq. (36) leads us to conclude that all the inclusive verteces are expressed in terms of only one function and we obtain in addition to the usual exchange degeneracy relation of Eq. (35) the relation

$$\frac{F_{A_2}(\gamma \to \pi)}{F_f\,(\gamma \to \pi)} = \frac{3}{5} . \tag{37}$$

This relation is reminiscent of the similar SU(3) relation between $\gamma\gamma A_2$ and $\gamma\gamma f$ verteces which can be written as

$$\frac{\beta^{\gamma}_{A_2}}{\beta^{\gamma}_{f}} = \frac{3}{5} , \qquad (38)$$

however, for the inclusive vertex functions, Eq. (37) can not be derived from SU(3) alone and we need another condition of absence of the diagram (2) of Fig. 11. Otherwise the ratio of Eq. (37) could have been any value. Regge-Mueller analysis of this process has been performed recently[12). Eq. (37) enables us to predict the inclusive cross sections of π^+ and π^- on neutron using the data of proton, because we are left with only two independent functions P and F for four cross sections. The results are shown in Fig. 12. The predictions are in very good agreement with experiment.

In this section we extended the ideas of the previous section to one particle inclusive processes. We have seen that our dual quark-parton model gives further constraints to the usual Regge-Mueller analysis of the inclusive process, which is satisfied by experiment.

CONCLUSIONS

We said at the beginning that dual quarks and parton quarks were introduced to describe very different phenomena in high energy physics. Duality, which is most vividly expressed in terms of dual quarks, has been abstracted from pure hadronic dynamics and the number of dual quarks within a nucleon is of course three. Partons, which we identify with quarks, were introduced to describe scaling in deep inelastic processes and an essential element of this model is incoherence of the individual current-parton interactions and we expect a great number of partons to exist within a nucleon. Thus we expect that very different phenomena compared with the usual strong interaction region are occurring in this deep inelastic region. However, we have seen that many of the ideas of strong interaction region like duality are applicable also in this region and we also know from experiment that a simple picture with only three quarks works surprisingly well[21)] in this region i. e. contributions from sea quarks are restricted to very small region of x. In a sense experiment forces us to change gradually the original picture of the parton model and we find that even in this new region of deep inelastic scattering the old ideas of strong interaction are very powerful.

There was one exeption. $SU(6)_W$ symmetry which was approximately valid in the usual strong interaction region turned out to be completely broken in a particular region of deep inelastic scattering, namely near $x \simeq 1$. The situation is not curable even by distinguishing constituent and current quarks. Thus we are lead either to break $SU(6)_W$ or to introduce a configuration mixing or even altogether to abandon it.

For single pion inclusive processes in the current fragmentation region we have obtained a relation between the inclusive vertex functions $F_{A_2}(\gamma \to \pi)$ and $F_f(\gamma \to \pi)$, which is in good agreement with experiment in deep inelastic region. It would be interesting to see whether this relation is satisfied or not in the photoproduction experiments, where there is no reason this to be the case.

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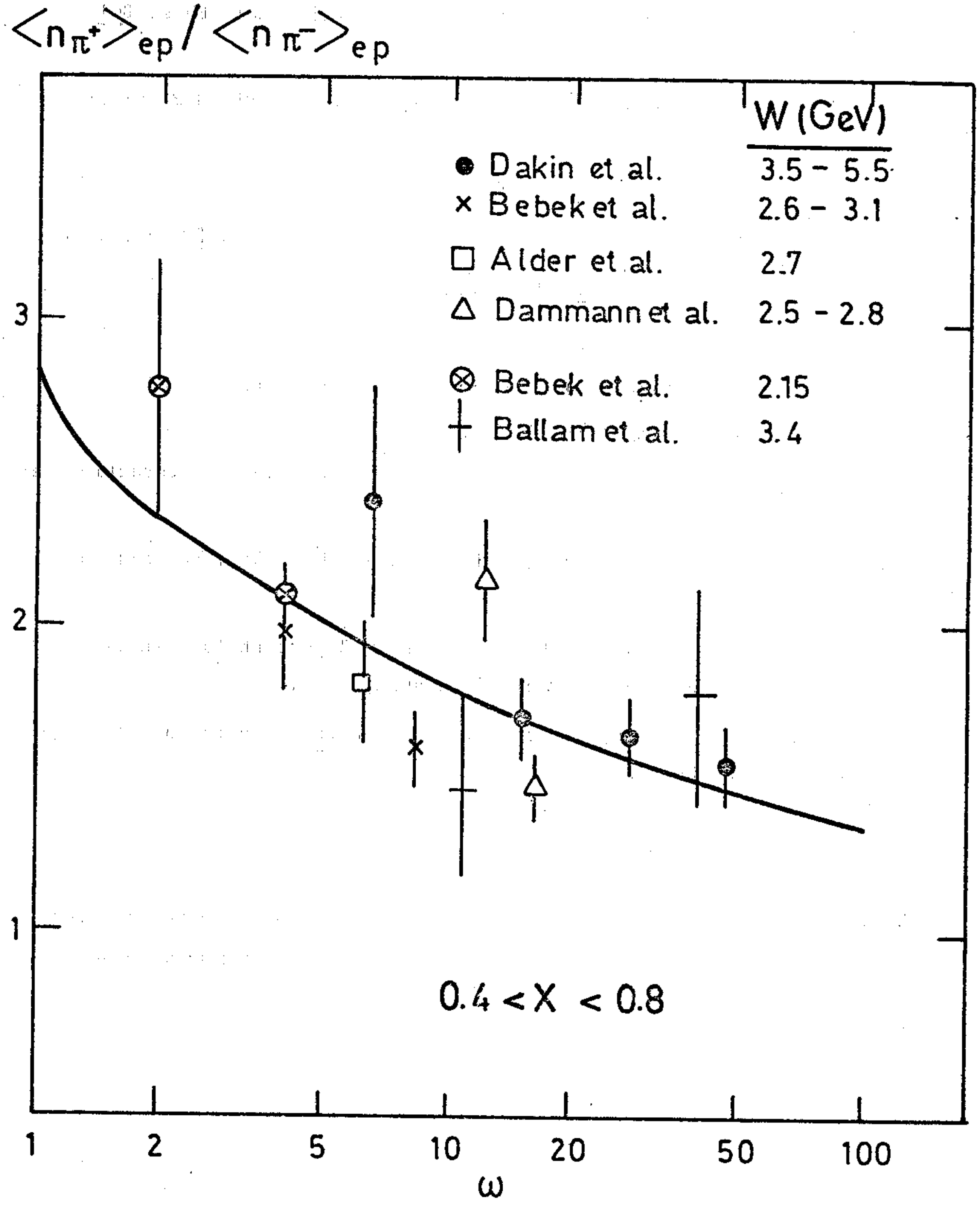

Fig. 8 Charge ratio for scattering on a proton versus ω

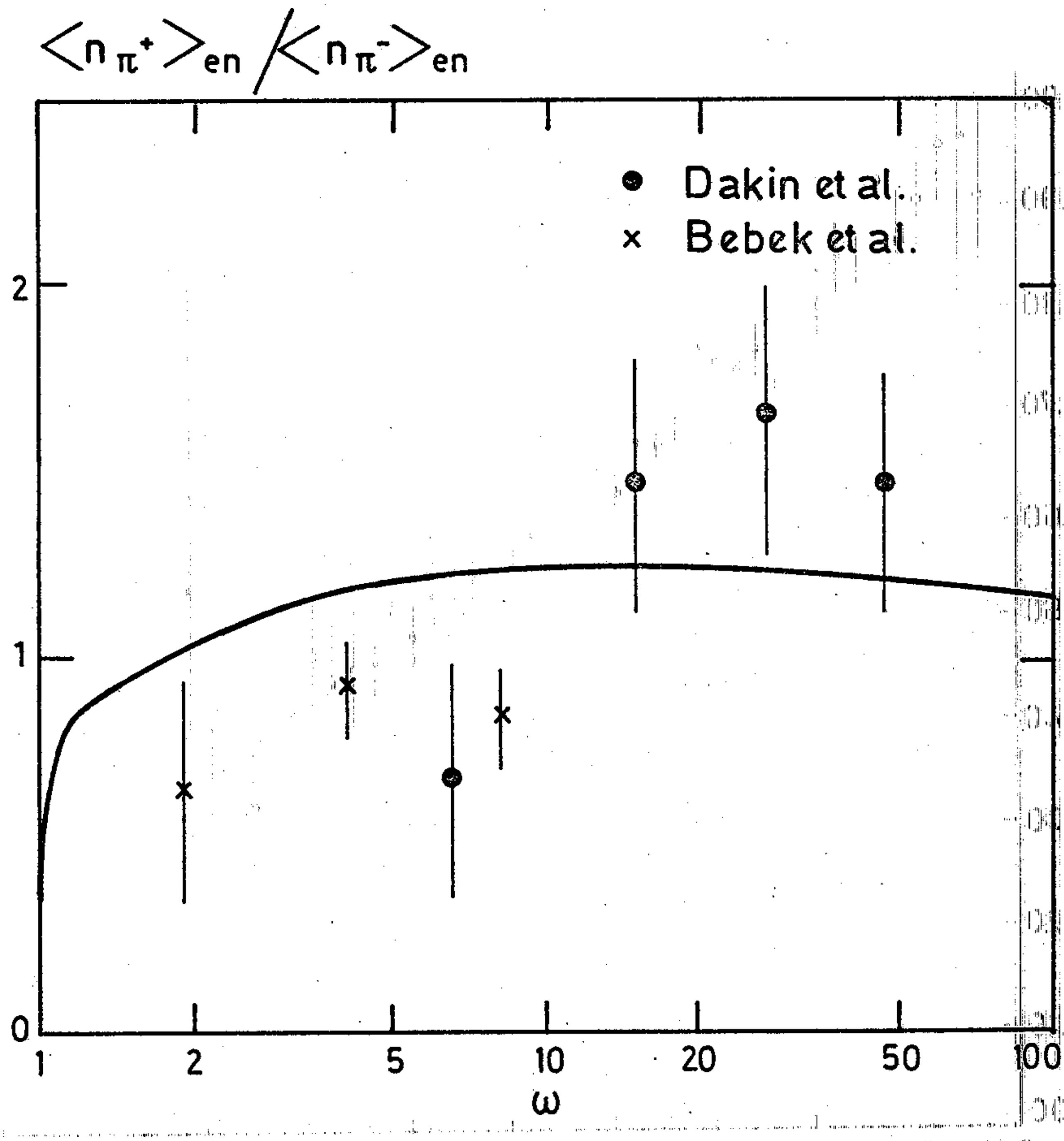

Fig.9 Charge ratio for scattering on a neutron versus ω

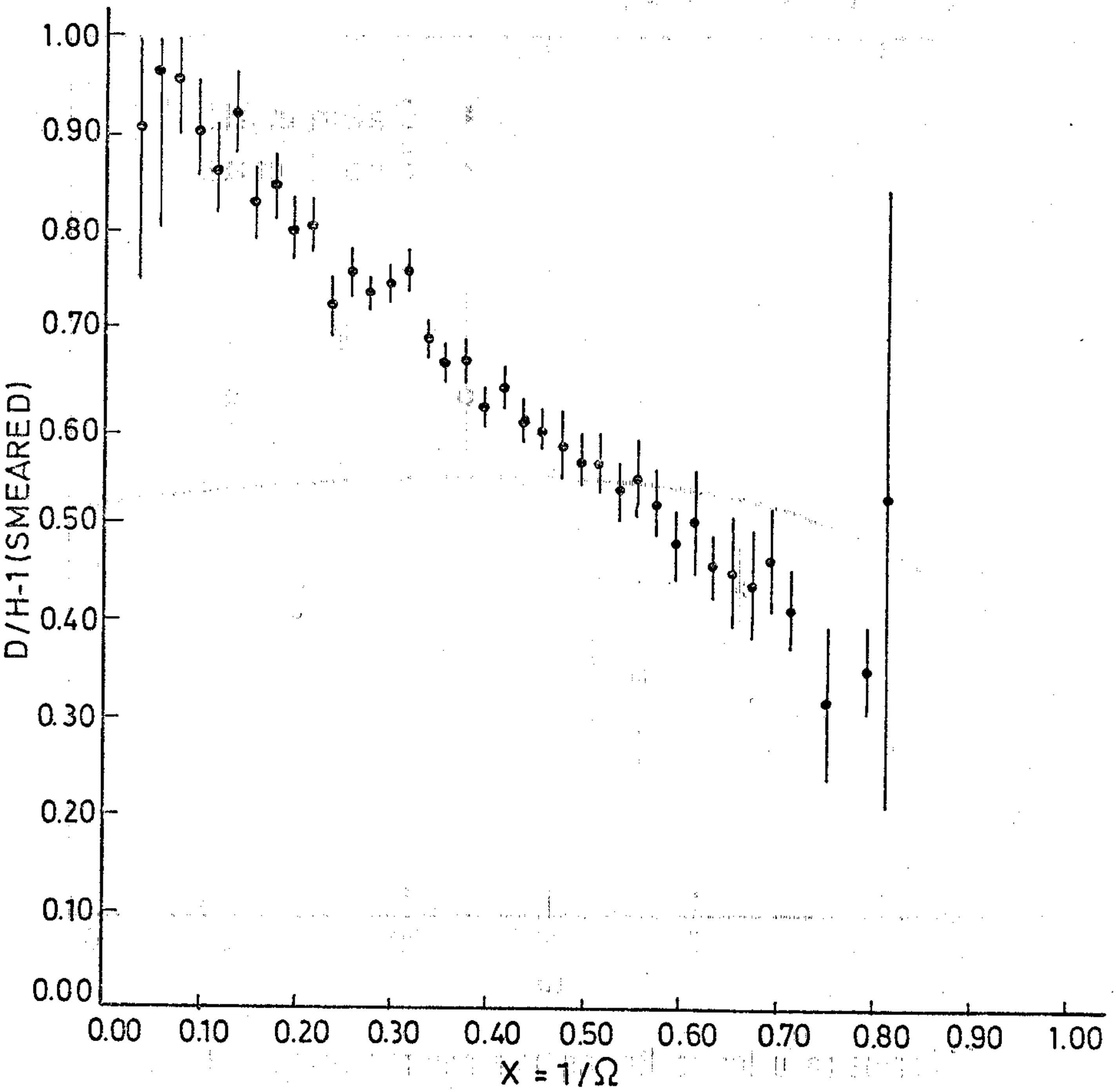

Fig.10 F_2^n / F_2^p

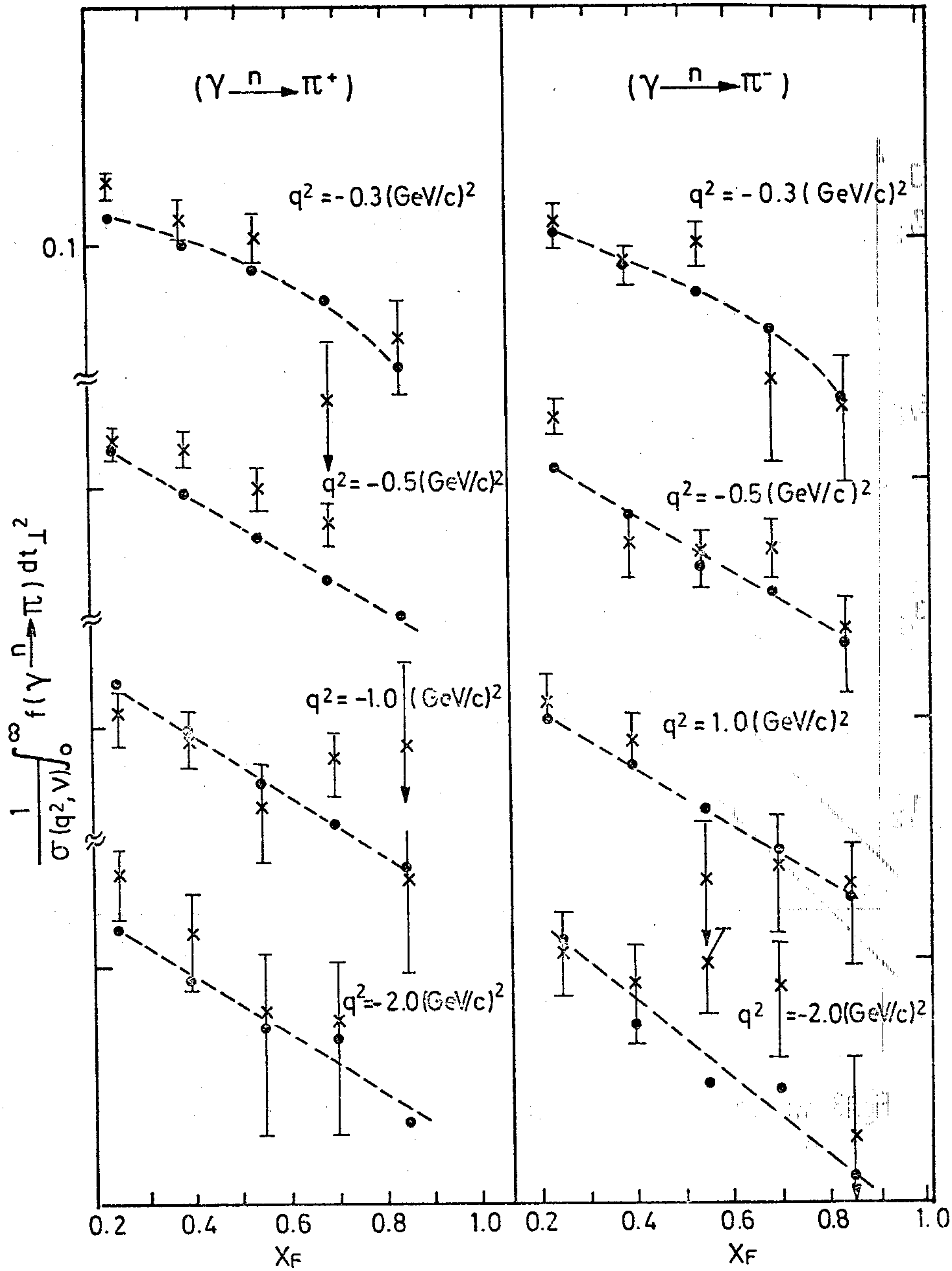

Fig. 12 The x_F distributions of the neutron target cross sections. The dotted lines are the predictions

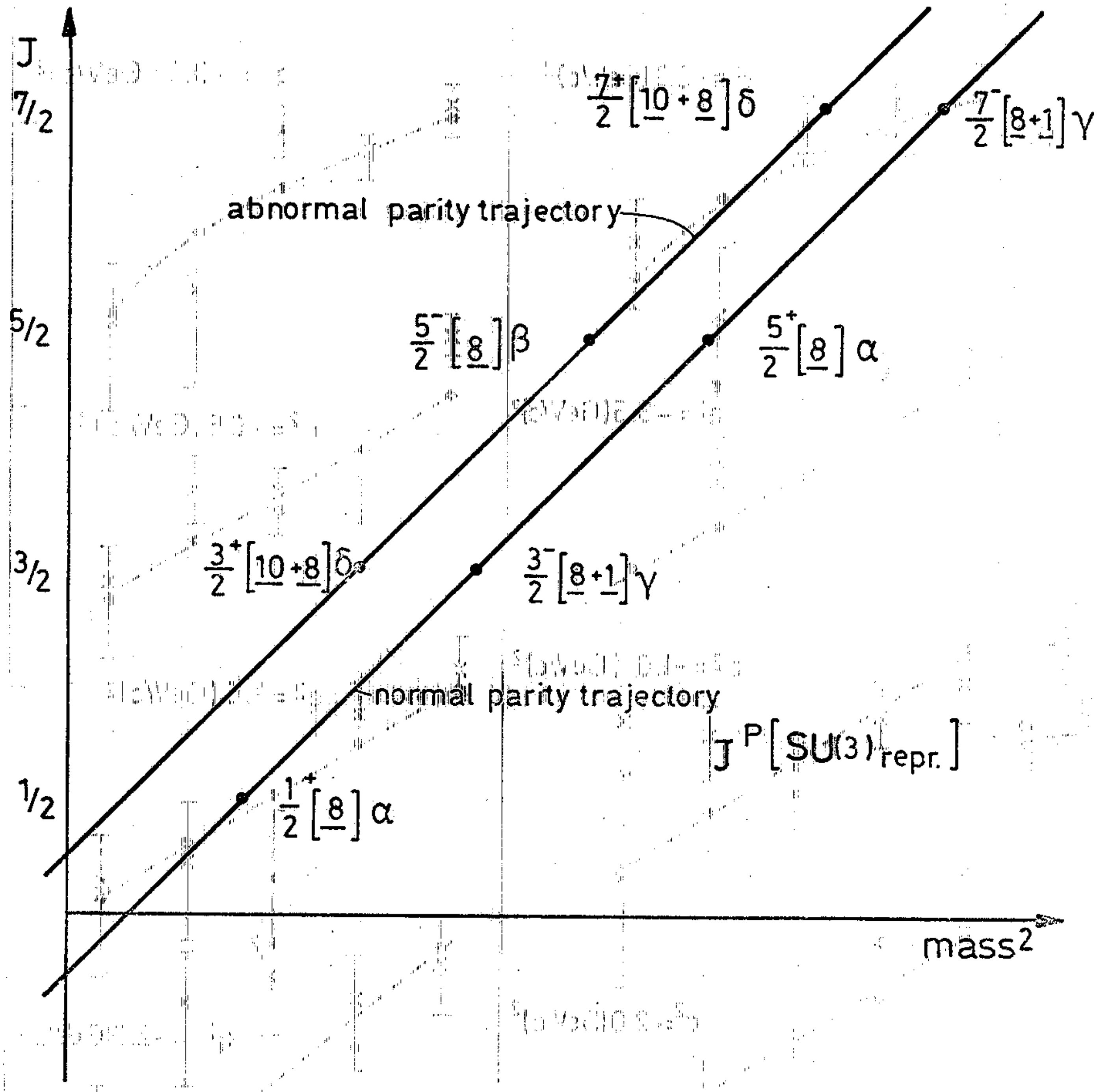

Fig. 13 Two sets of exchange degenerate $\alpha - \gamma$ and $\beta - \delta$ trajectories with normal and abnormal parity respectively.