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B.R. Kim, R. Rodenberg
and
H. Schlereth

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by

R. Rodenberg**

Centro Brasileiro de Pesquisas Físicas

and

B.R. Kim*, and H. Schlereth

III. Physikalisches Institut Technische Hochschule Aachen

Aachen, W-Germany

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ABSTRACT

Within the framework of Harari's conjecture and the constraints provided by duality we investigate the consequences of a non - SU(3) singlet Pomeron in deep inelastic lepton - hadron scattering. The corresponding distribution functions of the quark parton model are discussed.

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** On leave of absence from III. Physikalisches Institut der TH-Aachen, Jaegerstrasse Aachen, Germany (F.R.).

I - INTRODUCTION

The present experimental data[1] for deep inelastic electron nucleon scattering seem to the indicate that the sum rule of Gottfried is violated[2]. As in the Q.P.M. (Quark Parton Model) the Gottfried integral is a measure for the non diffractive contributions coming from the sea of quark - antiquark pairs ($q\bar{q}$ - sea) wich is supposed to make up the nucleon in addition to the three quarks of the low energy quark model, this can be interpreted such that the $q\bar{q}$ - sea does carry quantum numbers different from those of the vaccuum.

Harari[3] has put forward the hypothesis that the $q\bar{q}$ - sea is related to the Pomeron singularity known from high energy hadron - hadron scattering.

To formulate this hypothesis it is not necessary to invoke the Regge description, it suffices to make use of the fact that any amplitude can be decomposed into a resonant part and a background part. Utilizing strong interaction duality it is then clear that Harari's conjecture should apply to the background term in general, not only in a specific kinematical limit.

So adopting Harari's idea does not imply the validity of the Regge description of the deep inelastic data. Although the data indicate a Regge - like behaviour for the case where the parton momentum is a small fraction x of the nucleon momentum there are good reasons to believe that this will not remain true down to the wee region[11], at least not with a Pomeron being a simple pole wich leads to the by now well known 1/x behaviour of the parton distribution functions. One would expect this behaviour to be changed in the wee region such that the number of wee partons is finite[11]. Especially when the $q\bar{q}$ - sea is not neutral with respect to isospin[2] one must be prepared not to extrapolate the data to small and wee x with the usual Regge formula, because then the Gottfried integral does simply not exist.

So it is necessary to stress that Hararis conjecture although originally formulated in the framework of the Regge model can be expressed in a model independent general way.

Within the Q.P.M. this means especially that our formulae are valid for $x \to 0$, as well as for $x \to 1$.

When equations show up which are only correct in the Regge regime we will explicitly say this.

We will continue to use the term Pomeron, but it is understood that it means the background contribution defined in every kinematical region.

If we accept Hararis conjecture in conjunction with a $q\bar{q}$ - sea not neutral with respect to inner symmetries it has the consequence that the background and hence ã forteriori the usual Pomeron have an additional nondiffractive component. This conclusion fits well into what is known from strong interaction physics. There it turns out that the cross sections $\sigma_{tot}(\pi^+p)$ and $\sigma_{tot}(k^+p)$ are not equal at high energies, so that the Pomeron cannot be an SU(3) singlet[5].

Summing up what we are going to do in this paper is the following: In the framework of the duality constraints for deep inelastic lepton-hadron inclusive processes as formulated in ref. 4 we introduce non-SU(3)-singlet Pomeron terms (terms which are not constrained by duality) in addition to the SU(3) - - singlet Pomeron term.

We find (among others) that we can reproduce the sum rule of Adler,and we are able to fit the Gottfried sum rule to its experimental value.

We would like to mention that it can be concluded from the work of ref.

4 that there is a close connection between the sum rule of Gottfried and the SU(3) properties of the Pomeron. This conclusion works on a rather general basis without invoking the Parton Model.

Most easily this can be seen by leaving aside in ref. 4 the conditions coming from the light cone algebra of bilocal currents[10]. Then the Gottfried sum rule follows if the Adler sum rule is required from the SU(3) - singlet character of the Pomeron.

To have an easy way of speaking we introduce the following definitions:

P^V: SU(3)-singlet Pomeron part
(identical with the conventional Pomeron)

P^Q: Non-SU(3)-singlet Pomeron part

The paper is organized as follows:

In section II we fix our notations and kinematics. In section III we specify the properties of the P^Q so far, as we need for our purpose. Then using duality constraints we express the structure functions in terms of reduced matrix elements $a_i(FF)$, $a_i(DD)$ etc. In section IV we discuss our results.

II - NOTATION AND KINEMATICS

We study the absorptive parts of the scattering amplitudes of the process:

$$J^a(q) + h_\alpha(p) \to J^b(q) + h_\beta(p) \qquad (1)$$

($J^{a,b}$: currents, $h^{\alpha,\beta}$: hadrons; α,β, a,b: SU(3) indices) in the forward direction:

$$W_{\mu\nu} = \frac{1}{4\pi} \int e^{iqx} \langle h_\alpha(p)| \left[J^a_\mu(x), J^b_\nu(0)\right] | h_\beta(p) \rangle d^4x; \qquad (2)$$

Confining ourselves to those parts which are measured in electron resp. neutrino reactions modulo $O(m_{lepton})$ their expansion in Lorentz invariants reads (average over target spin is taken):

$$W_{\mu\nu} = W_1(\nu,q^2) \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}\right) + M_p^{-2} W_2(\nu,q^2) \left(p_\mu - \frac{pq}{q^2} q_\mu\right) \left(p_\nu - \frac{pq}{q^2} q_\nu\right) +$$

$$+ W_3(\nu,q^2) \left(\frac{-i\varepsilon_{\mu\nu\alpha\beta} p^\alpha q^\beta}{2M^2}\right) ; \qquad (3)$$

with $\nu = pq$, and the states covariantly normalized $\langle p|p'\rangle = (2\pi)^3 \, 2E.\delta^{(3)}(\vec{p}-\vec{p}\,')$ The metric is (1,-1,-1,-1).

Experimentally these functions seem to approach limiting forms for large values of ν and q^2 depending only on the dimensionaless variable: $\omega: = \frac{-2\nu}{q^2} = \frac{1}{x}$ ($\omega > 1$ by kinematics).

So in the limit $\nu \to \infty$, ω fixed (Bjorken Limit) we expect

$$W_1(\nu,q^2) \to F_1(\omega)$$

$$\nu W_2(\nu,q^2) \to F_2(\omega) \qquad (4)$$

$$\nu W_3(\nu,q^2) \to F_3(\omega)$$

In the following we assume the Callan-Gross relation

$$2F_1(\omega) = \omega F_2(\omega) \qquad (5)$$

which is valid for spin $\frac{1}{2}$ - Parton models. In these models there are further relations between $F_1(\omega)$ and $F_3(\omega)$ which are displayed explicitly in ref. 4) to which we refer the reader. We further generally ignore contributions from the Cabbibo angle.

III - DEFINITION OF THE POMERON PROPERTIES AND DUALITY CONSTRAINTS

Expressing the SU(3) structure of $F^{ab}_{i,\alpha\beta}$ in the form (Dictated by group theory);

$$F^{ab}_{i,\alpha\beta} = a_i(FF)\ if^{abc}\ if^{\alpha\beta c} + a_i(FD)\ if^{abc}\ d^{\alpha\beta c} + a_i(DF)\ d^{abc}\ if^{\alpha\beta c} +$$

$$+ a_i(DD)\ (1-\delta^{co})\ d^{abc}\ d^{\alpha\beta c} + a_i(SS)\ \delta^{co}\ d^{abo}\ d^{\alpha\beta o}\ , \qquad (6)$$

duality will give constraints for the functions $a_i(FF)$, $a_i(FD)$, $a_i(DF)$, $a_i(DD)$, $a_i(SS)$ which for the case of an SU(3)-singlet Pomeron are explicated in ref. 4).

Now we will introduce the basic properties of p^Q;

a) p^Q should dominate over all resonance contributions in the high energy region. (In Q.P.M. this means dominance for small and wee x)

b) p^Q is assumed to be non dual to resonances.

These are the non controversial properties of the conventional Pomeron p^V, which cannot be said of its SU(3) structure as already mentioned in the introduction.

We conjecture now that the Pomeron will not only contribute to the SU(3) singlet parts $a_1(SS)$ and $a_2(SS)$ in eq. (6) but also in some of the octet coefficents $a_i(FF)$. etc. To decide where such non-singlet Pomeron contributions are possible we proceed in the following way:

In order not to invalidade the relation

$$F^{ep}(x) - F^{en}(x) \to 0 \qquad \text{for } x \to 0 \tag{7}$$

which seems to be supported by the present data[1] the only possibility is that p^Q does not contribute to the electromagnetic structure functions at all. So p^Q can only show up in $F_i^{\nu,\bar{\nu}}$, $i = 1,2,3$, the data of which are presently not available.

With the usual SU(3) singlet Pomeron one has the following relation[4]

$$\begin{aligned} &F_i^{\nu p}(x) - F_i^{\bar{\nu} p}(x) \underset{x \to 0}{\to} 0 \qquad \text{for } i = 1,2 \\ &F_i^{\nu p},\ F_i^{\bar{\nu} p} \to \text{const} \neq 0 \\ &F_3^{\nu,\bar{\nu}}(x) \underset{x \to 0}{\to} 0 \end{aligned} \tag{8}$$

Now we assume that the non-neutrality of the $q\bar{q}$-sea with respect to quantum numbers of inner symmetries will show up in the structure functions in the form

$$\begin{aligned} &\lim_{x \to 0} (F_i^{\nu p}(x) - F_i^{\bar{\nu} p}(x)) = \text{const} \neq 0 \quad \text{for } i = 1,2 \\ &\lim_{x \to 0} F_3^{\nu p}(x) = \lim_{x \to 0} F_3^{\bar{\nu} p}(x) = \text{const} \neq 0. \end{aligned} \tag{9}$$

With this assumption and the relation (7) there is only one way of introducing p^Q, namely p^Q contributes to the following functions only:

$$a_i(SS) \qquad i = 1,2 \tag{10-a}$$

$$a_i(FF) \qquad i = 1,2 \tag{10-b}$$

$$a_3(SS) \tag{10-c}$$

As only the resonant parts are constrained by duality we separate out the pomeron contributions:

$$\begin{aligned} a_{1,2}(SS) &= a^P_{1,2}(SS) + a^R_{1,2}(SS) \\ a_{1,2}(FF) &= a^P_{1,2}(FF) + a^R_{1,2}(SS) \\ a_3(SS) &= a^P_3(SS) + a^R_3(SS) \end{aligned} \tag{11}$$

a^P_i denotes the non resonant continuum, hereafter refered as pomeron where as a^R_i denotes resonance contributions.

$a^P_1(SS)$ are the contributions from the conventional pomeron p^V whereas $a^P_{1,2}$ (FF) and $a^P_3(SS)$ are the now additionally appearing terms of p^Q.

Duality imposes the condition that the S-channel non pomeron part must be non exotic. This means that in the case of Baryon target no 27 and $\overline{10}$ representations are allowed to appear.

This implies the following set of constraints:

$$\text{(no 27)} \quad \frac{2}{3}\, a^R_i(SS) + \frac{1}{3}\, a_i(DD) + a^R_i(FF) = 0 \quad \text{for } i = 1,2,3 \tag{12}$$

$$\text{(no } \overline{10}\text{)} \quad \frac{2}{3}\, a^R_i(SS) - \frac{2}{3}\, a_i(DD) + a_i(FD) + a_i(DF) = 0 \qquad \text{for } i = 1,2,3 \tag{13}$$

The quark model rule that currents with (ϕ,f') - quantum numbers only couple to particles with strangeness give the further relation[4]

$$\frac{2}{3}\, a_i^R(SS) + a_i(DF) + \frac{1}{3}\, a_i(DD) = 0 \qquad i = 1,2,3 \tag{14}$$

Compatibility of the general form of eq. (6) with the quark parton model finally gives the following equations [4]:

$$\begin{aligned} a_1(FF) &= -\frac{1}{2}\, a_3(DF) \\ a_1(FD) &= -\frac{1}{2}\, a_3(DD) \\ a_1(DF) &= -\frac{1}{2}\, a_3(FF) \\ a_1(DD) &= -\frac{1}{2}\, a_3(FD) \end{aligned} \tag{15}$$

Using the eqs. (12) to (15) we can express the structure functions in the following way:

$$\begin{aligned} F_1^{ep} &= \frac{8}{9}\, a_1^p(SS) - \frac{2}{9}\, a_1(DD) - 2a_1^R(FF) \\ F_1^{en} &= \frac{8}{9}\, a_1^p(SS) - \frac{8}{9}\, a_1(DD) - \frac{4}{3}\, a_1^R(FF) \\ F_1^{\nu p} &= \frac{8}{3}\, a_1^p(SS) + 2\, a_1^p(FF) - 4\, a_1(DD) - 4\, a_1^R(FF) \\ F_1^{\bar{\nu} p} &= \frac{8}{3}\, a_1^p(SS) - 2\, a_1^p(FF) - 9\, a_1^R(FF) \\ F_3^{\nu p} &= \frac{8}{3}\, a_3^p(SS) + 8\, a_1(DD) + 8\, a_1^R(FF) + 4\, a_1^p(FF) \\ F_3^{\bar{\nu} p} &= \frac{8}{3}\, a_3^p(SS) + 16\, a_1^R(FF) + 4\, a_1^p(FF) \end{aligned} \tag{16}$$

In the rest of the paper we mainly discuss the consequence of this system of equations.

IV - DISCUSSIONS

IVa) Pomeron dominance region. ($x \to 0$)

For small x where the Pomeron terms dominate we have the following relations between the structure functions

$$F_1^{ep}(x) - F_1^{en}(x) = 0$$

$$F_3^{\bar{\nu}p}(x) - F_3^{\nu p}(x) = 0 \quad \text{and} \quad F_3^{\bar{\nu}p} \neq 0, \quad F_3^{\nu p} \neq 0 \tag{17}$$

$$F_1^{\nu p}(x) - F_1^{\bar{\nu}p}(x) = 4\, a_1^p(FF)$$

The expressions $F_1^{\nu p}$, $F_1^{\bar{\nu}p}$ from eq. (16) give us the possibility to make a statement about the x-dependence of $a_1^p(FF)$ in the small x-region:

$$F_1^{\nu p}(x) = \frac{8}{3}\, a_1^p(SS) + 2\, a_1^p(FF)$$
$$F_1^{\bar{\nu}p}(x) = \frac{8}{3}\, a_1^p(SS) - 2\, a_1^p(FF) \tag{18}$$

First we know that $F_1^{\nu p}$ and $F_1^{\bar{\nu}p}$ are positive semi-definite. On the other hand $a_1^p(FF)$ is supposed to dominate over any resonant contribution. Thus we are led to the following ansatz:

$$2\, a_1^p(FF) \underset{x\to 0}{=} \frac{c_1}{x}\left(1 + \frac{c_2}{(\log\ x)^{\varepsilon}}\right) \tag{19}$$

where $\varepsilon > 0$, c_1, c_2 are constant. $a_1^p(SS)$ has the usual Pomeron dependence:

$$\frac{8}{3}\ a_1^p(SS) \underset{x\to 0}{=} \frac{c}{x} \quad \text{and} \quad c > 0 \tag{20}$$

For simplicity we put $c_2 = 0$ *). Inserting eq. (19) and (20) into (18) we obtain:

$$\begin{aligned} F_1^{\nu p} &= \frac{c}{x} + \frac{c_1}{x} = \frac{1}{x}(c + c_1) \\ F_1^{\bar{\nu} p} &= \frac{c}{x} - \frac{c_1}{x} = \frac{1}{x}(c - c_1) \end{aligned} \tag{21}$$

The positivity condition for $F^{\bar{\nu}p}$ implies:

$$\begin{aligned} c + c_1 &> 0 \\ c - c_1 &> 0 \end{aligned} \tag{22}$$

Not knowing the sign of c_1 we can only predict that the difference $F_2^{\nu p} - F_2^{\bar{\nu}p}$ approaches the constant $4\ c_1$ for $x \to 0$. This difference is a direct measure for the assumed non neutrality of the $q\bar{q}$ - sea and can be measured experimentally.

In accordance with what has been stressed in the introduction eq. (19) is expected to hold in some region of small x, but not down to wee x.

IV b) Sum rules

From eq. (16) one can derive sum rules. In general one can say that sum rules which are independent of core structure must be the same in our approach as in the approach with SU(3) singlet Pomeron.

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*) This term would violate the Feynman scaling logarithmically.

For example the Ll. Smith sum rule:

$$12(F_1^p - F_1^n) = F_3^{\nu} - F_3^{\bar{\nu}} \qquad (23)$$

is also valid in our formalism which can easily by derived from eq. (16). Utilizing eqs. (28), (29) we also easily derive the Adler sum rule from eq. (16).

Any sum rule which depends on a neutral $q\bar{q}$ - sea must have a correction term, which amounts to separating out the deviations from neutrality.

We will bring here two examples:

$$F_1^{\nu}(x) - F_1^{\bar{\nu}}(x) + 6(F_1^p(x) - F_1^n(x)) = 4\, a_1^p(FF) \qquad (24)$$

$$F_3^{\nu} + F_3^{\bar{\nu}} - 36(F_1^p + F_1^n) + 12(F_1^{\nu} + F_1^{\bar{\nu}}) = \frac{16}{3}\, a_3^p(SS) + 8\, a_1^p(FF) \qquad (25)$$

Eqs. (24) and (25) can in principle be used to determine $a_3^p(SS)$ and $a_1^p(FF)$ experimentally.

Using the Adler sum rule which is true in every reputable model, and the Callan-Gross relation (5) we derive from eq. (24) the corrected Gottfried sum rule:

$$3\int_0^1 \frac{dx}{x}\,(F_2^p(x) - F_2^n(x)) = 1 + 4\int_0^1 a_1^p(FF)\, dx \qquad (26)$$

For $\int_0^1 a_1^p(FF)dx = 0$, eq. (26) reduces to Gottfried sum rule.

IV c) Estimate of $\int a_1^p(FF)dx$

The present experimental data for the left hand side of eq. (26) is about 0.5[1]. This fixes the value of the integral of the right hand side to

$$\int_0^1 a_1^p(FF)dx = -1/8 \tag{27}$$

To determine the integral of the respective resonant part we use the first of the following two relations:

$$-1/4 = \int_0^1 dx\ a_1(FF) \tag{28}$$

$$0 = \int_0^1 dx\ a_1(FD) \tag{29}$$

These relations follow from the identific tion of the bilocal current with the local corrents[4]. From eq. (27) and (28) we obtain

$$\int_0^1 a_1^R(FF)\ dx = -1/8 \tag{30}$$

The integral over $a_3^p(SS)$ can be determined from eq. (16), the quark sum rule

$$\int_0^1 (F_3^{\nu}(x) + F_3^{\bar{\nu}}(x))dx = -6 \tag{31}$$

and the duality relation $a_1(FD) = a_1(DD)$:

$$\int_0^1 a_3^p(SS)\ dx = -\frac{3}{8} \tag{32}$$

We should like to mention here that the (10b) implies (10c). Otherwise the sum rule (31), which is ontained in the quark parton model, cannot be fitted into our scheme.

From eqs. (30), (32) one observes that p^Q does not share an important property with p^V, namely it cannot be positive in the whole x-region. Of course, this does not come as a surprise because correcting the Gottfried sum rule amounts to lowering the value 1 to 0.5. This implies from eq. (24) that there must be a finite x intervall where the following inequality holds:

$$F_1^{\nu}(x) - F_1^{\bar{\nu}}(x) + 6\,(F_1^{p}(x) - F_1^{n}(x)) < 0 \qquad (33)$$

This relation can be tested experimentally.

IVd) Comparison with SU(3) Singlet Pomeron core

Utilizing the sum rules (27) - (30), (32) and eq. (16) we derive the following relations:

$$\int_0^1 (F_1^p - \frac{8}{9}\, a_1^p(SS)dx = 1/4 \quad = 1/4 \quad (1/2) \qquad (34\text{-a})$$

$$\int_0^1 (F_1^n - \frac{8}{9}\, a_1^p(SS))dx \quad = 1/6 \quad (1/3) \qquad (34\text{-b})$$

$$\int_0^1 (F_1^{\bar{\nu}p} - 3F_1^{en})dx \quad = 3/4 \quad (1) \qquad (34\text{-c})$$

$$\int_0^1 F_3^{\nu p}\, dx \quad = -5/2 \quad (-2) \qquad (34\text{-d})$$

$$\int_0^1 (F_3^{\bar{\nu}p}\, dx \quad = -7/2 \quad (-4) \qquad (34\text{-e})$$

$$\int_0^1 (F_3^{\nu p} - F_3^{\bar{\nu}p})\, dx = 1 \quad (2) \qquad (34\text{-}f)$$

The values in the brackets are the SU(3) singlet Pomeron case, for example in Ref.4.*

IVe) Some remark on Bloom-Gilman duality

If Bloom-Gilman duality[(6)] is believed to be the usual two component form, then we expect it to hold better for combinations of structure functions which have been chosen such that all Pomeron terms have been cancelled.

For example, we expect is to work better for $(F_1^{ep} - F_1^{en})$ than for $(F_1^{\nu p} - F_1^{\bar{\nu}p})$ and better for $(F_3^{\nu p} - F_3^{\bar{\nu}p})$ than for $F_3^{\nu p}$, $F_3^{\bar{\nu}p}$ seperately[(7)].

IVf) Parton Distribution Functions

From eq. (16) and the corresponding expressions from the Quark Parton Model we derive the following relations:

$$\frac{1}{2}\, u(x) = \frac{2}{3}\, a_1{}^P(SS) - a_1{}^P(FF) - 4\, a_1{}^R(FF) - \frac{1}{3}\, a_3{}^P(SS)$$

$$\frac{1}{2}\, \bar{u}(x) = \frac{2}{3}\, a_1{}^P(SS) + a_1{}^P(FF) + \frac{1}{3}\, a_3{}^P(SS)$$

*) The experimental evidence that $xF_3/F_2 \neq 0$ for all x and that F_3 is quite large can also be interpreted as being due to a Pomeron term in $F_3(a_3(SS)$ carries $G = -1)$.

$$\frac{1}{2}\,d(x) = \frac{2}{3}\,a_1{}^P(SS) - \frac{1}{3}\,a_3{}^P(SS) - 2\,a_1(DD) - 2\,a_1{}^R(FF)$$

$$\frac{1}{2}\,\bar{d}(x) = \frac{2}{3}\,a_1{}^P(SS) + \frac{1}{3}\,a_3{}^P(SS) \qquad (35)$$

$$\frac{1}{2}\,(s(x) + \bar{s}\,(x)) = \frac{4}{3}\,a_1{}^P(SS)$$

$u(x)$, $d(x)$ and $s(x)$ are distribution functions for the p- n- and Λ-quarks in the proton and $\bar{u}(x)$, $\bar{d}(x)$ and $\bar{s}(x)$ are those of the corresponding antiquarks.

For the limit $x \to 0$ where the Pomeron dominates we obtain from eq. (35)

$$u(x) + \bar{u}(x) = d(x) + \bar{d}(x) = s(x) + \bar{s}(x) \qquad (36\text{-a})$$

$$u(x) \neq \bar{u}(x) = d(x) \neq \bar{d}(x) = s(x) + \bar{s}(x) \qquad (36\text{-b})$$

The equality $\bar{u}(x)=\bar{d}(x)$ would imply $a_1^P(FF) = 0$,in the small x-region where it should dominate. We furthermore see that the Gottfried integral exists and has a known value, namely

$$\int_0^1 (\bar{u}(x) - \bar{d}(x))\,dx = -\frac{1}{16} \qquad (36\text{-c})$$

This is of course only the quark parton version of eqs. (26), (27). Furthermore the following pattern to realize a non-neutral $q\bar{q}$ - sea

$$u(x) = \bar{u}(x)$$

$$d(x) = \bar{d}(x) \qquad \text{for } x \to 0 \qquad (37)$$

and $u(x) \neq d(x)$ is excluded

is excluded in our approach as a direct consequence of pure FF and SS coupling. This model (37) would be valid in a pure DD-coupling. Here it leads to $a_1^p(FF)=0$, and $a_3^p(SS) = 0$ for $x \to 0$.

On the other hand the case

$$u(x) = \bar{d}(x) \qquad \text{for } x \to 0 \qquad (38)$$
$$\bar{u}(x) = d(x)$$

can be nontrivially realized if $a_1^p(FF)$ and $a_3^p(SS)$ are correlated by the equation $a_1^p(FF) + \frac{2}{3} a_3^p(SS) = 0$ for $x \to 0$. So we can imagine that the core is made up of $(p,\bar{n})$ and $(\bar{p},n)$ pairs which have the quantum numbers of $\rho^{\pm}$ mesons and $\pi^{\pm}$ mesons.

This core-structure is reminiscent of nuclear physics, where the hard core is believed to be made up from vector mesons.

V - POSSIBLE ROLE OF p^Q IN STRONG INTERACTIONS

We have used Harari's conjecture to justify (to some extend) the application of duality principles to current-hadron scattering, and to be able to identify those parts of the reduced matrix elements in eq. (6) which are not constrained by the duality equations (12), (13) and (14) with the Pomeron. Logically however our procedure can be completly decoupled from Harari's conjecture. We are then dealing purely with properties of current - hadron (resp. current - - quark) scattering, without any immediate connection to strong interaction physics. (See also ref. 4) Physically it would of course be more satiesfying to be able to maintain Harari's hypothesis.

That this in our case will be much more difficult to do than with an SU(3) singlet Pomeron is apparent from the result of the last chapter. There it turned out that as a consequence of our assumptions the $q\bar{q}$ - sea will in general carry charge.

Within the Q.P.M. this is acceptable because only the average charge of the nucleon is defined, and from eqs. (27), (30), (32) and (35) one easily derives $\int (u(x) - \bar{u}(x))\, dx = 2$, and $\int (d(x) - \bar{d}(x))dx=1$, and in addition we may put $s(x) = \bar{s}(x)$ in eq. (35), without running into any contradiction.

Thus we have a model where the proton's quantum numbers are well defined and the $q\bar{q}$ - sea carries charge.

On the other hand we do not know of a charged Pomeron in strong interaction physics. This may be a hint that Harari's conjecture does not apply in our case.

But we may be bold enough to speculate where p^Q could show up in strong interactions if we tentatively believe in Harari's hypothesis.

As the cross sections of exclusive hadronic reactions where quantum numbers different than those of the vacuum are exchanged decrease rapidly as the energy increases, p^Q cannot contribute to these because according to eq. 19 it would lead to a constant high energy cross section or if (19) is relaxed somewhat one would encounter to slowly a decrease with energy.

Hence we can expect p^Q to appear at most in inclusive processes. If this is the case then in order to keep it off exclusive processes, it must not factorize between strong vertices.

Feynman[8] has classified Regge poles according to wether they are exchanged between vertices (like virtual emissions in field theory) or appear through hadronic Bremsstrahlung (real emission). In this understanding p^Q is of Bremsstrahlung origin. Most easily p^Q should be dectable in the triple Regge limit (e.g. in $\pi^- p \rightarrow \pi^0 X$) when there is a triple Pomeron coupling $p^Q p^Q p^V$. (This is not completly unreasonable as there is some evidence for a $p^V p^V p^V$ coupling in $pp \rightarrow pX$). If p^Q has a actually properties like a $\rho^\pm$ meson it may couple to inclusive spin flip amplitudes(9), and thus have considerable influence on the spin dependence of inclusive processes.

Within the framework of this paper the spin properties of p^Q can also be investigated. But for the time being there is not much use doing this because no polarization data from deep inelastic lepton - hadron scattering are available.

We finally mention that due to duality principles by studying the missing massbackground of inclusive hadronic reactions one could in principle also detect p^Q.

So presently we can not come to a final conclusion concerning the validity of Harari's conjecture.

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