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A POSSIBLE DESCRIPTION OF BARYON DYNAMICS IN DUAL AND GAUGE THEORIESG.C. Rossi^{*)}

and

G. Veneziano^{**)}

CERN - Geneva

ABSTRACT

We present an attempt to generalize to baryons a framework recently proposed in order to unify gauge, dual and Regge-Gribov theories of mesons. We find it necessary to depart from the $1/N_{\text{colour}}$ expansion of quantum chromodynamics (QCD) and to replace it by a more general definition of a "dual" approximation of QCD, based on the zero-width limit.

Theoretical and phenomenological consequences of the scheme are derived. For $N_{\text{colour}} = 3$, the baryon resembles a Y shaped string; three families of new "baryonium" bound states are predicted and rough estimates of intercepts and slopes of the associated Regge trajectories are given.

A new type of Zweig-like selection rule is found to hold in leading order and its violations through higher-order topologies are discussed. Duality diagrams for baryons are ambiguous unless implemented with extra lines indicating the flow of certain colour indices. Duality between scattering and annihilation channels is found in $B\bar{B}$ scattering and its consequences are discussed. Some justification is given for the quark counting rule for total cross-sections. Finally, implications of our scheme for the Regge-Gribov calculus in processes involving baryons are discussed.

*) Address after 30 June 1977: Istituto di Fisica, Università di Roma, Roma, Italy.

**) On leave of absence from: The Weizmann Institute of Science, Rehovot, Israel.

1. INTRODUCTION AND OUTLINE OF THE PAPER

Recently it has been recognized that most of our present theoretical schemes for the description of hadrons, notably gauge theories, dual models and Regge-Gribov models, can be conceptually unified through the use of various $1/N$ expansions of QCD^{1,2)}. In spite of the lack of real quantitative results in four dimensions, due in particular to our present ignorance of the mechanism of colour confinement, this approach leads to interesting results in two dimensions³⁾ and allows to achieve a remarkable unification of hadronic concepts, putting on a broader and more fundamental perspective semiphenomenological schemes, such as the topological expansion (TE)⁴⁾ and the Reggeon field theory (RFT)⁵⁾.

Unfortunately, even at this qualitative level the scheme only works for mesons. The need to extend it to baryons hardly requires justification. The difficulties of such extension are related, on the one hand, to the old problems of duality for baryons⁶⁾ and, on the other, to the lack of a $1/N_c$ expansion of QCD when baryons are involved.

In spite of these difficulties various papers dealing with the TE (or dual unitarization) for baryons have appeared⁷⁾. The original ideas of Ref. 4) were extended to $B\bar{B}$ annihilation and to BB and $B\bar{B}$ scattering by several authors with a certain variety of (sometimes contradictory) conclusions.

More recently, effects of baryons and baryon loops in dual unitarization have been dealt with by Chan and co-workers⁸⁾. Finally, the question of "baryonium", i.e., of $B\bar{B}$ bound states with an approximate selection rule against their decay into mesons, has been revived by Rosenzweig⁹⁾ and also by Chew¹⁰⁾ in very stimulating papers.

In this work we try to get a general picture of all these phenomena by attempting to find also for baryons a link between gauge theories and dual models. We propose a possible way out of the above-mentioned difficulties and we discuss some immediate phenomenological consequences of the scheme, leaving a more thorough analysis of them to further work.

The outline of the paper is as follows. In Section 2 we recall some previous results on the mesonic problem which will be relevant for extending the scheme to baryons. We then discuss the difficulties associated with a straightforward generalization from mesons to baryons both in dual theories and in QCD.

In Section 3 we propose a definition of dual baryons in QCD which, instead of using $1/N$ expansion concepts, rests on the idea of a zero width (dual) approximation.

In Section 4 we couple colour singlet currents to baryons and through this we obtain meson-baryon-antibaryon vertices. The order of magnitude of these couplings is discussed.

In Section 5 we approach the problem of constructing $B\bar{B}$ and BB scattering amplitudes by a sewing-up procedure inspired by dual model techniques. We find various terms contributing to these amplitudes, more than expected from ordinary duality diagrams. The drawing of duality diagrams when baryons are present has to be implemented by showing the evolution of an extra line related to the flow of some colour indices, which together with the quark lines, accompanies each baryon. Duality between scattering and annihilation channels in $B\bar{B} \rightarrow B\bar{B}$ is then obtained.

The existence of new families of bound states, called baryonium¹⁰⁾, is predicted together with a new Zweig-like selection rule which forbids their decay into mesons. A comparison of our scheme with the string picture is made and duality in $B\bar{B}$ and MB scattering is discussed.

In Section 6 the bare Pomeron component of $B\bar{B}$ and BB scattering is introduced. The counting rule of the "naïve" quark model is rederived under some general assumptions and a comparison is made between $B\bar{B}$ and BB cross-sections.

In Section 7 we discuss some higher order effects. These include the mixing of baryonium states amongst themselves (leading to violation of the ordinary Zweig rule) and with ordinary mesons (leading to violations of the new Zweig rule). Finally, we discuss the implications of our results for the Regge-Gribov calculus in processes involving external baryons.

Our conclusions, as well as a summary of the results, are given in Section 8.

2. DUAL MESONS IN GAUGE THEORIES AND DIFFICULTIES WITH BARYONS

In this Section we briefly review some aspects of the mesonic problem in dual and gauge theories and point out the difficulties of applying similar ideas to baryons.

Let us first fix some of our notations. We shall consider a gauge theory -- QCD -- which describes the interaction of spin $\frac{1}{2}$ fermions (quarks) with spin 1 bosons (gluons) defined by the Lagrangian

$$\mathcal{L} = \frac{1}{4} G_{\mu\nu}^i G^{\mu\nu}_i + \bar{\Psi}^i \gamma_\mu D^\mu_i \Psi_i + \quad (2.1)$$

+ (mass term) + (gauge fixing term) + (ghost term) ,

where the flavour group is $SU(N_f)$ and the gauge group is the colour group $SU(N_c)$. In Eq. (2.1) all flavour indices are understood and the colour indices i and j run from 1 to N_c . $G_{\mu\nu}^j$ and D_μ^j are given, in terms of the gauge fields $A_{\mu i}^j$, by the equations:

$$G_{\mu\nu}^i = \partial_\mu A_\nu^i - \partial_\nu A_\mu^i + g [A_\mu, A_\nu]^i \quad (2.2)$$

$$D_\mu^i = \partial_\mu \delta^i + g \left(A_\mu^i - \frac{1}{N_c} \delta^i A_\mu^k \right).$$

Under $SU(N_c)$ the quark field ψ_i transforms like the fundamental representation, while the gauge field $A_{\mu i}^j$ transforms like the adjoint representation.

2.1 Mesons in dual and gauge theories

Under the fundamental assumption that colour is confined in the leading order of the $1/N_c$ expansion, one can draw some very suggestive analogies between dual mesons and mesons in QCD^{1,2)}. These are the following:

- 1) There is an infinite spectrum of zero width $q\bar{q}$ bound states in dual theories¹¹⁾. The same is supposed to be true in the leading $1/N_c$ approximation of QCD¹⁾, and has so far been verified in two dimensions³⁾.
- 2) We have a picture of dual mesons as open strings with two ends¹¹⁾. In QCD any intermediate state in the correlation functions of n colour-singlet currents has, in leading order of $1/N_c$, the same colour structure of the gauge invariant operator

$$O_{q\bar{q}}(P(x_1, x_2)) = \bar{\Psi}^i(x_1) [G(P(x_1, x_2))]^i_j \Psi_j(x_2) , \quad (2.3a)$$

$$[G(P(x_1, x_2))]^i_j = \left[\text{Tr exp} \left(g \int_{P(x_1, x_2)} A_\mu(x) dx^\mu \right) \right]^i_j \quad (2.3b)$$

when the exponential is expanded in powers of g . In Eqs. (2.3) $P(x_1, x_2)$ represents a path from x_1 to x_2 which looks like an open string with ends in x_1 and x_2 .

- 3) The analogy goes further if we consider quarkless states, sometimes referred to as gluonium. They are represented by closed strings in the dual resonance model (DRM). In QCD one finds²⁾ that the multigluon states which build up gluonium have the colour structure of the operator

$$O_{\text{gluonium}}(P_{\text{closed}}) = \text{Tr} \left[\exp \left(g \int_{P_{\text{closed}}} A_\mu(x) dx^\mu \right) \right]. \quad (2.4)$$

Notice that both the operators (2.3a) and (2.4) are "irreducible" colour singlets, in the sense that they cannot be written as a product of two colour singlets. This, under the assumption of colour confinement, guarantees that both kinds of mesons are stable in this approximation.

There are different equivalent ways of defining zero-width dual mesons in QCD. Consider, for simplicity, the case of a two-current correlation function

$$\Gamma_2^M = \langle 0 | T(\bar{\psi}^i(x) \gamma_\mu \psi_i(x) \bar{\psi}^j(y) \gamma_\nu \psi_j(y)) | 0 \rangle. \quad (2.5)$$

The QCD diagrams defining dual mesons as intermediate states of Γ_2^M can be characterized equivalently in one of the three following ways:

- a - The leading diagrams in $1/N_c$ (at $g^2 N_c$ and N_f fixed).
- b - The diagrams whose intermediate states have perturbatively the same global colour structure as that of the states

$$\bar{\psi}^{i_1} A_{i_1}^{i_2} A_{i_2}^{i_3} \dots A_{i_{m-1}}^{i_m} \psi_{i_m} | 0 \rangle, \quad (2.6)$$

where Lorentz indices have been suppressed.

- c - The set of diagrams which, once decomposed in time ordered diagrams of the old-fashioned perturbation theory, contain only vertices of the type shown in Fig. 1 *) for all **) time orderings. Notice that all creation vertices (Figs. 1a,b,c) are allowed, while annihilation vertices (Figs. 1d,e,f,g) are only allowed if the incoming colours, labelled by repeated indices, were already forced to be equal by previous interactions.

2.2 Difficulties with baryons in dual theories

For a long time baryons have been a nuisance to dual theorists. First of all, there is the old observation made by Rosner⁶⁾ that in $B\bar{B}$ scattering states with $qq\bar{q}\bar{q}$ quantum numbers have to be dual to ordinary $q\bar{q}$ mesons. An obvious question here is: do the $qq\bar{q}\bar{q}$ states have to be bound states, or can they be seen as a continuum of a pair of $q\bar{q}$ states? After all, in MM scattering the Pomeron states are dual to $qq\bar{q}\bar{q}$ states, but we know that this only means that the Pomeron is dual to background, as first suggested by Harari and Freund. We shall see in the following that solving the Rosner difficulty by interpreting $qq\bar{q}\bar{q}$ as a continuum is not self-consistent within our scheme. This, however, cannot be seen by simply drawing the usual duality diagrams for baryons.

A second difficulty is the pattern of exchange degeneracy for baryon states and trajectories which one gets from the duality constraints imposed by the absence of exotic resonances in $MB \rightarrow MB$ and $B\bar{B} \rightarrow MM$ scattering. One obtains an over constrained homogeneous system of equations whose solution seems to require baryonic multiplets not experimentally found. We shall see in Section 5.4 that in our approach certain constraints are no longer valid, so that the number of equations is reduced.

The third difficulty is within the string model of baryons¹¹⁾. There is a lot of ambiguity on what a baryonic string ought to be, e.g., should it have a Δ or a Y shape? Or should it have a linear structure, as in the meson case, with two quarks at one end replacing one antiquark, or perhaps with one of the three quarks sitting in the middle of the other two?

As a final remark we may add that in dual theories it is felt that the fermionic amplitudes which have been constructed in the Ramond-Neveu-Schwarz

*) Here, and in the following, we shall adopt the graphical conventions of Refs. 1) and 2).

**) The time orderings, in which a set of particles is annihilated into the vacuum or is created from the vacuum, need not be considered for the classification of the diagrams.

model¹¹⁾ do not describe actual baryons, but some sort of mathematical fermions which should not be interpreted as physical particles: Schwarz¹¹⁾ suggested that they may have something to do with the quarks themselves.

2.3 Difficulties with baryons in QCD

As already remarked by 't Hooft¹⁾ it is very difficult, if not conceptually impossible, to define a $1/N_c$ expansion for baryons. The problem is that baryons are supposed to be bound states of N_c quarks in a wave function completely anti-symmetric in colour [this is the way to form a singlet in $SU(N_c)$ with N_c ψ 's].

Another difficulty is that, iterating any reasonable N_c quark interaction (e.g., gluon exchange forces among the quarks) does not introduce in any obvious way factors of N_c : as a consequence, the $(n+1)^{th}$ iteration of the basic interaction is down with respect to the n^{th} iteration by powers of g , without compensating factors of N_c . This can be explicitly seen in a recent paper by Durgut¹²⁾ in which a two dimensional QCD model for baryonic bound states of N_c quarks is proposed, extending the ideas of Ref. 3)*).

Again this is not the situation for mesons: there the full set of planar diagrams is of the same order in the $1/N_c$ expansion.

The above considerations lead us to look for a definition of dual baryons which, although inspired by QCD, does not rely on the $1/N_c$ expansion.

3. A POSSIBLE DEFINITION OF DUAL BARYONS IN QCD

In order to find a definition of dual baryons in QCD which is not of the $1/N_c$ expansion type, we consider the two alternative definitions, b and c, of dual mesons in QCD, discussed in Section 2.1.

Definition b suggests that we look at gauge invariant operators with non-zero baryon numbers. The simplest one is, in $SU(N_c)$

$$\begin{aligned} & \bigcirc_{q_1 \dots q_{N_c}} (P(x, x_1), P(x, x_2), \dots, P(x, x_{N_c})) = \\ & = \varepsilon^{i_1 i_2 \dots i_{N_c}} [G(P(x, x_1)) \psi(x_1)]_{i_1} [G(P(x, x_2)) \psi(x_2)]_{i_2} \dots [G(P(x, x_{N_c})) \psi(x_{N_c})]_{i_{N_c}} \end{aligned} \quad (3.1)$$

where $G(P(x, y))$ is given in Eq. (2.3b).

*) Although we have not checked it completely, we believe that our proposal in Section 3 reduces to Durgut's one in the two-dimensional case.

We can interpret Eq. (3.1) in the following way: the "string operators" $G(P(x, x_k))$ acting on $\psi(x_k)$ make each term in the square brackets transform as a ψ field at x instead of at x_k . The ε tensor then builds up a local singlet out of N_c objects which transform as ψ 's at the same x .

Let us consider now the correlation function of two colour singlet baryonic currents

$$\Gamma_2^B = \langle 0 | T(j_B(x) \bar{j}_B(y)) | 0 \rangle, \quad (3.2)$$

where

$$j_B(x) = \varepsilon^{i_1 i_2 \dots i_{N_c}} \psi_{i_1}(x) \psi_{i_2}(x) \dots \psi_{i_{N_c}}(x). \quad (3.3)$$

The analogue of definition b of Section 2.1 which will enable us to define dual baryons as intermediate states of Γ_2^B becomes:

b' - The set of QCD diagrams whose intermediate states are all of the kind:

$$\varepsilon^{i_1 i_2 \dots i_{N_c}} (AA \dots A \psi)_{i_1} (AA \dots A \psi)_{i_2} \dots (AA \dots A \psi)_{i_{N_c}} | 0 \rangle. \quad (3.4)$$

In Eq. (3.4) we have suppressed the Lorentz indices of the gauge fields, showing only the global colour structure of the state.

Obviously the N_c colour labels $i_1, i_2 \dots i_{N_c}$ carried by the quark fields of one of the currents flow through the diagram, finally ending at the quark fields of the other current. We shall refer to these colour lines as "special" lines. One example is given in Fig. 2a for the case $N_c = 3$. The particles in the intermediate states in Γ_2^B are thus divided into N_c groups, each containing one of the quarks and carrying one of the special lines.

We want to emphasize two appealing properties of this definition of baryons. First, the states of Eq. (3.4) are "irreducible" singlets, exactly as in the case

of mesons. Hence they have to be thought of as single particle states and will have zero width within the set of diagrams we are considering^{*)}. Second, the gauge invariant operator (3.2) has a suggestive similarity to a star-like object, in which N_c strings with quarks at $x_1, x_2, \dots, x_{N_c}$ are joined by a "string junction" at x . With our definition then, a three-quark baryon has a Y shape and not a Δ shape or any other shape. As a matter of fact, we were not able to find gauge invariant operators for, say, a Δ shaped baryon. We shall see that the topological structure of the baryons will play an important role in the discussion of their scattering processes.

One can ask whether one can find an equivalent definition of dual baryons, similar to definition c for mesons. We have found that this is indeed possible and that b' is equivalent to the following statement:

c' - The diagrams which, in terms of time-ordered old-fashioned perturbation theory, contain the vertices described below for all time orderings^{**)} .

- i) All creation vertices are allowed.
- ii) Annihilation vertices amongst particles in the same group are allowed if they are of the type shown in Figs. 1d,e,f,g.
- iii) Annihilation vertices for particles in different groups are all allowed.

The appearance of these latter vertices, which are necessary in order to have any binding force at all, prevents baryonic diagrams from being of a definite order in the $1/N_c$ expansion.

In Fig. 2a we give examples of allowed diagrams, and forbidden diagrams are given in Figs. 2b and 2c where we show the intermediate state which violates our rules. One important point to be stressed here is that all fermion and gluon propagators are dressed by arbitrary planar insertions, because their addition to an allowed diagram always leads to another allowed diagram.

It is perhaps useful to visualize the wave function of the baryon in terms of quarks and gluons as consisting of three ($N_c = 3$ for a moment) valence quarks each one accompanied by a "plane of gluons". The last gluons of each plane interact amongst themselves through a colour exchange mechanism. When the baryon

*) While for mesons the $1/N_c$ expansion guarantees that the set of planar diagrams are gauge invariant in the leading order of the expansion, no such argument can be made to prove the gauge invariance of the set of diagrams by which we define the baryons. However, since we expect that a gauge invariant realization of the zero-width limit should exist, we are optimistic on this question. The fact that, in the two-dimensional case, our choice seems to reduce to that of Durgut¹²⁾ lends further support to this assertion.

) See note () on page 5.

evolves in time the three quarks describe the edges of three two-dimensional surfaces which tend to meet at a junction, although we cannot identify its position as that of any particular quark or gluon in the intermediate state: the emerging picture is that of a Y shaped baryon evolving in time.

4. COUPLINGS OF CURRENTS AND MESONS TO BARYONS

Even without a very specific model for the binding mechanism of the quarks inside the baryon, our approach allows to introduce couplings between baryons and currents and through them to mesons. This in turn will enable us to construct (formally) BB and B $\bar{B}$ scattering processes and to discuss some of their properties.

We first consider the coupling of a current like

$$J^k(x) = \beta \bar{\psi}_b^i(x) (\lambda^k)_a^b \Gamma \psi_i^a(x), \quad (4.1)$$

carrying zero baryonic number, by attaching the new vertex of Fig. 3 everywhere on the quark lines. In Eq. (4.1) a and b are flavour indices and i is a colour index. The λ^k 's are the usual Gell-Mann matrices, normalized to

$$\text{Tr} (\lambda^h \lambda^k) = 2 \delta_{hk} \quad (4.2)$$

and Γ is any Dirac matrix.

Since every quark propagator is completely dressed by planar insertions we have intermediate states of the q $\bar{q}$ meson variety in the current channel. By our rules, these states and their couplings to J^k will be exactly those found in the mesonic theory. In our approximation there will be poles in the current channel, so by factorization we can arrive to the definition of the baryon-antibaryon-meson vertex (Fig. 4).

At this point there is an interesting observation to make on the normalization of our couplings. Consider, for instance, a particular case of Eq. (4.1) with $\lambda^k \rightarrow \lambda^0 = \sqrt{2/N_f}$, $\Gamma \rightarrow \gamma_\mu$ and $\beta \rightarrow \sqrt{N_f}/N_c \sqrt{2}$ i.e., the baryon number current

$$J_{\mu}^B(x) = \frac{1}{N_c} \bar{\Psi}_a^i(x) \gamma_{\mu} \Psi_i^a(x), \quad (4.3)$$

normalized to give charge one on the single baryon state. Clearly, the coupling to each quark comes out to be $1/N_c$, because a quark has baryon number $1/N_c$.

The coupling of $J^k(x)$ to a vector meson was found in Refs. 1) and 2) to be

$$g_{vJ} \simeq \beta \sqrt{N_c}. \quad (4.4)$$

Since in our case $\beta \simeq \sqrt{N_f}/N_c \sqrt{2}$, we have

$$g_{vJB} \simeq \sqrt{\frac{N_f}{N_c}} \quad (4.5)$$

From the normalization of the baryonic charge

$$g_{vJB} \cdot g_{vB\bar{B}} = 1, \quad (4.6)$$

we deduce

$$g_{vB\bar{B}} \simeq \sqrt{\frac{N_c}{N_f}}. \quad (4.7)$$

One then gets a coupling of vector mesons to baryons which is N_c times larger than the coupling to mesons, for which one has²⁾

$$g_{vMM} \simeq \frac{1}{\sqrt{N_f N_c}}. \quad (4.8)$$

This is easily understood as coming from the fact that the vector meson can couple to each one of the N_c quark lines of the baryon.

The general result is actually the following:

1) for VMM coupling

$$g_{VMM_k}^{\pm} = \frac{C}{2\sqrt{N_c}} \left[\text{Tr}(\lambda^k \lambda^1 \lambda^k) \pm \text{Tr}(\lambda^k \lambda^2 \lambda^k) \right], \quad C = O(1), \quad (4.9)$$

where the sign is determined by the generalized Bose statistics;

2) for baryons in a totally symmetric flavour state (say the decuplet)

$$g_{V_k B\{n_\ell\} \bar{B}\{n'_\ell\}} = \frac{C'}{\sqrt{N_c}} \sum_{\ell, m=1}^{N_f} (\lambda^k)_\ell^m \sqrt{n_\ell n'_\ell}, \quad C' = O(1), \quad (4.10)$$

where n_ℓ is the number of quarks with flavour ℓ in $B\{n_\ell\}$.

This result, that couplings of mesons to baryons are in general of different order (in $1/N_c$) than couplings of mesons to mesons, seems to contradict an argument based on duality for meson-baryon scattering, according to which the two couplings should be of the same order. Indeed, naïve duality for the $MB \rightarrow MB$ scattering amplitude would give that the sum over s channel baryonic resonances is equal to the sum over t channel meson resonances. Writing schematically:

$$s \text{ channel contributions} \simeq \sum_{B^*} g_{MBB^*} g_{MBB^*} \quad (4.11)$$

$$t \text{ channel contributions} \simeq \sum_{M^*} g_{MMM^*} g_{M^*BB} \quad (4.12)$$

this would imply $g_{MBB} \simeq g_{MMM}$.

However, the argument is incorrect because there are, in $SU(N_c)$, N_c^2 different duality diagrams for $MB \rightarrow MB$, as shown in Fig. 5. Of these only the N_c contributions of type (I) can contribute to the duality relation

$$\sum_{B^*} g_{MBB^*} g_{MBB^*} = \sum_{M^*} g_{MMM^*} g_{M^*BB}, \quad (4.13)$$

because only they have ordinary mesons as intermediate states in the t channel and are planar. The remaining $N_c(N_c-1)$ diagrams [type (II) in Fig. 5] are non-planar and do not satisfy a duality constraint, such as Eq. (4.13). We shall see later in Section 5.4 what sort of t channel intermediate states they have.

The above argument can simply be rephrased as follows: the contribution of any intermediate baryonic state in Eq. (4.11) comes from the product of two couplings, each one having N_c components. However, only the N_c "diagonal terms" in which the two mesons are attached to the same quark line are given by Eq. (4.12) in which there are clearly only N_c contributions.

Similar considerations will arise when discussing the elastic coupling of the bare Pomeron to the baryons in Section 6. This will give us the clue for understanding the "counting rule" of the quark model¹³⁾ (see Section 6.1).

5. CONSTRUCTION AND SYSTEMATICS OF NON-DIFFRACTIVE $B\bar{B}$ SCATTERING

We shall approach the problem of constructing $B\bar{B}$ and BB scattering from a dual point of view. Due to our incomplete understanding of the baryon binding mechanism, this appears to us to be the only viable procedure at the moment. Hopefully one will be able to define eventually everything in terms of a complete classification of QCD diagrams.

5.1 $B\bar{B}$ scattering in the s channel and baryonium formation

The simplest object we can consider, once we are given (see Section 4) the couplings of $q\bar{q}$ mesons to baryons, is the $B\bar{B}$ scattering amplitude obtained by exchanging a complete tower of $q\bar{q}$ mesons in the t channel. One is thus led to the identification expressed by the equality in Fig. 6a, where the planar structure of the mesonic exchange has been evidentiated.

Since, by our construction, the mesonic intermediate state does not allow any exchange of special lines in the t channel, it is clear that the two groups of special lines, entering with the incoming baryon and antibaryon, have to come out with the outgoing baryon and antibaryon, respectively. In order not to forget this fact (which will turn out to be crucial), we redraw the duality diagram with two extra dotted lines as in Fig. 6b.

The t channel has been defined as the one in which ordinary mesons are exchanged. In order to gain some insight on the s channel hadronic intermediate state, it is useful to look perturbatively into the colour structure of that channel.

From the definition of the allowed diagrams in the baryon wave function and the construction of the scattering amplitude of Fig. 6, it follows that each s channel intermediate state has the global colour structure of the operator (from now on, unless differently stated, we shall take $N_c = 3$ for simplicity)

$$\varepsilon_{i_1 i_2 i_3} \varepsilon^{i_1 i_2 i_3} (\bar{\psi} A A \dots A)^{i_1} (\bar{\psi} A A \dots A)^{i_2} . \quad (5.1)$$

$$(A A \dots A)^{i_3} (A A \dots A \psi)_{i_2} (A A \dots A \psi)_{i_1} ,$$

where all the indices, but the colour ones, have been suppressed.

This can be directly verified in explicit examples, as the one shown in Fig. 7. The gauge invariant operator which, once expanded in powers of g , gives rise to terms like those of Eq. (5.1) is

$$O_{qq\bar{q}\bar{q}} (P(y_1, y), P(y_2, y), P(y, x), P(x, x_1), P(x, x_2)) =$$

$$= \varepsilon_{i_1 i_2 i_3} \varepsilon^{i_1 i_2 i_3} [\bar{\psi}(y_1) G(P(y_1, y))]^{i_1} [\bar{\psi}(y_2) G(P(y_2, y))]^{i_2} . \quad (5.2)$$

$$[G(P(y, x))]^{i_3} [G(P(x, x_1)) \psi(x_1)]_{i_2} [G(P(x, x_2)) \psi(x_2)]_{i_3} .$$

This operator is an irreducible colour singlet, in the sense we have used before. For consistency we have to assume that a new set of zero-width mesons with $qq\bar{q}\bar{q}$ quantum numbers appears in the s channel. We then find that the old problem of Rosner's exotics cannot be solved by interpreting these $qq\bar{q}\bar{q}$ states as a background made of a pair of $q\bar{q}$ mesons.

We now want to argue that in $B\bar{B}$ there is duality between scattering and annihilation processes. This is not obvious at this stage (Fig. 8a) because, after all, we have mesons in both s and t channels. But if we turn on fermion loops (Fig. 8b), as required by unitarity, we find that the s channel intermediate states are of the type $B\bar{B} + \text{mesons}$, i.e., they are scattering states. This is because the $qq\bar{q}\bar{q}$ bound states have now become resonances which, if kinematically possible, will decay, in the context of our graphs, into $B\bar{B} + \text{mesons}$, but not just into mesons.

In the t channel the zero-width ordinary mesons have been replaced by a multimeson annihilation state (see Section 5.2) similar to the one which generates the Reggeon in MM scattering.

Finally, there are no intermediate states in the u channel and therefore this amplitude gives a purely real contribution to $BB \rightarrow BB$ scattering.

At this stage a Zweig-type selection rule has emerged for the $qq\bar{q}\bar{q}$ states, in the same way as the planar approximation for mesons gives the ordinary Zweig rule¹⁴⁾. These states have, therefore, been named baryonium¹⁰⁾ in analogy with charmonium states.

In the following we shall denote them by the symbol M_4^J , where J stands for "junction" and the subscript is the total number of quarks and antiquarks. We shall discuss their properties further in Sections 6 and 7.

The next term to be considered is the exchange of a pair of $q\bar{q}$ mesons between B and $\bar{B}$. The amplitude for $B\bar{B} \rightarrow MM$ has already been discussed at the end of Section 4. When we sew two such amplitudes together, different topologies are generated according to which quark line emits and absorbs the two exchanged mesons. This is illustrated in Fig. 9.

The first term in the sum in Fig. 9 has already been included in the diagrams of Fig. 8b. The second term gives a new contribution which we shall discuss shortly. We first consider the last two terms: they are kind of interference terms which, however, have the same flavour structure of the basic diagram of Fig. 8a. Like the first term, they are of higher order in g , but without compensating factors of N_f . From our experience in the mesonic case, we are led to interpret them as absorptive corrections [$h = 1$ in the language of the TE⁴⁾] and we shall therefore drop them at this stage.

Let us go back to the second diagram of Fig. 9. It has a Reggeon-Reggeon cut in the t channel (or a meson-meson cut in the energy plane). In the s channel the structure is more complicated as is usual for the s channel content of a t channel Regge cut.

There are three kinds of intermediate states which correspond to the three topologically distinct ways of cutting the diagram, each one related to an s channel discontinuity of the two Reggeon cuts. In Fig. 10a, $B\bar{B}$ intermediate states are shown: they correspond to cutting no Reggeon. Fig. 10b shows M_4^J intermediate states which correspond to cutting one Reggeon. Finally, in Fig. 10c a new baryonium state of the type M_2^J is evidenced. It corresponds to cutting both Reggeons.

Although this new state has the quantum numbers of a $q\bar{q}$ pair, it cannot be thought of as an ordinary meson. This is most easily seen by analyzing, as we have done before, its colour structure, which corresponds to that of the gauge invariant operator

$$O_{q\bar{q}} [P(y_1, y), P(y, x), P'(y, x), P(x, x_1)] =$$

$$= \varepsilon_{i_1 i_2 i_3} \varepsilon^{i_1 i_2 i_3} [\bar{\psi}(y_1) G(P(y_1, y))]^{\delta_1} [G(P(y, x))]^{\delta_2}_{i_1}$$

$$\cdot [G(P'(y, x))]^{\delta_3}_{i_2} [G(P(x, x_1)) \psi(x_1)]_{i_3} . \quad (5.3)$$

Like M_4^J , these M_2^J mesons decay naturally into states of the type $B\bar{B} + \text{mesons}$ and not just into mesons. This becomes clear if one adds (at least) two quark loops to Fig. 10c, one in each of the planes of the two exchanged mesons (Fig. 11a).

The important point to be stressed here is that the contribution to $B\bar{B}$ scattering we are considering is not obtainable from the one of Fig. 8a by s-t crossing. We have a genuinely new set of diagrams (Fig. 11b), characterized by the fact that each group of special colour lines travels along with one quark in the s channel and not with the diquark in the t channel. The drawing of the dotted line is essential in order to show this property.

As before, we are led to interpret s channel intermediate states as scattering states and those in the t channel as annihilation states.

As the final case (with $N_c = 3$) we now consider the exchange of three mesons. Again omitting interference-type graphs, we are left with the diagrams of Figs. 12a and 12b. We have already drawn the dotted lines to indicate which is the scattering and which is the annihilation channel.

- In terms of t channel intermediate states we have three $q\bar{q}$ mesons or a three Reggeon cut. Analyzed in terms of s channel intermediate states, we find
- 1) $B\bar{B}$ states. They can be seen by cutting Fig. 12a so that none of the Reggeons are cut. There are six distinct contributions of this type.
 - 2) M_4^J states. They correspond to cutting a single Reggeon in Fig. 12a. This can be done in six different ways.
 - 3) M_2^J states. They correspond to cutting two Reggeons and there are again six possible contributions.

- 4) Finally we encounter a new family of exotic mesons, which we denote by M_0^J . They have the flavour quantum numbers of the vacuum and correspond to cutting all three Reggeons. The gauge invariant operator they correspond to is

$$O[P(x,y), P'(x,y), P''(x,y)] = \varepsilon_{\delta_1 \delta_2 \delta_3} \varepsilon^{i_1 i_2 i_3} [G(P(x,y))]_{i_1}^{j_1} \cdot [G(P'(x,y))]_{i_2}^{j_2} [G(P''(x,y))]_{i_3}^{j_3} \quad (5.4)$$

M_0^J is another meson of the baryonium family in the sense that it decays into $B\bar{B}$ + mesons and not just into mesons in our approximations. This again can be seen by allowing a quark loop in each one of the three planes of the exchanged mesons (Fig. 13).

Summarizing this subsection, we have identified three (for simplicity we have considered the case $N_c = 3$) distinct contributions to $B\bar{B}$ scattering whose s channel absorptive part is entirely given by either $B\bar{B}$ continuum or by single mesons of the baryonium family. A fourth contribution, the bare Pomeron, will be discussed in Section 6.

Within the class of diagrams we have discussed so far -- diagrams of Figs. 8a, 11b and 12b with any quark loop insertion -- there is one selection rule which holds exactly. It is a new kind of Zweig rule¹⁴⁾ forbidding the M^J mesons formed in the s channel, to decay into ordinary mesons, i.e., to go into purely annihilation channels. Violations of this selection rule will be discussed in Section 7, together with the question of mixing of baryonium states amongst themselves and with ordinary mesons.

Let us now turn our attention to the contributions of our diagrams to the physical region of the t channel.

5.2 $B\bar{B}$ annihilation in the s channel and baryonium exchange

We have already emphasized that our diagrams of Figs. 8a, 11b and 12b are neither crossing symmetric nor related to each other by crossing. Instead, in order to get the complete crossing symmetric amplitude, one has to add the s - t crossed diagram to each of them. One then gets the new diagrams of Figs. 14a,b,c, respectively. Their s channel absorptive parts represent $B\bar{B}$ annihilation, as can be seen by adding arbitrarily quark loops, and correspond to the following processes:

a) Annihilation into a single $q\bar{q}$ jet of mesons (Fig. 14a). We expect the multiplicity for this process to be the same as the one associated with the Reggeon. The s dependence of the cross-section will be $s^{\alpha_4^J(0)-1}$, where $\alpha_4^J(0)$ is the intercept of the leading M_4^J trajectory.

b) Annihilation into two $q\bar{q}$ jets (Fig. 14b). It exhibits Pomeron-like multiplicity (i.e., twice the multiplicity of the Reggeon) with an energy-dependence of the cross-section of the type $s^{\alpha_2^J(0)-1}$, where $\alpha_2^J(0)$ is the intercept of the leading M_2^J trajectory.

c) Annihilation into three $q\bar{q}$ jets (Fig. 14c). The multiplicity here is three times the Reggeon multiplicity and the cross-section behaves as $s^{\alpha_0^J(0)-1}$, i.e., is controlled by the intercept of the leading M_0^J trajectory.

Of these three contributions only the third one is a direct annihilation process, in the sense that there are no intermediate states through which annihilation actually proceeds. Instead in the two others one can find intermediate states of the baryonium type (M_2^J and/or M_4^J) in the s channel. Thus to the extent that the new Zweig rule for the baryonium family holds, the contributions a) and b) to $B\bar{B}$ annihilation may be suppressed as compared to c).

This is indeed what happens at high energies in our approach (for a discussion of the low-energy region, see Section 7). Let us in fact study the energy and impact parameter dependence of these three contributions by giving an estimate of the intercepts and slopes of the leading baryonium trajectories.

For the intercepts this can be done using a simple multiperipheral-like approximation similar to the one proposed by H. Lee and one of the authors¹⁵⁾ in the meson-meson scattering and generalized to $B\bar{B}$ annihilation by several authors⁷⁾. In this way one easily gets

$$\begin{aligned}\alpha_4^J(0) &\simeq 2\alpha_B(0) - 1 + (1 - \alpha_R(0)) \simeq -1/2 \\ \alpha_2^J(0) &\simeq 2\alpha_B(0) - 1 + 2(1 - \alpha_R(0)) \simeq 0 \\ \alpha_0^J(0) &\simeq 2\alpha_B(0) - 1 + 3(1 - \alpha_R(0)) \simeq 1/2.\end{aligned}\tag{5.5}$$

In the above numerical estimates we have set the baryon intercept $\alpha_B(0)$ equal to 0, and the Reggeon intercept $\alpha_R(0)$ equal to $\frac{1}{2}$ ^{*)}. Since in any case we have

$$\alpha_4^J(0) < \alpha_2^J(0) < \alpha_0^J(0),$$

the three-jet component is seen to dominate at high energies, giving asymptotically, for the average multiplicities:

$$\frac{\bar{n}_{B\bar{B} ann.}}{\bar{n}_{B\bar{B} att.}} = \frac{3}{2}, \quad (5.6)$$

a result known for some time ⁷⁾.

An amusing comparison can be tried between e^+e^- and $p\bar{p}$ annihilation assuming that the latter is dominated by the three $q\bar{q}$ jet process. If the "hadronization" of a $q\bar{q}$ pair is weakly dependent on its total angular momentum, one could expect

$$\bar{n}_{e^+e^-}(s) \simeq \frac{1}{3} \bar{n}_{p\bar{p} ann.}(s), \quad (5.7)$$

where we have assumed that one third of the total available energy goes into each $q\bar{q}$ jet in $p\bar{p}$. Within the errors, Eq. (5.7) seems not to conflict with the data ¹⁷⁾.

For the slopes we use the known fact that a multiperipheral chain gives a random walk in impact parameter space ¹⁸⁾ and we extend this idea to the case of more than one chain (jet).

*) A phenomenological analysis of the inclusive reaction

$$\pi^- p \rightarrow p X^-$$

where the p is in the fragmentation region of the π^- (see Fig. 27b) has been done in Ref. 16) with the idea of relating, through FESR, the meson production in the p -virtual Δ^{++} annihilation to the intercept of the trajectory exchanged in the crossed scattering channel. They found -0.5 ± 0.3 in good agreement with our estimate of α_4^J (for a further discussion, see Section 7.2).

Consider first the case of planar MM scattering, which corresponds to one jet in the s channel and to Reggeon exchange in the t channel. Denoting by $P_1(s,b)$ the impact parameter profile of the imaginary part of the elastic amplitude, one has

$$\begin{aligned} P_1(s,b) &= \exp \left[-b^2 / (2 B_R^{MM} + 4 \alpha'_R \log s) \right] \approx \\ &\approx \exp \left[-b^2 / (R^2 \bar{n}_1(s)) \right], \end{aligned} \quad (5.8)$$

where R is the length of the elementary step in the random walk and $\bar{n}_1(s)$ is the average number of steps, i.e., the average multiplicity associated with one jet. B_R^{MM} is normalized so that $d\sigma(MM)/dt = e^{-B_R^{MM}|t|} + \text{real part contributions}$ and α'_R is, as usual, the slope of a typical mesonic Regge trajectory.

Since R and $\bar{n}_1(s)$ are expected to be universal for a $q\bar{q}$ jet, we can use Eq. (5.8) also for the single-jet contribution to $B\bar{B}$ annihilation of Fig. 14a, obtaining

$$\alpha'_4 = \alpha'_R \approx 1 \text{ GeV}^{-2}. \quad (5.9)$$

For a process with two $q\bar{q}$ jets, such as the $B\bar{B}$ annihilation of Fig. 14b or the bare Pomeron, P_0 , we have two random walks each with average multiplicity $\sim \bar{n}_1(s/4)$ (for simplicity we assume equipartition of the energy between the two jets). The two random walks have to end at the same point in $\vec{b}$ space, so if we suppose that they are independent and denote by $P_2(s,b)$ the profile function of two jets, we get

$$P_2(s,b) \approx (P_1(s/4,b))^2 \approx \exp \left[-2b^2 / (R^2 \bar{n}_1(s/4)) \right]. \quad (5.10)$$

Similarly for the three-jet annihilation process of Fig. 14c we expect

$$P_3(s,b) \approx (P_1(s/9,b))^3 \approx \exp \left[-3b^2 / (R^2 \bar{n}_1(s/9)) \right]. \quad (5.11)$$

Relating (as before) P_2 and P_3 to the corresponding Regge exchanges in the t channel, we obtain

$$\alpha'_2{}^J \simeq \alpha'_{P_0} \simeq \frac{1}{2} \alpha'_R, \quad (5.12)$$

$$\alpha'_0{}^J \simeq \frac{1}{3} \alpha'_R. \quad (5.13)$$

Note that with our approximation, apart from possible fragmentation region effects, one really has the more detailed prediction

$$\begin{aligned} B_R^{MM} &= B_4^J = 2 B_2^J + 4 \alpha'_R \log 2 = \\ &= 3 B_0^J + 4 \alpha'_R \log 3 = 2 B_{P_0}^{MM} + 4 \alpha'_R \log 2, \end{aligned} \quad (5.14)$$

where the B 's are defined by equations analogous to (5.8).

The resulting baryonium trajectories, before the mixing effects discussed in Section 7, are shown in Fig. 15^{*)}.

Following our (admittedly) crude arguments, we can give general formulae for intercepts and slopes. For the leading trajectory corresponding to hadrons with n_i quarks plus antiquarks of type i and n_J junctions, one has

$$\alpha_{(n_i, n_J)}(0) = 1 - \sum_i c_i n_i - c_J n_J, \quad (5.15)$$

$$\alpha'_{(n_i, n_J)}(0) = \frac{2 \alpha'_R}{4 + (N_c - 2) n_J - \sum_i n_i}. \quad (5.16)$$

*) The shape of the trajectories shown in Fig. 15 follows from extrapolating our slope predictions for $M^2 \leq 0$ to $M^2 > 0$. We cannot exclude, however, the existence of $M_{0,2}^J$ trajectories of higher slope which become leading for $M^2 > 0$.

where, comparing with the known ρ , K^* and B trajectories, the numerical constants c_i and c_J are approximately

$$c_u = c_d \simeq 1/4, \quad c_s \simeq 1/2, \quad c_J \simeq 1/4. \quad (5.17)$$

Eqs. (5.15) and (5.16) apply also to the Pomeron ($n_i = n_J = 0$), giving $\alpha_{P_0}(0) \simeq 1$, $\alpha'_{P_0}(0) \simeq \frac{1}{2}$.

5.3 String model picture

The results of these two subsections can be rephrased within a dual string model¹¹⁾, where the baryon consists of three ($N_c = 3$) quarks at the end of three strings, which are then joined at a single point called the string "junction".

Such a model has been considered by several authors in recent years¹⁹⁾. It was found that for a complete specification of dual diagrams involving baryons, one has to consider the evolution of the string junction and this introduces the extra dotted lines in the diagrams.

In this picture the baryonium states have two junctions and either zero, one or two $q\bar{q}$ pairs. They are represented together with ordinary mesons and baryons in Fig. 16.

As for ordinary mesons, the preferred formation and decay processes occur through string fusion and breaking, respectively. In QCD this corresponds to the introduction of fermion loops. In analogy, $B\bar{B}$ tend to form baryonium and baryonium tends to decay into $B\bar{B}$ + mesons through the breaking of one, two or three strings. Thus a Zweig-like selection rule arises, whose violations occur only through junction creation or annihilation.

Such a mechanism is not included in the diagrams considered so far, but is introduced by higher-order corrections (see Section 7).

Note that our selection rule differs somewhat from the one in which, assuming point-like diquarks, M_4^J decays into ordinary mesons occur only through a diquark-antidiquark annihilation⁹⁾, because in this scheme a M_4^J with non-vacuum quantum numbers would be stable. It also differs from the Freund-Waltz-Rosner rule²⁰⁾ which would allow $M_2^J \rightarrow MM$.

Let us finally remark that, although in our model the baryon does not look like a diquark-quark system, this could be the approximate picture for the lowest

lying states. Indeed, the fact that the baryon slope seems to be the same as that of ordinary mesons may indicate a Y shaped baryon in which two of the strings are much shorter than the third one.

5.4 Duality in baryon-antibaryon and in meson-baryon scattering

We now discuss some of the implications of our scheme for duality in $B\bar{B}$ and in MB scattering.

As we have already mentioned, our model states that the old Rosner problem in $B\bar{B}$ scattering can only be solved by introducing a new family of exotic mesons, that we have denoted by M_4^J . Such states cannot be regarded as M-M continuum of the type which is dual to the Pomeron in MM scattering because of the different topology of the two diagrams.

We even claim that the need for new bound states does not exhaust itself with M_4^J mesons. M_2^J and M_0^J states have to exist for consistency and are dual to two and three meson continuum, respectively.

The duality pattern of M^J trajectories in the t channel can be derived by demanding a vanishing amplitude when the annihilation channel is exotic. One thus gets exchange degeneracy of $C = \pm 1$, $I = 0, 1, 2$ M_4^J states, of $C = \pm 1$, $I = 0, 1$ M_2^J states and of $C = \pm 1$, $I = 0$ M_0^J states.

These degeneracies will be somewhat broken at small masses by mixing effects, but we may expect them to be still quite good near $B\bar{B}$ threshold (see Section 7).

Finally, by a standard application of duality, the usual exchange degeneracy results⁶⁾ are obtained for f, ω, ρ, A_2 trajectories and their couplings to $B\bar{B}$.

It is interesting to consider the implication of baryonium on $MB \rightarrow MB$ duality. Here, historically, no problem of exotics of the type encountered in $B\bar{B}$ scattering was found. Yet, the exchange degeneracy pattern for baryon trajectories, as derived from the exoticity constraints, in either $MB \rightarrow MB$ and in $B\bar{B} \rightarrow MM$ led to difficulties with experiments⁶⁾. Typically, the homogeneous system of equations one gets is overconstrained.

In our scheme, whereas an exotic MB channel is always associated with a real amplitude, an exotic $B\bar{B} \rightarrow MM$ channel is not, because, as one can easily check, there are always M_4^J intermediate states.

Thus, in our scheme, results coming from exoticity in the annihilation channels are not necessarily valid. They could, however, be approximately right, if the coupling of M_4^J to the MM system is small.

6. THE BARE POMERON IN $B\bar{B}$ AND IN BB

In the previous section we have considered the non-diffractive contribution to $B\bar{B}$, showing that $B\bar{B}$ annihilation is dual to baryonium exchange and $B\bar{B}$ scattering to ordinary meson exchange, i.e., that annihilation and scattering are dual to each other.

We now go to discuss the most important high-energy contribution to $B\bar{B}$ and BB cross-sections, i.e., the bare Pomeron (P_0).

This is easily done if we realize that, as in the mesonic case, Pomeron exchange has to be interpreted in QCD as gluonium (M_0) exchange.

6.1 Pomeron exchange in $B\bar{B}$ scattering

There are several equivalent ways of drawing P_0 exchange in the t channel in $B\bar{B}$ collisions. They are shown in Figs. 17a,b,c, and d.

Figure 17a represents a simple QCD graph which contributes to the Pomeron. Figure 17b is the corresponding dual diagram which puts in evidence that the s channel intermediate states are those of a continuum made out of an ordinary $q\bar{q}$ meson and an M_4^J meson. In the t channel we have gluonium which couples to the external lines through ordinary mesons. Inclusion of quark loops leads to the diagrams of Figs. 17c and d. They correspond in QCD and in dual theories, to multiparticle intermediate states, consisting of two jets: a standard $q\bar{q}$ jet of mesons and a $qq\bar{q}\bar{q}$ jet consisting of $B\bar{B}$ + mesons (Fig. 18).

The multiplicity of P_0 in $B\bar{B}$ is the same as in MM or in MB scattering. This is expected from the fact that the bare Pomeron is a factorizable pole [see, for instance Ref. 4)].

If we compare the production mechanism of Fig. 18, whose shadow defines the bare Pomeron, with that of Fig. 8b associated with the Reggeon, we see that in the Pomeron graphs there is an extra active $q\bar{q}$ pair which emits mesons. This difference makes the intercept of the Pomeron higher than that of the Reggeon and near to 1, as discussed in Section 5.2.

We may ask what happens if we take the diagram of Fig. 11b and let each of the two annihilating $q\bar{q}$ pairs produce a jet, as shown in Fig. 19a.

When we sum over the jets (Fig. 19b), a Pomeron-like singularity emerges whose intercept is again found to be near to 1, even though this diagram has a topology which is different from that of the bare Pomeron. For instance, it has intermediate states of twice the multiplicity associated with P_0 . This already suggests that it has to be thought of as contributing to the Pomeron-Pomeron cut.

In order to discuss the situation in more detail, we have to make a digression on the important question of the magnitude of the P_0 coupling to baryons.

This coupling can be estimated if we remember that in the TE the Pomeron is f dominated (when it is not already the $f^{(2)}$) and that in turn the coupling of the f to the baryons is related by exchange degeneracy (see Section 5.4) to that of the ρ . Assuming, as usual, vector meson dominance for the $t = 0$ non-spin flip coupling of the isospin current, one gets $c = c'$ in Eqs. (4.9) and (4.10) and furthermore the validity of Eq. (4.10) also for baryon octets.

For the couplings of the flavour singlet M^1 to mesons and baryons we then find the ratio

$$\frac{g_{M^1 M_i M_i}}{g_{M^1 B_\kappa \bar{B}_\kappa}} \simeq \frac{2C/\sqrt{N_c N_f}}{C \sqrt{N_c/N_f}} = \frac{2}{N_c}, \text{ for } \begin{matrix} i=j \\ \kappa=\kappa \end{matrix} . \quad (6.1)$$

The same formula holds if we consider the couplings of an ideally mixed f to mesons and baryons with only u and d quarks

$$\frac{g_{f M_i M_i}}{g_{f B_\kappa \bar{B}_\kappa}} \simeq \frac{2C/\sqrt{2N_c}}{C \sqrt{N_c/2}} = \frac{2}{N_c}, \text{ for } \begin{matrix} i=j \\ \kappa=\kappa \end{matrix} . \quad (6.2)$$

Under our assumption, Eq. (6.2) implies

$$g_{P_0 M M} \simeq \frac{2}{N_c} g_{P_0 B \bar{B}} \quad (6.3)$$

for the bare Pomeron, so that from the known result²⁾

$$g_{P_0 M M} = O(1/N_c) , \quad (6.4)$$

we finally have

$$g_{P_0 B \bar{B}} = O(1) . \quad (6.5)$$

In our scheme Eq. (6.5) is easily seen to follow from the fact that the Pomeron couples with strength $1/N_c$ to each of the N_c quarks of the baryon (see Fig. 17). This will have important consequences on the relation between RFT for baryons and $1/N$ expansion (see Section 7.3).

Since factorization holds for the bare Pomeron, the results of Eq. (6.3) lead, for the (inelastic, unabsorbed, non-diffractive) total cross-section σ_0 related to the bare Pomeron, to the relation

$$\sigma_0(\pi\pi) : \sigma_0(\pi p) : \sigma_0(pp) = 4 : 2N_c : N_c^2. \quad (6.6)$$

These are, for $N_c = 3$, the predictions of the naïve quark model which are experimentally good to within 10%¹³⁾. This kind of "quark counting rule" has been generalized to broken SU(3) by assuming different amplitudes for processes involving strange and non-strange quarks, i.e., by writing

$$A(h_1 h_2 \rightarrow h_1 h_2) = \sum_{\ell \in h_1} \sum_{m \in h_2} a_{\ell m}, \quad (6.7)$$

$$a_{\ell m} = a(q_\ell q_m \rightarrow q_\ell q_m).$$

A similar result can be obtained easily here. In the TE the bare Pomeron at $t = 0$ is expected to be somewhere between a pure singlet and an ideal f. Introducing the mixing angle θ , one can associate with P_0 the λ matrix

$$\lambda_{P_0} = \begin{pmatrix} \cos \theta & 0 & 0 \\ 0 & \cos \theta & 0 \\ 0 & 0 & \sqrt{2} \sin \theta \end{pmatrix}. \quad (6.8)$$

So that for a tensor dominated P_0 , Eqs. (4.9) and (4.10) give

$$g_{P_0 M_\ell M_\kappa} = \frac{\tilde{e} \delta_{\ell\kappa}}{\sqrt{N_c}} \left[\cos \theta (n_u^M + n_d^M + n_{\bar{u}}^M + n_{\bar{d}}^M) + \sqrt{2} \sin \theta (n_s^M + n_{\bar{s}}^M) \right], \quad (6.9)$$

$$g_{P_0 B_\ell \bar{B}_\kappa} = \frac{\tilde{e} \delta_{\ell\kappa}}{\sqrt{N_c}} \left[\cos \theta (n_u^B + n_d^B) + \sqrt{2} \sin \theta n_s^B \right], \quad (6.10)$$

where $n_f^h(n_{\bar{f}}^h)$ is the number of quarks (antiquarks) of flavour f in the hadron h . Using the factorization properties of the $|P_0\rangle$, we have for MB scattering

$$\begin{aligned} \sigma_0(MB) = & \frac{\tilde{c}^2}{N_c} \left[\cos^2 \theta (n_u^M + n_d^M + n_{\bar{u}}^M + n_{\bar{d}}^M) (n_u^B + n_d^B) + \right. \\ & + 2 \sin^2 \theta (n_s^M + n_{\bar{s}}^M) n_s^B + \sqrt{2} \sin \theta \cos \theta \cdot \\ & \left. \left((n_u^M + n_d^M + n_{\bar{u}}^M + n_{\bar{d}}^M) n_s^B + (n_s^M + n_{\bar{s}}^M) (n_u^B + n_d^B) \right) \right], \end{aligned} \quad (6.11)$$

and similar formulae for $\sigma_0(MM)$ and $\sigma_0(BB)$. In this way we get the quark counting rule of Eq. (6.7) with amplitudes

$$a_{uu} = \frac{\tilde{c}^2}{N_c} \cos^2 \theta, \quad a_{ss} = \frac{2\tilde{c}^2}{N_c} \sin^2 \theta, \quad a_{us} = \frac{\sqrt{2}\tilde{c}^2}{N_c} \sin \theta \cos \theta$$

for uu , ss and us scattering, respectively. An estimate of θ can be obtained experimentally by the relation

$$\frac{\sigma^T(Kp)}{\sigma^T(\pi p)} \simeq \frac{1 + \sqrt{2} \tan \theta}{2}. \quad (6.12)$$

Although we have recovered the quark model result, we see that the physical picture here is very different. The process is not dominated by two-body quark scattering, when the remaining quarks act as spectators: all quarks take part in the reaction and form two jets in the s channel. The quark model result appears for us as a consequence of factorization and of assuming f dominance of the Pomeron, exchange degeneracy and vector meson dominance of the isospin charges.

We can now go back to the discussion of the diagram of Fig. 19b. Consider the structure of the two-Pomeron cut in $B\bar{B}$ scattering. We can identify two components in the two-Pomeron-two baryon amplitude. The first one consists of N_c terms in which the two Pomerons couple to the same quark line. The second one has $N_c(N_c-1)$ "interference" terms in which the two Pomerons are attached to different quark lines. The diagram of Fig. 19b belongs to this second class.

The two types of contributions are characterized by the behaviour of the intermediate diffractive states for large masses. In fact, one can convince oneself that only the diagrams of the first class can have triple Pomeron behaviour (Figs. 20a and b). This is consistent with the result

$$g_{P_0 P_0 P_0} = O(1/N_c), \quad (6.13)$$

as can be seen from the comparison of the magnitude of the two diagrams in Fig. 20a, which gives

$$\left[N_c \left(\frac{g_{P_0 B \bar{B}}}{N_c} \right)^2 \right]^2 = g_{P_0 P_0 P_0}^2 g_{P_0 B \bar{B}}^2. \quad (6.14)$$

6.2 Pomeron exchange in BB scattering

Following the graphical notations introduced for $B\bar{B}$ (Fig. 17), the bare Pomeron diagrams for BB scattering can be drawn in the various ways indicated in Figs. 21.

We see that, as in the usual dual picture, the bare Pomeron in BB is the shadow of the inelastic production of two jets (Fig. 22), each one consisting of a leading baryon plus ordinary mesons.

In QCD language the production mechanism can be described in the following way. Two quarks, one in each of the initial baryons, swap their colour through the exchange of soft gluons. This leads to two colour singlet states of high mass, which then "hadronize" into two jets (Fig. 23). A similar mechanism with soft gluons replaced by hard gluons could be relevant in large-angle processes.

An argument analogous to the one made for $B\bar{B}$ relates the diagrams in which more than one quark line radiates mesons to multipomeron cuts.

By inspection of Figs. 17c and d, and 21 c and d it can be seen that, apart from the leading baryons, the rest of the diagrams have the topology of the bare Pomeron in MM scattering. To be more precise the Pomeron in $BB(B\bar{B})$ scattering contains the graphs with an odd (even, but not zero) number of twisted Reggeon exchanges in the corresponding multiperipheral production. In other words, given the contribution with n twists, the one with $n + 1$ twists will contribute to the amplitude with the same sign, with respect to the former, or with the opposite, according to whether the state in the t channel has charge conjugation $C = +1$ or $C = -1$.

The term with zero twist has already been discussed in Section 5.1 and is depicted in Fig. 8a. It gives the degenerate ω, f, ρ, A_2 trajectory and, as in the mesonic case, gets renormalized by the "cylinder" diagram. As a consequence the $f-\omega$ degeneracy is removed, because of their different charge conjugation, with the ω intercept pushed below $\frac{1}{2}$, while ρ and A_2 remain unchanged.

This analysis shows that, apart from annihilation contributions to $B\bar{B}$, one has, with obvious notations

$$\sigma_{n.a.}^T(B\bar{B}) = P_0 + f + \omega + \rho + A_2 = P_0 + f + \omega + 2\rho \quad (6.15a)$$

$$\sigma^T(BB) = P_0 + f - \omega - \rho + A_2 = P_0 + f - \omega. \quad (6.15b)$$

6.3 Comparison between BB and $B\bar{B}$ cross-sections

Let us first compare total cross-sections for BB and $B\bar{B}$, including annihilation, at the level of approximation at which we have described these reactions so far, i.e., neglecting absorptive corrections and other higher order effects, discussed in Section 7.

Apart from terms which drop faster than $s^{\alpha_M-1} \approx s^{-\frac{1}{2}}$, we only have to keep the Pomeron, the Regge exchange (Fig. 8a) and the leading annihilation diagram of Fig. 14c. We then obtain

$$\sigma^T(B\bar{B}) = P_0 + f + \omega + 2\rho + f_0^J + \omega_0^J \quad (6.16a)$$

$$\sigma^T(BB) = P_0 + f - \omega + f_0^J - \omega_0^J, \quad (6.16b)$$

where we have denoted by $f_0^J(\omega_0^J)$ the annihilation contribution with $C = +1$ (-1).

At this order, having exact ρ - A_2 degeneracy, we get immediately the famous Harari-Freund result that $\sigma^T(BB)$ is isospin independent [$\sigma^T(pp) = \sigma^T(pn)$]. Indeed, also f_0^J and ω_0^J are exchange degenerate (they are built up by annihilation which is absent in BB) and cancel in BB scattering, giving

$$\sigma^T(BB) = P_0 + f - \omega. \quad (6.17)$$

In the Harari-Freund scheme one also has $f = \omega$ and hence a flat total cross-section, if $\alpha_{P_0}(0) = 1$. As we said, for us the f - ω degeneracy is generally broken; however, in some simple models²¹⁾, the flatness of $\sigma^T(BB)$ is regained as a consequence of f and ω decoupling at $t = 0$.

Taking the difference between Eqs. (6.16a) and (6.16b), we have

$$\Delta\sigma = \sigma^T(B\bar{B}) - \sigma^T(BB) = 2(\omega + \rho + \omega_0^J), \quad (6.18)$$

where only the last term, being a cross-section, is positive-definite. If the coupling $\rho B\bar{B}$ is small, as suggested by phenomenology and universality, we finally get

$$\Delta\sigma = \sigma_{ann.}^T(B\bar{B}) + 2\omega. \quad (6.19)$$

We see in this way that the difference between $\Delta\sigma$ and $\sigma_{ann.}^T(B\bar{B})$ is not necessarily zero, but rather measures the importance of the ω trajectory.

In the standard picture $\Delta\sigma$ is identified with the ω exchange, giving rise to the following problem. Experimentally it looks that $\Delta\sigma \simeq \sigma_{ann.}^T(B\bar{B})$ ²²⁾, with the consequence that a central process as annihilation²³⁾, seems to be dominated by a trajectory which, like the ρ , is expected to be built up by peripheral processes. We avoid this difficulty by having two ω like trajectories, which mix -- as we shall discuss in Section 7.2. Since the leading output trajectory is near the ω_0^J at $t = 0$, we can get $\Delta\sigma \simeq \sigma_{ann.}^T(B\bar{B})$ without running into problems with the non-peripherality of annihilation.

We conclude this section by comparing our results with the ones obtained in other schemes.

Eylon and Harari⁷⁾ identify the "three-jet annihilation" with a piece of the Pomeron that they assume to compensate for a missing term in the $B\bar{B}$ Pomeron graphs. Therefore they have

$$\sigma^T(BB) = P_0, \quad (6.20a)$$

$$\begin{aligned} \sigma^T(B\bar{B}) &= P_0 - (\text{missing term}) + (\text{annihilation}) + (\text{Regge exchange}) = \\ &= P_0 + R, \end{aligned} \quad (6.20b)$$

where R stands for the ordinary exchange degenerate ω, f, ρ, A_2 trajectories of Fig. 8a. Hence $\Delta\sigma = R$ vanishes, for instance, for a super exotic initial state, such as $\Delta^{++}\bar{\Delta}^-$. In our case annihilation has nothing to do with Pomeron and corresponds to a new exchange degenerate $\omega_0^J - f_0^J$ trajectory. Consequently, even for a super exotic reaction, $\Delta\sigma$ does not necessarily vanish and the total $B\bar{B}$ cross-section is not flat. This result also contradicts the so-called "Lipkin rule"²⁴⁾.

Our approach is more close to Webber's analysis⁷⁾ and gives a systematic generalization of his ideas. In particular, we are confirming on general grounds his statement that a multiperipheral chain with twisted or untwisted baryon exchanges does not remove $C = \pm 1$ exchange degeneracy.

7. HIGHER-ORDER EFFECTS

We have briefly mentioned in Section 5.1 some higher order corrections associated with Pomeron-Pomeron cuts. Before discussing this type of effects in more detail, we address ourselves the question of mixing baryonium states among themselves through quark loops and with ordinary mesons through junction loops.

7.1 Mixing among baryonium states.

Extension and violations of the ordinary Zweig rule

As we said, the mixing of baryonium states among themselves occurs through ordinary planar loops and is therefore already included at the lowest order of our (topological) expansion.

The mixing of M_2^J states with the $I = 0, 1$ M_4^J states is illustrated in Fig. 24a. In Fig. 24b an example is shown where all baryonium states get mixed in the $I = 0$ sector.

The situation is similar to that of ordinary mesons when loops are included in the Pomeron graphs²¹⁾. Thus the following discussion of the mixing of baryonium trajectories will parallel very much the mesonic case.

At high baryonium masses (M^2) the "mixing" diagrams (Figs. 24) are depressed as compared to the one of Fig. 8a (asymptotic planarity²¹⁾), because the latter is dominated by Regge pole exchange, whereas the others have either a two-Reggeon or a three-Reggeon cut. We then expect the leading baryonium trajectories to approach the unmixed M_q^J trajectory^{*)} at large M^2 .

At $M^2 = 0$ the arguments of Section 5.2 give roughly

$$\alpha_{I=0}(0) - \alpha_{I=1}(0) \simeq \alpha_{I=1}(0) - \alpha_{I=2}(0), \quad (7.1)$$

while, as M^2 becomes negative, we expect the splitting to increase, in agreement with our estimate of the slopes. In conclusion the qualitative behaviour of the leading $I = 0, 1, 2$ baryonium trajectories will look like in Fig. 25.

We now turn to the discussion of the broken $SU(3)$ [or $SU(4)$] quantum number of the baryonium states. This question is related to the form the ordinary Zweig rule takes for them and to its violations.

The situation is again analogous to the one encountered for ordinary mesons. Consider, for example, the M_q^J family. The ordinary Zweig rule for such states is exact if we limit ourselves to the diagrams of Figs. 8 and can be stated in the following way:

- a) The physical M_q^J states are "pure" quark states. (This is the analogue of ideal mixing for ordinary mesons.)
- b) M_q^J decays in those $B\bar{B} + \text{meson}$ states which can be reached through one or more string breakings and $q\bar{q}$ creations.

A consequence of the rule is that the quarks and antiquarks of M_q^J have to be found in the final state.

When one takes into account diagrams such as the one of Fig. 24a, mixing will occur between various M_q^J mesons: e.g., $M_q^J(u\bar{u}s\bar{s})$ will mix with $M_q^J(u\bar{u}d\bar{d})$.

*) Following a Chew-Rosenzweig approach²¹⁾ one would only get a splitting of the $I = 0, 1, 2$ M_q^J , without new trajectories being produced (as in the $f-\omega$ case). However, on the basis of QCD, of the string picture, and the arguments of Ref. 2), we do expect non-leading trajectories related to $M_{q,2}^J$ families to appear.

We can think of this mixing as going through an $M_2^J(u\bar{u})$ state, or through a $B(u\bar{s})-\bar{B}(\bar{u}s)$ state, exactly as we can think of $M(ss)$ mixing with $M(u\bar{u})$ as going through either gluonium or a $K(s\bar{u}) - \bar{K}(\bar{s}u)$ intermediate state.

The conclusion is that, through diagrams such as those of Fig. 24, the $I = 0$ and $I = 1$ trajectories with no strange quarks get shifted and in broken $SU(3)$ mixed with the corresponding trajectories with strange quarks.

Since the magnitude of shifting and the magnitude of $SU(3)$ mixing are related, we expect, as for ordinary mesons, to find at large-positive M^2 small deviation from "pure" quark content (ideal mixing) and a well satisfied Zweig rule; whereas at large-negative M^2 the trajectories will tend to approach pure $SU(3)$ representations. At $M^2 = 0$ the situation will be sort of mid-way with, for instance, $\sigma_{ann.}^T(\Lambda\bar{p})$ smaller, but not too different from $\sigma_{ann.}^T(p\bar{p})$.

7.2 Mixing of baryonium states to ordinary mesons and violations of the new Zweig rule

In order to obtain mixing effects between baryonium and ordinary mesons, we have to add to the diagrams considered so far, junction (or baryon) loops in arbitrary numbers. An example is given in Fig. 26a, where mixing of M_2^J and ordinary $q\bar{q}$ states is shown. M_4^J states with $I = 0,1$ mix to single $q\bar{q}$ states only through a violation of both the ordinary and the new Zweig rule. By violating only the new Zweig rule they can, however, couple to two-jet annihilation channels, as shown in Fig. 26b.

If such a coupling is too strong, it could also wipe out the baryonium states as resonances and make them disappear in the continuum. This is how we can understand, in our framework, the effect found by various authors²⁵⁾ in potential scattering calculations of low-energy $p\bar{p}$ interactions. Before the effects of annihilation are added, there is a whole spectrum of states²⁶⁾ which disappear as a consequence of the annihilation contribution to the effective $p\bar{p}$ potential, if this is too strong.

In our scheme the question of how good the new Zweig rule is for baryonium is related to the magnitude of annihilation as compared to scattering, near threshold.

We have already pointed out that at large M^2 each annihilation contribution is down by a factor $\sim (M^2)^{\frac{1}{2}}$ with respect to the corresponding scattering diagrams, having the same flavour structure. But we know experimentally that annihilation is large at threshold. This does not necessarily imply, however, a large $M_4^J \rightarrow MM$

transition^{*)}. In fact, it may well happen that the total energy discontinuity of the two-jet annihilation is small, or that the three-jet annihilation, in which M_4^J does not appear as an intermediate state, dominates at all energies.

Even in this case, however, one cannot exclude in principle the possibility of a coupling of M_4^J to three-jet annihilation, through intermediate $B\bar{B}$ states, due to non-planar unitarity corrections of the type shown in Fig. 25c. Clearly an estimate of how big $B\bar{B}$ annihilation is in the channels with M_4^J quantum numbers is extremely difficult. However, in our approach one can argue in the following way. If a rather central process (see Section 5.2) like the three-jet annihilation dominates, only the low spin daughters of the baryonium family will communicate with annihilation channels, whereas the states on the leading trajectories, which have spins up to 3 or 4 (Fig. 25), would still show a remarkably good selection rule, forbidding their decay into mesons.

In our scheme, then, at least if the annihilation is dominated by the three-jet process^{**)}, there is certainly room for highly elastic resonances near the $p\bar{p}$ threshold, as narrow as the ϕ or the f' . They could be identified with the recently discovered²⁸⁾ narrow structures in $p\bar{p}$ total and elastic cross-sections.

It is interesting to note that there could also be narrow bound states just below threshold if the new Zweig rule is not too fastly spoiled when M^2 decreases. Experimentally they could be seen as narrow bumps in the missing mass in inclusive reactions such as

$$\bar{p}n \rightarrow \pi^+ X^{--}, \quad \pi^+ \text{ in the fragmentation region of the } \bar{p}, \quad (7.2)$$

where a virtual Δ^{++} annihilates with the neutron through an exotic M_4^J state (Fig. 27a^{***)}). In a similar way one can analyze the reactions (Figs. 27b and c):

$$\pi^- p \rightarrow p X^-, \quad p \text{ in the fragmentation region of the } \pi^- \quad (7.3)$$

$$\bar{p} p \rightarrow \pi^- X^+, \quad \pi^- \text{ in the fragmentation region of the } p. \quad (7.4)$$

As we mentioned, the mixing of baryonium with annihilation channels is expected to increase as we move to lower M^2 .

*) We discuss in detail the case of M_4^J but similar arguments can be constructed for M_2^J .

**) This assumption was first made in the naïve quark model in Ref. 27).

***) See the note on page 18.

Let us discuss the situation at $M^2 \lesssim 0$. Clearly we have no effect on $I = 2$ baryonium and, for $I = 1$, the ρ - A_2 trajectory is quite far from $I = 1$ baryonium. For $I = 0$, however, the leading baryonium trajectory (see Fig. 25) lies near the ω , whose intercept after the renormalization effects due to the Pomeron (Section 6.2) is lowered below $\frac{1}{2}$. These two trajectories are expected to cross at small positive M^2 and strongly mix. The outgoing leading trajectory ω_{ph} will be again around $\frac{1}{2}$ at $M^2 = 0$ but with a small slope $[\alpha'_{\omega_{ph}}(0) \approx \frac{1}{3} \alpha'_\rho]$. As M^2 increases it turns into a normal slope trajectory which approaches the ρ and goes through the physical ω particle. The situation is illustrated in Fig. 28.

7.3 Reggeon calculus for baryons

It was recognized some time ago⁴⁾ that the topological expansion for mesons is connected to the Gribov Reggeon calculus. A term of order $(1/N)^h$ in TE, where h is the number of handles, gives the sum of all enhanced and non-enhanced Gribov diagrams of the order g^{2h} , where one takes the n particle m Pomeron vertex to be of the order g^{n+m-2} .

The fact that, as we have explained in Section 5.1, the coupling of the Pomeron to the baryons is about $N_c/2$ times larger than the one to mesons, modifies the relationship between TE and RFT, when baryons are present.

It is important to realize that this modification depends on the behaviour of the $P_0 B \bar{B}$ vertex as a function of impact parameters, as we vary N_c . Such behaviour in turn depends on the spatial wave function of the baryon as a function of N_c , which of course we do not know.

In the following, for the sake of illustration, we shall consider the case of an N_c independent impact parameter shape of the baryon. The order of a graph in RFT will then be $(1/N_c)^{v_3+2v_4+\dots}$, where v_3, v_4 is the number of trilinear, quadrilinear, ... vertices not involving baryons. Any number of $P_0 B \bar{B}$ couplings can be included at each order in $1/N_c$. Thus in BB or $B\bar{B}$ diffractive scattering in the energy region

$$\frac{1}{N_c} \log s \ll 1 \quad (7.5)$$

the leading diagrams (of order 1) are the s channel eikonal graphs of Fig. 29a (the so-called Froissaron²⁹⁾). The next corrections are given by graphs with just one triple Pomeron vertex (Fig. 29b) and so on. If we consider the region in which

$$g_{P_0 MM} g_{P_0 P_0 P_0} \log S \approx g_{P_0 P_0 P_0}^2 \log S \approx \frac{1}{N_c^2} \log S \ll 1 \quad (7.6)$$

but

$$g_{P_0 P_0 P_0} g_{P_0 BB} \log S \approx \frac{1}{N_c} \log S \approx 1, \quad (7.7)$$

we find that BB (or $B\bar{B}$) diffractive scattering is dominated by the eikonalized tree diagrams in which an arbitrary number of Pomerons is coupled to the baryon lines (Fig. 29c), while MB scattering by those in which only one Pomeron in each tree is coupled to the meson line (Fig. 29d).

This reminds us of the dominance of tree diagrams advocated by Schwimmer³⁰⁾ in particle-nucleus scattering, where the Pomeron-nucleus coupling at fixed impact parameter is taken to be proportional to $A^{1/3}$ (A is the atomic number), as a consequence of the fact that the nuclear radius grows as $A^{1/3}$.

8. CONCLUSIONS AND SUMMARY OF THE RESULTS

The purpose of this paper has been to show how it is possible to generalize to baryons the successful schemes^{1,2)} based on $1/N$ type expansions employed in the mesonic interactions.

We have seen (Section 2) that a trivial extension of $1/N_c$ expansion to baryons is made difficult, if not impossible, by the fact that the quark content of a baryon does itself change with N_c . Our procedure has been, then, to abstract some crucial properties of the $1/N_c$ expansion for mesons, notably the colour structure of the intermediate states leading to the zero width dual limit, and to generalize those properties, rather than the $1/N_c$ expansion itself (Section 3).

The resulting prescriptions have allowed us to discuss the coupling of currents and ordinary mesons to baryons (Section 4) and through them to construct $B\bar{B}$ and BB scattering amplitudes (Sections 5 and 6).

The main results coming out of our approach can be summarized as follows:

- 1) Non-diffractive $B\bar{B}$ reactions exhibit a new type of dual structure between states appearing in the s and t channels (u being the $BB \rightarrow BB$ channel). In the absence of quark loops this duality is between single $q\bar{q}$ mesons, pairs of $q\bar{q}$ mesons, triplets of $q\bar{q}$ mesons and three distinct families of new zero width bound states (called baryonium states) with a $qqq\bar{q}$, $q\bar{q}$ and no quark content, respectively.
- 2) As soon as quark loops are turned on (i.e., in the leading order of the $TE^{(4)}$), baryonium states only decay into $B\bar{B} + \text{mesons}$ or into lower baryonium states + mesons but not just into mesons. This we call the new Zweig rule for baryonium. At this order baryonium satisfies the ordinary Zweig rule as well -- if suitably defined. Thus in the TE $B\bar{B} \rightarrow BB$ duality is between annihilation and scattering channels.
- 3) Baryonium trajectories can be estimated by exploiting their dual relationship to annihilation and using for the latter the dual multiperipheral picture, already successfully employed for the Reggeon and the Pomeron. The impact parameter structure of annihilation is roughly predicted and a new, flat, exchange degenerate $I = 0$ trajectory is found to dominate at high energy. Annihilation is expected to be more central than scattering and to consist essentially of three $q\bar{q}$ jets giving asymptotically $\bar{n}_{\text{ann.}}/\bar{n}_{\text{scatt.}} = 3/2$.
- 4) Our results on the topological structure of amplitudes involving baryons can be rephrased in a dual string picture in which the baryon (for $N_c = 3$) is a Y shaped object with quarks sitting at the ends and with a string junction in the middle. For the string junction we have a QCD interpretation in terms of certain colour indices flowing through the Feynman graphs.
- 5) Diffraction, or Pomeron exchange, can be added to both $B\bar{B}$ and BB scattering via gluonium exchange. In certain approximations we can recover the successful quark counting rule for the total MM , MB , BB cross-sections in spite of the fact that our process is not mediated by single qq or $q\bar{q}$ elementary interactions with the remaining quarks acting as spectators.
- 6) One can compare $B\bar{B}$ and BB cross-sections and discuss the relation of their difference to annihilation. It turns out that $\Delta\sigma = \sigma^T(B\bar{B}) - \sigma^T(BB)$ does not coincide (necessarily) with annihilation. A sort of paradox, connected with the believed central nature of annihilation and the peripherality of ω exchange, is neatly avoided.

7) We are able to discuss (Section 7) the violations of both the old and the new Zweig rule for baryonium, which are related to the question of mixing of baryonium states amongst themselves and with ordinary mesons. We find that annihilation could still let the baryonium picture survive (especially for high spin states near $B\bar{B}$ threshold), at least if it is dominated at all energies by the central three-jet process. The importance of precise $B\bar{B}$ annihilation experiments just above or even below the $B\bar{B}$ threshold is emphasized.

8) There are implications of our results for what concerns the RFT for baryons (Section 7.3). The fact that the Pomeron coupling to baryons comes out of $O(1)$ and not out of $O(1/N_c)$, as in the meson case, makes the kind of TE for baryons we have defined in Section 7.3, more like the RFT expansions recently employed for scattering on nuclei.

Our picture clearly still needs much refinement. On the more formal and theoretical side a systematic classification of QCD diagrams for baryons into topologically characterized gauge invariant classes is certainly needed. At the phenomenological level too we have been leaving quite a lot of possible directions unexplored in which one could go in order to test these ideas.

We hope to be able (or that somebody will be able) to fill in those gaps in the near future.

NOTE ADDED

While this work was being completed, we received an interesting paper by M. Imachi, S. Otsuki and F. Toyoda (Kyushu University preprint 76-HE-4, 1976), where many of the results obtained here are derived within the string model picture mentioned in our Section 5.4.

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Figure captions

- Fig. 1 : Vertices entering in the definition of zero-width dual mesons in QCD. (a), (b) and (c) are creation vertices; (d),(e),(f) and (g) are annihilation vertices.
- Fig. 2 : (a) example of an allowed diagram for the baryon wave function: the flow of the "special colour lines" (heavy solid line, dashed-dotted line, dashed line) is explicitly shown. (b) and (c) examples of forbidden diagrams. The time-ordering, which gives rise to the forbidden intermediate state, is indicated. In (b) the quark paths are indicated by the heavy solid lines.
- Fig. 3 : The point-like vertex $J^k q \bar{q}$.
- Fig. 4 : Definition of the meson-baryon-antibaryon vertex.
- Fig. 5 : Diagrams contributing to $MB \rightarrow MB$ scattering. Type (I) diagrams have meson resonances in the t channel and obey the duality constraint (4.13). Type (II) diagrams are "interference terms" in which the two mesons are not emitted by the same quark.
- Fig. 6 : Construction of the first contribution to $B\bar{B}$ scattering, as a sum of t -channel $q\bar{q}$ meson exchanges. The evolution of the two groups of N_c special colour lines is indicated in (b) by two horizontal dotted lines following the flow of the baryonic number.
- Fig. 7 : A QCD graph corresponding to the diagram of Fig. 6b for $N_c = 3$.
- Fig. 8 : (a) The diagram of Fig. 6b with $N_c = 3$. (b) The same but with insertion of quark loops: the intermediate states in the s and t channel are also indicated.
- Fig. 9 : Dual diagrams relative to the exchange of two $q\bar{q}$ mesons in $B\bar{B}$ scattering.

- Fig. 10 : The possible s channel intermediate states in the second diagram of Fig. 9. (a) $B\bar{B}$ intermediate state; (b) M_4^J intermediate state; (c) M_2^J intermediate state.
- Fig. 11 : (a) Insertion of quark loops, allowing the decay $M_2^J \rightarrow B\bar{B}$.
(b) The second contribution to $B\bar{B}$ scattering.
- Fig. 12 : (a) t channel exchange of three mesons in $B\bar{B}$ scattering;
(b) the third (and last if $N_c = 3$) contribution to $B\bar{B}$ scattering.
- Fig. 13 : Insertion of quark loops in the previous diagram allowing the decay $M_0^J \rightarrow B\bar{B}$
- Fig. 14 : Annihilation contributions to $B\bar{B}$ scattering for $N_c = 3$.
- Fig. 15 : Unmixed M^J trajectories.
- Fig. 16 : String-like representation of M_0^J, M_2^J, M_4^J, M and B .
- Fig. 17 : The Pomeron in $B\bar{B}$ scattering: (a) a simple QCD graph;
(b) dual-type representation of P_0 ; (c) insertion of quark loops in a QCD graph; (d) an artistic view of the dual P_0 when quark loops are present.
- Fig. 18 : The two-jet structure which builds up the Pomeron in $B\bar{B}$.
- Fig. 19 : (a) The four-jet structure which builds up a piece of the two-Pomeron cut; (b) a dual picture of it, showing that the two Pomerons are attached to different quark lines.
- Fig. 20 : (a) The piece of the two-Pomeron cut in which the two P_0 's are attached to the same quark line. The corresponding Regge-Gribov diagram is also shown. (b) Another way of drawing the piece of the two Pomeron cut in which the two P_0 's are attached to different quark lines. The corresponding Regge-Gribov diagram is also shown.

- Fig. 21 : The Pomeron in BB scattering: (a) a simple QCD graph; (b) dual representation of $|P_0\rangle$; (c) insertion of quark loops in a QCD graph; (d) an artistic view of the dual $|P_0\rangle$ when quark loops are present.
- Fig. 22 : The two-jet structure which builds up the Pomeron in BB scattering.
- Fig. 23 : The basic colour exchange mechanism responsible for the two-jet formation in the Pomeron diagram.
- Fig. 24 : Example of diagrams violating the usual Zweig rule: (a) mixing between M_2^J and M_4^J ; (b) mixing amongst all baryonium states.
- Fig. 25 : The leading $I = 0, 1, 2$ baryonium trajectories after mixing amongst baryonium states.
- Fig. 26 : Example of diagrams violating the new Zweig-like selection rule through mixing of baryonium with ordinary $q\bar{q}$ mesons: (a) coupling of M_2^J to one-jet annihilation; (b) coupling of M_4^J to two-jet annihilation; (c) coupling of M_4^J to three-jet annihilation through non-planar loops.
- Fig. 27 : Dual diagrams for the one-particle inclusive reactions:
 (a) $\bar{p}n \rightarrow \pi^+ X^{--}$ with the π^+ in the fragmentation region of $\bar{p}$
 (b) $\pi^- p \rightarrow p X^-$ " " p " " " " π^-
 (c) $p\bar{p} \rightarrow \pi^- X^+$ " " π^- " " " " p
- Fig. 28 : The mixing pattern of the $C = -1$ $I = 0$ trajectories.
- Fig. 29 : (a) Leading diagrams in BB scattering for $(1/N_c)\log s \ll 1$;
 (b) next to leading diagrams in BB scattering for $(1/N_c)\log s \ll 1$;
 (c) leading diagrams in BB scattering for $(1/N_c)\log s \approx 1$;
 (d) leading diagrams in MB scattering for $(1/N_c)\log s \approx 1$.

FIG.1

FIG. 2

FIG.3

FIG.4

FIG.5

FIG.6

FIG. 7

FIG. 8

FIG. 9

FIG. 10

FIG. 11

FIG. 12

FIG. 13

FIG. 14

FIG. 15

FIG. 16

FIG.17

FIG.18

FIG. 19

(a)

(b)

FIG. 20

FIG. 21

FIG. 22

FIG. 23

FIG. 24

FIG. 25

(a)

(b)

(c)

FIG. 26

FIG.27

FIG. 28

FIG. 29