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A 7 Colour Chemistry and Mini-Review on Hadron Dynamics

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This report consists of two logically disconnected parts. The first is a contribution to the spectroscopy session A6 requested by the session organizer Professor Igi; the second is a mini-review on hadron dynamics for the Session A7 organized by myself.

A. Colour Chemistry

Colour is now popularly believed to be the basis of strong interactions, and if so, it is as fundamental as electric charge. Unfortunately, the empirical evidence for colour remains as yet limited. The two often quoted examples are: $R=\sigma(e^+e^-\to \text{hadrons})/\sigma(e^+e^-\to\mu^+\mu^-)$ and $\pi^0\to\gamma\gamma$, the first being a factor 3 and the second a factor 9 bigger than expected in the case of no colour. This is not much evidence for a new degree of freedom as fundamental as colour. Instead, one ought to expect an entirely new class of phenomena qualitatively different from those in a colourless world.

One obvious place to seek such manifestations is in spectroscopy, where a new degree of freedom should lead to a richer spectrum. This richness is not previously noticed because:

(i) Hadrons are colour singlets; colour is distinguishable only at subhadronic levels;

(ii) It cannot be seen in qq mesons and qqq baryons since in each case the quarks can combine in only one way to form a colour singlet, namely $(q^3q^{\bar 3})^1$, $((q^3q^3)(^{\bar 3}q^3))^1$, where colour is denoted by a superscript. To see colour, one has to examine multiquark states, $(qq)(\bar q\bar q)$. Here the number of states are doubled, $((q^3q^3)^3(\bar q^3\bar q^3)^3)^1$ and $((q^3q^3)^6(\bar q^3\bar q^3)^{\bar 6})^1$.

(iii) Even in multiquark states, different coloured combinations are in general mixed and experimentally indistinguishable. In any case, the spectrum is so rich that it is neither feasible nor desirable simply to enumerate the levels. One must select special, readily identifiable configurations for study.

The following configurations are believed to be particularly suitable for observation: $(A)^x \overset{L}{\cdot} (B)^{\bar x}$, where A and B are each aggregates of quarks and/or antiquarks of colour x and $\bar x$ respectively, which are spatially separated by some (high) orbital angular momentum L. The reasons for this belief are as follows:

(a) Colour mixing (dipole) forces, in contrast to confining forces, decrease with distance, so that for large L, states with different x become pure (unmixed) and experimentally distinguishable.

(b) States with higher colour x are more stable and therefore easier to observe. For example, for ordinary $q\bar q$ meson decays such as $f\to\pi\pi$ where $x=3$, only one $q\bar q$ pair need be created on cutting the colour bond, namely $q^3+\bar q^{\bar 3}=(q\bar q)+(q\bar q)$, whereas for the decay $(qq)^6+(\bar q\bar q)^{\bar 6}$ at least two $(q\bar q)$ pairs need be created in order to neutralize the colour 6.

(c) These states often have unusual decay characteristics. For example, in the case of $(qq)-(\bar q\bar q)$:

(i) The decay into mesons: $(q\bar q)+(q\bar q)$ are suppressed for both $x=3$ and 6 since to effect the quark interchange an angular momentum barrier has to be overcome.

(ii) For $x=3$, the decay into baryon-antibaryon is effected by cutting the colour bond and creating a $q\bar q$ pair as in $f\to\pi\pi$, leading to widths ~100 MeV.

(iii) For x=6, the decay into baryon-antibaryon cannot be obtained by cutting the colour bond as explained above, which requires the creation of 2 $q\bar q$ pairs. The preferred decay is by cascade via meson emission into another resonance of the same type, namely $(qq)^6 \overset{L}{-} (\bar q\bar q)^{\bar 6}\to(qq)^6 \overset{L'}{-} (\bar q\bar q)^{\bar 6}+(q\bar q)$. Such unusual decay patterns make these states easy to identify.

(d) The spectrum is largely calculable. For example:

(i) quark-spin (hyperfine) splitting between states can be estimated by one-gluon exchange as for $q\bar q$ and qqq spectroscopy.

(ii) L-dependence (Regge slope α') is calculable from, *e.g.*, the MIT bag model, where $\alpha'_x=\text{const}/\sqrt{C_x}$, C_x being the quadratic Casimir operator for colour x. This fact also helps in their identification.

We shall henceforth focus our attention on these elongated configurations. The simplest example is $(qq)-(\bar{q}\bar{q})$ which we call diquoniums. They are characterized experimentally by suppressed decays into meson modes and relatively large branching ratios into $B\bar{B}$ channels (hence the misnomer 'baryonium'). There are two types.

(A) $(qq)^{\bar{3}}-(\bar{q}\bar{q})^{3}$ which we call "T"-diquoniums. They are broad in general with widths $\Gamma\sim 100$–200 MeV except near or below the $N\bar{N}$ threshold. Because of their larger coupling to $B\bar{B}$, they are best looked for in formation experiments in $N\bar{N}$ collisions, *e.g.*, $N\bar{N}\rightarrow$ "T"

(B) $(qq)^{6}-(\bar{q}\bar{q})^{\bar{6}}$ which we call "M"-diquoniums. They are in general narrow with widths $\sim$10–20 MeV, preferring wherever possible to decay by cascade via π-meson emission into other M-diquoniums. A further identifying feature is that they are relatively easy to produce in reactions such as $\pi+p\rightarrow$"M"$+$X, but are hard to form in reactions such as $\bar{p}p\rightarrow$"M". This has to do with the fact that their mixing with ordinary hadron is strongly /-dependent.

The spectrum, as well as the decay pattern and production mechanism of diquoniums have been worked out in some detail. Experimentally also, a fair amount of information has been accumulated. We summarize here the essential features.

The predicted spectrum for "T" diquonium with u, d (p, n) quarks only is shown in Fig. 1. This has been calculated with one free parameter, the effective diquonium intercept, which was fixed by normalizing with respect to one of the experimental candidates for baryonium;

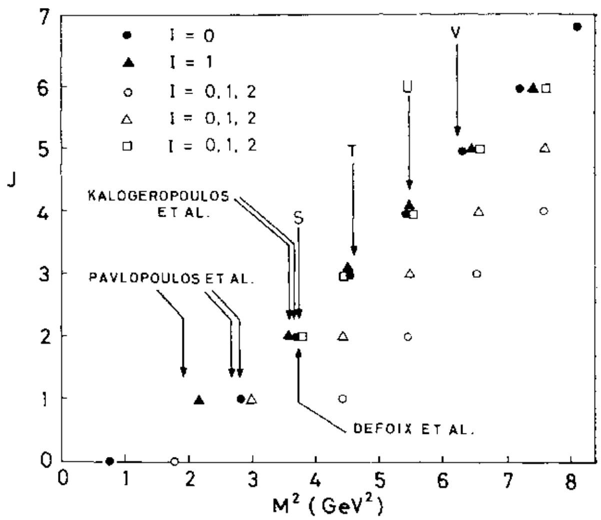

Fig. 1-A

all other parameters having previously been determined from $q\bar{q}$ and qqq spectroscopy. Existing empirical information comes mainly from $N\bar{N}$ formation experiments. Some further states below $N\bar{N}$ threshold are reported in $\bar{p}n$ annihilation using virtual neutrons in deuterium as target and in $N\bar{N}\rightarrow\gamma$ inclusive. All the structures seen in this way and known to us, are shown in Fig. 1, where they are seen to correspond quite closely in mass with the predicted states. Most of these states are known to have large couplings to baryon-antibaryon and small couplings to mesons, and therefore good candicates for T-baryoniums. Above the $N\bar{N}$ threshold, a fair amount of data exist for their decay into various channels which permit some determination of their quantum numbers. In Table I we list the information accumulated from the following reactions: $p\bar{p}\rightarrow p\bar{p}$, $p\bar{p}\rightarrow n\bar{n}$, $p\bar{p}\rightarrow\pi^{+}\pi^{-}$, $p\bar{p}\rightarrow\pi^{0}\pi^{0}$, $p\bar{p}\rightarrow K^{+}K^{-}$, $p\bar{p}\rightarrow\pi^{0}\eta$, $p\bar{p}\rightarrow\rho\omega$. The masses and quantum numbers of the signals observed compare quite well with the dominant peripheral states in the predicted spectrum.

Table I.

Energy region		S			T			U			V		
Predicted states	Traj.	[0, 9]	[2, 36]		[0, 9]	[2, 36]		[0, 9]	[2, 36]		[0, 9]	[2, 36]	
	L	2	0		3	1		4	2		5	3	
	$I^G(J^P)$	$0^+(2^+)$	$0^+(2^+)$	$1^-(2^+)$	$0^-(3^-)$	$0^-(3^-)$	$1^+(3^-)$	$0^+(4^+)$	$0^+(4^+)$	$1^-(4^+)$	$0^-(5^-)$	$0^-(5^-)$	$1^+(5^-)$
	Mass	1.92	1.94	1.94	2.13	2.11	2.11	2.33	2.34	2.34	2.51	2.56	2.56
Exptal states	$I^P(J^G)$	$0^+(2^+)$	$0^+(2^+)$	$1^-(2^+)$	$0^-(3^+)$		$1^+(3^-)$	$0^+(4^+)$		$1^-(4^+)$	$0^-(5^-)$		$1^+(5^-)$
	Mass	1.936	2.0	1.940	2.15		2.15	2.34		2.32	2.5		2.5

Our main interest however is in M-diquoniums whose very existence would be a verification for colour. Again, the spectrum can be calculated up to only one parameter. As explained previously, M-diquoniums prefer to decay by cascades, which are known to be strongly governed by angular momentum barriers. For example, they favour those modes in which a light meson (π) is emitted, while the change in angular momentum ΔL is small and mass Δm large between the initial and final diquonium states. These conditions together with the constraints on isospin for n emission allow one then to estimate rough branching ratios for various cascade modes. In particular, one can deduce in which final states a certain resonance is expected to be found. Thus for example for diquonium made of only u, d (p, n) quarks:

(I) The configurations $(qq)_3^6-(\bar{q}\bar{q})_3^{\bar{6}}$ (where a superscript denotes the colour representation and a subscript 2s+1 for total quark spin .y) with *I=0* are predicted mostly in $B\bar{B}$ final states.

(II) The configurations $(qq)_3^6-(\bar{q}\bar{q})_1^{\bar{6}}$ and $(qq)_1^6-(\bar{q}\bar{q})_3^{\bar{6}}$ with *I=1* are expected mostly to cascade once, into $(qq)_3^6-(\bar{q}\bar{q})_3^{\bar{6}}+\pi$, thus ending up in final states $B\bar{B}\pi$.

(III) The configurations $(qq)_1^6-(\bar{q}\bar{q})_1^{\bar{6}}$ with *I=0,* 1, 2 tend to cascade twice ending up in *BBTZTC.*

The predicted spectrum is summarized in Table II, where each entry represents a near-degenerate multiplet whose members differ only in the relative orientation of the diquark spins. This degeneracy is lifted when spin-dependent interactions between diquarks are introduced. For *L=l₉* mass splittings in each multiplet can be estimated from similar splittings of $q\bar{q}$ mesons. For *L>1* the splitting becomes probably too small to be resolved by present experiments and the whole multiplet will probably appear as a single structure with an apparent width given by the splitting.

Unfortunately, the experimental situation on M-diquoniums is not yet very clear. There are indeed some narrow structures seen with the expected properties of M-diquonium, whose masses agree well with the predicted spectrum, including the *L-S* splitting of the *L=l* states. (Table II) Most of these peaks, however, have yet to be confirmed and there is little information on their quantum numbers. We note that the structure at 2.95 GeV in category 2 of Table II was seen to cascade into 2.020, 2.20 and 2.62(?) GeV, exactly as expected. Also, the fact that all the listed peaks have been observed only in production and not in $p\bar{p}$ formation experiments (formation cross-section of 2.020 and 2.204 is $\sigma_T \lesssim .4$ mb compared with ~ 10 mb for T-diquonium S(1936)) is encouraging for their interpretation as M-diquonium. No candidate yet exists for the interesting *I=2* exotic states in category III. Recently, however, one (perhaps two) exotic state(s) in $\Lambda\Delta^{++}$ and Σ^-p has been reported, well into the strange diquonium spectrum calculated previously by Tsou without introducing any additional parameter.

In Fig, 2 we exhibit the predicted spectrum of diquoniums (both "T" and "M") with $J^{PC}=1^{--}$ which can couple directly to the photon. Below the *NN* threshold, both "T" and "M" diquoniums are narrow. Above the threshold, while "M" diquoniums remain narrow, "T" diquoniums are broad but both

Table II.

	Prepicted peaks in $B\bar{B}$ ($I=0$)		Exptal peaks in $p\bar{p}$			Predicted peaks in $B\bar{B}\pi$ ($I=1$)		Exptal peaks in $p\bar{p}\pi^-$ ($I>0$)		Pre. peaks $B\bar{B}\pi\pi$ ($I=0,1,2$)	Exptal $\bar{p}p\pi\pi$
L	Central mass	$L\cdot S$ Spread	Mass	Quoted width	$I^G(J^P)$	Central mass	$L\cdot S$ Spread	Mass	Quoted width	Mass	Mass
1	2.23	.30	2.204 2.020	Resolved	$0^-(3^-)$? $I=0$?	2.34	.20			2.45	
2	2.56	.11	2.62?	~100	?	2.46	.06			2.77	
3	2.84	.05	2.85	39	$I=0$?	2.95	.03	2.95	.02	3.06	
4	3.10		3.05	~10	$I=0$?	3.21		3.19?		3.32	

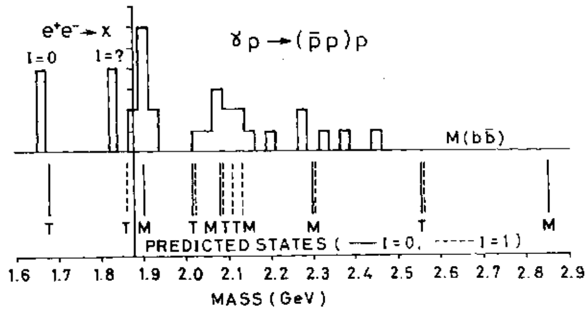

Fig. 2–A

can be indentified by a large branching ratio into $B\bar{B}+n\pi$. Experimentally, these diquoniums can be looked for in e^+e^- annihilation or diflfractive photoproduction. The existing data is summarized in Fig. 2. Two states found by DCI and Frascati at 1.66 and 1.82 GeV respectively, whose narrow widths defy assignment to radial excitation, are seen to fit well with our predicted diquoniums below $p\bar{p}$ threshold. Above the threshold, the $p\bar{p}$ spectrum obtained recently by Bonn *et al* in photoproduction is seen to show exactly the characteristics expected from our theoretical prediction.

Generalization to other multiquark states such as $(q\bar{q})^8$—$(q\bar{q})^8$, $(q\bar{q})^8$—$(qqq)^8$, $(qq)^6$—$(qqq)^6$, $(qqq)^8$—$(qqq)^8$ etc. is fairly straightforward. Experimentally some unusual baryon states have newly been reported which look like appropriate candidates for $(q\bar{q})^8$—$(qqq)^8$ mesobaryons. Now, all these generalizations share the common feature of having two quark aggregates of conjugate colour charges held together and neutralized by a tube of colour flux. This reminds one of a parallel situation in chemical compounds such as Na^+–Cl^- where two ions are neutralized by an ionic bond. The analogy is only superficial since the dynamics is very different, due presumably to the different gauge groups in QED and QCD. Thus whereas an ionic bond breaks easily at large separation, a "chromionic" bond becomes more stable as its length increases. There is, nonetheless, one great advantage in adopting the chemical language. In chemistry, the properties of a molecule are determined to a large extent by the properties of its constituent atoms or ions and of the bonds which hold them together. The same may apply to colour molecules, which would then greatly simplify their study. One may therefore begin by compiling a list of "chromions" and "chromionic bonds" together with their properties, such as masses, stability parameters etc. (a tentative version of this already exists). From such a list, one can then construct one's own molecules and predict their general properties. One need not step at the diatomic molecules so far studied; one could equally well construct triatomic molecules or even perhaps the equivalent of a benzene ring (glue ring). It looks thus as if one has just opened a new area of study which may aptly be labelled "colour chemistry". As physicists, however, we are probably not interested in complication for complication's sake. Our main concern perhaps is to establish colour as a degree of freedom and to learn about colour dynamics. For this purpose, we can restrict ourselves to the simplest molecules such as diquoniums. There it seems that just a few good experiments in perhaps one or two years will settle the issue.

References for Part A

For references, see Chan Hong-Mo and H. Hogaasen: Nucl. Phys. **B136** (1978) 401 and Chan Hong-Mo CERN preprint TH 2540 (1978); Rutherford Laboratory preprint RL-78-089 (1978). Also contributions to this Conference No. 418, 928, 1103.

The previous talks were all on spectroscopy. On behalf of Professor Igi then let me now declare the spectroscopy session formally closed, and switch to my own session on hadron dynamics, beginning with my mini-review.

B. Mini-Review on Hadron Dynamics

This is intended to cover what is now known as 'soft hadron collision' or 'small p_T physics', which account for the bulk of hadronic cross sections. There are a fair number of contributions to this conference, but no new trend is detected. The revolutions by QCD *et al.* so prominent elsewhere in particle physics is not strongly felt here since we are in the infrared region which is particularly difficult to attack. Though some of the contributed papers are quite interesting, they are mainly elaborations in ideas already current for a few years. The best I can do therefore is simply to go over the usual list, noting any new developments. I shall arrange the models roughly in order of the degree of sophistication in dynamics; this has no reflection on the merit of the individual papers which is usually more a junction of the authors than of the subjects they deal with.

At the bottom of this scale are statistical, thermodynamic and hydrodynamic models. Some sophisticated methods are suggested for studying multiparticle correlations etc. Although they may be important for describing and analyzing data, they are clearly not a direct means for probing the dynamics which is seen only as deviations from the models. I shall therefore not spend more time on them.

A little above these are the geometrical models in which the properties of hadronic reactions are ascribed only to the shape of hadrons or to their interaction region. Here, there are two contributions from Schrempp and Schrempp (613, 614) in which the interaction region is pictured as an elongated spheroid with length proportional to $\sqrt{s}$ and width around 1 fm. Some neat scattering theory is employed leading to some interesting predictions, *e.g.*,

$$E\frac{d^3\sigma}{dp^3}(ab\to hX)$$
$$=f_{ab}(p_T)f_{ab}\left(A_h\frac{p_T}{M}\right)\cdot g_h(X_R). \qquad (1)$$

($M=M_x$, $X_R\sim 1-M^2/s$, A_h=constant) deduced from the factorisation of ^-channel Regge pole residues which occur in the theory. This is seen to agree quite well with data over a wide range of the kinematical variables. Such models however have probably the most interesting question unanswered, namely how or why the hadron interaction region acquires this particular shape?

Half-way up the scale comes Regge theory in its various forms. Though not as popular as before it still receives the bulk of the contributions. First, there are some new results on standard Regge phenomenology. I shall quote just three somewhat randam examples:

(i) Geneva-Lausanne (Cleland *et al)* (1081) presented new data on the reactions $K^{\pm}p\to K^{*\pm}(890)p$ and $K^{\pm}p\to K^{*\pm}(1420)p$ at 10 and 50 GeV/c. They noted cross-overs between $d\sigma/dt$ for $K^{\pm}p$ induced reaction at $t=-.3$ GeV[2] for all four cases. Regge fits to 10 GeV data gave the correct shape in *da/dt* at the higher energy.

(ii) SLAC (Ballam *et al)* (604) presented new data at 11.6 GeV on two pairs of reactions related by line-reversal, namely:

$$\begin{cases}\pi^+p\to K^+\Sigma^+, & \text{(a)}\\ K^-p\to\pi^-\Sigma^+, & \text{(b)}\end{cases}$$
$$\begin{cases}\pi^+p\to K^+Y^{*+}(1385), & \text{(c)}\\ K^-p\to\pi^-Y^{*+}(1385), & \textbf{(d)}\end{cases} \qquad (2)$$

measured m the same set-up. The cross sections are equal within experimental errors in the first pair, and after kinematical corrections for t_{min} effects also in the second, in agreement with the prediction from exchange degeneracy of the K^* and K^{**} trajectories. For Σ-production, this result agrees m essence with the earlier data of Berglund *et al.*,[1] although the earlier experiment, having higher statistics, saw deviations at 10-20 percent level. The discrepancy between the two experiments for Y^* production however, has yet to be resolved. Measured polarisations are consistent with weak exchange degeneracy.

(iii) Caltech-LBL (Bonnes *et al)* (1065) presented some new data on $\pi^-p\to(\pi^0,\eta)X$ at 100 GeV/c. These are paritcularly suitable for triple-Regge analysis since in each case only one trajectory is exchanged, namely ρ and A_2 respectively for π^0 and η-production, and the energy is sufficiently high. The ρ and A_2 trajectories extracted from such an analysis of the data is in good agreement with those obtained from the exclusive reactions $\pi^-p\to(\pi^0,\eta)n$. The π^0 production cross section even shows a dip at $t\sim-.5$ GeV^2 corresponding to the wrong signature nonsense zero point for the ρ-trajectory.

We conclude that the standard Regge model for ordinary quantum number exchanges continues to be useful and approximately valid where better data become available.

The recent interest on the spectroscopy of multiquark states as reported by Igi (mini-rapporteur, session A6) is reflected in the search for Regge exchanges of multiquark trajectories, *e.g.*, $qq\bar{q}\bar{q}$. The data relevant for these exchanges accumulated through the years were recently collected and analysed by Nicolescu.[2] Although his conclusions are often too optimistic, they become somewhat more convincing when combined with earlier indirect evidences and the critical prejudices. The intercept of the leading $qq\bar{q}\bar{q}$ trajectory with exotic quantum numbers is probably of an order $\alpha \sim -.7$.

The same interest on multiquark states extends to the study of duality. We recall first an old idea due, I believe, to Rosner based on quark diagrams, which suggests that in $B\bar{B}$ scattering $qq\bar{q}\bar{q}$ resonances in the direct channel are dual to $q\bar{q}$ trajectory exchanges and *vice versa.* All $qq\bar{q}\bar{q}$ resonances are then thought by some to couple mainly to $B\bar{B}$ channels and little to mesons. This last assumption has recently been subjected to a direct phenomenological test by Pennington.[3] He examined the reaction $\pi^- p \to (N^*, \Delta)X$. This satisfies the duality condition illustrated in Fig. 1, where $qq\bar{q}\bar{q}$ resonances in the direct channel are dual to ρ–f exchange. Knowing the ρ–f couplings (estimated by factorisation from $K^+ p \to \Lambda X$), we can calculate the total inclusive cross section of any mass X. This is then compared with the experimental cross section of $\pi^- p \to (N^*, \Delta)$ $(p\bar{p})$ measured by Benkheiri *et al.* at $M(p\bar{p}) \sim 2.2$ GeV where some sharp resonances were observed. It was found that the predicted inclusive cross section is 10-100 times bigger than the $(p\bar{p})$ production cross section where already both the resonant and background contributions have been included. One draws then the conclusion that the bulk of the $qq\bar{q}\bar{q}$ states m the direct channel dual to ρ–f exchange do decay into mesons and not into $p\bar{p}$, in con-

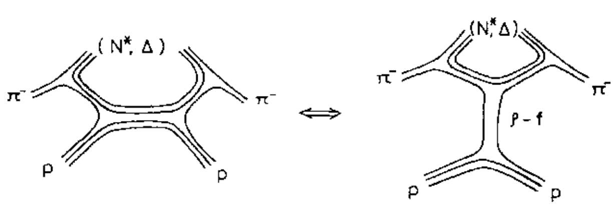

Fig. 1-B

trast to earlier expectations from duality. Such a result, however, would be intuitively obvious in a quark bag picture where, except under very special circumstances as the case considered in Part A of this paper, a $qq\bar{q}\bar{q}$ system can easily recombine into $(q\bar{q})(q\bar{q})$ and decay into mesons.

Theoretically, there have been several attempts to incorporate baryons and $qq\bar{q}\bar{q}$ states explicitly into duality. In my opinion, none of them is perfect. For example:

(i) Inami *et al* (284, 936) assumed that one of the 3 quark lines in a baryon is inactive so that quark diagrams for baryon reactions are planer as for meson reactions. Although the scheme has some attractive features the assumption of an inactive quark line is quite arbitrary and can be justified only by subjection to extensive phenomenologidal tests.

(ii) The Kyushu group (Imachi *et al.* 537) and effectively also Rossi and Veneziano[4] introduce a new constituent to a baryon in addition to the usual three quarks, which can be viewed as a junction joining together three strings. This has some motivation from QCD where the colours of the 3 quarks in a baryon are neutralised at the junction by an ε_{ijk} symbol, as illustrated in Fig. 2. Quark diagrams are now drawn treating the junction line effective as another quark line. Difficulties in this scheme are that: (a) in diagrams for *p-f* exchange, as seen in Fig. 3, the direct channel carries two junction lines and contains necessarily $B\bar{B}$ final states, thus contradicting the observation of Pennington summarised above, (b) There is no explanation, except by injecting new dynamical assumptions, for the apparent existence of narrow $B\bar{B}$ states at high masses, (see, *e.g.*, contribution by Six, session A5).

There are however quite a number of contributions especially from Japan based on the junction model. Therefore, to counteract my own lack of sympathy for this approach, I have asked Professor Uehara later in this session to represent the opposite point of view. In my opinion, the dual properties of baryons, if any exist, are still an open question.

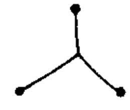

Fig. 2-B

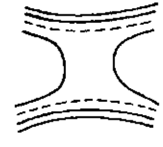

Fig. 3-B

In the last several years, a fair amount of effort was devoted to reconciling duality with unitarity, based on a general framework now known as dual topological unitarisation (DTU).[5] Investigations were carried out in detail only in the approximation where exchanges other than $q\bar{q}$ are neglected. In the zeroth order approximation, planar diagrams are considered. From the ansatz that planar loops should leave planar quantities unrenormalised one deduces certain consistency conditions through which some Regge parameters can be determined. These are then applied to calculate the non-planar loop correction required by unitarity. At the planar level, the following properties are all exact by definition: exchange degeneracy (*e.g.*, between ρ and A_2), isospin degeneracy (*e.g.*, between f and A_2), purity of $q\bar{q}$ states, (*e.g.*, $\omega = n\bar{n} + p\bar{p}$, $\phi = \lambda\bar{\lambda}$) and the Okubo-Zweig-Iizuka (OZI) rule. 1 hey are all broken however by the calculable nonplanar corrections. Hence, both the existence of these approximate symmetries previously taken as independent assumptions, and the magnitude and pattern of their violations are now predictable in dual topological unitarisation.

The main result in this scheme so far is the calculation of the crossed meson loop with the topology of a cylinder, the most detailed version of which was carried out by Tsou.[7] This correction renormalises and mixes $q\bar{q}$ trajectories with zero flavors (*e.g.*, f and ω) but leaves others unchanged (*e.g.*, ρ, A_2, K^*). At $t \sim 0$, the effect was found to be large giving renormalisation of the Regge intercept $\Delta\alpha_f \sim .5$ and hence approximately constant σ_T for intermediate energies, say s< 100 GeV^2. This effect was ascribed in the past to a separate vacuum trajectory called by phenomenologists the Tomeron' (this should be distinguished from *"the"* singularity governing asymptotic behavior considered by some theorists),—hence the confusing terminology: y-Pomeron identity' used to discribe this renormalisation of the /-trajectory. The renormalisation effect decreases as t increases, giving thus a small[6] Pomeron' slope (renormalised $\alpha_f' \sim .3\ \mathrm{GeV}^{-2}$) and becomes quite small in the resonance region implying approximate validity of exchange degeneracy ($m_{A_2}^2 - m_f^2 \sim .1\ \mathrm{GeV}^2$) and the OZI rule, in good agreement with experiment. The deduction from a systematic approximation scheme of all these effects which were taken as independent postulates in the past, represents a major step forward in Regge phenomelogy.

In this Conference, there are some contributions (32, 75, 261, 568) but not much progress reported. Instead some aspects of the scheme came under close scrutiny:

(i) The procedure allows calculation also of the renormalisation in reggeon couplings. Hence, starting from exchange degenerate couplings, the ratio of couplings for the renormalised (*i.e.*, physical) f and ω can be calculated. The results of such calculations have passed previous phenomenological tests, but when applied recently to $K^\pm p \to K^{*\pm}$ (890) p, (32) it was found that the predicted cross section has a much slower energy dependence than the data. In other words the predicted ratio in coupling of P/ω is too large. These analyses however have neglected the considerable mixing between ω and $qq\bar{q}\bar{q}$ states and the contributions of ϕ and f' trajectories which may improve the agreement. Much depends also on one data point above 30 GeV, which may soon be checked by the Geneva-Lausanne experiment (1081). I regard the problem as unsettled at present.

(ii) Pennington *et al* (615) showed that the DTU scheme for '*f*-Pomeron identify' is inconsistent with the usually assumed analytic properties for the trajectory function $\alpha(t)$. The proof is interesting but the result is no obvious objection to the DTU scheme, and may be even expected. We note first that power behaviour in t of the amplitude is expected asymptotically only at fixed s or *l(t)* but not necessarily in the limit where *l(t)* increases with t. Next, we note that in contrast to the planar loop diagram of Fig. 4(a) the spins of the intermediate resonances in the crossed loop of Fig. 4(b) cannot increase with total / in the t-channel because of the twists; any increase in *J(t)* must be at the cost of increasing the orbital angular momentum *1(E).* Now, along the trajectory, $J = \alpha(t)$ linearly with t, hence also *l(t).* For short-ranged interactions (Yukawa type, say) the amplitude must damp exponentially. The usual analytic behaviour assumed for a cannot therefore be valid.[8]

(iii) A recent experiment by Geneva-Lausanne (1082) confirms earlier evidence for the existence of an $I^G(J^P) = I^-(4^+)$ resonance at mass 1.95 GeV. A natural assignment is to the recurrence of A_2 and the isospin partner of h with $I^G(J^P) = 0^+(4^+)$ at 2.04 GeV. If so the mass difference $\Delta m^2 = m_{A_4}^2 - m_h^2 \sim -.36\ \mathrm{GeV}^2$ is of the wrong sign and too large to be explained by DTU. (*cf.* $m_{A_2}^2 - m_f^2 \sim .1\ \mathrm{GeV}^2$).

Some further investigations on these questions are needed.

fa) (b)

Fig. 4-B

The extension of the DTU scheme to include baryon exchanges is difficult because:

(i) duality for baryons undefined; there is no obvious equivalent to planar diagrams.

(ii) there are more quarks in a baryon which means more higher loop diagrams than in the meson case, making the convergence problem worse. Some older calculations which exist involving baryons are good only for crude estimates. They give for example the mixing angles of f and ω with $qq\bar{q}\bar{q}$ trajectories of the order $\tan\theta_f \sim .15$, $\tan\theta_\omega \sim 1$, which would make a big difference to the previous dicussion of $K^{\pm}p \rightarrow K^{*\pm}p$ as regards the problem of 'f-Pomeron identity.' The best way to check this mixing is in the production of $qq\bar{q}\bar{q}$ state,[9] which can proceed in this case by /-Pomeron or ω exchange and not merely by the exchange of $qq\bar{q}\bar{q}$ trajectories with much lower intercept, as illustrated respectively in Fig. 5(a) and (b). Similar considerations apply also to ρ, π and other trajectories. They will make a big difference to the planning of experiments for the production of multiquark states in general. For instance, the diffractive production of a $qq\bar{q}\bar{q}$ states from meson induced reactions may be estimated as $\sin^2\theta_f$ times the cross section of an ordinary $q\bar{q}$ meson state of similar mass and spin. For a mass of around 3 GeV, one obtains a value of some fraction of a μb, quite consistent with the fragmentary information available for the production of some existing $qq\bar{q}\bar{q}$ candidates.

(a) (b)

Fig. 5-B

Difficulties encountered in attempts to account for baryon exchanges are perhaps less serious at high energy where we need consider only 'Pomeron' and multi-Pomeron exchanges. With the bare 'Pomeron' defined by the renormalised f or the topological cylinder, one can effect a formal link-up with the Gribov reggeon field theory (RFT). However, practical calculations of higher order exchanges are difficult, the main impedance being the imprecise formulation of the planar consistency condition which become compounded in the higher order corrections. We note that attempts at a reformulation entirely within the framework of S-matrix theory are being made by Chen, Sursock and collaborators.

Of course, asymptotic problems can be studied in Reggeon field theory without reference to DTU. There has been some recent progress but some confusion remains. The subject has become highly sophisticated and specialised, and I feel incompetent to give you a proper account of it. I have therefore asked Professor Le Bellac later in this session to review for you the situation.

With the recent success of the quark-parton model and QCD in hard processes, it is natural to attempt applying similar ideas to soft collisions also. Indded many constituent models have been suggested in contributions to this meeting. They are based on various assumptions some of which are extremely elaborate (*e.g.*, 219). Since the assumptions are often quite arbitrary, they can only be justified eventually only by comparison with experiment, unfortunately, tests are usually neither extensive nor systematic and it is not easy to judge their merit. I shall quote here only one simple amusing example contributed by Wakaizumi (51,1112). He used an eikonal model to consider the scattering of two protons each made up of N constituents. Assuming a Gaussian form for both the scattering amplitude between constituents and the proton wave functions, he calculated the ratio between slopes of the single and double scattering terms, $R = b_2/b_1$ as a function of N and the relative size of the constituent a to that of the proton r. For $a \ll r$ he finds that:

$$R = b_2/b_1 = (N-2)/2(N-1) \tag{3}$$

Experimentally, from the slopes of the primary and secondary peaks in pp elastic $d\sigma/dt$, estimates the ratio $R \sim 1/5$, which can be obtained in (3) by setting $N \sim 3$. The limit $N \to \infty$, which in this language corresponds to the Chou-Yang case, gives a much smaller value of R= 1/2. He further checked that the same conclusion still holds for somewhat different parameterisations of the constituent amplidute and proton wave function. He was able to get a good fit to experiment by giving the constituent a finite but still small size relative to the proton. It would be amusing to see if similar ideas apply at higher *t* values and to other processes.

There have been some attempts also to relate the dual Regge approach to QCD. In a hypothetical, purely mesonic world this can be formally achieved by the $1/N_c$ expansion of 't Hooft, and its generalisation to $1/N_c$, $1/N_f$ expansions by Veneziano.[10] Calculations can be done however only in the 2-dimensional level which is of doubtful physical significance. Also, a controversy with T. T. Wu has yet to be resolved.[11] A subsequent attempt by Rossi and Veneziano,[4] generalising the considerations to a world with baryons, leads essentially to the junction model of Imachi *et al.* quoted earlier, and shares both its advantages and disadvantages. One may point out also a logical difference here with the original 't Hooft $1/N_c$ expansion. Whereas the usual $1/N_c$ expansion is merely an ordering of all Feynman diagrams obtained from QCD, the baryon model chose a basic set of Feynman diagrams and generate the rest by unitarity. It is not obvious that the set of diagrams so generated is the same as the original set of all Feynman diagrams.

Instead of attempting a link-up with the dual-Regge model, we can try to deduce physical consequences directly from gauge field theories, for example in the asymptotic region. There is an interesting contribution here by Hung Cheng (33) who will speak about it himself later this session. He will also cover an application by Bourrely *et al.* (80) of similar ideas to phenomenology.

On the whole, it is probably fair to say that the progress made in this Conference on soft hadron dynamics is not dramatic. Although, this is due partly to the fact that a fair portion of the general effort has been siphoned off to more exciting fields. One cannot help feeling that it is also partly due to the fact that the phenomenology S-matrix theory which dominated this field for a decade, has finally run out of stream. We have pushed it as far as it can go, and to progress further, some new theoretical injection may be needed.

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A 7 Comment on Chan Hong-Mos Talk: Quasinuclear States in Baryon-Antibaryon Systems

Presented by O. D. DALKAROV

We would like to emphasize that the bound and resonant states of the quasinuclear type in baryon-antibaryons systems have been predicted about ten years ago[1] (for review see ref. 2).

We can resume the main predictions 01 theory as follows: (a) the theory predicts the existence of a large number of heavy mesons *(M~2M,* m-mass of baryon) strongly coupled to the *BB* channel. The latter means that the ratio of the elastic $B\bar{B}$ channel width (real or virtual) to the total width must be close to the unitary limit, *i.e.,* $\Gamma_{B\bar{B}}/\Gamma\sim 1$: (b) the radii at the quasinuclear baryonium states must be large enough: the inequality $R\gg 1/m$ needs to be satisfied: (c) $B\bar{B}$ interaction provides (mainly due to spin-orbit forces) the existence of a rich spectrum of quasinuclear states of baryonuim. Their number must be of about 10 near each $B\bar{B}$ threshold. Thus, the theory predicts several tens of $B\bar{B}$ states in the 1.5-3 GeV mass range, in particular, exotic states, *e.g.,* $\bar{\Sigma}N$ (Isospin 7=3/2) and $\Sigma\bar{\Sigma}$ $(I=2)$; (d) the annihilation widths and level shifts of the baryonium can be reliably estimated[2] (in order of magnitude) provided the theory contains a certain smallness parameter. It has been already mentioned that the ratio of the annihilation radius to the average distance between *B* and $\bar{B}$ $(\tau_a/R\sim 0.1)$ may serve as a smallness parameter. This means that for bound states the probability to find *B* and $\bar{B}$ in the annihilation region is estimated as $|\psi(0)|^2/m^3$ and should be considerably lower than unity. At the same time the annihilation of slow *B* and $\bar{B}$ from the continuum may be large (compared to the unitary limit) due to strong attraction between baryons (the dimensionless enhancement factor is in this case proportional to $|\psi_v(0)|^2$, where $\psi_v(0)$ is the wave function of the continuum state). According to theoretical estimates the annihilation width of the quasinuclear baryonium level may vary from 0.1 to 100 MeV (in order of magnitude) and mainly depends upon the relative orbital momentum of the $B\bar{B}$ pair (roughly, the higher the orbital momentum, the smaller is the width), (e) existence of the baryonium quasinuclear spectrum must manifest itself in a lot of nearthreshold phenomena, for instance, discrete *y*-spectrum in $p\bar{p}$ annihilation at rest, large *p*-wave annihilation from atomic p-states etc, which were predicted theoretically. Most of them have been observed in recent years; (f) the theory predicts that besides the simplest two-particle $B\bar{B}$ (baryonium) system there exist more complicated three- and four-particle quasinuclear systems—the baryons $2B\bar{B}$ bosons $2B2\bar{B}$. The masses of these systems must be about 3 and 4 GeV, respectively.

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A 7 Developments in Reggeon Field Theory

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Although there has not been a tremendous activity in this field since the past conference, there are still a number of open problems, some of which I'll enumerate below.

First, in my view at least, Reggeon field theory (RFT) is an intermediate step in a theory of high energy (H_E) collisions, which is taken in order to sum classes of graphs which are believed to be the dominant ones.[1] For example in the (unrealistic!) $\lambda\varphi^3$ model, one begins by building reggeon from ladders, reggeon interactions from ladder splitting etc. Then the question is the following: how can we derive a Reggeon calculus from a "fundamental" theory of strong interactions (may be *via* some other intermediate step like the topological expansion)? Then, can we compute, or estimate the main parameters of RFT, for example

$$\Delta = 1-\alpha \quad \alpha = \text{Pomeron intercept}$$

r=triple Pomeron coupling etc.

In the framework of gauge theories, which is of course what we would like to take as "fundamental", some progress has been made by a number of people, for example Bartels,[2] to whom I refer for a complete list of references.

The second question, to which I hope to bring a firm answer, is the following: given the RFT hamiltonian, what is the asymptotic behaviour of the total cross section σ_T (for example).

For $\alpha \leq 1$, or more precisely for $\alpha \leq \alpha_c$, where α_c is the critical intercept, everyone agrees. However for $\alpha > \alpha_c$ (supercritical case) the situation is still controversial. According to Amati *et al.*[3] the cross-section rises with energy:

$$\sigma_T \sim (\log s)^2 \quad \alpha > \alpha_c,\ s \to \infty$$

while White[4] claims that:

$$\sigma_T \to 0 \quad \alpha > \alpha_c,\ s \to \infty$$

I'll come back later on this controversy.

Finally my third question will be: since the mathematical apparatus of RFT is rather involved, can we find an intuitive picture in which the results appear to be physically plausible? And also: if the pomeron is supercritical, is RFT the best intermediate step?

I don't have time to develop the first point, which is very technical, and in order not to get stuck at once in mathematical details, let me begin by the third point.

Chemical analogy and parton interpretation

Grassberger[5] has given an interesting interpretation of RFT; following Feynman, he assumes that hadron cross sections at high energy are governed by the interactions of wee partons. If a hadron is boosted, the wee partons can split and recombine, they can be converted into hard partons, and thus are lost from the point of view of the collision, or they can recombine to give hard partons:

wee+wee wee (1a)

wee hard (1b)

wee+wee hard (1c)

This suggests an analogy with chemical reactions. One can take 5 kinds of molecules, *A*, *B*, M, *N* and *R* which undergo chemical reactions in a vessel: we shall be interested in the evolution of the molecules of type *R* (which will correspond to the wee partons), while the chemically inert molecules *A* and *B* correspond to the hard partons (once they are formed, they are lost from the point of view of the chemical reaction, exactly as the hard partons are lost from the point of view of the hadron collision). The chemical reactions corresponding to (1a)-(1c) are thus

$$R+N \underset{\mu}{\overset{\sigma}{\rightleftarrows}} 2R$$

$$R+M \overset{\rho}{\rightarrow} A$$

$$R+R \overset{\nu}{\rightarrow} B$$

where σ, μ, ν, ρ are reaction rates. Other rules of correspondence are:

rapidity ↔ time

impact parameter $\boldsymbol{b}$ ↔ space coordinate x

Pomeron propagator ↔ density of molecules of type

$G(y, b)$ ↔ $R^\dagger$: $n_R(\boldsymbol{x}, t)$

Furthermore, the parameters of the RFT hamiltonian:

$$H=\int \mathrm{d}^{D_t}b\left\{\Delta\bar{\psi}\psi+\frac{ir}{2}\bar{\psi}(\bar{\psi}+\psi)\psi + \frac{\lambda}{4}\bar{\psi}^2\psi^2+\alpha'(\nabla\bar{\psi})(\nabla\psi)\right\} \quad (3)$$

can be put in correspondence with the reaction rates:

$1-\alpha=\Delta\leftrightarrow\rho-\sigma$

$\alpha'\leftrightarrow$ diffusion coefficient D etc. . .

Now, suppose that at $t=0$ I put one molecule of type R at the point $x=0$, and I look at the density $n_R(\boldsymbol{x}, t)$ at a fixed point x, and at large time t. If the production rate α is small $(\alpha\ll 1)$, this density will tend to zero, probably in an exponential way $(\mathrm{e}^{-\delta t})$. If I increase σ, the time constant δ will increase, and for α very large, I can expect an equilibrium situation: $n_R(\boldsymbol{x}, t)\underset{t\to\infty}{\to}\text{constant}$. There will be a critical value σ_c, corresponding to a critical value α_c at which this phenomenon begins to take place:

$\sigma<\sigma_c(\alpha<\alpha_c)$ $n_R(\boldsymbol{x}, t)\to 0$

$t\to\infty$, $\boldsymbol{x}$ fixed

$\sigma>\sigma_c(\alpha>\alpha_c)$ $n_R(\boldsymbol{x}, t)\to\text{constant}$

$t\to\infty$, $\boldsymbol{x}$ fixed.

The region in space in which n_R has reached this constant will expand with some velocity v,

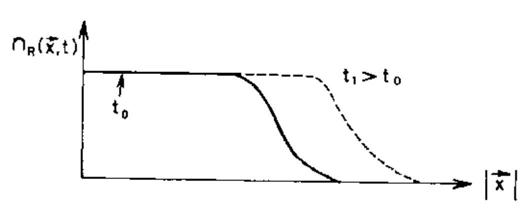

Fig. 1.

Coming back to the parton picture, we see that the density of wee partons will tend to a constant in an expanding disc of radius $vy=v$ log s. This leads (in two transverse dimension: $D_t=2$) to a total cross-section increasing like $(\log s)^2$.

On the contrary, according to White,[4] the density of wee partons will again fall to zero where $\sigma>\sigma_c$. This goes certainly against physical (or rather chemical!) intuition, but we must go a little further into the mathematics in order to ascertain this conclusion.

Reggeon field theory in zero transverse dimension

Fortunately everyone agrees that the issue can be settled by studying what should be a simple exercise in mathematics, but which turns out to be a rather involved problem: compute the evolution operator in zero dimension: $D_T=0$. In this case we can replace the Pomeron field operators (ψ and $\bar{\psi}$) by simple harmonic oscillator annihilation and creation operators:

$\psi\to a$ $\bar{\psi}\to a^+$

$$H=\Delta a^+a+\frac{ir}{2}a^+(a^++a)a \quad (4)$$

We are interested in the following object:

$$U(\bar{z}, z;\ Y)=\langle\bar{z}^*|\mathrm{e}^{-HY}|z\rangle \quad (5)$$

where $|z\rangle$ is a coherent state: $|z\rangle=\mathrm{e}^{a^+z}|0\rangle$. Notice that for $z=-if$, $\bar{z}=-ig$, U is nothing but the S-matrix if f and g (real) are the Pomeron couplings to the external particles:

$$S(f, g;\ Y)=U(-ig, -if;\ Y) \quad (6)$$

In ref. (6) we used a differential equation satisfied by U:

$$-\frac{\partial U}{\partial Y}=\bar{z}\left(\Delta\frac{\partial}{\partial\bar{z}}+\frac{ir}{2}\left(\bar{z}\frac{\partial}{\partial\bar{z}}+\frac{\partial^2}{\partial\bar{z}^2}\right)\right)U \quad (7)$$

with boundary conditions (suggested by the eigenvalue problem $H_E=E$)

$$U(\bar{z}, z;\ Y)\underset{\bar{z}\to -i\infty}{\to}\text{Constant} \quad (8)$$

With this boundary condition U is well-defined, can be proved to be the Borel sum of the exponential series:

$$U=\sum_N\frac{(-Y)^N}{N!}\langle\bar{z}^*|H^N|\bar{z}\rangle \quad (9)$$

and can be computed explicitly (up to the numerical solution of a Schrodinger equation). The main result is that:

$$U(\bar{z}, z;\ Y)\sim\mathrm{e}^{-\delta Y}\quad Y\to\infty \quad (10)$$

$$\delta\cong\Delta \text{ if } \Delta\to\infty\quad \delta\cong e^{-\frac{\Delta^2}{2r^2}},\ \Delta\to-\infty$$

The value of δ for Δ large and negative is reminiscent of a tunnelling effect. The results can be summarized in the figure below:

[T] Assuming that at time $t=0$ one molecule of type R has been put at the point $X=0$.

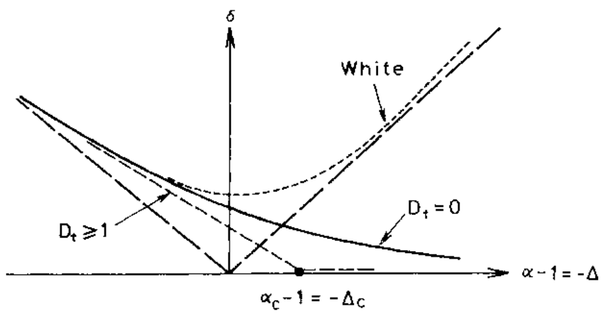

Fig. 2.

In non-zero transverse dimension ($D_t > 1$), the gap δ vanishes for a finite value $\alpha = \alpha_c$ of the Pomeron intercept, leading to a phase transition. As one may expect, the phase transition disappears, being pushed at infinity, when the transverse degrees of freedom are frozen.

I won't describe White's arguments,[4] since he seems to have changed his mind recently,[7] but I'll only mention new arguments which indicate that the CLR solution[6] is what is needed from a physical point of view. In fact one can try the following things:

1) Derive a Lagrangian PI formulation equivalent to the CLR-Schrodinger equation, from the hamiltonian form. This has been attempted by Ciafaloni,[8,9] but no complete proof seems to be possible if there is only a triple Pomeron coupling, since the hamiltonian PI itself is not well defined. However, if one adds a $\lambda a^{+2} a^{2}$ term, this term acts as a regulator and a complete proof seems to be possible

2) Prove that U_{CLR} gives the Borel sum of the perturbative series. This has been proved for $\Delta = 0$, but it seems difficult to extend the proof to the case $\Delta \neq 0$

3) Prove the $D_t = 0$ version of Reggeon unitarity:

$$\sum_n \langle \bar{z}^* | e^{-HY_1} | n \rangle \langle n | e^{-HY_2} | z \rangle = \langle \bar{z}^* | e^{-H(Y_1+Y_2)} | z \rangle \qquad (13)$$

$|n\rangle$ = bare Pomeron states.

A formal proof of (13) is easy to find, starting from ref. 6 and it is probably possible to make it rigorous.

To summarize: it is very likely that U_{CLR} give the correct physical solution, (remember that Raykin and Ryskin had also to withdraw their criticisms[10]) and that the supercritical Pomeron behaves in the way advocated first in ref. 3

Conclusion

There are many other developments which I could not examine in this talk. Among them I would like to quote:

—the contribution of Alessandrini *et al.*[11] who exhibited the states responsible for the expanding disc

—the calculation by Brower *et al.*[12] of the critical indices in the spin model which checks nicely with those of the continuous theory

—Finally applications to phenomenology, inclusive processes[13,14] etc.

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A7 Critical and Super-Critical RFT

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RFT is both a method of summing Reggeon graphs and a solution of *t* and s-channel unitarity at high energy. In my path-integral formalism the RFT phase transition is described directly in terms of Feynman (Reggeon) diagrams and the unitarity of the theory is explicit. The phase-transition is unambiguously described by a field shift to a new vacuum in the path-integral formalism provided the creation and destruction operator character of the Reggeon fields is respected. I have space only to describe the results.

The super-critical Reggeon expansion satisfies Reggeon unitarity and (hence) s-channel unitarity, with a(0)<l. This implies *the total cross-section rises only at the (very special) critical point.* The new phase is distinguished from the sub-critical phase by the appearance of singular (in momentum transfer) Pomeron interaction vertices shown in Fig. 1. Cut

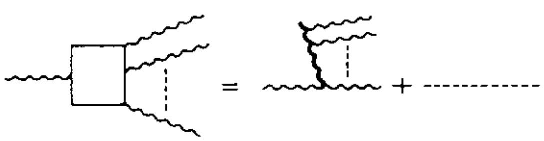

Fig. 1.

RFT (AGK cutting rules) shows the .v-channel unitarity property of the theory. From this it can be seen that the singular Pomeron interactions are due to *degenerate odd-signature* partners to the Pomeron which decouple at the critical point, and away from the critical point look very like the Reggeised "gluons" of a spontaneously broken gauge theory. Therefore the critical Pomeron has a "dual" property which looks very appropriate for strong interactions — a_Q -> a_{Qc} (—) => Pomeron constructed from usual bare expansion using hadrons. a_0 a_{0c}(+)=> Pomeron constructed from Reggeised, gluons which become massless and (confine) at the critical point.

There results give the (optimistic) hope that studying a_0 $a_{0c}(+)$ in RFT derived directly from (broken) gauge theories (*e.g.,* Bartels work) will show that the critical Pomeron occurs in unbroken gauge theories including QCD. It is important that RFT can be derived directly in this situation from the multiparticle dispersion relations derived by Stapp and myself.

RFT with zero transverse dimensions (D=0)

The RFT spin model and the expanding disc solution to super-critical RFT are based on the *D=Q* theory with just a triple Pomeron coupling. The theory with $a_0>l$ can be defined either directly from the Fock-space path integral or (see Bellac's talk) by analytic continuation (out of the Fock space) from $a_0<l$. The energy spectrum is shown in Fig. 2.

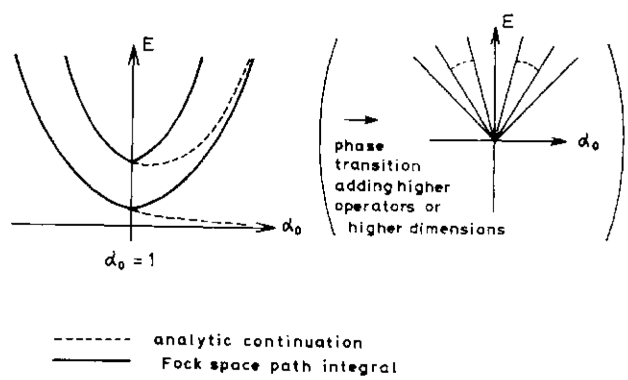

Fig. 2.

The degeneracy with the vacuum (tunnelling effect) which occurs when a_0-> — oo (in the analytically continued theory) is argued to be analogous to that in $\lambda\varphi^4$ in one dimension. The phase transition in higher dimensions is then studied by spin-model methods successful for However, RFT is *very different* to Critical RFT is *not* infra-red singular as Z 0 whereas massless $\lambda\varphi^4$ is extremely singular. The RFT phase-transition should *not disappear* as Z 0. Instead all higher-order couplings acquire the same canonical dimension and become "relevant" operators for the phase-transition, which will involve an infinite number of degenerate zero-energy states, and not just two as in the analytically continued triple Pomeron theory. In fact $a_0 = 1$ (not $a_0 = -$) is most closely analogous to a phase-transition in the triple Pomeron theory since there is a "change of vacuum" at this point leading to non-analytic behaviour of the energy levels. This suggests that the tunnelling effect and similarly the spin-model expanding disc are solutions to RFT which occur far from the critical point and also lie outside of the bare Pomeron Fock space.

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A7 The Pomeron in Gauge Theories

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There are various ways to study the Pomeron which dominates scattering processes of *s*Z *o* and small$p_{\pm}$. One way is to construct models. There exist two general classes of models for the Pomeron:

(i) *The eikonal models:* In these models, one studies the S-matrix as a function of the impact parameter *b*. For this S-matrix, one puts

$$S(s, b) \sim e^{i\chi(b)}, \tag{1}$$

where % has a non-negative imaginary part. In this way, the t-channel unitarity has been incorporated. The central question in such models is in the choice of the eikonal *%(b)*. If we treat potential scattering, one finds that *X* is simply related to the potential. In the droplet model, χ is taken to be purely imaginary and is related to the distribution of the constituents in the hadrons

(ii) The Regge models: In these models, one utilizes the fact that the elastic scattering amplitude due to the exchange of a Regge pole in the t-channel is

$$\beta(t)s^{\alpha(t)}. \tag{2}$$

The central question in such models is how to treat other exchanges, such as the exchange of two or more Regge poles. The Reggeon field theory represents one attempt to answer this question.

One notices that while these two classes of models both incorporate certain physical aspects of high-energy scattering, they appear irreconcilably different.

Instead of constructing models, we may try to derive everything from the most fundamental level. In these days, this means deducing high-energy scattering from gauge field theories.

In the late sixties, a great deal of work in this direction has been done. Most of such work was, however, based on perturbation or Feynman diagrams. Such an approach has two intrinsic difficulties: (a) It cannot yield non-perturbative effects such as confinement; (b) Almost all individual diagrams for elastic scattering yield terms larger than the Froissart bound. Consequently, extensive cancellations occur among these terms.

Because of these difficulties, it appears that one has to sum almost all diagrams to be assured of the correct answer.

Since we cannot sum all diagrams, we

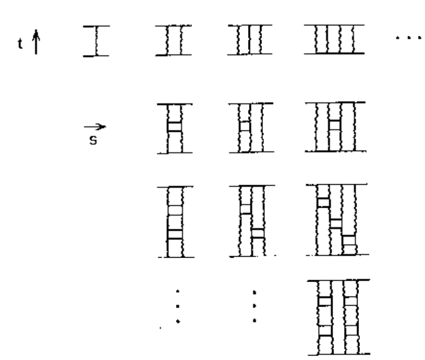


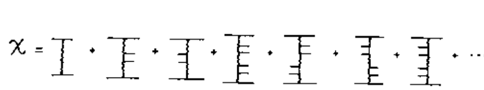
Fig. 1.

started with a program which is a great deal more modest. We summed the following set of diagrams for elastic scattering in QED.

Remarks

1. All permutations of the gluon vertices are included in the set.
2. For each diagram, we calculate only the leading term.
3. For the pair creation process, for example, we attach a pair to any of the virtual gluon lines.

We have no space here to explain in detail how these diagrams are chosen. It suffices to say that the above set of diagrams, together with their counterparts for inelastic processes, are generated by unitarity and crossing from the following multi-peripheral diagrams:



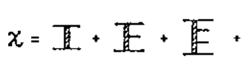
(3)

The calculation of such a large class of diagrams is tedius, but the result is surprisingly simple. We get

$$S = e^{i\chi} \tag{4}$$

where x is an operator the matrix elements of which are represented by the diagrams in (3). We wish to remark that: (i) eq. (4) gives exactly all the leading terms calculated, (ii) It is valid not only for the elastic amplitude but also for all inelastic amplitudes; (iii) Unitarity is explicit as x is hermitian; (iv) If we expand e^{ix} into a Taylor series of then, for elastic scattering, the χ^4 term represents the diagrams in the 2nd column (exchange of a Regge pole or Regge branch point), while the χ^2 term represents the diagrams in the 4th column (exchange of two Regge poles plus others). Thus (4) unites the eikonal models and the Regge pole models.

Next we turn to Yang-Mills theories. Here we have calculated only the elastic amplitude and up to only the 10th perturbative order. We again get the eikonal form (4), with



(5)

where ▨ represents a Reggeized gluon. In other words, the propagator of ▨ is $s^{\alpha(q_1)-1}/(\boldsymbol{q}_\perp^2+\lambda^2)$, where X is the mass of the gluon. For the scattering of colorless hadrons, we may set $X=0$ without encountering infrared divergences.

There remain the two big questions: (i) How to take care of confinement? (ii) How about the diagrams we have not included? Since (4) was obtained rather unexpectedly, we suspect that its validity exceeds the derivation above. We therefore try to formulate the problem without perturbation.

Let us consider the scattering of a hadron by a potential. The Lagrangian is

$$\mathcal{L} = \mathcal{L}_0 + \mathcal{L}_I + g\rho(x)V(\boldsymbol{x}), \tag{6}$$

where $\mathcal{L}_I$ is the strong interaction. The scattering matrix is

$$S_{fi} = \left\langle \phi_f^{(-)} \middle| \left\{ \exp\left[-ig\int \rho(x)V(x)\mathrm{d}^4x \right] \right\}_+ \times \middle| \phi_i^{(+)} \right\rangle,$$

where $\phi_i^{(+)}$, $\phi_f^{(-)}$ are the dressed incoming and outgoing states corresponding to $\mathcal{L}_0+\mathcal{L}_I$, and { }+ denotes the product ordered with respect to $x_+ \equiv t+z$. We now *assume* that all particles created in the scattering process have large and positive z-momentum. Then in *pipe),* the scale of x_+ is very large and the scale of x_- is very small. Since

$$\lim_{L\to\infty}\int f\left(\frac{x}{L}\right)g(x)\mathrm{d}x = f(0)\int g(x)\mathrm{d}x,$$

we may set x_+ in $p(x)$ to be zero and the ordered integral in (7) becomes an ordinary integral. Thus we may integrate the exponent in (7) directly and obtain

$$S_{fi} \cong \langle \phi_f^{(-)} | \exp(i\chi_0) | \phi_i^{(+)} \rangle, \tag{8}$$

where

$$\chi_0 = -g\int \rho(x_+=0)V(x_-=0)\mathrm{d}^4x.$$

The dressed states and the bare states are related by

$$\phi_f^{(-)} = \Omega\phi_f,\ \phi = \text{bare state},\ \Omega \equiv U(0,\infty) \tag{9}$$

If i is a single particle state, we have also

$$\phi_i^{(+)} = \Omega\phi_i. \tag{10}$$

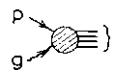
Substituting (9) and (10) into (8) gives

$$S_{fi} \cong \langle \phi_f | \exp(i\chi) | \phi_i \rangle,\ \chi = \Omega^{-1}\chi_0\Omega. \tag{11}$$

This is the operator eikonal form (4).

The matrix elements of x are given by

$$\langle\phi_m|\chi|\phi_n\rangle=\langle\psi_m^{(-)}|\chi_0|\psi_n^{(-)}\rangle. \qquad (12)$$

If n is a single particle state, (12) is the particle-gluon scattering amplitude 

m. And if *m* is also a single particle state, (12) is the form factor 

It is seen that (3) is the lowest-order approximation of (12). We note that since the gluon has small energies, $s_1=(p+q)^2$ is much smaller than ***s***. Thus (3) is a good approximation of (12) as long as s_1 is not very large and the coupling constant is small. When ***s*** becomes extremely large, may also be large, then we must express the particle-gluon scattering amplitude, in turn, by exponentiation. By repeating this process of exponentiation, we may obtain good approximation of the amplitude for arbitrarily high energies.

Equation (8) also suggests that, to treat confinement, one simply put in the confinement wavefunction as in the Glauber form or the impact factor form. It also suggests that, for hadron-hadron scattering, impact factors for both hadrons must be used.

In conclusion, we have succeeded in extracting high-energy amplitudes from gauge field theories and much phenomenological work awaits us in the future.

Finally, I like to report the phenomenological works done by Bourrely, Soffer and Wu as well as by Fujisaki and Tsukahara. Both of these works use models of the eikonal class. Both give good description of the p-p differential cross section up to $|t| = 10$ GeV^2 at ISR energies. They also show the absence of the second dip even when $\sqrt{s}$ is as high as 10^3 GeV. Some of their results are illustrated in the following figures.

These works are in contrast with the phenomenological analysis of Chou and Yang, who concluded that a second dip will develop for p-p scattering.

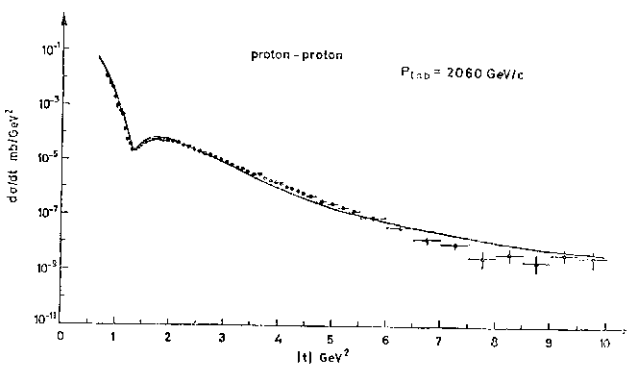


Fig. 2.

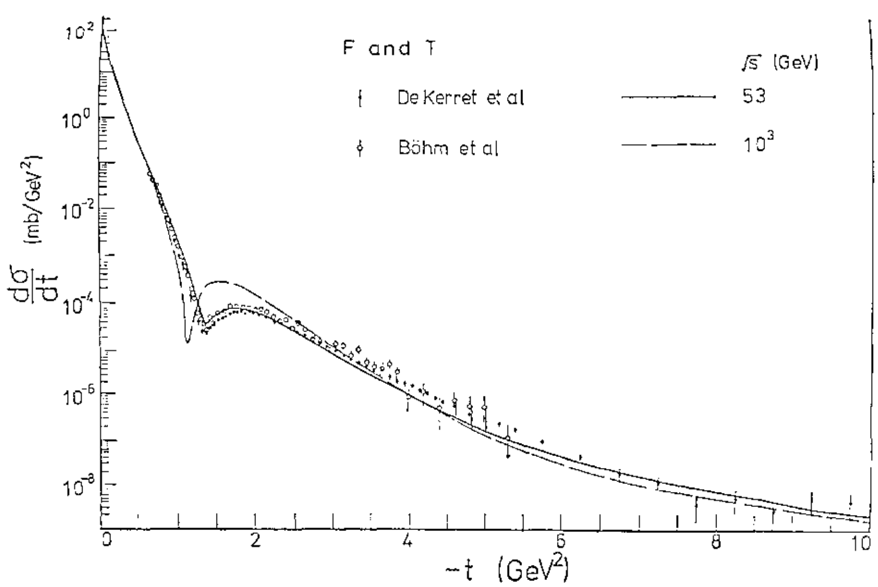


Fig. 3.

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A 7 Duality and Baryonium in the String-Junction Model

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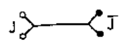
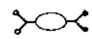
The string-junction model[1] leads to the expectation of massive exotic mesons with junction-antijunction pairs, which we call baryonium as it is a popular name. The leading family of them has two junctions and three different topologies depending on the number of quarks, $M_4^2=$  $M_2^2=$  and $M_0^2=$  .

One of the experimentally mysterious aspects of baryonium candidates is that there are observed narrow states in baryon-antibaryon channels with masses far above the $p\bar{p}$ threshold, but they are not observed in $\sigma_{\bar{p}p}^{tot}$ and $\sigma_{\bar{p}p}^{el}$. In order to resolve the problem we propose a new scheme of baryonium assignment within the string-junction model, that is, observed narrow states are assigned to M_2^2 and wide ones to M_4^2, since an M_2^2 (with $l>0$ at least) turns out to be narrow because of its non-planar structure and of the small EXD breaking of baryons.

Adopting a linear trajectory for M_2^2 given by generalized dual topological unitarization,[2] $\alpha_2^2(s)\sim 1/2\alpha's$ $(\alpha'\sim 0.9)$, we predict the leading spin states, made of non-strange quarks, as shown in Table. The phenomenological implications are summarized: (1) Narrow states corresponding to our M_2^2 would not be observed as clear bumps in $\sigma_{p\bar{p}}^{tot}$ and $\sigma_{p\bar{p}}^{el}$ because of the weak coupling to $p\bar{p}$ and 01 relatively low spins. (2) These states can occupy only a very small fraction of an inclusive spectrum in *MB-^B'X,* since M_2^2 is not dual to the leading M and M_0^2 reggeon exchanges, the latters being dual to wide M_4^2 and continuous meson productions, respectively. (3) Narrow

J^{PC}	1^{--}	2^{++}	3^{--}	4^{++}
M(GeV)	1.5	2.1	2.6	3.0
Expt.	1.49	2.02	2.62?	2.85 } ?
Narrow	*via* e^+e^-	2.20	*via* cascade	2.95 } ?
Bumps			from 2.95	3.05 } ?

exotic flavor states, if they are confirmed, have to be assigned to the next family with 4 junctions such as  or  , but their produciton would be suppressed compared to the production of the leading family because of weak couplings to conventional hadrons and of their low lying regge trajectories.

The narrowness of an M_2^2 width is seen as follows: Let's take a discontinuity of a *BB* elastic amplitude saturated by M_2^2 poles, which is dual to an MM-cut. The vertices ot the cut connected to external baryons contain twisted lines of produced baryons. The amplitude corresponding to twisted lines of baryons is suppressed compared to the one corresponding to untwisted lines of baryons owing to the cancellation between baryons with opposite signatures, similarly to the suppression of the OZI breaking decays in the meson sector. Thus we have $\Gamma(M_2^2\rightarrow B\bar{B})\propto \varepsilon_B^2 s^{2\alpha_M-1}$, where ε_B is the ratio of the twisted baryon lines to the untwisted ones, a small number. Mesonic decay as well as cascade decay is also suppressed by the similar argument to suppress mesonic decay of M_4^2. (See ref. 3 for details.)

The string-junction model forms a sharp contrast to the bag model or the color chemistry approach[4]: Our baryonium states can communicate with conventional hadrons through unitarity and duality both in the time- and space-like regions, while the mock states are made of completely new entities which are not connected to conventional hadrons through unitarity and duality.

Our scheme is based on the viewpoint that there are hadrons with closed rings as well as tree-type hadrons. There is an alternative viewpoint that there should be only tree-type hadrons. In the latter viewpoint there are no M_2^2 distinguished from M_4^2, that is, we are led to an M_0^2–M_2^2–M_4^2 identity hypothesis like the Pomeron-/' identity hypothesis. The baryonium spectroscopy will shed a new light on the problem concerning the Pomeron-/ identity hypothesis.

References

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2. G. Rossi and G. Veneziano: *loc. cit.;* M. Uehara: Progr. theor. Phys. **59** (1978) 1587.
3. M. Uehara: Contributed paper No. 561 and preprint Saga-78-2.
4. Chan Hong Mo: Mini-raporteur talk and preprint RL-78-027.

Session A9: High Energy Hadron Reactions, Large p_T and Jets

Chairman: R. COOL

Organizer: P. DARRIULAT

Scientific Secretaries: K. NAKAMURA
T. INAGAKI

1. (i) Quark-Quark Scattering with Gluon Corrections
(ii) Results from E260

a Fox

2. Results from the Fermilab-Lehigh-Penn-Wisconsin Collaboration

W. SELOVE

3. Scaling of Symmetric Hadron Pairs

R. L. MCCARTHY

(Friday, August 25, 1978; 16: 35-17: 20)

1. Results from the CERN-Columbia-Oxford-Rockefeller Collaboration

R. COOL

2. Results from the CERN-Saclay Collaboration

A. CLARK

3. Results from the Split Field Magnet Detector

K. H. HANSEN

4. Results from the Athens-BNL-CERN-Syracuse-Yale Collaboration

K. NAKAMURA

5. Direct Photon Production

B. BORGIA

(Saturday, August 26, 1978; 9: 00-10: 30)