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CONFINEMENT THROUGH REGGE-LIMIT INFRA-RED ANALYSIS OF QCD

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A B S T R A C T

It is shown that the usual mass-shell infra-red analysis of QED can be carried out indirectly in the Regge limit and that a more extensive analysis of QCD can similarly be performed using multi-dispersion theory based multi-Regge theory. It is argued that the exponentiation of reggeization "screens" all but a special class of mass-shell singularities. This class is shown to be due to the soft gluon content of the vacuum and argued to produce confinement with chiral symmetry breaking. In addition the resulting pomeron is strongly dependent on both the gauge group and the number of fermions. Experiment favours QCD with SU(3) colour, and saturated with quarks !

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INTRODUCTION

In this talk I shall describe an approach to solving a central part of QCD which is not amongst the presently most popular but which I believe may eventually prove to be very powerful. The basic aim of the formalism[1)] which I shall outline is to understand and calculate the pomeron (diffraction scattering) in QCD. However, the deeper that we probe into the pomeron the more we find that it is intimately related to the central problems of confinement and chiral symmetry breaking in QCD. In Fig. 1 I have made a qualitative "Chew-Frautschi plot" for QCD separating the "resonance region", to which most "non-perturbative" work is directed, from the "$t \sim 0 \sim m_\pi^2$ strip" where my formalism can be applied and where I believe an understanding of the pomeron is essential.

Many people would probably agree that to believe the pomeron can really be described within a particular formalism it is certainly necessary to also understand confinement, at least qualitatively, within the same context. (It is, of course, much more unconventional to suggest, as I shall do, that the Pomeron is itself necessary for a complete understanding of chiral symmetry in a confining theory.) In this talk therefore, I shall first describe how I see the confinement hadronic spectrum emerging from my formalism and then outline the very beautiful results, which then follow, relating the pomeron in a very detailed way to both the structure (centre) of the gauge group and the fermionic content of the theory. I would like to emphasize that while it is possible that with very much effort my "derivation" of the hadronic spectrum could be elevated to the status of a "proof of confinement" (if this can have any meaning short of giving a complete expression for a unitary QCD hadronic S matrix), I regard it rather as a necessary confirmation that I am handling the pomeron correctly.

The technical tool that I shall use is very highly developed "S matrix theory" in which multiparticle dispersion theory[2)] is used as a basis for multi-Regge theory[1),2)]. I shall argue that many technical points and points of principle combine to allow me to apply this formalism to first extrapolate existing perturbative calculations and then to analyze simultaneously the mass-shell infra-red problem of QCD and the high energy Regge behaviour. A comprehensive development of the formalism is given in Ref. 1) and a non-technical account is given in Ref. 3). Here I wish to confront technical problems directly but aim at a level of presentation which can be followed by a general ("well informed") reader.

Since I have to analyze the infra-red problem of QCD it will clarify my procedure if I first review the infra-red problem in QED and discuss its significance in the Regge limit. Of course, since I am to arrive at a confining theory my analysis must go way beyond the infra-red structure found in QED. Beginning from a perturbation theory where there are massless gluons I must show that all infra-red divergences cancel or exponentiate themselves (to finite answers) except those which can be interpreted as due to a "vacuum" background of soft gluons. [It is now commonly accepted that the infinite volume limit for QCD generates vacuum expectation values for many gauge invariant gluonic operators[4),5)]. We shall find an S matrix analogue of this effect.] The nature of this "vacuum" is central in what I shall describe and so if the reader is to believe the results I shall quote for the Pomeron he must clearly believe that I have adequately located the confining vacuum.

QED INFRA-RED ANALYSIS

The most well-defined way to handle the infra-red divergences of QED is to give the photon a mass M and construct the Fock space S matrix $S_F(M^2)$. We can then write

$$S_F(M^2) = Z(M^2)\ \tilde{S}(M^2)\ Z^+(M^2) \tag{1}$$

where

$$Z = \exp\left[R + i\phi\right] \tag{2}$$

R contains the real "radiation" divergences while ϕ is conventionally referred to as the "coulomb phase operator". Z contains all the divergences as $M^2 \to 0$ and so potentially we can identify $\tilde{S}(0)$ as the QED "S matrix" between coherent asymptotic states[6)-9)] of the form

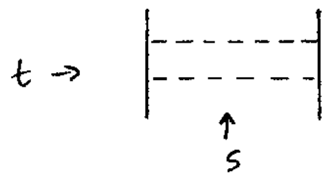
$$|\text{coherent state}\rangle = Z^+ \,|\text{Fock space state}\rangle \tag{3}$$

The need to define coherent physical states can (qualitatively) be seen as required by gauge invariance[10)]. Local gauge invariant operators cannot be defined for charged particle states. That is $\bar{\psi}(x)$ cannot be a gauge invariant (bare) electron creation operator but $\bar{\psi}(x)\ \exp\left[\int_x^\infty dx'_\mu\ A_\mu(x')\right]$ could be. The exponen- exponentiation of infra-red divergences can be directly described by operators of the latter form[11)] although momentum space coherent state operators are usually more convenient.

In perturbation theory, the factorizability of the infra-red divergences, as in (1), is well known to be due to the external line rule for soft photons[12)-14)]. That is all soft photon lines (causing infra-red divergences) are attached only to external (on mass-shell) charged particle lines as illustrated in Fig. 2.

The question we need to consider is, how much of the above infra-red singularity structure is retained in the Regge limit ? Specifically the Regge limit in QED (and QCD) is described by transverse momentum diagrams which are infra-red divergent as $M^2 \to 0$. What is the significance of the divergences ?

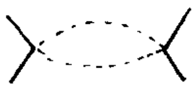
In fourth order (leading log) electron elastic scattering, two photon-exchange gives



$$\underset{s \to \infty}{\sim} \quad s \ln s \; K_{M^2}(q^2), \quad q^2 = t \tag{4}$$

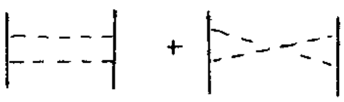
where $K_{M^2}(q^2)$ is a transverse momentum integral associated with the simplest transverse momentum loop diagram

$$K_{M^2}(q^2) = \int \frac{d^2 \underline{k}}{[\underline{k}^2 - M^2][(\underline{q}-\underline{k})^2 - M^2]} \equiv \tag{5}$$



$$\underset{M^2 \to 0}{\longrightarrow} \int \frac{d^2 \underline{k}}{\underline{k}^2 (\underline{q}-\underline{k})^2} = \infty \tag{6}$$

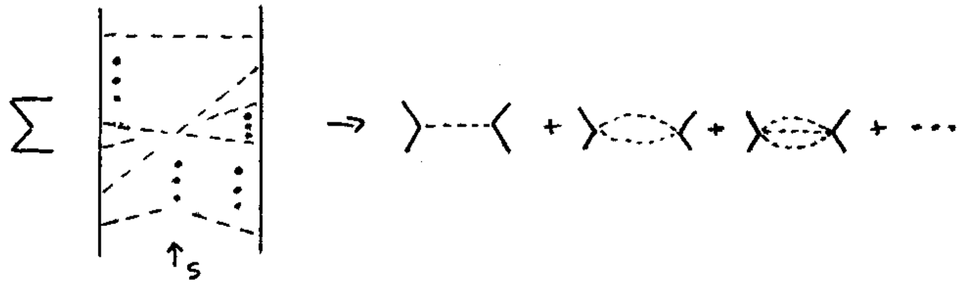
Adding the crossed box diagram, the real part of (4) cancels giving



$$\underset{s \to \infty}{\sim} \quad i \pi s \; K_{M^2}(q^2) \tag{7}$$

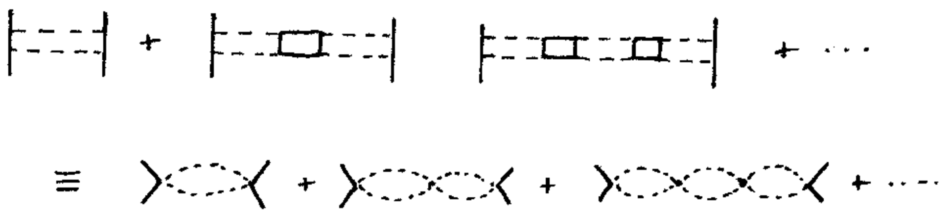
$$\underset{M^2 \to 0}{\longrightarrow} \quad i \infty \quad \text{the infinite Coulomb phase !} \tag{8}$$

(That this is indeed the Coulomb phase is confirmed by noting that the internal electron lines are placed on-shell by the high energy limit.) The exponentiation of the Coulomb phase comes from the usual eikonalization which in the Regge limit is written as a sum of transverse momentum diagrams[15].



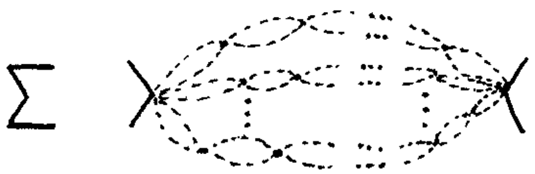
(9)

The higher order leading logs, however, come from a different class of diagrams[16),17]



(10)

The sum of these diagrams gives a j plane cut at $j = 1 + 11\pi\alpha^2/32$, thus violating the Froissart bound at infinite energy. This cut originates from $\underline{k}^2 \sim \infty$ and so occurs additively with the infra-red divergence from $\underline{k}^2 \sim 0$. It is possible[18),19] to eikonalize the leading log diagrams and restore the Froissart bound

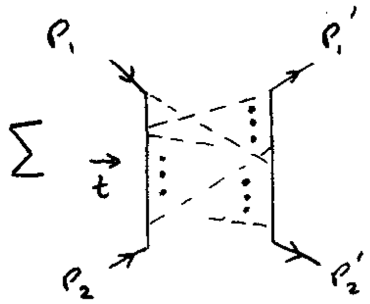


$$\rightarrow \sigma_T \sim [\ln s]^2 \tag{11}$$

Many people have argued[20),21] that adding additional diagrams required by unitarity will destroy this result. (In fact since QED is not well defined at large $\underline{k}^2$, where the high energy behaviour originates, there may not be a sensible result for the sum of all diagrams.) We note only that the infra-red divergences can also be seen as indicating the inadequacy of eikonalizing the leading log diagrams since they *do not multiply the* $[\ell n\ S]^2$ *behaviour as they should.*

We conclude that transverse momentum diagrams keep the Coulomb phase divergences explicitly and that analyzing them correctly is vital for high energy behaviour. (The real radiation divergences do not appear explicitly.)

But is the Coulomb phase simply a trivial phase factor in S matrix elements or is it related to the need to use coherent states to eliminate divergences ? Fortunately it is. In the t channel the sum of diagrams (9) contains only real radiation divergences which exponentiate as usual[12)-14)].



$$\to \exp\left[\sum_{i,j=1,2} e_i e_j' \, p_i \cdot p_j \int \frac{d^4q}{(q^2+M^2)(p_i\cdot q)(p_j'\cdot q)}\right] \tag{12}$$

$$\times \text{ finite part}$$

$$\equiv \exp\,[R] \times \text{finite part} \tag{13}$$

Using external coherent states will, of course, absorb the $\exp[R]$ factor. The point we wish to make is that analytic continuation of $[R]$ from the t channel to the s channel gives an imaginary part which is in fact[8),12)] the Coulomb phase

$$[R] \underset{t<0}{\to} [R] + [i\phi] \tag{14}$$

That is the Coulomb phase is the *s channel imaginary part* of the "non-local" factors needed to define *t channel states.* The high energy limit, in which t channel states are exchanged, retains through the imaginary part divergence, essential information on how to redefine physical t channel states as $M^2 \to 0$.

Qualitatively, that bare gauge invariant physical states can be created by operators of the form $\bar{\psi}(x)\,\exp\left[-\int_x^\infty dx'\,A(x')\right]$ is indirectly equivalent to the exponentiation of transverse momentum divergences in the sum of diagrams (9). Consequently by analyzing transverse momentum divergences in QCD we can hope to analyze the full infra-red problem indirectly provided we look at the implications of the divergences for *t channel states.*

PHASE STRUCTURE IN QCD

Moving on to QCD the first point of principle we encounter is - can we hope to add a mass to the gluons and remove it smoothly by simply factorizing off infra-red divergences. By now well known, lattice results[22),23)] tell us that if we use the Higgs mechanisn to make the gluons massive *there is no phase transition from the confining phase if the Higgs scalars are in the fundamental representation.* The relevant phase diagram is shown in Fig. 3. As illustrated we can, *on the lattice*, go analytically from a "Higgs theory" with massive gluons to a confining pure gauge theory. The underlying reason for this analyticity is that with fundamental representation scalars string-like states can form smoothly from local states[1),22),23)].

This analyticity is unlikely to be preserved in the continuum, assuming the continuum theories are defined by the continuum limits shown, because of the intervening line of phase transitions. (In fact the ultra-violet problems of the Higgs theory throw considerable doubt on its existence.) The underlying reasoning implies, however, that the analyticity should be preserved provided any form of ultra-violet cut-off is used. Of course, a major reason for introducing the lattice is that it provides the only known gauge invariant ultraviolet cut-off. Fortunately *in the (multi-) Regge limit a transverse momentum cut-off in the massive S matrix is a gauge invariant momentum cut-off.*

In practice this means that if we describe the S matrix by transverse momentum diagrams we should keep an ultra-violet cut-off in the corresponding integrals. We can then hope, perhaps, to write in analogy with (1) for the (multi-) Regge limit S matrix

$$S_{F\Lambda}(M^2) = Z_\Lambda(M^2)\, \tilde{S}_\Lambda(M^2)\, Z^+_\Lambda(M^2) \tag{15}$$

where Λ is the transverse momentum cut-off, $S_{F\Lambda}(M^2)$ is the S matrix of a massive Higgs theory, $Z_\Lambda(M^2)$ contains all the infra-red divergences and $\tilde{S}_\Lambda(0)$ is the S matrix of the confining unbroken gauge theory. We can then take $\Lambda \to \infty$ *after taking* $M^2 \to 0$.

At first sight (15) may seem inconsistent with the absence of massless particles in the confining theory since this would imply there can be no infra-red singularities to factorize off from the S matrix. In fact *(15) can only be consistent with confinement if* Z_Λ *contains infra-red singularities only from configurations of gluons which are uniformly soft.* Such configurations can be regarded as just a modification of the vacuum. Thus the Z factors of QCD and QED must be quite different. We shall elaborate on this further later.

If we calculate directly on a lattice we do not expect[22] to encounter massless particles at any stage during the continuation from Higgs to confinement shown in Fig. 3. However, it is well known that the fate of confined particles in a theory is very dependent on the regularization procedure used[24]. We are certainly not on a lattice but rather are starting from the perturbation expansion for the massive Fock space S matrix. Note that using fundamental representations of Higgs scalars implies that the gauge symmetry is completely broken by the Higgs mechanism, including the centre of the gauge group. Consequently the functional integral contains no topologically non-trivial contributions and the cut-off perturbation expansion should be Borel summable and *define the cut-off theory*. In this expansion the gluon mass is the (gauge) symmetry breaking parameter and setting it to zero at the same time restores gauge invariance and introduces infra-red singularities. Clearly we can hope to handle this only if we can write an equation of the form of (15). Actually since we are trying to go round the line of phase transitions in Fig. 3 and we are starting with a perturbation expansion valid near the line $g = 0$, we should try to shrink the line as much as possible by taking the lattice spacing (that is the inverse of our transverse momentum cut-off) as large as possible. This might imply that we can only hope to describe properly the $\underline{k} = 0$ part of Z_Λ, that is the "vacuum part" !

As we emphasized in the introduction we have very powerful machinery[1),2)] at our disposal to calculate $S_{F\Lambda}(M^2)$ and so if we can extract $\tilde{S}(0)$ through an equation of the form of (15) we are in a very good position. We shall begin the introduction of this machinery in the next section. First we note that in addition to allowing the (completely broken) cut-off Higgs theory to be defined by perturbation theory, fundamental scalars also prescribe the symmetry restoration. Decoupling the scalar representations one at a time increases the gauge symmetry through the sequence[25]

$$\text{"}SU(1)\text{"} \rightarrow SU(2) \rightarrow SU(3) \rightarrow SU(4) \rightarrow \cdots \tag{16}$$

until the full symmetry is restored. The expected smooth formation of states created by (fixed time) Wilson loop operators (the gauge theory analogue of closed string operators) from local states as this symmetry restoration takes place is illustrated in Fig. 4. The topological complexity[26] of "closed strings" as the centre of the gauge group increases in size is provided by adding local singlet vectors (and open strings) to a basic "non-orientated" SU(2) string.

We shall find that infra-red singularities from the SU(2) symmetry restoration produce the confinement vacuum and that the additional singlet vector Regge trajectories build up the complexity of the Pomeron as the gauge group enlarges. This gives a correspondence

$$\text{centre of group} \leftrightarrow \text{complexity of closed strings} \leftrightarrow \text{complexity of Pomeron spectrum} \tag{17}$$

MULTI-REGGE THEORY OF NON-ABELIAN HIGGS THEORIES

This is a very big subject which if it is reviewed at all should really be treated in depth. Bartels has already reviewed the subject at this meeting. As we noted in the introduction it is comprehensively reviewed in Ref. 1) and briefly reviewed for our present purpose in Ref. 3). Here we note only the really major points as follows.

1) Perturbative calculations[27)-31)] up to tenth order show the high energy behaviour of all amplitudes is described by reggeon diagrams[32)] involving reggeized gluons and quarks. (Reggeon diagrams are transverse momentum diagrams with $\ln s$ dependence incorporated through reggeon propagators which we give explicitly below.) At tenth order, in "order of magnitude"

$$\simeq 500 \text{ Feynman diagrams} \leftrightarrow 5 \text{ Reggeon diagrams} \tag{18}$$

This illustrates the power of reggeon diagrams to analyze the infra-red behaviour of sums of Feynman diagrams.

2) Multi-dispersion theory based multi-Regge theory[1),2)] can be used (in principle) to build perturbative calculations into a complete reggeon diagram description of an arbitrary multiparticle amplitude.

Bartels has described the use[30)] of the general formalism to build the weak coupling calculations, as far as possible, into an explicit solution of s channel unitarity. Relations between reggeon interaction vertices, important for the infra-red analysis we shall describe below, emerge directly as s channel unitarity constraints. Alternatively the same relations can be derived from Ward identities, as a reflection of gauge invariance, in the unbroken gauge theory[33)].

As will become clear we need to analyze the infra-red singularities of reggeon diagram amplitudes which are still far outside the explicit weak coupling calculations. We need therefore to rely heavily on the general formalism to extrapolate the weak coupling results as logically as possible. It is vital for this that the reggeon interactions in reggeon diagrams are simple, in particular they must not be singular in the transverse momentum variables. Several people[29)-31)] have given reggeon diagram rules for elastic amplitudes with signature conserving interactions which are singular. Fortunately, we have shown[1)] that the perturbative calculations can also be written in terms of reggeon diagrams in which

3) all quarks and gluons have reggeon propagators $[E-\Delta_i(\underline{k}^2)]^{-1}$ - E being the Mellin conjugate variable to $\ell n\ s$ - where the trajectory functions $\Delta_i(k^2)$ are non-singular when gluon masses go to zero and, of course, vanish (or = 1/2) at $\underline{k}^2 = M_i^2$ where M_i is the mass of the gluon (or quark). Also all interactions are regular functions of transverse momenta with signature conservation resulting from (non-sense) zeros required by the general Regge theory !

With this last form of the reggeon rules all infra-red singularities arise from gluon particle poles. That is the complete propagator for a reggeized gluon is

$$- - - - - - = \frac{1}{[E-\Delta_i(\underline{k}^2)][\Delta_i(\underline{k}^2)]} \quad \text{where} \quad \Delta_i(\underline{k}^2) \underset{\underline{k}^2 \to M_i^2}{\sim} \alpha'(\underline{k}^2 - M_i^2) \tag{19}$$

There is an integration $\int_{-i\infty}^{i\infty} dE \int d^2\underline{k}$ for each reggeon (gluon or quark), so that when $M_i^2 \to 0$ infra-red singularities clearly arise from the $[\Delta_i]^{-1}$ factors. The reggeon interactions always conserve "momentum" $\underline{k}$ and in general also conserve "energy" E. In multiparticle amplitudes, however, energy non-conserving vertices play an important role[1)]. We shall not give explicit expressions for reggeon interactions in what follows since we have not yet worked out the most compact form for the reggeon diagram rules in the form 3) above. The reader will therefore have to accept the cancellations of vertices described below and take on trust that we are confident how to handle multiparticle amplitudes knowing only that rules of the form 3) do exist.

INFRA-RED ANALYSIS OF SU(2) GAUGE SYMMETRY

We now begin to tackle the central problem, that is the mass-shell infra-red problem in a non-Abelian gauge theory. The triple gluon coupling allows soft gluons to cascade so that there is no "external line soft gluon rule" as there is for photons (except in leading logs[34),35)]). Weinberg[9)], many years ago, called the unravelling of such cascading divergences "a Herculean task" but suggested it could explain confinement. We shall try to analyze the problem in terms of reggeon diagrams believing that we are indirectly looking at the full problem for t channel states. Clearly we shall be pushing our knowledge of reggeon diagrams to the limit. Nevertheless we believe the attraction of the results should be sufficient to at least suggest to the reader that the problem can be unravelled with the techniques we are using, even if he is not convinced that the solution is as we shall present it.

We consider the first stage symmetry restoration in (16) and use the notation (19) for the massive $SU(2)$ gluons which are to be made massless. Because of the equivalence of the 2 and $\bar{2}$ representations of $SU(2)$ the Higgs potential will have a symmetry which gives the completely massive theory, from which we start, a global $SU(2)$ "colour" symmetry which we refer to as I. This is *the colour that will be confined*. The massive gluons form a triplet $(I=1)$ while $SU(2)$ quarks are doublets $(I=1/2)$.

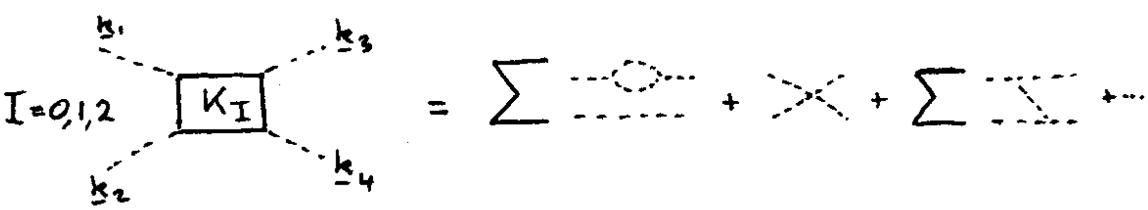
We begin our infra-red analysis by noting that the *major effect of cascading soft gluons is the exponentiation of reggeization*[27)-33)], or in our formulation reggeon self-energy effects,

$$= \left[E - \Delta(t) - \Delta_R(t)\right]^{-1} \qquad (20)$$

which in rapidity space is an exponentiation of infra-red divergences

$$s^{1-\Delta(t)-\Delta_R(t)} \underset{M^2\to 0}{\sim} s^{-\Delta_R(t)} \sim e^{-g^2 \ln s \ln t/M^2} \to 0 \qquad (21)$$

At first sight this exponentiation sends *all* reggeon diagrams to zero ! However, other reggeon interactions also produce divergences. To analyze this we look at interaction "kernels" defined for fixed numbers of gluons with definite t channel colour, e.g., for two gluons[28)-31)]

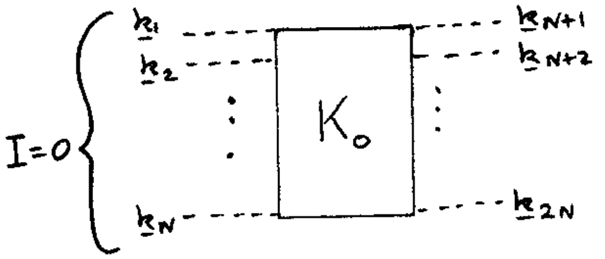


(22)

A central result from the existing weak coupling calculations is that because of the dominance of reggeization ( )

$$K_I \underset{M^2 \to 0}{\longrightarrow} \infty \qquad \text{for } I \neq 0 \tag{23}$$

This result seems to generalize[30),33)] to an arbitrary number of gluons. That is reggeization (or self-energy exponentiation) dominates all kernels except those with $I = 0$ so that only



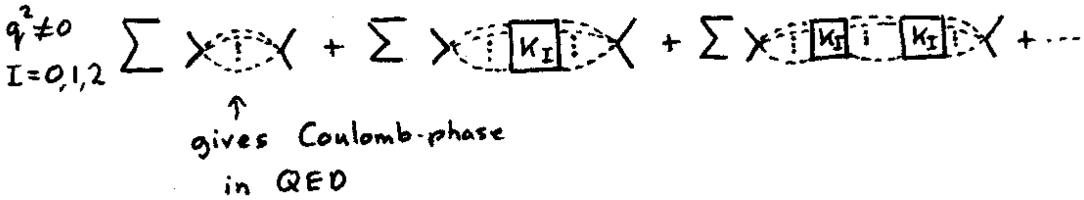
(24)

is finite and therefore (in the infra-red region) "scale invariant" as $M^2 \to 0$. The scale invariance property suggests that $K_0 \to \infty$ also if any subset of an $I = 0$ set of gluons goes to zero transverse momentum. Weak coupling results give that in this case reggeization of the remaining gluons with finite transverse momentum dominates[31),35)]. *The only infra-red limit in which* K_0 *is finite and reggeization does not exponentiate the infra-red divergence is all gluons taken uniformly to zero transverse momentum.*

$$\underline{k}_1 \sim \underline{k}_2 \sim \cdots\cdots \sim \underline{k}_{2N} \to 0 \tag{25}$$

This is just the configuration that we anticipate could give a divergence which could be absorbed into a Z factor as a modification of the vacuum. (The infra-red behaviour of the gluon kernels is pointing us towards a vacuum modification !) However, the most important point to emphasize at this stage is that the above description, in terms of reggeon self-energy effects dominating the effects of other interactions in every infra-red region except (25), is one that we believe can very safely be extrapolated from the weak coupling calculations to quite general higher order amplitudes.

We consider first the infra-red behaviour of elastic quark and gluon amplitudes which results from the behaviour of the gluon kernels. The configuration (25) does not appear in these amplitudes (for reasons that we shall explain shortly) and so all infra-red singularities are exponentiated by the sum



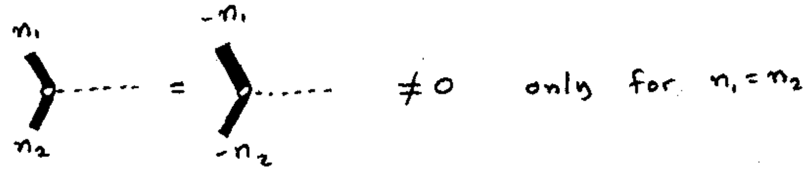
(26)

For $I \neq 0$, amplitudes are exponentiated to zero identically. For $I = 0$, reggeization exponentiates the infra-red divergences but leaves finite amplitudes from the transverse momentum regions where the K_0 kernels are finite. The finite amplitudes retain a singularity at $q^2 = 0$, as do, for similar reasons, the amplitudes that many authors[28)-30),36),37)] have calculated from sums of the form (26) with external colour zero states. Such amplitudes cannot therefore represent a confining theory since the q^2 singularities would necessarily imply the existence of massless particles in the theory.

In effect the internal infra-red exponentiation of reggeization has "screened" the Coulomb phase divergences occuring in the first term of (26), leaving only $I = 0$ amplitudes finite. This is a "suppression" of colour which is analogous to that in form factors[38)] and also to the more extensive phenomenon of "preconfinement"[39),40)]. If we stopped our infra-red analysis at this stage we might well conclude that simple $I = 0$ bound states of quarks are suitable external states. This is really the framework in which many calculations of the Pomeron have been done[28)-30),36),37)]. The Pomeron appears in leading logs as a fixed-cut violating the Froissart bound which shows little sign of improving as more complicated diagrams are summed[41)]. We believe that the bad behaviour of the pomeron is simply an indication, in addition to the zero q^2 singularities, that we have not carried the infra-red analysis far enough. In fact we must expose confinement if we are to hope to have the Pomeron treated correctly.

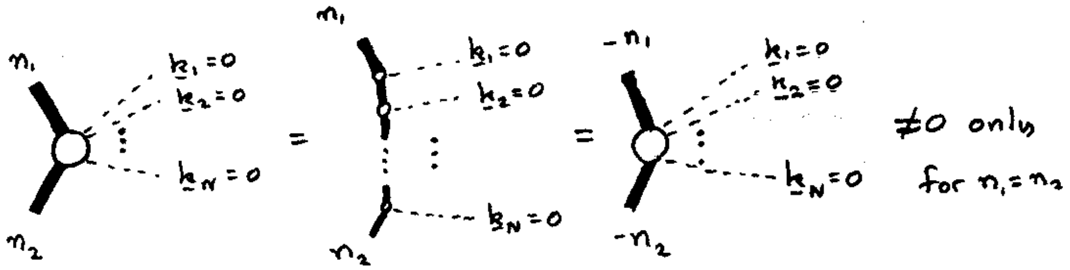
CONFINEMENT FROM MULTIPARTICLE AMPLITUDES

Confinement must involve some further infra-red phenomenon. This will be the occurrence of the configuration (25) accompanying bound state Regge poles in multiparticle amplitudes. To discuss this and also to explain why this configuration does not occur in elastic scattering we need some further properties of the theory. Firstly *a colour charge parity* C can be defined from the global SU(2) symmetry so that $I = 0$ combinations can have $C = \pm 1$. Also as a direct consequence of gauge invariance, a *massless gluon with zero transverse momentum couples only elastically* to physical states, with the coupling both *helicity conserving* and *helicity independent*. That is



(27)

Because all amplitudes are Regge-behaved asymptotically it follows (from a contour closing argument) that elastic couplings to arbitrary numbers of gluons can be computed through dispersion relations[1)] and so (27) extends to such couplings

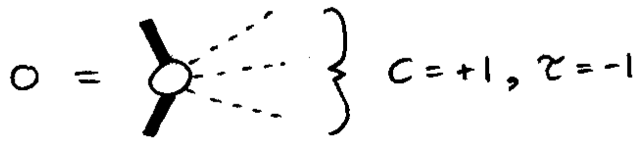


(28)

This result implies the known weak coupling result that gluons couple elastically only to t channel states with even parity (P). Since the signature of a Regge singularity is given (with appropriate conventions) by[42),43)]

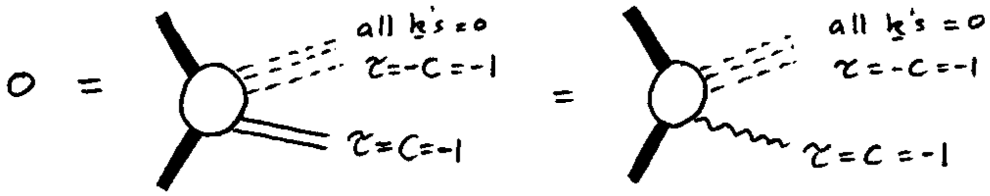
$$\tau = CP \tag{29}$$

in elastic amplitudes (at $q = 0$) signature and charge parity coincide (again a familiar fact perturbatively). Consequently combinations of gluons for which the Regge cut signature rules[2)] [$\tau = +(-1)$ for an even (odd) number of gluons] require that $\tau = -C$, cannot couple elastically at $q^2 = 0$. For SU(2), all $I = 0$ combinations have $C = +1$ and so it is only odd signature combinations of gluons that cannot couple, e.g.



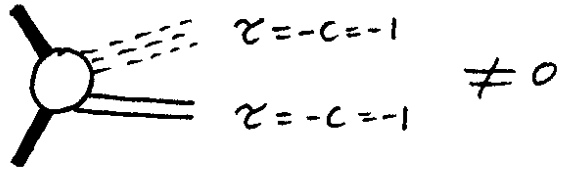
(30)

If massive gluons (〰) or quarks (——) are also involved in the coupling then unless they compensate for the $\tau \neq C$ mismatch of the massless gluons there will again be no coupling. Therefore



(31)

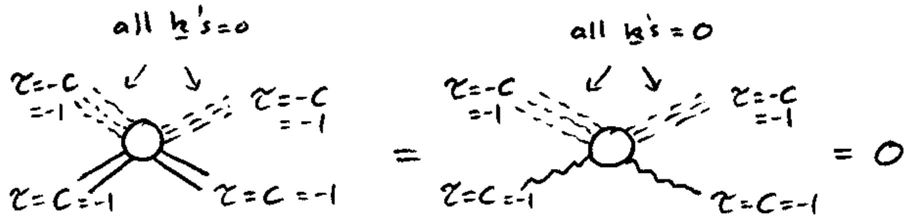
but



(32)

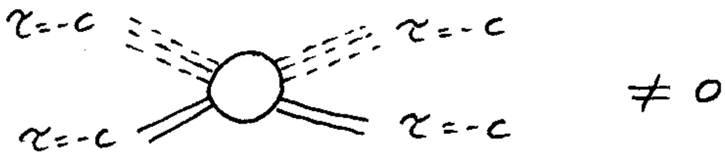
Note that the Regge cut rules for a pair of reggeized quarks fix $\tau = -1$, but they do not fix C (or equivalently P).

Since *local* reggeon interactions can similarly be computed by dispersion relations we expect (31) to generalize to



(33)

while (32) will generalize to



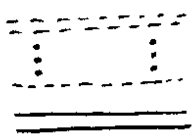
(34)

This structure for the reggeon interactions, together with the behaviour of the massless gluon kernels K_I, leads directly to the confinement spectrum with chiral symmetry breaking, in the following. Therefore, our derivation of this spectrum could perhaps be elevated to the status of "a proof" if, in particular, the above structure for the reggeon interactions could be confirmed by explicit weak coupling calculations, rather than justified by general arguments, as above.

We are now in a position to discuss when the configuration (25) will occur and more importantly, when the infra-red divergence will *not* be exponentiated. Potentially this configuration can occur in any $I = 0$ t channel, accompanying any $I = 0$ combination of massive gluons or quarks. However, if there is an interaction between such massive reggeons and the zero transverse momentum gluons, such as (34), then reggeon unitarity[1)] implies that summing iterations of the interaction will invert the associated E plane singularity. This is equivalent to exponentiation of the infra-red singularity of the phase space. Consequently, only when the gluons have $\tau = -C$ and the massive reggeons have $\tau = C$, as in (33), so that such interactions are absent, will the soft gluon divergence not exponentiate. The uniformly soft gluons will then behave like the soft photons of QED and simply accompany the massive reggeons, which must therefore carry all the transverse momentum. Note that no matter how many gluons are involved the divergence will remain the same [a single logarithmic divergence from the region (25)] . Hence summing over numbers of gluons simply produces a gauge invariant result which we summarize as *a gauge invariant uniformly soft flux of gluons with* $I = 0$ *and* $\tau = -C = -1$ *will accompany an* $I = 0$ *combination of massive quarks or gluons, with* $\tau = C$, *without screening its own singularity.*

Before considering how we factorize such divergences off from amplitudes let us consider the possible t channel reggeon states that are picked out by the divergence. Flavour quantum numbers play the essential role of distinguishing $I = 0$ channels in which reggeized quark singularities are leading. [Note that with $SU(2)$ flavour C becomes G parity.]

e.g.



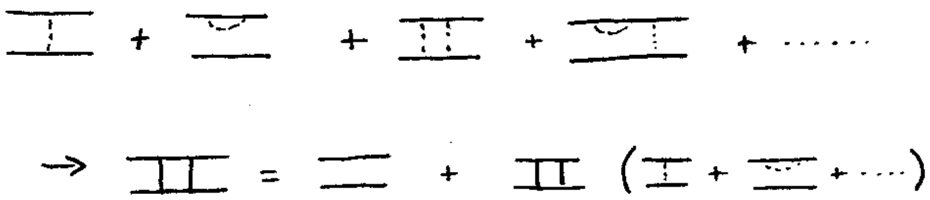
$\} \tau = -1, C = +1$

$\} \tau = -1, C = -1$

$\} \tau = +1, = -1 \Rightarrow P = -1$ (35)

$\Rightarrow$ a pseudoscalar pion ?

Can this $\tau = C$ quark/antiquark state, which is an odd signature Regge cut without the gluons, become an even signature Regge pole giving the pion ? The gluon flux will indeed change the signature but cannot directly enhance the Regge cut to become a pole. However, if asymptotic freedom is restored in the massless limit, we can take $\Lambda \to \infty$ after $M^2 \to 0$. The large $\underline{k}^2$ interactions give an integral equation in the *odd signature* quark/antiquark channel with an infra-red finite kernel analogous to that[28),29)] in the two gluon channel. For both $\tau = C$ and $\tau = -C$ we have



(36)

If asymptotic freedom can be used to justify using the running coupling constant in the kernel the solution of this equation is a sequence of Regge poles[44]. More investigation of this is required, but the following is what we anticipate.

1) $M^2 \to 0$, with the "physical" quark mass m_q kept fixed, adds the soft gluon flux to the quark/antiquark channel with $\tau = C$ and converts it to even signature. The leading singularity is the two reggeon cut (with a cut-off dependent trajectory) at

$$\alpha_{q\bar{q}}(t) = 2\,\alpha_q\left(\frac{t}{4}\right) - 1 \Rightarrow \alpha_{q\bar{q}}(4m_q^2) = 2\left(\tfrac{1}{2}\right) - 1 = 0 \tag{37}$$

2) $\Lambda \to \infty$ introduces large $\underline{k}^2$ interactions which replace this cut by a sequence of Regge poles. However, since the dominant large $\underline{k}$ interaction [that shown in (36)] vanishes at $t = 4m_q^2$ we expect the leading trajectory (the pion !) to still satisfy

$$\alpha_\pi(0) \underset{m_q^2\sim 0}{\simeq} \alpha_{q\bar{q}}(4m_q^2) = 0 \quad !! \tag{38}$$

giving (because of the *even signature*) a light mass particle, with $P = -1$ from (35). That is we have a physical *pseudoscalar pion* trajectory with a pion mass determined by the chiral symmetry breaking parameter - the light quark mass.

3) The trajectory slope will be set by the large $\underline{k}$ scale (the usual QCD scale μ). So for hadrons made out of N light quarks we similarly have

$$\alpha(0) \simeq N\,\alpha(m_q^2) - N + 1 = -\frac{N}{2} + 1 \tag{39}$$

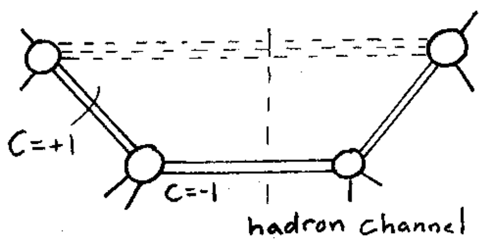
giving for the nucleon

$$\alpha_N(0) \doteq -\tfrac{1}{2} \quad ! \qquad (40)$$

The nucleon mass will clearly be determined by the scale of the slope, that is μ and not m_q !

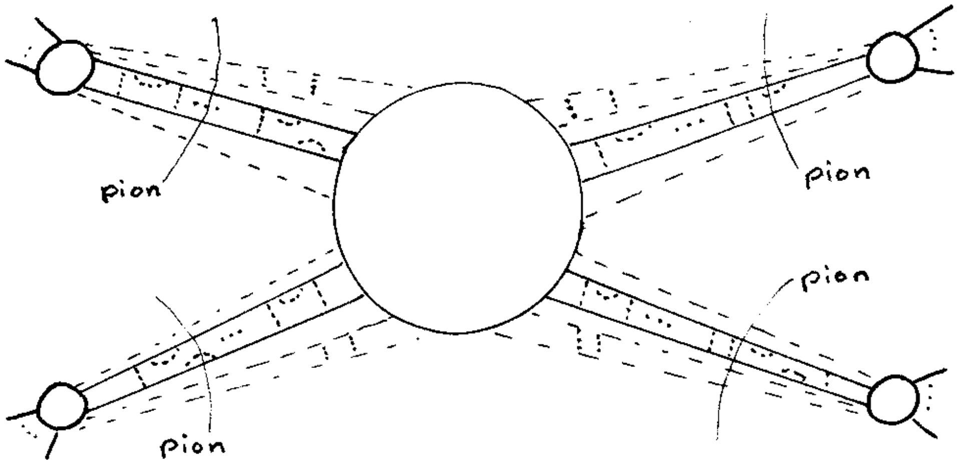
We conclude that, with asymptotic freedom, we have a bound state hadronic spectrum with chiral symmetry as $m_q^2 \to 0$ which, in effect, has the wrong signature (for Regge trajectories) until the infra-red divergent soft gluon flux is added - changing the signature for one parity only. Thus if we can properly understand this flux as a vacuum modification we have that the small mass pseudoscalar pion is a direct consequence of the vacuum and we have true Nambu-Goldstone chiral symmetry breaking. This we now discuss.

First we note that it follows from (31) and its generalization that the infra-red divergent configurations analogous to (35) do not occur in elastic quark and gluon amplitudes. Instead they can occur only in multiparticle amplitudes where the massive reggeons and the massless gluons originate at a different rapidity. The simplest example is probably



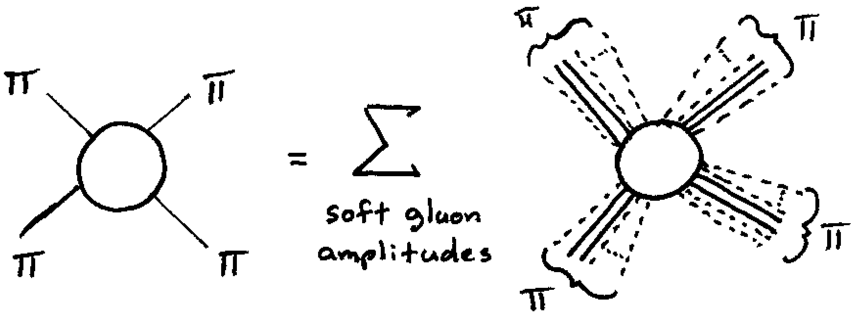
(41)

To extract hadronic scattering amplitudes we, in principle, go to much more complicated amplitudes. The four pion amplitude we extract from an amplitude of the form



(42)

In principle we continue to positive t in each of the pion channels and go to the particle pole on the trajectory generated by the sum of diagrams (36). The infra-red divergences of the gluon fluxes can be factorized off on to external residues which we can drop. [In fact we have recently realized that computing directly the Regge cut discontinuities in all reggeon diagrams in the foregoing would give finite answers for the infra-reddivergent exchanges we consider and zero for all other reggeon amplitudes in the infra-red limit. This suggests that the pole residues in (42) may in fact be finite but we shall not elaborate on this here.] Denoting the "perturbative" bound states generated by (36) as $=\!=\!=$ we see that the pion amplitude is described by a sum of amplitudes



(43)

where the sum is over the soft gluons needed to build up a gauge invariant flux with $\tau = -1$, $C = +1$ for each hadron. This is (superficially at least) the S matrix reggeon analogue of the vacuum expectation value for gluonic operators referred to in the introduction.

The individual amplitudes in (43) can be computed perturbatively. The role of the gluon flux is, of course, to effectively produce new couplings as would be anticipated from a modified vacuum. But is this the expected confinement vacuum[45] that we have arrived at ? We believe that it is for several reasons. Firstly the soft gluon flux can be regarded as due to the subtraction of all infra-red divergences resulting from two-dimensional (transverse momentum) configurations of gluons going uniformly to large (transverse) distance and giving a radial (equivalently a perimeter) divergence. This could well be the analogue, in our formalism, of requiring that a Wilson loop expectation value have no perimeter dependence when quarks are not present. Secondly as we shall discuss further below, the hadronic couplings that appear in (43) originate from (multiple) fixed-pole residues whose structure is indirectly related to the anomaly structure of chiral currents. This structure is essential to show that a consistent set of amplitudes having the form of (43) can be defined (and hence is centrally involved in breaking chiral symmetry as anticipated). Finally, as is discussed extensively in Ref. 1), there is a close parallel between the properties of the pomeron and those of Wilson loops in our formalism, as we would expect in a confining string-like solution of QCD. In fact we have emphasized that the pomeron is the real goal of our formalism and that our view of the above "derivation" of the hadronic spectrum is that it is needed to justify the treatment of the pomeron that we now outline.

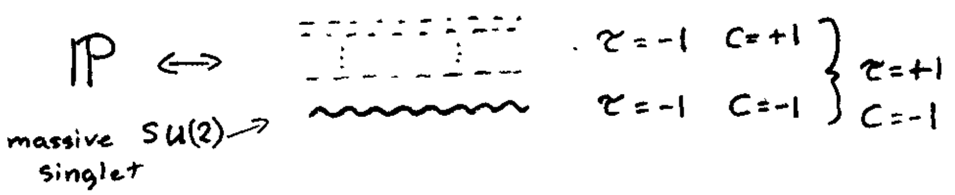
THE POMERON

Our first analysis[1] of the pomeron in QCD was simply based on looking for infra-red divergent multiparticle amplitudes as discussed above. However, we now believe that the analysis is best phrased in terms of looking for finite amplitudes of the form (43). We so far have only a very incomplete understanding of the following reasoning;

1) if the soft gluon flux is to couple, it must be to vertices which contain internal on mass-shell quarks, because of (28);
2) the soft gluon flux couples like an "axial-vector" in that it carries $\tau = -1$ but $C = +1$;
3) axial current anomalies are traceable[46] to higher order Landau singularities (in which internal quarks are on-shell) as chiral symmetry is restored;
4) higher order Landau singularities appear outside the straightforward Regge dispersion theory used above and might be expected to manifest themselves in higher order reggeon couplings.

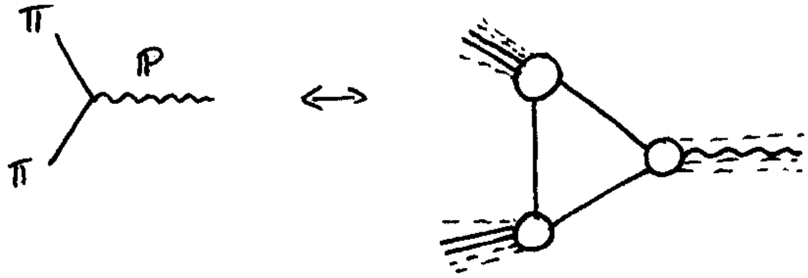
We think that these facts combine to give that hadronic couplings follow the pattern of chiral axial anomalies and that pomeron couplings are similarly associated with flavour singlet anomalies. This will then imply that only odd numbers of gluon fluxes can couple and immediately gives that SU(2) gauge theory contains no combination of gluons which can couple in (43) and give an even signature pomeron.

A pomeron with intercept close to one appears for the first time when we consider QCD but with the SU(3) gauge symmetry broken to SU(2) by one triplet of scalars. In a first approximation we then have



(44)

This pomeron can potentially couple to pions through a triangle Landua singularity involving internal quarks (such couplings could be the origin of the additive quark model !).

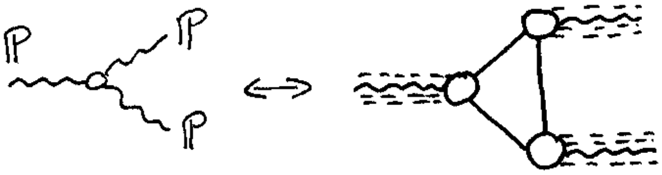


(45)

The most important property of (44) is that it gives a pomeron trajectory exchange degenerate with that on which the massive singlet vector lies. Consequently restoring the full gauge symmetry of QCD potentially places the intercept of the pomeron at one ! We will be unable to exploit this, however, if we subsequently have to remove the transverse momentum cut-off Λ. But to remove the cut-off before we restore the SU(3) symmetry, we must have asymptotic freedom of the Higgs theory. This requires[47),48)] that we have 16 flavours of quarks (if all quarks are colour triplets), that is we must *saturate QCD with quarks*.

We refer to pure QCD saturated with quarks as QCD_M, and QCD_M with the gauge symmetry broken to SU(2), by one triplet of scalars, as QCD_{MB}. We can construct QCD_{MB} directly by restoring the SU(2) gauge symmetry as in the last sections and subsequently taking $\Lambda \to \infty$. We will have a confined spectrum of hadrons, with, probably, a sequence of Regge poles for each hadronic channel [coming from the integral equation (36)]. The pomeron will have intercept less than one and so will have little impact on the hadronic spectrum.

We shall need that QCD_{MB}, as *a continuum theory*, can when the singlet vector mass goes to zero, give QCD_M. We believe this is very likely a priori, but has to be shown to be self-consistent. Since the pomeron intercept potentially goes to one we must consider the nature of the pomeron interactions in QCD_{MB} to determine what kind of "critical limit" is involved. In fact the "vacuum" soft gluon flux also plays a vital role in determining the pomeron interactions. The $\tau = +1$, $C = -1$ property of (46) implies that there must be a charge parity transition to give a triple pomeron vertex. This implies firstly we must have quark loops involved and secondly the vacuum is essential. There is no C non-conserving triple vertex for the singlet vector, even with quark loops. With the soft gluon flux, however, we can have an anomaly structure for the vertex as we discussed above. It could be that a simple quark loop will suffice

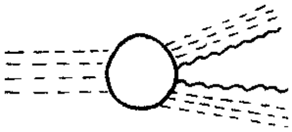
$$\gamma_0 = $$  (46)

or a more complicated non-planar structure involving the pion vertex (45) could be involved for reasons discussed in Ref. 1). However, the fundamental consequence of the need to use quark loops is that

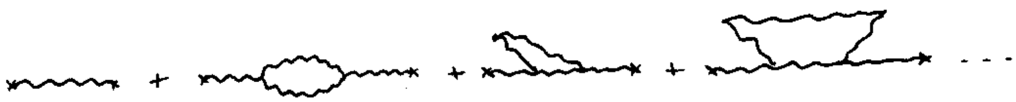
$$\gamma_0 \underset{m_q^2 \to 0}{\sim} \frac{1}{m_q^2} \tag{47}$$

which implies that *the chiral symmetry limit* $m_q^2 \to 0$ *is closely related to the limit of infinite pomeron coupling.* We shall return to this very shortly.

The soft gluon flux also allows the production of pairs of pomerons from the vacuum which, because of the elastic scattering constraint (28) for soft gluons, must have the gluon mass scale, that is

 $\sim M^2$ (48)

Given the existence of the vertices (46) and (48) and the exchange degeneracy of the pomeron we have all the features of the super-critical pomeron theory[49] which is obtained by making a vacuum shift in a theory which contains only pomerons but in which formally the intercept is above one. After the vacuum shift the theory is described by pomeron graphs

 (49)

There are many reasons[1], both general and particular, to believe that the QCD_{MB} pomeron is, in fact, described completely by the super critical pomeron formalism, and we are confident that we shall be able to demonstrate this by detailed infra-red analysis in the near future. Establishing this connection will, for the first time, *connect the pomeron intercept to a controllable physical parameter in a continuum theory* - namely the singlet vector mass in QCD_{MB}. (Such a connection can only be seen as remarkable by anybody who has studied the pomeron seriously, either in the abstract or in field theory calculations.) The following conclusions can then be drawn

1) Setting $M^2 \to 0$ gives $\alpha_{\mathbb{P}}(0) = 1$ and simultaneously decouples the odd signature vector trajectory from the physical states. This gives the necessary consistency for the limiting theory to be QCD_M.

2) The high energy behaviour of QCD_M is described by the reggeon field theory critical pomeron[50),51)] with its explicitly calculable scaling asymptotic diffraction peaks.

3) Repeating the arguments with N_F (number of flavours) < 16 but taking $M^2 \to 0$ before $\Lambda \to \infty$ shows[1)] the pomeron is then sub-critical and so asymptotically the cross-section falls.

4) Since $\alpha_{\mathbb{P}}(0) = 1$ the reggeon-pomeron cuts can shield the non-leading hadron trajectories [in the sequence generated by (36)] giving a simple spectrum, with chiral symmetry breaking, which is consistent with a string-like spectrum in the resonance region of Fig. 1.

5) Most dramatic of all, perhaps, (47) gives that the chiral limit $m_q^2 \to 0$ implies $r_0 \to \infty$, which within the pomeron formalism, is unlikely to exist unless $\alpha_{\mathbb{P}}(0) = 1$. In this *special case* $r_0 \to \infty$ gives the critical diffraction peak at finite energy. Clearly the scaling critical pomeron theory can give a sensible, one scale, strongly interacting theory - which the chiral limit must be.

The strongest conclusion from this last point would be that the existence of the chiral limit actually requires $\alpha_{\mathbb{P}}(0) = 1$, that is QCD must contain sixteen flavours (less[1)] if some quarks are not colour triplets). Before dismissing this conclusion as too strong we note firstly that as a *necessary* constraint for a unitary QCD with massless quarks, it has been argued[52)] that there must be at least ten flavours. Secondly sixteen is a particularly easy number of fermions to put on a lattice in four dimensions[53)]. To put less is already difficult and to ask whether the chiral limit exists with less is much more complicated. The sub-critical pomeron theory with less than 16 flavours is very non-perturbative with respect to the pomeron theory of QCD_{MB} and so it could be that the $m_q^2 \to 0$ limit does not involve $r_0 \to \infty$ in the sub-critical theory. This would, however, be very difficult to see.

The above analysis of QCD can be extended[1] to higher SU(N) groups by restoring the symmetry as in (16). From Fig. 4 it is clear that for each increase in N a new singlet "pomeron" trajectory is added to the theory. This leads to the conclusion that producing rising cross-sections in SU(4) or a higher group involves a more complicated critical phenomenon with more "pomeron" Regge trajectories. The method of calculation makes it clear that there is a correspondence between the structure of Wilson loops in transverse space and the pomeron spectrum as illustrated in Fig. 5, implying the correspondence between the pomeron and the centre of the group anticipated in (17). The need for asymptotic freedom after the first SU(2) symmetry is restored again leads to the conclusion that a rising cross-section requires close to the maximum number of fermions[48] allowed by asymptotic freedom of the unbroken gauge theory considered.

CONCLUSION

I shall finish by summarizing my results for the pomeron, so that they are clearly understood, and then compare some recent calculations of the critical pomeron with experiment. This allows me to be very optimistic about the future for the comparison of theory and experiment in diffraction scattering. I would like also to emphasize that the logical presentation of this talk is actually back to front from the historical development of my understanding. I began with the supercritical pomeron[49], looking for its significance in gauge theories. It has only gradually dawned on me that, if the development I have presented has even a grain of truth in it, getting the strong interaction pomeron right is essentially equivalent to understanding almost entirely what the strong interaction force is and is very likely to be the most fundamental way of approaching a large part of the basic conceptual problems involved.

SUMMARY OF THE EMERGING RESULTS FOR THE POMERON

SU(2) gauge theory

$$\sigma_T \xrightarrow[s \to \infty]{} 0 \qquad \forall \ N_F$$

SU(3) gauge theory

$$\sigma_T \xrightarrow[s \to \infty]{} 0 \qquad \forall \ N_F < 16$$

<u>$N_F = 16$</u> → RFT critical pomeron
→ many predictions

$$\sigma_T \underset{s \to \infty}{\longrightarrow} \beta_i \beta_j [\ln s]^{\eta} \qquad \eta \simeq \tfrac{1}{3} \quad (= -\gamma)$$

$$\frac{d\sigma}{dt} \longrightarrow \beta_i^2 \beta_j^2 [\ln s]^{2\eta} \, f_1 \left(t (\ln s)^{\nu} \right)$$

$$M^2 \frac{d^2\sigma}{dt\,dM^2} \underset{M^2, \frac{s}{M^2} \to \infty}{\longrightarrow} \beta_i^2 \beta_j^2 \beta_k^2 \left[\ln s/M^2 \right]^{\alpha_1} \left[\ln M^2 \right]^{\alpha_2} \times$$

$$\times f_2 \left(t \left(\ln \tfrac{s}{M^2}\right)^{\nu_1}, \; t \left(\ln M^2\right)^{\nu_2} \right)$$

$$\left[\frac{d\sigma}{dt}\right]_{pp} - \left[\frac{d\sigma}{dt}\right]_{p\bar{p}} \longrightarrow 0$$ for any particle and its anti-particle

η, ν, α_1, α_2, ν_1, ν_2 are calculable critical exponents, while f_1 and f_2 are calculable scaling functions.

<u>SU(4) gauge theory</u>

$$\sigma_T \underset{s \to \infty}{\longrightarrow} 0 \qquad N_F < 20$$

$$\sigma_T \underset{s \to \infty}{\longrightarrow} \infty \;, \quad \left[\frac{d\sigma}{dt}\right]_{pp} - \left[\frac{d\sigma}{dt}\right]_{p\bar{p}} \not\longrightarrow 0 \qquad N_F = 21$$

That is there is an odd signature component of the pomeron so that probably there is no simple factorization property for differential cross-sections. We anticipate that the relevant critical phenomenon involves real parts and imaginary parts scaling with different powers of log s [54].

<u>SU(N) gauge theory</u>

There is an increasing complexity of the pomeron with N (that is more Regge trajectories of both signatures). A rising cross-section requires close to the maximum number of fermions allowed by asymptotic freedom and in general if

$$\sigma_T \to \infty \quad \text{then} \quad \left[\frac{d\sigma}{dt}\right]_{pp} - \left[\frac{d\sigma}{dt}\right]_{\bar{p}p} \not\to 0$$

and there is no factorization.

So QCD with 16 flavours (much less if non-triplet quarks are present - as supersymmetry might suggest) *seems to be the essentially unique theory with factorizing, rising, total cross-sections, and equal particle → antiparticle differential cross-sections. All the scaling laws for the critical pomeron are predicted.*

Experimentally the situation is very encouraging. Recent comparison of the "$O(\varepsilon^2)$" critical pomeron with existing ISR results, by Bourrely and Dash[55], is shown in Figs 6 and 7. First, maximizing the agreement of the data with the scaling law is used to extract a simple parametrization of the residue functions (β_i,...). Then the calculated scaling function is compared with ISR data at the energy where the largest t range is available. The agreement is spectacular (for an approximate, but no parameter, fundamental calculation) and makes it very exciting to look for detailed verification of critical Pomeron behaviour at the $p\bar{p}$ colliders soon to come into being. Explicit predictions can be found in the recent paper of Baumel, Feingold and Moshe[56].

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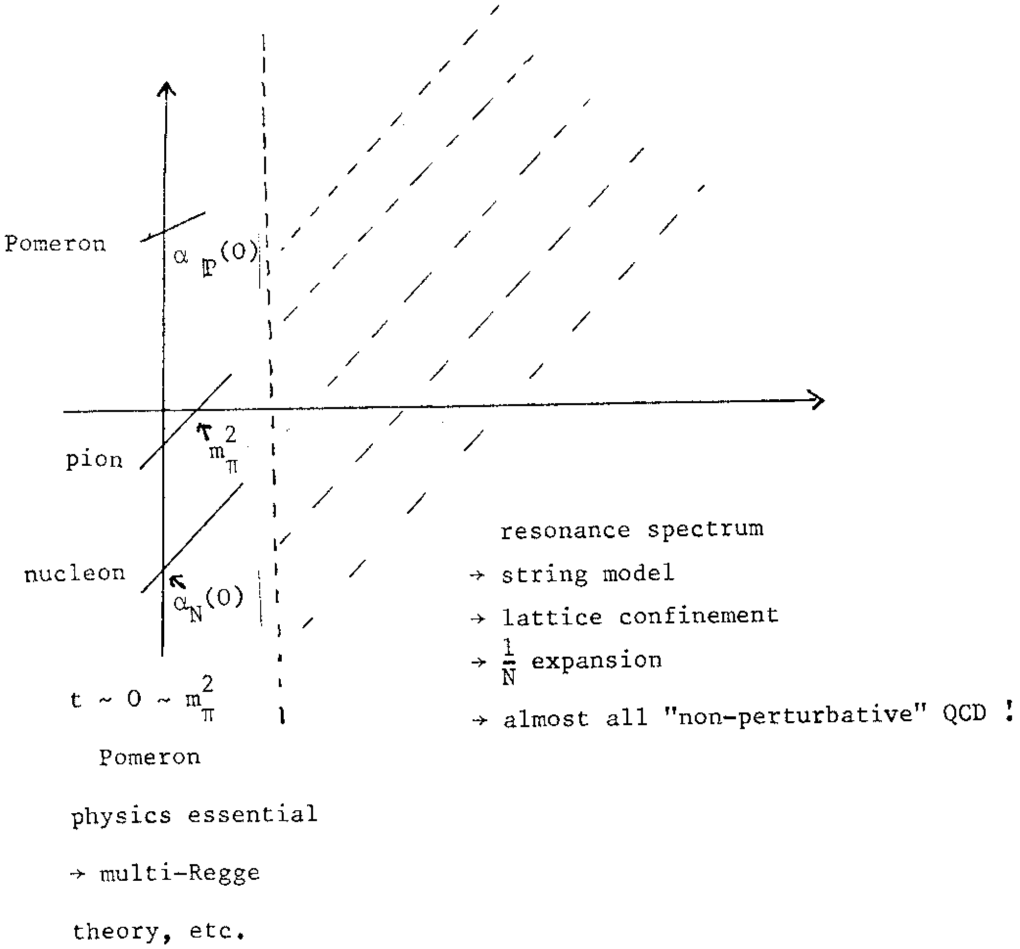

Fig. 1 : A "Chew-Frautschi plot" for QCD.

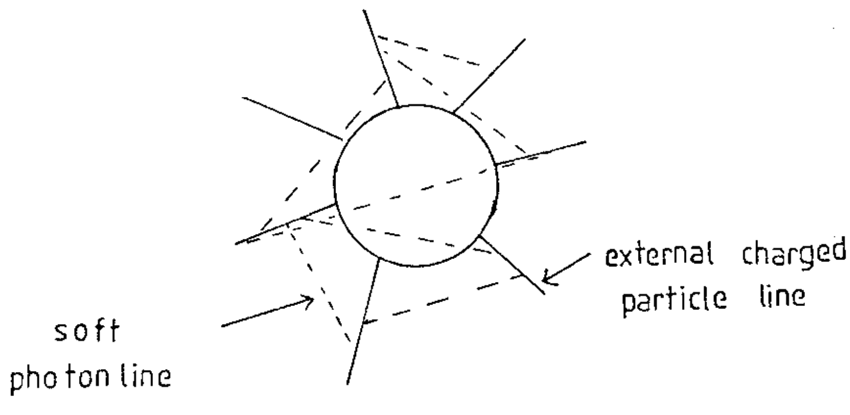

Fig. 2 : The external line rule for soft photons

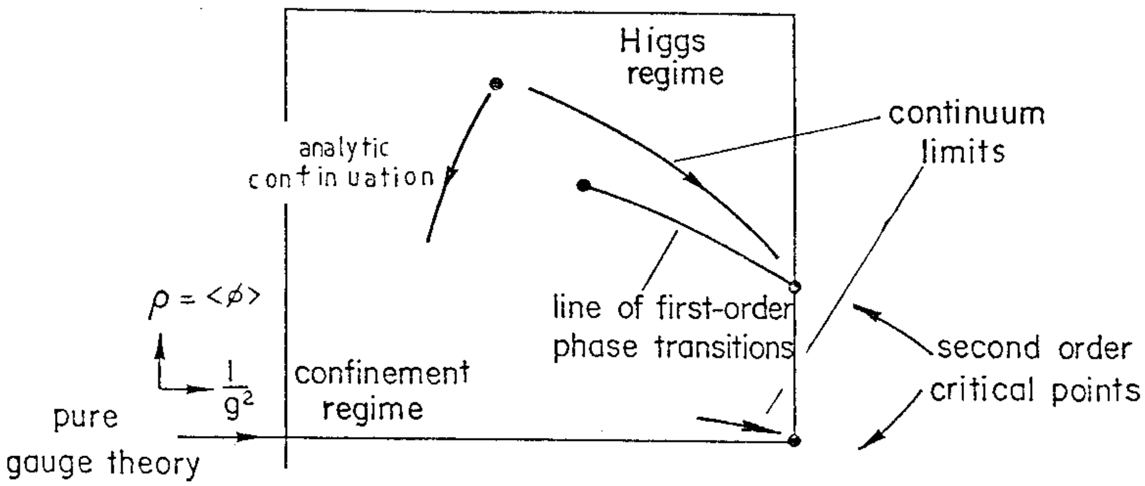

Fig. 3 : The phase diagram for fundamental representation scalars

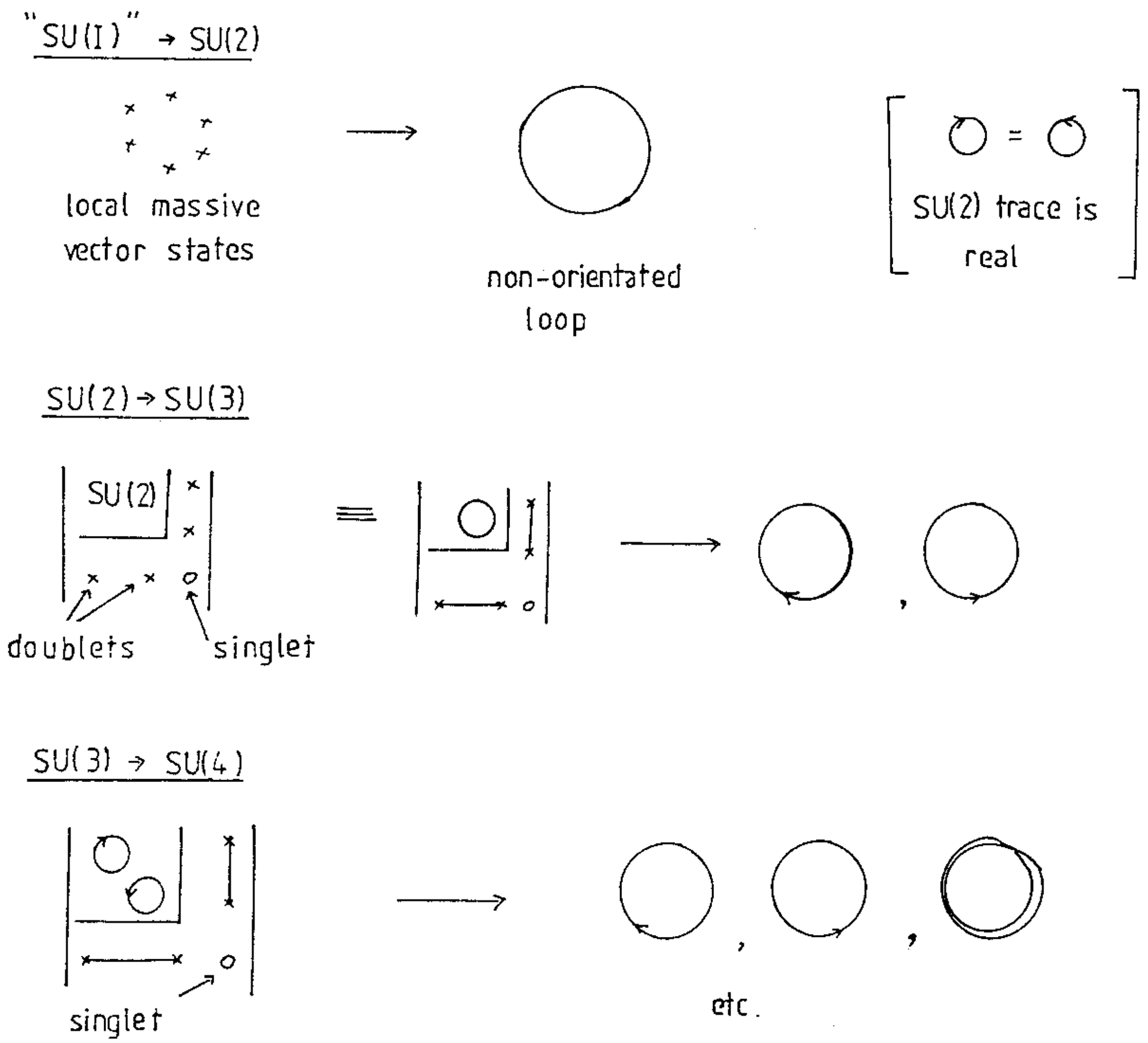

Fig. 4 : Formation of Wilson loops as the gauge symmetry is restored

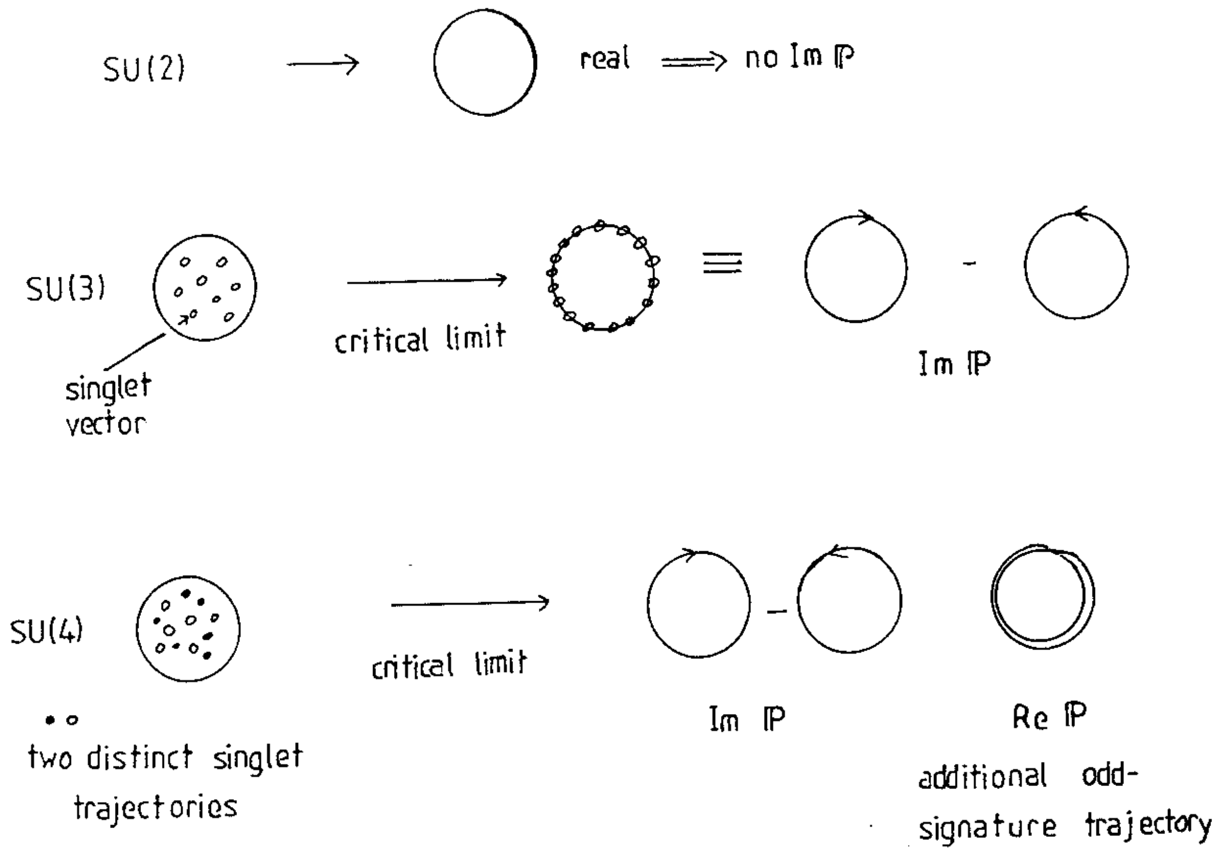

Fig. 5 : The correspondence between Wilson loops and the Pomeron

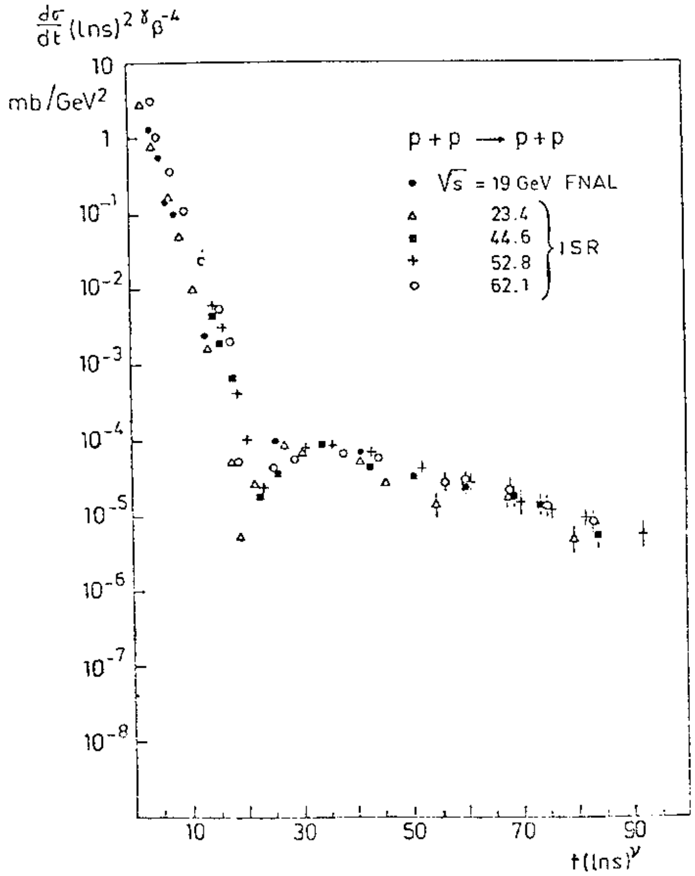

Fig. 6 : Experimental evidence for critical Pomeron scaling

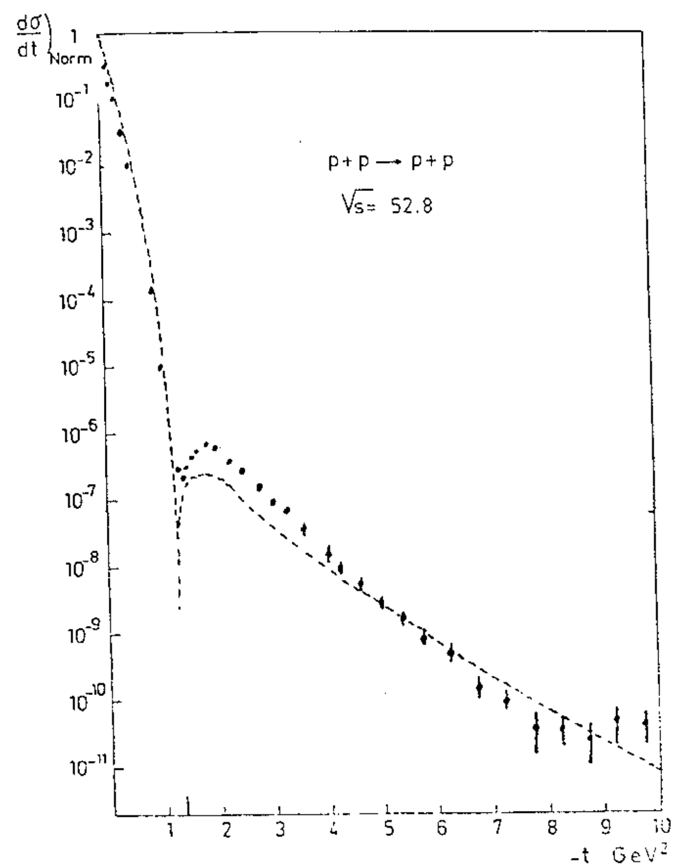

Fig. 7 : Comparison of the critical Pomeron scaling function with ISR data