

SOLVING QCD USING MULTI-REGGE THEORY*Alan. R. White[†]High Energy Physics Division
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9700 South Cass, Il 60439, USA.**Abstract**

This talk outlines the derivation of a high-energy, transverse momentum cut-off, solution of QCD in which the Regge pole and “single gluon” properties of the pomeron are directly related to the confinement and chiral symmetry breaking properties of the hadron spectrum. In first approximation, the pomeron is a single reggeized gluon plus a “wee parton” component that compensates for the color and particle properties of the gluon. This solution corresponds to a supercritical phase of Reggeon Field Theory.

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1 Introduction

To solve QCD at high-energy we must find both the hadronic states and the exchanged pomeron giving unitary scattering amplitudes. Experimentally (in first approximation) the pomeron appears to be a Regge pole at small Q^2 and ¹ a single gluon at larger Q^2 . Neither property is present in QCD perturbation theory. In this talk I will outline a high-energy, transverse momentum cut-off, “solution” of QCD in which these non-perturbative properties of the pomeron are directly related to the confinement and chiral symmetry breaking properties of hadrons.

The arguments have taken me a long time to assemble. They involve

- i) the techniques of multi-regge QCD calculations,
- ii) the dynamics of the massless quark U(1) anomaly,
- iii) Reggeon Field Theory phase-transition analysis.

The emphasis will be on ii) in this talk. A major outcome of the results² is a demonstration of how confinement and chiral symmetry breaking, normally understood as consequences of the vacuum, can instead be produced by a “wee-parton” distribution. This is a very non-trivial property that provides, I hope, a deeper basis for the parton model (and even³ the constituent quark model) in QCD !

The framework for my analysis is multi-regge theory. By using reggeon unitarity equations^{4,5}, well-known Regge limit QCD calculations^{6,7,8} can be extended to obtain multiparticle amplitudes involving multiple exchanges of reggeized gluons and quarks in a variety of channels. In particular we can study amplitudes, such as those of the form illustrated in Fig. 1, in which Regge pole bound states (e.g. the pion) and their scattering amplitudes (pomeron exchange) appear. Presently the simultaneous study of bound states and their scattering amplitudes is impossible in any other formalism.

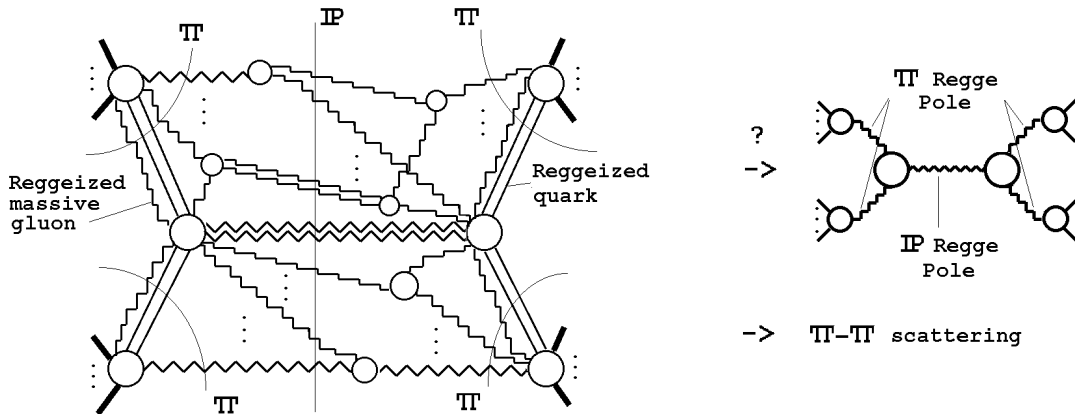


Fig. 1 The Anticipated Formation of Pion Scattering Amplitudes

In general, limits wrt all mass, gauge symmetry, and cut-off parameters are crucial. The most important feature, however, is the dynamical role played by new “reggeon helicity-flip” vertices that appear in the amplitudes we discuss. Hadron amplitudes are initially isolated via a (“volume”) infra-red divergence that appears when SU(3) gauge symmetry is partially broken to SU(2) and the limit of zero quark mass is also taken. The divergence is produced by quark loop helicity-flip vertices involving chirality violation (c.f. instanton interactions). The chirality violation survives the massless quark limit because of an infra-red effect closely related to the triangle anomaly⁹. The divergence produces (what we call) a “wee parton condensate” which is directly responsible, when the gauge symmetry is partially broken, for confinement and chiral symmetry breaking.

The pomeron, in first approximation, is a reggeized gluon in the wee parton condensate and so is obviously a Regge pole. Although we will not give any description of supercritical RFT⁵ in this talk we do find that all the essential features of this phase are present. We briefly discuss the restoration of SU(3) gauge symmetry. It is closely related with the critical behaviour of the pomeron¹⁰ and the associated disappearance of the supercritical condensate. We note that the large Q^2 of deep-inelastic scattering provides a finite volume constraint that can keep the theory (locally) in the supercritical phase as the full gauge symmetry is restored. A single gluon (in the background wee parton condensate) should then be a good approximation for the pomeron. Finally we discuss the (very special) circumstances under which our solution can be realized in QCD.

2 Multi-Regge Theory

This is an abstract formalism based on the existence of asymptotic analyticity domains for multiparticle amplitudes derived^{5,11} via “Axiomatic Field Theory” and “Axiomatic S-Matrix Theory”. All the assumptions made are expected to be valid in a completely massive spontaneously-broken gauge theory. Since we begin with massive reggeizing gluons, this is effectively the starting point for our analysis of QCD. We can very briefly list the key ingredients as follows.

i) Angular Variables

For an N-point amplitude we can introduce variables corresponding to any Toller diagram, i.e. any tree diagram, drawn as in Fig. 2, that involves only three-point vertices. The result is that we can write

$$M_N(P_1, \dots, P_N) \equiv M_N(t_1, \dots, t_{N-3}, g_1, \dots, g_{N-3})$$

where $t_j = Q_j^2$ and g_j is in the little group of Q_j , i.e. for $t_j > 0$, $g_j \in \text{SO}(3)$, and for $t_j < 0$, $g_j \in \text{SO}(2,1)$. A set of $3N - 10$ independent variables is obtained, $N-3$ t_i variables, $N-3$ z_j ($\equiv \cos \theta_j$) variables and $N-4$ u_{jk} ($\equiv e^{i(\mu_j - \nu_k)}$) variables.

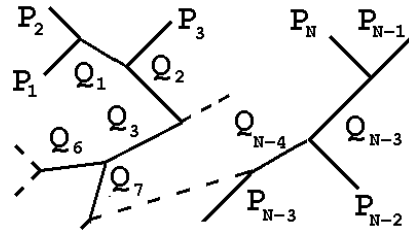


Fig. 2 A Tree Diagram with Three Point Vertices.

ii) *Multi-Regge Limits*

These limits are defined by $z_j \rightarrow \infty$, $\forall j$. We will also be interested in *Helicity-Pole Limits* in which some $u_{jk} \rightarrow \infty$ and some $z_j \rightarrow \infty$. In a helicity-pole limit a smaller number of invariants is taken large. In a “maximal” helicity-pole limit the maximal number of u_{jk} are taken large.

iii) *Partial-wave Expansions*

Using $f(g) = \sum_{J=0}^{\infty} \sum_{|n|, |n'| < J} D_{nn'}^J(g) a_{Jnn'}$, for a function $f(g)$ defined on $SO(3)$, leads to

$$M_N(\underline{t}, g) = \sum_{\underline{J}, \underline{n}, \underline{n}'} \prod_i D_{n_i n'_i}^{J_i}(g_i) a_{\underline{J}, \underline{n}, \underline{n}'}(\underline{t})$$

iv) *Asymptotic Dispersion Relations*

By dispersing in all z_j variables simultaneously, and applying the Bargman-Weil formula, we can write $M_N = \sum_C M_N^C + M^0$ where

$$M_N^C = \frac{1}{(2\pi i)^{N-3}} \int \frac{ds'_1 \dots ds'_{N-3} \Delta^C(\dots t_i, \dots u_{jk}, \dots s'_i)}{(s'_1 - s_1)(s'_2 - s_2) \dots (s'_{N-3} - s_{N-3})}$$

and \sum_C is over all sets of (N-3) Regge limit asymptotic cuts. M^0 is non-leading in the multi-regge limit. The resulting separation into spectral components, which can be described using a “hexagraph” notation^{2,5}, is crucial for the development of multiparticle complex angular momentum theory.

v) *Sommerfeld-Watson Representations of Spectral Components*

For each spectral component a multiple transformation of the partial-wave expansion can be performed, e.g.

$$M_4^C = \frac{1}{8} \sum_{N_1, N_2} \int \frac{dn_2 dn_1 dJ_1 u_2^{n_2} u_1^{n_1} d_{0, n_1}^{J_1}(z_1) d_{n_1, n_2}^{n_1 + N_2}(z_2) d_{n_2, 0}^{n_2 + N_3}(z_3)}{\sin \pi n_2 \sin \pi(n_1 - n_2) \sin \pi(J_1 - n_1)} a_{N_2 N_3}^C(J_1, n_1, n_2, \underline{t})$$

$$+ \sum_{\underline{J}, \underline{n}} d_{0, n_1}^{J_1}(z_1) u_1^{n_1} d_{n_1, n_2}^{J_2}(z_2) u_2^{n_2} d_{n_2, 0}^{J_3}(z_3) a_{\underline{J}, \underline{n}}(\underline{t})$$

These representations give the form of the asymptotic behaviour in both multi-Regge and helicity-pole limits. In particular, in a “maximal” helicity-pole limit, in which the maximal number of $u_{jk} \rightarrow \infty$, only a single (analytically-continued) partial-wave amplitude appears.

vi) *t-channel Unitarity in the J-plane*

After the hexagraph separation, multiparticle unitarity in every t -channel can be projected and continued to complex J as an equation for partial-wave amplitudes, i.e.

$$a_J^+ - a_J^- = i \int d\rho \sum_{\underline{N}} \int \frac{dn_1 dn_2}{\sin \pi(J - n_1 - n_2)} \int \frac{dn_3 dn_4}{\sin \pi(n_1 - n_3 - n_4)} \dots a_{J \underline{N} \underline{n}}^+ a_{J \underline{N} \underline{n}}^-$$

Regge poles at $n_i = \alpha_i$, together with the phase-space $\int d\rho$ and the “nonsense poles” at $J = n_1 + n_2 - 1, n_1 = n_3 + n_4 - 1, \dots$ generate multi-reggeon thresholds, i.e. Regge cuts.

vii) *Reggeon Unitarity*

In ANY J -plane of any partial-wave amplitude, the “threshold” discontinuity due to M Regge poles with trajectories $\underline{\alpha} = (\alpha_1, \alpha_2, \dots, \alpha_M)$ is given by the reggeon unitarity equation

$$\text{disc}_{J=\alpha_M(t)} a_{\underline{N}\underline{n}}(J) = \xi_M \int d\hat{\rho} a_{\underline{\alpha}}(J^+) a_{\underline{\alpha}}(J^-) \frac{\delta \left(J - 1 - \sum_{k=1}^M (\alpha_k - 1) \right)}{\sin \frac{\pi}{2} (\alpha_1 - \tau'_1) \dots \sin \frac{\pi}{2} (\alpha_M - \tau'_M)}$$

Writing $t_i = k_i^2$ (with $\int dt_1 dt_2 \lambda^{-1/2}(t, t_1, t_2) = 2 \int d^2 k_1 d^2 k_2 \delta^2(k - k_1 - k_2)$), $\int d\hat{\rho}$ can be written in terms of two dimensional “ k_\perp ” integrations, anticipating the reggeon diagram results of direct s -channel high-energy calculations^{6,7,8}. The generality of reggeon unitarity makes it particularly powerful when applied to the partial-wave amplitudes appearing in (maximal) helicity-pole limits.

3 Reggeon Diagrams in QCD

Leading-log Regge limit calculations of elastic and multi-regge production amplitudes in (spontaneously-broken) gauge theories show^{6,7,8} that both gluons and quarks “reggeize”, i.e. they lie on Regge trajectories. Non-leading log calculations are described by “reggeon diagrams” involving reggeized gluons and quarks. Reggeon unitarity implies that a complete set of reggeon diagrams arise from higher-order contributions.

Gluon reggeon diagrams involve a reggeon propagator for each reggeon state and also gluon particle poles e.g. the two-reggeon state

$$\begin{array}{c} \text{~~~~~} \\ \text{~~~~~} \end{array} \longleftrightarrow \int \frac{d^2 k_1}{(k_1^2 + M^2)} \frac{d^2 k_2}{(k_2^2 + M^2)} \frac{\delta^2(k'_1 + k'_2 - k_1 - k_2)}{J - 1 + \Delta(k_1^2) + \Delta(k_2^2)}$$

The BFKL equation⁶ corresponds to 2-reggeon unitarity, as illustrated in Fig. 3

$$\left[\text{diagram} \right] + \left[\text{diagram with } R_{22} \right] + \left[\text{diagram with } R_{22}^2 \right] + \dots$$

Fig. 3 Iteration of the 2-Reggeon State.

with $R_{22} = [(\underline{k}_1^2 + M^2)(\underline{k}_2'^2 + M^2) + (\underline{k}_2^2 + M^2)(\underline{k}_1'^2 + M^2)]/[(\underline{k}_1 - \underline{k}_1')^2 + M^2] + \dots$

We assume that two leading-order properties of the limit $M \rightarrow 0$ generalize to all orders. The first is that infra-red divergences exponentiate to zero all diagrams that do not carry zero color in the t -channel. The second property is that infra-red finiteness implies canonical scaling ($\sim Q^{-2}$) for color zero reggeon amplitudes when all transverse momenta are simultaneously scaled to zero (this requires $\alpha_s(Q^2) \not\rightarrow \infty$ when $Q^2 \rightarrow 0$).

4 Reggeon Diagrams for Helicity-Pole Limit Amplitudes

For our purposes, “maximal” helicity-pole limits of multiparticle amplitudes are the most interesting to study. Because the Sommerfeld-Watson representation involves only a single partial-wave amplitude, reggeon unitarity implies that reggeon diagrams again appear. Although we will not discuss it here, the physical significance of such diagrams is subtle². In particular, “physical” k_\perp planes in general contain lightlike momenta !

As an example, we introduce variables for the 8-pt amplitude corresponding to the tree diagram of Fig. 4. We consider the “helicity-flip” limit $z, u_1, u_2^{-1}, u_3, u_4^{-1} \rightarrow \infty$. The behavior of invariants is

$$\begin{aligned} P_1.P_2 &\sim u_1 u_2^{-1}, & P_1.P_3 &\sim u_1 z u_3, \\ P_2.P_4 &\sim u_2^{-1} u_4^{-1}, & P_1.Q_3 &\sim u_1 z, \\ Q_1.Q_3 &\sim z, & P_4.Q_1 &\sim z u_4^{-1} \dots \\ P_1.Q, P_2.Q, P_3.Q, P_4.Q & \text{finite} \end{aligned}$$

($u_1, u_2^{-1} \rightarrow \infty$ is a “helicity-flip” limit,
 $u_1, u_2 \rightarrow \infty$ is a “non-flip” limit.)

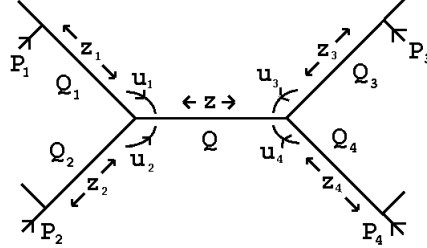


Fig. 4 Variables for the 8-pt Amplitude

Reggeon unitarity determines that the helicity-flip limit is described by reggeon

diagrams of the form shown in Fig. 5. The amplitudes \mathcal{A} contain all elastic scattering reggeon diagrams. The T^F are new “reggeon helicity-flip” vertices that play a crucial role in our QCD analysis. (These vertices do not appear in elastic scattering reggeon diagrams).

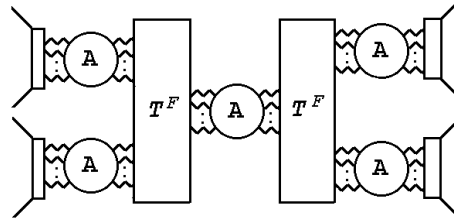


Fig. 5 Reggeon Diagrams for the 8-pt Amplitude

5 Reggeon Helicity-Flip Vertices

The T^F vertices are most simply isolated kinematically by considering a “non-planar” triple-regge limit which, for simplicity, we will define by introducing three distinct light-cone momenta. (This limit actually gives a sum of three T^F vertices of the kind discussed above², but in this talk we will not elaborate on this subtlety.) We use the tree diagram of Fig. 6(a) to define momenta and study the special kinematics

$$\begin{aligned}
P_1 &\rightarrow (p_1, p_1, 0, 0), & p_1 &\rightarrow \infty \\
P_2 &\rightarrow (p_2, 0, p_2, 0), & p_2 &\rightarrow \infty \\
P_3 &\rightarrow (p_3, 0, 0, p_3), & p_3 &\rightarrow \infty
\end{aligned}$$

$$\begin{aligned}
Q_1 &\rightarrow (0, 0, q_2, -q_3) \\
Q_2 &\rightarrow (0, -q_1, 0, q_3) \\
Q_3 &\rightarrow (0, q_1, -q_2, 0)
\end{aligned}$$

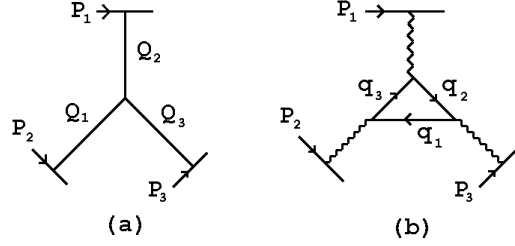


Fig. 6 (a) A Tree Diagram and (b) a quark loop coupling for three quark scattering.

Consider, first, three quarks scattering via gluon exchange with a quark loop coupling as in Fig. 6(b). The non-planar triple-regge limit gives

$$\rightarrow g^6 \frac{p_1 p_2 p_3}{t_1 t_2 t_3} \Gamma_{1+2+3+}(q_1, q_2, q_3) \longleftrightarrow g^3 \frac{p_1 p_2 p_3}{t_1 t_2 t_3} T^F(Q_1, Q_2, Q_3)$$

where $\gamma_{i+} = \gamma_0 + \gamma_i$ and $\Gamma_{\mu_1 \mu_2 \mu_3}$ is given by the quark triangle diagram i.e.

$$\Gamma_{\mu_1 \mu_2 \mu_3} = i \int \frac{d^4 k \text{Tr}\{\gamma_{\mu_1}(\not{q}_3 + \not{k} + m)\gamma_{\mu_2}(\not{q}_1 + \not{k} + m)\gamma_{\mu_3}(\not{q}_2 + \not{k} + m)\}}{[(q_1 + k)^2 - m^2][(q_2 + k)^2 - m^2][(q_3 + k)^2 - m^2]}$$

where m is the quark mass. We denote the $O(m^2)$ chirality-violating part of T^F ($\equiv g^3 \Gamma_{1+2+3+}$) by T^{F,m^2} and note that the limits $q_1, q_2, q_3 \sim Q \rightarrow 0$ and $m \rightarrow 0$ do not commute, i.e.

$$T^{F,m^2} \underset{Q \rightarrow 0}{\sim} Q i m^2 \int \frac{d^4 k}{[k^2 - m^2]^3} = R Q$$

where R is independent of m . This non-commutativity is an “infra-red anomaly” due to the triangle Landau singularity⁹.

T^F is one of a general set of quark loop reggeon interactions that have ultra-violet divergences. To maintain the reggeon Ward identities that ensure gauge invariance², we introduce Pauli-Villars fermions as a regularization. (Note that we take the regulator mass $m_\Lambda \rightarrow \infty$ after $m \rightarrow 0$. This implies that the initial theory with $m \neq 0$ is non-unitary for $k_\perp \gtrsim m_\Lambda$.) For the regulated vertex, $T^{\mathcal{F},m^2}$, we obtain (for $m \neq 0$)

$$T^{\mathcal{F},m^2}(Q) \sim T^{F,m^2} - T^{F,m_\Lambda^2} \underset{Q \rightarrow 0}{\sim} Q^2$$

However, since $T^{F,0} = 0$, we also have

$$T^{\mathcal{F},0}(Q) \sim -R Q$$

implying that imposing gauge invariance for $m \neq 0$ gives a slower vanishing as $Q \rightarrow 0$ when $m = 0$.

After color factors are included and all related diagrams summed, $T^{\mathcal{F},0}(Q)$ survives only in very special vertices coupling reggeon states with “anomalous color parity”. We define color parity (C_c) via the transformation $A_{ab}^i \rightarrow -A_{ba}^i$ for gluon color matrices and say that a reggeon state has anomalous color parity if the signature τ (i.e. whether the number of reggeons is even or odd) is not equal to the color parity. (Note that the reggeized gluon and the BFKL two reggeon state both have normal color parity.) We will be particularly interested in the “anomalous odderon” three-reggeon state with color factor $f_{ijk}A^iA^jA^k$ that has $\tau = -1$ but $C_c = +1$ (c.f. the winding-number current $K_\mu = \epsilon_{\mu\nu\gamma\delta}f_{ijk}A_\nu^iA_\gamma^jA_\delta^k$). $T^{\mathcal{F},0}(Q)$ appears in the triple coupling of three anomalous odderon states as in Fig. 7.

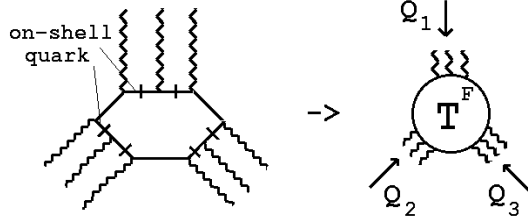





Fig. 7 An Anomalous Odderon Triple Coupling.

6 A Quark Mass Infra-Red Divergence

A vital consequence of the “anomalous” behavior of $T^{\mathcal{F},0}$ as $Q \rightarrow 0$ is that an additional infra-red divergence is produced (as $m \rightarrow 0$) in massless gluon reggeon diagrams. The divergence occurs in diagrams involving the T^F where $Q_1 \sim Q_2 \sim Q_3 \sim 0$ is part of the integration region. This requires that T^F be a disconnected component of a vertex coupling

distinct reggeon channels, as in Fig. 8. In this diagram an anomalous odderon reggeon state (\equiv ) is denoted by  while  denotes any normal reggeon state. Fig. 8 is of the general form illustrated in Fig. 5, except that we are allowing the vertices V_i to involve more complicated external states than a single scattering quark.

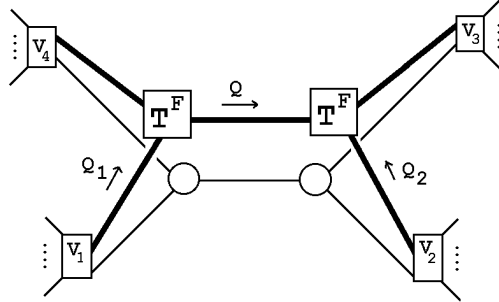


Fig. 8 A Divergent Reggeon Diagram

The canonical scaling of the anomalous odderon states gives the infra-red behaviour

$$\int \cdots \frac{d^2 Q_1 d^2 Q_2 d^2 Q}{Q_1^2 Q_2^2 Q^2 (Q - Q_1)^2 (Q - Q_2)^2} V_1(Q_1) V_2(Q_2) V_3(Q - Q_2) V_4(Q - Q_1) \\ \times T^{\mathcal{F}}(Q_1, Q) T^{\mathcal{F}}(Q, Q_2) \times [\text{regular vertices and reggeon propagators}]$$

for Fig. 8. Depending on the behaviour of the V_i , it is clear that a divergence may indeed occur when $Q \sim Q_1 \sim Q_2 \rightarrow 0$. In general we can show² that gauge invariance produces a cancelation involving similar divergences of diagrams related to that of Fig. 8 by

reggeon Ward identities for the reggeons within the anomalous odderon states. However, the divergence of Fig. 8 is preserved and a cancelation eliminated if we partially break the SU(3) gauge symmetry to SU(2). In this case, a divergence can occur in any diagram of the form of Fig. 8 in which --- is any SU(2) singlet combination of massless gluons with $C_c = -\tau = +1$ (i.e. a generalized SU(2) anomalous odderon) and --- is any normal reggeon state containing one or more SU(2) singlet massive reggeized gluons (or quarks).

A-priori reggeon Ward identities imply $V_i \sim Q_i$ when $Q_i \rightarrow 0$, $\forall i$. This would actually be sufficient to eliminate any divergence in Fig. 8. However, if we impose the “initial condition” that $V_1, V_2 \not\rightarrow 0$, the divergence is present in a general class of diagrams, including those having the general structure illustrated in Fig. 9. In this diagram there are $n + 3$ multi-reggeon states of the form --- . Imposing $V_1, V_2 \not\rightarrow 0$ and assuming that reggeon Ward identities are satisfied by the remaining vertices, i.e.

$$V_i(Q_i) \sim V(Q_i) = Q_i$$

$i \neq 1, 2$, gives that Fig. 9 has the infra-red behavior

$$\int \frac{d^2 Q}{Q^2} \left[\int \frac{d^2 Q}{Q^4} \right]^n [V(Q) T^{\mathcal{F}}(Q)]^n$$

giving (as $m \rightarrow 0$) an overall logarithmic divergence. In general, this divergence occurs in just those multi-reggeon diagrams which contain only SU(2) color zero states of the form --- coupled by regular and $T^{\mathcal{F},0}$ vertices, as in the examples we have discussed.

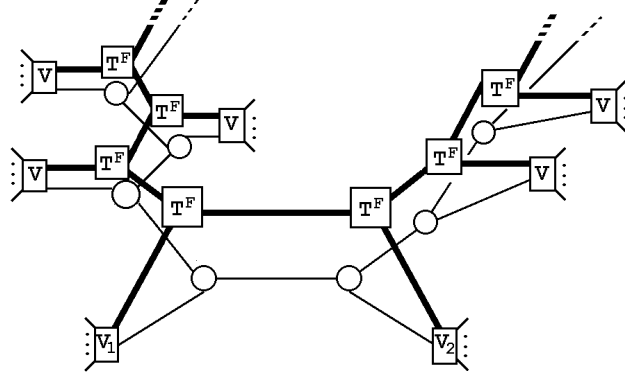


Fig. 9 A General Divergent Diagram

7 Confinement and a Parton Picture

We define physical amplitudes by extracting the coefficient of the logarithmic divergence. There is “confinement” in that a particular set of color-zero reggeon states is selected that contains no massless multigluon states and has the necessary completeness property to consistently define an S-Matrix. That is, if two or more such states scatter via QCD interactions, the final states contain only arbitrary numbers of the same set of states. Since $k_{\perp} = 0$ for the anomalous odderon component of each reggeon state, an “anomalous odderon condensate” is generated.

The form of physical amplitudes is illustrated in Fig. 10. In addition to the $k_{\perp} = 0$ (“wee-parton”) component, each physical reggeon state has a finite momentum “normal” parton component carrying the kinematic properties of interactions. We emphasize that the “scattering” of the $k_{\perp} = 0$ condensate is directly due to the infra-red quark triangle anomaly.

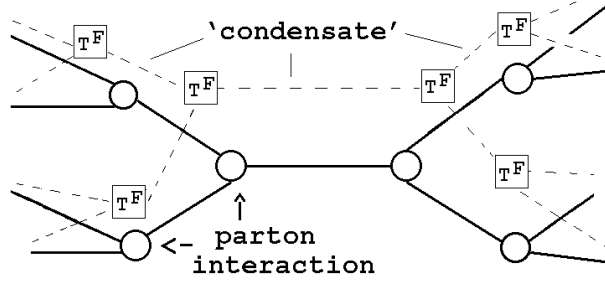


Fig. 10 A Physical Amplitude

The breaking of the gauge symmetry has produced physical states in which the “partons” are separated into a universal wee-parton component and a normal reggeon parton component which is distinct in each distinct physical state. However, the condensate has the important property that it switches the signature compared to that of the normal parton component. The following are a direct consequence.

- The “pomeron” has a reggeized gluon normal parton component, but is a Regge pole with $\tau = -C_c = +1$ and intercept $\neq 0$.
- There is a bound-state reggeon formed from two massive SU(2) doublet gluons, giving an exchange-degenerate partner to the pomeron. The SU(2) singlet massive gluon lies on this trajectory.
- There is chiral symmetry breaking -

studies of $\tau = -1$ quark-antiquark exchange¹² can be used to demonstrate the “reggeization cancelation” shown in Fig. 11(a). Because of this cancelation the sum of quark-antiquark diagrams shown in Fig. 11(b) generates a Regge pole with zero intercept and with $\tau = -1$, $C_c = +1$ and $P = -1$. In the condensate we obtain a Regge pole with $\tau = +1$, $P = -1$, giving the massless pion associated with chiral symmetry breaking. (We use the reggeon diagram notation that $\text{wavy line} = \text{a reggeized gluon}$ and $\text{solid line} = \text{a reggeized quark/antiquark}$).

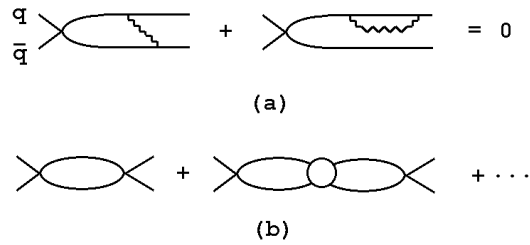


Fig. 11 (a) The Reggeization Cancelation and (b) the Sum of Quark-AntiQuark Diagrams.

We will not discuss Reggeon Field Theory, except to note that all the features of my supercritical RFT solution⁵ are present. (This solution was very controversial 20 years ago - although it was supported by Gribov !)

8 Restoration of SU(3) Gauge Symmetry

We make only a few brief comments on this, obviously important, subject. To discuss it in detail requires extensive use of Reggeon Field Theory and, in this talk, we are avoiding this. Because of complementarity¹³, restoring SU(3) symmetry (which involves decoupling a color triplet Higgs scalar field) should be straightforward if we impose a transverse momentum cut-off $k_\perp < \Lambda_\perp$. Restoring the symmetry involves removing the mass scale that produces the reggeon condensate and distinguishes normal (finite momentum) partons from wee (zero momentum) partons. If the (partially) broken theory can be mapped completely onto supercritical RFT then the condensate and the odd-signature partner for the pomeron will disappear simultaneously and the result will be the critical pomeron¹⁰. The wee-parton condensate will be replaced by a universal, small k_\perp , wee parton, critical phenomenon that merges smoothly with the large k_\perp normal (or constituent) parton component of physical states, just as originally envisaged by Feynman¹⁴. (Note that, because of the odd SU(3) color charge parity of the pomeron, the two-gluon BFKL pomeron will not contribute.)

Mapping partially-broken QCD onto supercritical RFT has further consequences. In particular, it implies that the Λ_\perp scale mixes with the symmetry breaking mass scale and becomes a “relevant parameter” for the critical behavior. It then follows that, after the symmetry breaking scale is removed, there will (for a general number of quark flavors) be a $\Lambda_{\perp c}$ such that $\Lambda_\perp > \Lambda_{\perp c}$ implies the pomeron is in the subcritical phase, while $\Lambda_\perp < \Lambda_{\perp c}$ will give the supercritical phase. This implies that the supercritical phase can be realized with the full gauge symmetry restored if Λ_\perp is taken small enough. However, $\alpha_{\mathbb{P}}(0)$ and the mass of the exchange degenerate, composite, “reggeized gluon” will be functions of Λ_\perp . We can also anticipate that in deep-inelastic diffraction large Q^2 will act as an additional (local) lower k_\perp cut-off and produce a “finite volume” effect that can keep the theory supercritical as the SU(3) symmetry is restored.

To remove Λ_\perp requires $\Lambda_{\perp c} = \infty$. As we briefly elaborate in the next Section, this requires a specific quark flavor content. It is interesting that, for any quark content, we can take $\Lambda_\perp \ll \Lambda_{\perp c}$, and go deep into the supercritical phase. We obtain a picture in which constituent quark hadrons interact via a massive composite “gluon” (and an exchange degenerate pomeron). Confinement and chiral symmetry breaking are realized via a simple, universal, wee parton component of physical states. This is remarkably close to the realization of the constituent quark model via light-cone quantization that has been advocated by light-cone enthusiasts³.

9 When is this Solution Realized in QCD ??

We have found a high-energy S-Matrix via a transverse momentum infra-red phenomenon involving massless gluons and quarks. At first sight, it would appear that this could not occur in QCD since non-perturbative effects should eliminate massless gluons for $k_\perp < \lambda_{QCD}$! Our solution requires that massless QCD remain weak-coupling at $k_\perp = 0$. This is

the case only if there are a sufficient number of massless quarks in the theory to give an infra-red fixed point for α_s .

With the maximum number of flavors allowed by asymptotic freedom, there is such an infra-red fixed point and we can also break SU(3) symmetry to SU(2) with an asymptotically-free scalar field. This can be used to show that $\Lambda_{\perp c} = \infty$. This, in turn, implies that critical pomeron scaling occurs for all k_{\perp} and allows a smooth match with perturbative QCD.

The above arguments suggest that if “single gluon” supercritical pomeron behavior is actually observed at HERA then new QCD physics, in the form of a new fermion sector, remains to be discovered above the (diffractive) Q^2 range presently covered. Everything is consistent if the electroweak scale is a QCD scale, i.e. the “Higgs sector” of the Standard Model, that is yet to be discovered, is composed¹⁵ of higher-color (sextet) quarks. A special definition of QCD is necessarily involved, but we will not discuss this here.

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