HIGHER-ORDER LIPATOV KERNELS AND THE QCD POMERON

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Abstract

Three closely related topics are covered. The derivation of $O(g^4)$ Lipatov kernels in pure glue QCD. The significance of quarks for the physical Pomeron in QCD. The possible inter-relation of Pomeron dynamics with Electroweak symmetry breaking.

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1. Introduction

I will spend most of my time in this talk describing my recent derivation[1] of higher-order Lipatov kernels, for pure-glue QCD, directly from $t$-channel unitarity. However, I also want to put this analysis in the context of my general study[2] of the physical, or soft, Pomeron in full QCD with quarks. I shall outline why quarks play such an important part in the emergence of confinement in the low transverse momentum region. I will also describe how study of the full unitarity properties of the Pomeron leads naturally to the introduction of a new higher-color quark sector which can provide a very attractive picture of dynamical electroweak symmetry breaking.

2. Higher-order Lipatov Kernels

Currently there is much excitement about the small-x behavior of structure functions. In particular it seems that at HERA the “Lipatov Pomeron”[3] may have been seen i.e.

$$F_2(x, q^2) \sim x^{1-\alpha} \sim x^{-\frac{1}{2}}$$  \hspace{1cm} (1)$$

The Lipatov equation was originally derived from extensive leading and next-to-leading log calculations in the Regge limit of (massive) Yang-Mills theories[4, 5] and is applied as an evolution equation for parton distributions at small-x i.e.

$$\frac{\partial}{\partial (\ln 1/x)} F(x, q^2) = \tilde{F}(x, q^2) + \frac{1}{(2\pi)^3} \int \frac{d^2 k}{k^4} K(k, q) F(x, k^2)$$  \hspace{1cm} (2)$$

$K(k, q)$ is given in terms of the $O(g^2)$ Lipatov kernel by

$$K(k, q) = K^{(2)}(k, -k, q, -q)$$  \hspace{1cm} (3)$$

where

$$\frac{2}{3g^2} K^{(2)}_{2,2}(k_1, k_2, k_3, k_4) = \sum_{1<\cdots<2} \left( (2\pi)^3 k_1^2 J_1(k_2^2) k_2^2 \left( k_3^2 \delta^2(k_2 - k_4) + k_4^2 \delta^2(k_2 - k_3) \right) \right)$$

$$- \frac{k_1^2 k_4^2 + k_2^2 k_3^2}{(k_1 - k_3)^2} - (k_1 + k_2)^2$$  \hspace{1cm} (4)$$

To obtain non-leading corrections to the $O(g^2)$ kernel, it appears that (very complicated) very non-leading log Regge limit calculations are required.

Back in the days when Pomeron Reggeon Field Theory (RFT) was studied intensely, it was understood that to satisfy multiparticle $t$-channel unitarity in the Regge limit a theory must be describable in terms of reggeon diagrams that satisfy reggeon unitarity[6]. For an even signature Pomeron, reggeon unitarity is very simple - as we now briefly review.

Introducing RFT variables, $E = 1 - \ell$ and $k^2 = -t$ , and writing $\alpha_R(t) = 1 - \Delta(k^2)$, a “partial-wave amplitude” $a^+(\ell, t)$ will satisfy reggeon unitarity if we can write

$$a^+(\ell, t) \equiv F \left( E, k^2 \right) = \sum_{n,m=1}^{\infty} F_{nm}(E, k^2),$$  \hspace{1cm} (5)$$
with
\[ F_{nm} (E, k^2) = \frac{1}{(2\pi)^{3m+3n}} \int \left[ \prod_{i,j} d^2k_i d^2k_j \delta^2 \left( k - \sum_{i=1}^{n} k_i \right) \delta^2 \left( k - \sum_{j=1}^{m} k'_j \right) \right] \] (6)

\[ g_n (k_1, \ldots, k_n) g_m (k'_1, \ldots, k'_m) G_{nm} (E, k_1, \ldots, k_n, k'_1, \ldots, k'_m), \]

where the \( G_{nm} \) are reggeon “Green’s functions” satisfying a “unitarity equation” -

\[ G_{nm} (E + i\epsilon, k) - G_{nm} (E - i\epsilon, k) = \frac{1}{(2\pi)^{3M}} \sum_M (-1)^{M-1} \int \prod_s d^2k_s \]
\[ \delta^2 \left( k - \sum_{s=1}^{M} k_s \right) \delta \left( E - \sum_{s=1}^{M} \Delta(k_s^2) \right) G_{nr} (E + i\epsilon, k) G_{rm} (E - i\epsilon, k). \] (7)

This equation is satisfied by writing Pomeron reggeon diagrams. For an \( n \)-Pomeron state we write the (“non-relativistic”) propagator

\[ \Gamma_n = \frac{1}{(E - \sum_{i=1}^{n} \Delta(k_i^2))} \] (8)

A minimal unitary set of diagrams is obtained by iterating this propagator with transverse momentum conserving \( 1 \rightarrow 2 \) and \( 2 \rightarrow 1 \) triple Pomeron interactions and then integrating over all transverse momenta. The arbitrariness arising from the possible existence of an infinite number of unknown higher-order Pomeron couplings is overcome by demonstrating that (when \( \alpha_P (0) = 1 \)) there is a fixed point where only the triple coupling is relevant. The result is the well-known Critical Pomeron scaling theory[7].

For an odd-signature Regge pole the corresponding reggeon diagrams are much more complicated - due to the presence of vector particle poles. It has been known for some time that the Lipatov equation can be written as a reggeon “Bethe-Salpeter equation” with the \( O(g^2) \) kernel as a singular reggeon interaction. (The singularity of the kernel is due to gluon poles). Nevertheless, the complete set of reggeon diagrams should determine (and be determined by) all Regge limit logs. Therefore, if we could construct such diagrams directly we could predict the results of arbitrary higher-order log calculations.

At first sight, such a construction looks impossible both because of singular reggeon interactions due to gluon poles and undetermined parameters in higher-order interactions. The first problem is resolved in general by the construction of reggeon loops via multiple discontinuities - this explicitly involves no singular interactions. The second problem is resolved for Yang-Mills theories by the imposition of Ward identity constraints directly on reggeon amplitudes. The result is a powerful technique for constructing higher-order reggeon diagrams, and therefore higher-order Lipatov kernels, without going through the very, very, complicated underlying Feynman diagram calculations.

We will illustrate the method by outlining a derivation of the Lipatov equation directly from reggeon diagrams, then present the \( O(g^4) \) kernels. We first discuss general properties of the diagrams.

A general reggeon amplitude is gauge-invariant and (in a suitable multi-Regge limit) can be embedded in a multigluon S-Matrix element as illustrated in Fig. 1. If \( k_\perp \rightarrow 0 \) for
one reggeon, then we obtain an amplitude $A_\nu$ for a gluon with zero four-momentum $k_\mu$ to couple to a physical multigluon state. This gluon amplitude satisfies the Ward identity

$$k_\mu A_\mu = 0$$  \hspace{1cm} (9)

which, provided there are no massless fermions in the theory, implies $A_\nu$ vanishes at $k_\mu = 0$. Consequently, the reggeon amplitude must also vanish when $k_\perp = 0$.

To use multiple discontinuity theory we introduce an $\alpha'$ for gluons from the outset - we will show that $\alpha' \to 0$ gives perturbative results. The gauge group is manifest via group factors in reggeon vertices and gauge invariance is imposed by the Ward identity constraint we have just discussed. In addition to reggeon propagators, diagrams will now contain a particle pole or “signature factor” $[\alpha' k^2]^{-1}$ for each uncut reggeon line. The triple reggeon vertex contains a “nonsense zero” i.e.

$$\Gamma_{12} \sim g \sqrt{\alpha'} [E - \alpha' k^2_1 - \alpha' k^2_2] \sim [k^2 - k^2_1 - k^2_2] \sim k^2$$  \hspace{1cm} (10)

The simplest reggeon diagrams are the set of “self-energy” diagrams. These give gluon reggeization. In the cut diagrams - all two reggeon propagators are cancelled by nonsense zeroes, leaving one zero per loop to provide the reggeization. The result is a series which sums up to

$$[E - \alpha' k^2 - g^2 k^2 J_1(k^2)]^{-1} \rightarrow [E - g^2 k^2 J_1(k^2)]^{-1}$$  \hspace{1cm} (11)

where

$$J_1(k^2) = \frac{1}{(2\pi)^3} \int d^2q \frac{1}{q^2(k - q)^2}$$  \hspace{1cm} (12)

which is the the perturbative reggeization result.

To construct multi-loop diagrams we have to proceed loop-by-loop, utilising multiple discontinuities. E.g. for the two-loop diagram shown in Fig. 2, we first construct the triple discontinuity of a three-reggeon vertex diagram. As illustrated in Fig. 3, nonsense zeroes at the vertices cancel the reggeon propagators $\Gamma_2$ and $\Gamma_3$ (and produce an external zero factor). The cuts also remove an internal signature factor. Only a transverse momentum integral remains which is then used as a vertex in the one-loop bubble diagram to obtain the two-loop diagram.

The two-loop diagrams in the color zero, even-signature, channel are shown in Fig. 4. Now only the three reggeon propagators are cancelled by nonsense zeroes. With external couplings $\alpha'g^2$, the sum of such diagrams gives

$$\frac{g^4}{(2\pi)^6} \int \frac{d^2k_1}{k_1^2} \frac{d^2k_2}{k_2^2} \delta^2(k - k_1 - k_2) \int \frac{d^2k_3}{k_3^2} \frac{d^2k_4}{k_4^2} \delta^2(k - k_3 - k_4) \frac{1}{(E - \alpha' k_1^2 - \alpha' k_2^2)(E - \alpha' k_3^2 - \alpha' k_4^2)} K^{(2)}_{2,2}(k_1, k_2, k_3, k_4)$$  \hspace{1cm} (13)
where \( K^{(2)}_{2,2}(k_1, k_2, k_3, k_4) \) is a sum of transverse momentum diagrams. The relative weight of the diagrams is uniquely determined by demanding both infra-red finiteness of the (integral) kernel and the Ward identity vanishing when \( k_i \to 0, \ i = 1, ..., 4 \). We thus obtain, without calculating a single Feynman diagram, the Lipatov kernel, to \( O(g^2) \), given in (4) above. The limit \( \alpha' \to 0 \) of \( (13) \) gives directly the sixth-order perturbative result (if we identify \( g \) with the gauge coupling). Iteration of the construction procedure to obtain the full Lipatov equation is straightforward.

The \( O(g^4) \) contribution to the (2-2) Lipatov kernel originates from three-loop reggeon diagrams of the form shown in Fig. 5 and the resulting kernel is a sum of the corresponding transverse momentum diagrams. If the vanishing at \( k_i \to 0, \ i = 1, ..., 4, \) is imposed together with infra-red finiteness, the relative weights of the distinct transverse momentum diagrams is uniquely determined and the result is

\[
g^{-4}K^{(4)}_{2,2}(k_1, k_2, k_3, k_4) = \sum_{i < j < k < l} \left( \frac{2}{3}(2\pi)^3 k_1^2 J_2(k_1^2) k_2^2 \left( k_3^2 \delta^2(k_2 - k_3) + k_4^2 \delta^2(k_2 - k_3) \right) \right) - \left( k_1^2 J_1(k_1^2) k_2^2 k_3^2 + k_1^2 k_2^2 J_1(k_2^2) k_3^2 + k_1^2 k_2^2 k_3^2 J_1(k_3^2) \right) / (k_1 - k_3)^2 + J_1((k_1 - k_3)^2) \left( k_2^2 k_3^2 + k_1^2 k_4^2 \right) + k_1^2 k_2^2 k_3^2 k_4^2 I(k_1, k_2, k_3, k_4) \right)
\]

with

\[
I = \frac{1}{(2\pi)^3} \int d^2q \frac{1}{q^2(q + k_1)^2(q - k_3)^2(q + k_1 - k_4)^2}, \quad J_2 = \frac{1}{(2\pi)^3} \int d^2q \frac{1}{(k - q)^2} J_1(q^2)
\]

Similiarly a new (2-4) kernel is generated which is also uniquely determined by the Ward identity and infra-red finiteness constraints

\[
K^{(4)}_{2,4}(k_1, ..., k_6) = \sum_{i < j < k < l < m} 2\pi^3 k_1^2 \left( \delta^2(k_2 - k_6) K^{(4)}_{1,3}(k_1, k_3, k_4, k_5) + \delta^2(k_2 - k_5) K^{(4)}_{1,3}(k_1, k_3, k_4, k_6) + \delta^2(k_2 - k_4) K^{(4)}_{1,3}(k_1, k_3, k_5, k_6) + \delta^2(k_2 - k_3) K^{(4)}_{1,3}(k_1, k_4, k_5, k_6) \right) - K^{(4)}_{2,4}(k_1, ..., k_6)
\]

where \( K^{(4)}_{1,3}(k_1, k_3, k_4, k_5) \) and \( K^{(4)}_{2,4}(k_1, ..., k_6) \) are given by

\[
K^{(4)}_{1,3}(k_1, k_3, k_4, k_5) = \frac{1}{(2\pi)^3} \int d^2k_1 d^2k_2 k_1^2 k_2^2 \delta^2(k - k_1 - k_2) K^{(4)}_{2,3}(k_1, k_2, k_3, k_4, k_5)
\]

with

\[
g^{-4}K^{(4)}_{2,3}(k_1, ..., k_5) = \sum_{i < j < k < l < m} \left( (k_1 + k_2)^2 - \left( k_1^2 (k_4 + k_5)^2 / (k_1 - k_3)^2 + k_1^2 (k_3 + k_5)^2 / (k_1 - k_4)^2 \right) + \frac{k_1^2 (k_3 + k_5)^2}{(k_1 - k_4)^2} \right) + \frac{1}{3} \left( k_1^2 k_2^2 k_3^2 (k_2 - k_5)^2 / (k_2 - k_4)^2 + k_1^2 k_2^2 k_4^2 (k_2 - k_3)^2 / (k_2 - k_4)^2 \right) \]

\[
+ 2 \left( k_1^2 k_2^2 k_3^2 k_4^2 / (k_1 - k_3)^2 (k_2 - k_5)^2 + k_1^2 k_2^2 k_3^2 k_4^2 / (k_1 - k_4)^2 (k_2 - k_5)^2 + k_1^2 k_2^2 k_3^2 k_4^2 / (k_1 - k_5)^2 (k_2 - k_4)^2 \right)
\]

4
and

\[
g^{-4} K_{2,4}^{(d)}(k_1, k_2, k_3, k_4, k_5, k_6)_c = \sum_{i<j=0} \left( (k_1 + k_2)^2 - \frac{k_1^2(k_4 + k_5 + k_6)^2}{(k_1 - k_3)^2} \right) - \frac{1}{4} \left( \frac{k_1^2k_2^2}{(k_2 - k_3)^2} \right)
\]

This technique explicitly combines the defining Ward identities of the theory with unitarity in the Regge limit. Since this should be sufficient to define the theory it appears that Yang-Mills reggeon theories can be constructed directly. That is, the Feynman diagram expansion can be bypassed completely. It certainly looks straightforward to construct even higher-order kernels. There are, of course, many issues to be studied already with the \(O(g^4)\) kernels, including the new Pomeron intercept, consequences for the anomalous dimensions of two and four gluon operators, consequences for deep-inelastic high-mass diffraction etc.\[10\].

3. Confinement and Quarks

The two defining properties we have used for pure glue reggeon theories are actually sufficient to show that such a theory can not produce a confining Pomeron RFT. While there is an extensive exponentiation of infra-red divergences in color non-zero channels, which removes all color non-zero amplitudes from the theory, all color zero reggeon amplitudes are infra-red finite. Also, the Ward identity constraint that reggeon amplitudes vanish when any \(k_\perp \rightarrow 0\) guarantees that phase-space integrals for multi-reggeon intermediate states are finite. Consequently color zero multigluon reggeon states necessarily survive and give singularities (but not divergences) in all amplitudes at zero transverse momentum. This clearly implies there is no confinement.
As we now discuss, this conclusion is avoided, when quarks are present, in a manner which does lead to confinement. To argue that the Ward Identity (9) requires that $A_\nu$ vanishes at $q = 0$ we simply differentiate

$$A_\nu + \frac{\partial A_\mu}{\partial q_\nu} q_\mu = 0 \implies A_\nu \to 0 \quad \text{unless} \quad \frac{\partial A_\mu}{\partial q_\nu} \to \infty$$

The Ward identity follows directly from gauge invariance and can not be violated, but massless quark infra-red divergences of $A_\nu$ - arising from the Regge limit - can prevent the conclusion that $A_\nu$ vanishes at $q = 0$. There are quark reggeon diagrams for which this is the case. In the color zero two gluon channel, quark loops of the form shown in Fig. 6 satisfy the Ward Identity by “$A_\mu \equiv A_-$” and “$q_\mu \equiv q_-$”, but produce a vertex (for quark mass $m$)

$$V_c(q, k, k') \sim \int d^2 p \frac{Tr \left[ (-\not p + \not q + m)\not k (-\not p - \not q + m)\not k' \right]}{((p + q)^2 + m^2)((p - q)^2 + m^2)}$$

which satisfies

$$V_c \sim \frac{8\pi}{m \sim 0} \frac{2(q \cdot k)(q \cdot k')}{q^2} - k \cdot k' \quad q \to \pm k, \pm k'$$

and so the Ward identity constraint is not satisfied. (This is consistent with (20) because $\frac{\partial V_c}{\partial q_0} \sim 0 (\equiv \frac{\partial V_c}{\partial m} \sim 0) \sim \frac{1}{q_0} \to \infty$ due to quark propagator infra-red divergences).

$V_c$ (as a partial-wave amplitude) is distinct from the Lipatov kernel, but “$V_c^2$” does contribute to this kernel. Consequently a new infra-red divergence is generated which produces the exponentation to zero of the two-gluon reggeon state. In fact, in SU(2) gauge theory all color zero reggeon states produce divergences which are similarly exponentiated[4], apart from a scaling divergence associated with the “anomalous Odderon” three gluon state. This divergence couples only via a “triangle quark loop anomaly” and so does not exponentiate. Instead it produces a reggeon (“winding number”) condensate which gives a confinement spectrum. In SU(3) gauge theory, a reggeized gluon in the background of this condensate gives a Regge pole Pomeron.

4. Pomeron Dynamics and Electroweak Symmetry Breaking

A Regge pole Pomeron gives $\sigma_T \to 0$ unless $\alpha_\not{F}(0) = 1$. Therefore, (confining) Pomeron RFT is, asymptotically, inconsistent with the validity of perturbative QCD at large $k_\perp$ unless the Pomeron is Critical !! My detailed analysis[2] of QCD Pomeron RFT suggests that $\alpha_\not{F}(0) = 1$ if, and only if, the number of flavors $N_f$ is a maximum. This requires that $N_f = 16$ or $N_f = 6$ for color triplet quarks and $N_f = 2$ for color sextet quarks. It is a remarkable coincidence that two flavors of color sextet quarks can provide a natural form of dynamical symmetry-breaking[11] for the electroweak interaction which meshes perfectly with the observed experimental features.

Add a massless doublet $(u_6, d_6)$ with the usual quark quantum numbers to the Standard Model (with no scalar Higgs sector). The axial $U(2) \otimes U(2)$ chiral symmetry breaks
spontaneously, producing four pseudoscalar Goldstone bosons ($\pi_6^+, \pi_6^-, \pi_6^0$ and $\eta_6$) which couple “longitudinally” to the sextet axial currents i.e.

$$<0|A_{\mu}^\tau|\pi_6^\tau(q)> \sim F_{\pi_6}q_{\mu}, \quad <0|a_{\mu}|\eta_6(q)> \sim F_{\eta_6}q_{\mu}$$

(22)

The $SU(2)$ gauge fields $B_{\mu}^\tau$ couple via

$$\mathcal{L}_I = g W^{\tau\mu}(V_{\mu}^\tau - A_{\mu}^\tau)$$

(23)

As a result $\pi_6^+, \pi_6^-$ and $\pi_6^0$ are “eaten” and provide masses for the $W^+, W^-$ and $Z^0$, with

$$M_W \sim g F_{\pi_6}$$

(24)

where $F_{\pi_6}$ is a QCD scale. The “Casimir Scaling” rule\[11\] gives

$$C_6 \alpha_s(F_{\pi_6}^2) \sim C_3 \alpha_s(F_{\pi}^2)$$

(25)

which is consistent with $F_{\pi_6} \sim 250$ GeV! Consequently, the electroweak scale is naturally explained as a second QCD scale. The restriction to a sextet flavor doublet, necessarily gives

$$\rho = (M_W^2/M_Z^2\cos^2\theta_W) = 1$$

(26)

as also required by experiment.

Finally we note that the sextet sector may also be deeply tied to the issue of Strong CP conservation and the origin of CP violation. The $\eta_6$ is a Goldstone boson associated with the $U(1)$ axial chiral symmetry and may be the axion. It is a conventional (Peccei-Quinn) axion\[12\] except that it can acquire an electroweak scale mass as a result of enhanced electroweak scale color instanton interactions\[13\].

Not only may the sextet sector be the explanation of Strong CP conservation in the triplet sector (via the $\eta_6$ axion), but it may also be responsible for CP violation at the weak scale. Because the sextet sector has no axion the QCD interactions at this scale will naturally be “Strong CP-violating”. The normal triplet quark hadrons will contain a small admixture of sextet quark states which could provide their CP violating interactions.

5. Conclusion

The QCD Pomeron may be the key to many of the remaining puzzles of the Standard Model.

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Fig. 1 A reggeon amplitude gives a gluon amplitude as $k_{\perp} \to 0$.

Fig. 2 A two-loop reggeon diagram.

Fig. 3 The reduction of a three-reggeon vertex diagram to a transverse momentum diagram.
Fig. 4 Two-loop even-signature reggeon diagrams.

Fig. 5 Three loop diagrams giving the $O(g^4)$ kernel.

Fig. 6 A one loop quark reggeon diagram.
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