QCD REVIEW

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This review is focused on QCD theoretical issues and their phenomenological relevance specially for LHC. It is incomplete and mostly neglects the phenomenology of long distance physics.

1 Introduction

In 1973, with the discovery of asymptotic freedom $^1$, QCD was at the frontier of particle physics studies. Now, in 2006, QCD is strongly installed (returning) at the center of particle physics researches. This, not only because of the abundance of data on jet emission at HERA and Tevatron, but especially for the necessity of preparing tools for discoveries at LHC. Indeed, events with large $E_T$, as in the decay of new massive particles, are accompanied by intense hadron emission. So the interpretation of a new elementary process requires an accurate description of short distance hadron physics, and this is the domain of perturbative QCD. Therefore calculations have been performed in recent years in order to produce accurate hard physics predictions for LHC.

The results of high order QCD studies go beyond their phenomenological importance, they are exposing various new features of Feynman graphs which may point to new general properties of quantum gauge theories. Moreover the possibility that QCD can be viewed as a solvable string theory is getting strength after the discovery of the AdS/CFT correspondence. In the last two or three years this fact encouraged studies of “solvable” string theories leading to phenomenological results in QCD.

QCD studies are developing in many directions and I can discuss only a limited number of points. Before describing NLO results with their relevance for phenomenology and for understanding new features of QCD in general, I describe new avenues of the phenomenological attempts to connect QCD with string theory and its impressive list of results.

2 Phenomenology of QCD as a solvable string theory?

String theory originated from the Veneziano model $^2$ for strong interactions even before the QCD era. The possible relation of QCD with string theory is based on various key observations. First, QCD Feynman graphs can be embedded $^3$ on the topological expansion of string theory (sphere, torus, etc.). Second, a consistent string theory must have more than four dimensions. Recently, Polyakov $^4$ and Maldacena $^5$ suggested a holographic correspondence between a four dimensional gauge theory and a string theory in higher dimensions. That is, the gauge theory observables correspond to the observables on a four dimensional boundary of the string theory in higher dimensions. This correspondence was explicitly shown by Maldacena for the large $N_c$ four dimensional supersymmetric Yang-Mill theory with four supercharges and the ten dimensional string theory with $\text{AdS}_5 \times S_5$ metric (AdS/CFT correspondence). How to go from SYM theory to QCD? This question is one of the major issues in recent string theory studies $^6$, one tries to move away from conformal symmetry and reduce the symmetries to approach QCD.

In the last two or three years this prob-
lem has been approached also directly form the QCD side (AdS/QCD correspondence). The question is: what are the characteristics of a string theory that holographically would reproduce in the four dimensional boundary the key proprieties of QCD. Results in this direction are quite abundant. There are many approaches with a common starting point: a string theory with a AdS modified metric. In five dimensions (the minimum allowed for string theory (by semiclassical approximation) and in the maximal helicity violation (MHV) configuration can be reduced to a small single line formula. The formula was generalized to the case with any number of gluons. The multi-gluon amplitude can be written (a, λ, p = colour, helicity, momentum) as

\[ M_n = g_n^{\alpha-2} \sum_{\text{perm}} \text{Tr}(t^{a_1} \cdots t^{a_n}) \times M(p_1, \lambda_1, \ldots, p_n, \lambda_n) \]

with \( \text{Tr}(\cdots) \) the colour order factors and \( M(\lambda_1, \lambda_2 \cdots, \lambda_n) \) the helicity and momentum ordered amplitudes. For the MHV configuration (all positive helicities \( \lambda_k = +1 \) but two \( \lambda_i = \lambda_j = -1 \)) one has the single line formula

\[ M_{\text{MHV}}^{\text{tree}}(1 \cdots n) = \frac{(p_1 p_2)^4}{(p_1 p_2) \cdots (p_n p_1)} \]

with \( \langle p_i p_j \rangle = \sqrt{2p_i p_j} e^{i \phi_{ij}} \) and \( \phi_{ij} \) a phase proportional to the relative transverse momentum. This simple formula, although valid only for the MHV tree amplitudes, indicates that Feynman diagrams have structures which are surprisingly simple. One is then induced to find the origin of this simplicity. An attempt in this direction was made by Witten who suggested that perturbative gauge theories could be viewed as a string theory in the twistor space. The ground for this proposal is simple. Using the 2-spinors (massless Dirac equation solutions) one has

\[ \langle pp' \rangle = \bar{u}_-(p) u_+(p'), \quad p_\mu = \frac{1}{2} \bar{u}_\pm(p) \gamma_\mu u_\pm(p). \]
This shows that, while the momentum conservation $\delta^4(\sum p_i)$ depends on both $u_-$ and $u_+$ (together with the corresponding conjugate spinors $u_+, \bar{u}_-$), the amplitude $M_{\text{tree}}$ depends only on $u_+$ (together with its conjugate one $\bar{u}_-$), see (1). It is then natural to perform a “Fourier transform” in $u_-$ and in the “Fourier” space, the twistor space, one finds that the amplitude with momentum conservation is described by a line in which the gluons are attached in an ordered way. The problem is then, how to describe tree amplitudes which are not in MHV configurations and how to go beyond tree level. The results in these two directions have been reported by S.Moch $^{16}$ and Z.Berm $^{17}$.

The prescription on how to go beyond MHV in the tree level have been simplified and generalized $^{18,19}$. It consists of sewing together MHV amplitudes to construct the non-MHV ones. To do this one needs to use off-shell MHV amplitudes which are obtained by analytical continuation in the complex momentum space. Tree amplitudes have been constructed for various processes: massless fermions $^{20}$, Higgs boson $^{21}$, EW vector bosons $^{22}$. Results at one loop are also obtained $^{19}$. The analytical expression for the one-loop $gg \rightarrow gggg$ amplitude in all helicity configurations will soon be obtained while its numerical evaluation is available $^{23}$.

### 4 High order QCD results

Tree level amplitudes involving many QCD partons and next-to-leading order (NLO) amplitudes are very important for LHC studies. Crucial studies are the search for the Higgs meson and for all signals which could indicate a way to complete the Standard model. The relevant events involve large mass particles accompanied by intense hadron emission. Thus it is crucial, for the interpretation of the events, to have accurate QCD predictions. This requires calculations of a large variety of many parton NLO matrix elements and distributions needed for: i) direct studies of new physics signals; ii) merging exact matrix elements with QCD resummation results for jet shape distributions $^{24}$ and with Monte Carlo simulations of QCD jet emission $^{25}$.

There are various numerical programs to compute many-leg amplitudes at tree level $^{26}$. The most common $t\bar{t}$ decay ($t\bar{t} \rightarrow b\bar{b}W^+W^- \rightarrow bbq\bar{q}$) involves 6 final state jets. Various processes have been computed at NLO and some presented at this conference. The LHC “priority” wishlist $^{27}$ for NLO calculations includes the following processes:

| $V \in \{Z,W,\gamma\}$ | background to |
|----------------------|--------------|
| $pp \rightarrow VV\_\text{jet}$ | $t\bar{t}H$ |
| $pp \rightarrow t\bar{t}bb$ | $t\bar{t}H$ |
| $pp \rightarrow t\bar{t} + 2\text{jets}$ | $t\bar{t}H$ |
| $pp \rightarrow VV\_\text{bb}$ | VBF $\rightarrow H \rightarrow VV$, $t\bar{t}H$ |
| $pp \rightarrow VV + 2\text{jets}$ | VBF $\rightarrow H \rightarrow VV$ |
| $pp \rightarrow V + 3\text{jets}$ | new physics signatures |
| $pp \rightarrow VV\_\text{V}$ | SUSY trilepton |

The procedure for NLO calculation of an inclusive distribution is in principle “simple”. Typically, one starts from the tree level amplitude for the process under consideration with $n$-partons (Born approximation). Then one computes the $(n+1)$-tree amplitude (real contribution) and the one-loop correction to the Born amplitude (virtual correction). Finally one puts together both contributions to construct the distribution and check that collinear and infrared divergences cancel (for regular observables). Since in general the inclusive sum is done numerically, the cancellation needs to be controlled analytically first, a difficult issue which requires the understanding of the physics of the problem (see for instance $^{28}$). Results and discussion on this issues have been presented at this conference $^{29}$. Various techniques are used. On one hand there are seminumerical approaches. An example is the $gg \rightarrow gggg$ one loop amplitude $^{23}$ (some of the helicity configurations have been computed analytically by twistor techniques). On the other hand there are direct analytical approaches. Powerful methods $^{17}$ are based on Cutkosky
rules (unitarity), on the use of Passarino-Veltman reduction of any one-loop amplitudes in terms of a basis of scalar integrals and on recurrence relations.

The results of these studies are usually obtained after profound understanding of the general structure of Feynman diagrams. Often one finds general properties in gauge theories which point toward simple structures.

Contributions to these conferences on phenomenological studies at Lep, Hera, Tevatron and LHC of high order results results have been presented. They are:
i) running coupling measurements at Hera and Lep. Lattice calculations can be used to reduce the theoretical errors;
ii) jet emission studies at Lep, Hera, Tevatron and LHC. In particular at the Tevatron jet-finding algorithms start to be used;
iii) $W/Z$ and $W/Z+jets$ at Tevatron;
iv) $Higgs$ and $W/Z$ production to NNLO at LHC.

The general comment for these analyses is that NLO corrections improve the accuracy and the description of the data.

5 High order parton splitting

Parton density functions enter DIS and, due to QCD factorization, hadron collider distributions. Fragmentation functions, which describe inclusive final state emission, enter all collider studies. Their $Q^2$-evolution is governed by the corresponding space- and time-like parton splitting functions (anomalous dimensions) which have been computed in 1980 at two loops both for the singlet and non-singlet cases. Recently the anomalous dimensions have been computed at three loops: the singlet and non-singlet ones for the space-like case; the non-singlet ones for the time-like case. These very important results obtained in $\overline{MS}$ scheme have been already used for various phenomenological studies: Sudakov resummations, lepton pair and Higgs boson production, quark form factor, threshold resummation, DIS by photon exchange, longitudinal structure function, non-singlet analysis of deep inelastic world data.

High order anomalous dimensions are also important to understand general features of QCD and gauge theories in general. I discuss here two examples: relation between DGLAP and BFKL evolution in SYM theory and relation between space- and time-like anomalous dimensions.

**DGLAP and BFKL evolutions.** It has been shown that in the $\mathcal{N} = 4$ Supersymmetric Yang-Mills theory there is a deep relation between the BFKL and DGLAP evolution equations. In this theory, the eigenvalues of the space-like anomalous dimension matrix are expressed in terms of a universal function constrained (obtained) from the BFKL equation. This was checked at two loops by direct calculations and at three loop with the anomalous dimensions obtained for SYM from the QCD ones. It is important to explore to what extent the relations between BLKF and DGLAP can be extended to QCD.

**Relating S- and T-evolution.** The search for a relation between S- and T-anomalous dimension (space-like $\gamma_-(N)$ and time-like $\gamma_+(N)$) has a long story: Drell-Levi-Yan relation, Gribov-Lipatov relation, the analytical continuation. Consider DIS with $q$ the large space-like momentum transferred from the incident lepton to the target nucleon $P$ and $e^+e^-$ annihilation with $q$ the time-like total incoming momentum and $P$ the final observed hadron. The Bjorken and Feynman variables in DIS and $e^+e^-$ are

$$x_B = \frac{-q^2}{2Pq}, \quad x_F = \frac{2(Pq)}{q^2}.$$  

These variables are mutually reciprocal: after the crossing operation $P \rightarrow -P$ one $x$ becomes the inverse of the other (although in both channels $0 \leq x \leq 1$ thus requiring the analytical continuation). This fact was
the basis for the search of reciprocity relations between $\gamma_-(N)$ and $\gamma_+(N)$ (here $N$ is the Mellin moment conjugate to $x$).

Recently it has been noticed that new information on the relation between $\gamma_-(N)$ and $\gamma_+(N)$ could be obtained by taking into account that such a reciprocity property $x \rightarrow 1/x$ ($x = x_B$ or $x_F$) can be extended to the Feynman diagram for the two processes and, in particular, to the contributions from mass-singularities described by multi-parton splitting. Consider the three-parton vertex kinematics of the decay $k_0 \rightarrow k + k'$ in the DIS situation: $k_0^2 < 0$, $k^2 < 0$, $k'^2 > 0$. To change to the annihilation kinematics, $-k \rightarrow -k_0 + k'$, one has to change signs of $k_{+0}$ and $k_+$ and of the corresponding virtualities, $k_0^2 > 0$, $k^2 > 0$. The virtuality $k^2$ enters the denominators of the Feynman diagrams. In order for the transverse momentum integration produce a logarithmic enhancement, the conditions must be satisfied

$$|k_0^2| < \frac{k_{+0}^2}{k^2} = z^2 \cdot \kappa^2, \quad \sigma = \pm$$

(2)

with $\sigma = -1$ and $\sigma = 1$ for S- and T-case respectively. This kinematical fact has a strong impact on the relation between the S- and T-probability $D_\sigma(N, \kappa^2)$ to find a parton with virtuality up to $\kappa^2$. From (2) one directly deduces the following reciprocity respecting equation (RRE)

$$\kappa^2 \partial_\sigma D_\sigma(N, \kappa^2) = \gamma_\sigma(N) D_\sigma(N, \kappa^2)$$

(3)

$$= \int_0^1 \frac{dz}{z} z^N P(z, \alpha_S) D_\sigma(N, z^2 \kappa^2).$$

The difference between the two channels is simply in the fact that the virtuality of the integrated parton distribution is $\kappa^2 z^\sigma$, see (2). The splitting function $P(z, \alpha_S)$ does not depend on the S- or T-channel (its Mellin moments are not the anomalous dimensions). The running coupling in the splitting function depends on the virtuality in a reciprocity respecting form. This equation (in general a matrix equation) is non-local: for $\sigma = -1$ ($\sigma = 1$) the right hand side involves the parton distribution with all virtualities larger (smaller) than $\kappa^2$. So RRE is not suitable for explicit calculations of the anomalous dimensions, but for relating them.

Is eq. (3) correct? It is the result of the vertex kinematical ordering (2) for mass singularities. However, when dimensional regularization is used, the implication of the vertex kinematical ordering gets mixed with the fact that the S- and T-channel phase space differ by a factor $z^{-2\kappa}$. However corrections coming from this $z$-factor do not lead to really new structures but are essentially related to the anomalous dimensions at lower orders. This suggests (see discussion in 44) that these corrections could be an artifact of dimensional regularization so that, at the end of the calculation, reciprocity is actually restored leaving RRE unmodified.

The reciprocity relation (3) can be tested for higher order S- and T-anomalous dimensions 38,39. To do that one has to account also for the arguments of the running coupling in S- and T-cases which give contributions proportional to the beta-function. However beta-function contributions arise also from the factorization scheme used for S- and T-case and how to account for their reciprocity has not been studied yet. Then RRE have been tested only for the fixed $\alpha_S$ contribution. In this case RRE can be written as

$$\gamma_\sigma(N) = P(N + \sigma \gamma_\sigma),$$

(4)

$$P(N) = \int_0^1 \frac{dz}{z} z^N P(z),$$

with $P(N)$ a universal function depending on $\alpha_S$. Eq. (4) has been tested to high order in three cases: the non-singlet case, the large $x$ and the small $x$ behaviour.

1) Non-singlet case. From (4) one has

$$\gamma_\sigma^1(N) = P^1(N)$$

$$\gamma_\sigma^2(N) = P^2(N) + \sigma \gamma_\sigma^1(N) \hat{\gamma}_\sigma^1(N) + \gamma_\sigma^3(N) + \hat{\gamma}_\sigma^3(N) + \gamma_\sigma^4(N) \hat{\gamma}_\sigma^2(N)$$

$$\cdots$$

$$+ \sigma \gamma_\sigma^4(N) \hat{\gamma}_\sigma^1(N) + \gamma_\sigma^5(N) \hat{\gamma}_\sigma^2(N)$$

with $\gamma_\sigma^n$ and $P^n$ the $n$-th expansion coeffi-
coefficients in $\alpha_S$ of non-singlet $\gamma_\sigma(N)$ and $\mathcal{P}_\sigma(N)$. Dots are derivatives with respect to $N$. The first is the Gribov-Lipatov relation (independence of $\sigma$ valid only to one loop). The second is the two loop relation which has been pointed out in \ref{37} and has been one of the important elements used \ref{45} to derive RRE. The last one has been verified \ref{39} at three loop order.

2) Large $x$ behaviour. The dominant channels are the diagonal ones $gg$ and $qq$. Denoting by $\gamma_\sigma^{(a)}(x)$ the $x$-space anomalous dimension with $a = gg$ or $q\bar{q}$, one has

$$\gamma_\sigma^{(a)}(x) = \frac{A^{(a)} x}{(1-x)_+} + B^{(a)} \delta(1-x) + C^{(a)} \ln(1-x) + D^{(a)} + \cdots$$

(5)

The various coefficients are functions of $\alpha_S$. The two most singular terms do not depend on $\sigma$. In particular the first term corresponds to the classical soft radiation \ref{46} which is universal and depends only on the charge $A^{(a)}$ of the source. It can be expressed in terms of a physical coupling \ref{47} as $A^{(a)} = (C_a/\pi)\alpha_S^{\text{phys}}$ with $C_a = C_A, C_F$ for $= gg, q\bar{q}$. Using (4) one deduces

$$C^{(a)} = -\sigma A^{(a)}, \quad D^{(a)} = -\sigma A^{(a)} B^{(a)}.$$ 

These two relations are verified at three loop level, apart for a beta-function contribution entering $D^{(a)}$ which is not considered in (4).

Very recently the expansion of the leading coefficient $A$ has been evaluated \ref{48} in SYM with $\mathcal{N} = 4$ to all order in $\alpha_S$ (in dimensional regularization scheme suited for supersymmetric theories) in the S-case. All coefficients satisfy the “trascendentality principle” (i.e. are given in terms of derivative of the Euler psi-function). The first three terms agree with the calculation reported in \ref{41}. According to (5) the same result should be obtained for the T-case.

3) Small $x$ behaviour. The leading order contributions for $N \to 0$ (corresponding to small $x$) are given by ($\bar{\alpha}_S = C_A \alpha_S/\pi$)

$$\gamma_-(N) = \frac{\bar{\alpha}_S}{N}, \quad \gamma_+(N) = \frac{1}{4}(\sqrt{N^2 + 8\bar{\alpha}_S} - N)$$

These two expressions satisfy (4). They are the result of cancellations in the phase space due to coherence of soft radiation: angular ordering for the T-case and transverse momentum ordering for the S-case. Thus cancellations in the mass singularity phase space are reciprocity related so that RRE incorporates them into the universal function $P(z)$. RRE can be tested to higher order. One of the most interesting outputs is that the “accidental” absence in the leading BFKL anomalous dimension of the $\alpha_S^2/N^2$ and $\alpha_S^3/N^3$ terms implies, via reciprocity (4), the fact that exact angular ordering is valid beyond leading order up to NNLO.

6 Additional problems in hadron-hadron collisions

Hard processes in hadron collider are initiated by elementary hard cross sections with 2 incoming and $n$ outgoing partons

$$p_1 p_2 \to 0 \quad \text{DY, WW, ZZ} \cdots$$
$$p_1 p_2 \to 1 \quad (p_t\text{-Higgs}, \text{Higgs+jet}) \cdots$$
$$p_1 p_2 \to 2 \quad \text{dijet-distributions, jet-shape} \cdots$$
$$p_1 p_2 \to n \quad n > 2, \text{ multi-jet distributions}$$

For $n \geq 2$ one has a new QCD challenge. Consider inclusive distributions with two different hard scales ($Q \gg Q_0 \gg \Lambda_{QCD}$) for which logarithmic resummations are needed. An example is the out-of-event-plane energy distribution. Because of the complex structure of colour matrices for $n \geq 2$, soft gluon at large angles contribute to single logarithmic accuracy. Consider the hard vertex with four partons

$$p_1 \ p_2 \to p_3 \ p_4$$

(6)

Resummation of collinear and infrared logarithmic contributions coming from radiation emitted off the four primary QCD partons gives rise to four Sudakov form factors. Each factor has the charge of the emitting parton and the hard scale identified (as in $\gamma^+\gamma^-$)
by the single logarithmic analysis. To single log accuracy one has additional contributions coming from large angle soft QCD emission which are resummed into a fifth form factor. To understand why these large angle soft terms enter only for \( n \geq 2 \) consider the hard process (6) with matrix colour charges \( T_i \). Emission of a soft gluon \( q \) off these four partons is given by the square of the eikonal current. Using charge conservation \( (T_1 + T_2 = T_3 + T_4) \) one has

\[
j^2(q) = \left( \sum_i T_i \frac{\mu}{\mu q} \right)^2 = \sum_i T_i^2 W_i(q) + T_t^2 A_t(q) + T_a^2 A_a(q)
\]

with \( T_i^2 = C_i \) the square colour charge and \( W_i(q) \) the collinear and infrared divergent emission factor leading to the standard Sudakov form factor of the primary parton \( i \). The emission factors \( A_t \) and \( A_a \) are infrared but not collinear singular. While \( T_i^2 \) are proportional to unity, the exchanged charges \( T_t^2 = (T_1 - T_3)^2 \) and \( T_a^2 = (T_1 - T_4)^2 \) are colour matrices (6 \( \times 6 \) in \( SU(N) \) for \( gg \to gg \)). These last two terms, together with the corresponding Coulomb phases, give rise to the matrix fifth form factor (which needs to be diagonalized). From this discussion it is clear that these additional soft, but non-collinear, contributions are absent for less than four colour particles since there isn’t any exchanged channel.

For \( gg \to gg \) there is a puzzle. The relevant eigenvalues of the soft distribution (7) are symmetric under the exchange between external and internal space variables:

\[
\frac{\ln \frac{\mu^2}{\mu q} - 2 \pi i}{\ln \frac{\mu}{\mu}} \iff N_c.
\]

Another surprising result concerns the contribution beyond single-log (next-to-next-to-leading) to the fifth form factor: it has been found that it is proportional to the corresponding one-loop contribution and again one reconstructs the physical coupling \( A^{(a)} \) in (5). This result makes possible a variety of resummations at next-to-next-to leading order.

7 Jets, small-\( x \) and all that

Jets shape distributions. Examples in hadron hadron collisions are the out-of-event-plane energy distribution and the energy-energy azimuthal correlations. Both reliable predictions and experimental data are difficult to obtain. The large number of hadrons emitted at LHC adds a further algorithmic difficulty. A typical \( k_t \) jet-finder algorithms scales as \( n^3 \) with \( n \) the number of particles. A good news is that a jet-finder algorithms was found, using techniques developed in computational geometry, that scales as \( n \ln n \). The techniques required to obtain reliable jet distribution predictions are well understood: single logarithmic resummation, matching with known fix-order NLO results and non-perturbative power corrections. Reliable predictions have been obtained when no more than four jet are involved, counting also incoming jets that is one in DIS and two in hardon-hadron collisions. In this last case most events have four jets and analytical calculations become very laborious (see previous section) and the use of automate resummation becomes essentially unavoidable. An additional difficulty in the study of hadron emission associated to new physics events is that one may want to describe QCD radiation in given geometrical regions. In such a case one needs to resum non-global logs and this requires solving non-linear equations.

Small-\( x \) physics enters a large number of processes and gives rise to many interesting problems. The connection between BFKL and DGLAP equations has been already mentioned. Important developments are on linear and non-linear small-\( x \) regimes.

In the linear regime there have been extensive studies of the higher-order corrections which could stabilize the NLO poor perturbative convergence and construct a framework.
useful for phenomenological applications. The general method consists of matching small-$x$ resummation with collinear singularity resummations (running $\alpha_S$, anomalous dimensions, factorization scheme...).

The non-linear regime, a major issue in small-$x$ physics, is characterized by saturation which should be seen in heavy ion collisions. A major point is how to go beyond the Balitsky-Kovchegov equation viewed as a mean field approximation of a more general equation taking into account unitarity in perturbative QCD. A number of new formulations have been proposed, but it seems that a complete and consistent formulation is still lacking. An important attempt to account for general non-linear QCD corrections at small-$x$ is based on the observation that the BK equation is in the same universality class as the Fisher-Kolmogorov-Petrovsky-Piscounov equation introduced in statistical physics to discuss reaction-diffusion phenomena. Therefore methods of stochastic physics are used to study small-$x$ QCD phenomena (and viceversa).

Perturbative QCD studies of multiple interactions in DIS and hadron-hadron collisions are important for various reasons: they are needed to set up the theoretical framework for the small-$x$ non-linear equations as the BK equation and to perform phenomenological studies including multi-jet final states (background to new physics), heavy flavour jets (near forward direction) and underlying event.

Diffraction is phenomenologically important since even hard events have diffractive components. Moreover, diffraction could be used to study Higgs production. Diffraction is a non-perturbative fields and phenomenological models are needed. Their basis is unitarity and interactions between Pomeron.

8 Final remarks

This (theoretical) review of QCD is incomplete and mostly neglects the phenomenology of long distance physics. The list of issues here discussed and of the ones not discussed (NLO matching with Monte Carlo simulations, PDF and structure functions, photon emission, two photon scattering, prompt photons, power corrections, Bjorken and Gross-Llewellyn Smith sum rules, underlying events...) speaks for itself of the importance of QCD for high energy physics in general. QCD for LHC physics poses new problems and gives new information on emission such as rotation in colour space and consequent non-planar corrections.

Very laborious calculations, needed for LHC accurate predictions, reveal simple structures and properties. This fact points to the possibility of a new more efficient formulation of QCD. Is it possible that this could be provided by ideas developed in string theory? There are indications in this directions although not yet convincing: the abundant phenomenological results obtained within the framework of AdS/CFT as a way to “solve” QCD; Feynman diagrams in twistor formulations. However, for the moment, all these developments are only formal and, to make a decisive step in the understanding of QCD, one needs to find their physics basis.

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