STATUS OF STRUCTURE FUNCTIONS AND PARTONS

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We briefly review some of the developments in the study of parton distributions which have occurred since DIS2000, including discussion of uncertainties, shadowing, unintegrated and generalized distributions.

1 Data and parton distribution functions

The situation is summarised in Fig. 1, which shows the kinematic regions in the $(x, Q^2)$ plane covered by (i) the experiments at HERA, (ii) the fixed target deep-inelastic scattering experiments and (iii) the single jet inclusive experiments at the Tevatron. Roughly speaking, the fixed target experiments determine the distributions of the $u, d, \bar{u}, \bar{d}, s, \bar{s}$ quarks for $x \sim 0.1$ (and $Q^2 \sim 10 \text{ GeV}^2$). The HERA $F_2$ measurements determine the sea quark distributions, and their $\partial F_2/\partial \ln Q^2$ data determine the gluon, for $x \sim 10^{-3}$ (and $Q^2 \sim 10 \text{ GeV}^2$). The Tevatron jet data determine the quarks and gluon distributions in a region around $x \sim 0.1$ (and $Q^2 \sim 10^4 \text{ GeV}^2$).

In the past year new data have been available in each of the three domains. The H1 and ZEUS experiments have measured $F_2$ over a larger domain (see Fig. 1) with much improved precision. The D0 and CDF collaborations have measured the inclusive single jet $E_T$ distribution and provided correlated statistical and systematic errors. In particular, the D0 experiment has measured the $E_T$ distribution in five different $\eta$ bins, out to $\eta = 3$ and so samples partons in a much wider range of $x$ (see Fig. 1). The NuTeV collaboration have used $\nu, \bar{\nu}$ beams to measure $F_2, x F_3$ with increased precision, and also, by observing $\mu^+\mu^-$ production, to obtain information on the $s$ and $\bar{s}$ quark distributions; CCFR have made a model-independent re-analysis of their $\nu, \bar{\nu}$ data, which removes the discrepancy for $x < 0.1$ with NMC. The E866 collaboration have observed both $pp$ and “$pm$” Drell-Yan production and further constrained the difference of the $\bar{u}$ and $\bar{d}$ distributions.

2 Uncertainties in parton distribution functions and observables

The parton distributions are determined from NLO DGLAP fits to data from about 14 diverse experiments, with typically 1500 data points and more than...
Figure 1. The kinematic domains probed by the various experiments, shown together with the partons that they constrain. The new data that have become available since DIS2000 are also indicated.

20 parameters. There are many sources of uncertainty. First the statistical and systematic experimental errors. The latter are often not randomly distributed and, moreover, may depend on theory. The theoretical uncertainties come from higher order QCD contributions, the choice of factorization and renormalization scales, resummation corrections (from resumming \( \ln 1/x, \ln(1 - x) \) terms), power law contributions, nuclear target corrections, the choice of the particular parametrization of the starting distributions and, in hadronic processes, from the treatment of the underlying event.

To date, the best attempt to determine the errors on parton distributions is due to Botje. He includes the statistical and systematic covariance matri-
cesses and allows for some of the theory uncertainties. However, the analysis is based on five massless flavours and so is only applicable for $x > 10^{-3}$. Moreover there is no constraint on the gluon at large $x$ values (as would come from including Tevatron jet data in the fit) and so the gluon is undetermined for large $x$, which feeds down to small $x$ via the momentum sum rule. A similar analysis, also based on the Hessian method, has been performed by CTEQ.

Recently there have been several contributions concerning the errors on partons and observables. Giele et al. have obtained 100,000 ‘optimized’ sets of partons, and expressed the uncertainties as a density measure in pdf functional space. In principle it is ‘easy’ to propagate errors given complete data information, but the question of theory uncertainties has to be addressed. So far only $F_{2}^{ep}$ data have been considered; Giele et al. conclude that NMC and (old) ZEUS data are incompatible with the other $F_{2}^{ep}$ data sets. This type of analysis is in its infancy. Clearly it is necessary to include other types of data, e.g. neutrino and jet data.

It is advantageous to determine the uncertainties on observables directly, rather than to go through the intermediate stage of determining the parton errors. The disadvantage is that global analyses have to be carried out for each observable separately. In view of the importance of $W$ and $Z$ boson hadroproduction as luminosity monitoring processes at the Tevatron and the LHC, attention has focused on $\sigma_{W,Z}$. MRST(1999) found an uncertainty $\Delta \sigma_{W}$ of $\pm 3\%$ at the Tevatron and $\pm 5\%$ at the LHC. In the latter case the uncertainty is mainly due to the (conservative) error of $\pm 0.005$ assigned to $\alpha_{S}(M_{Z}^{2})$. More recently CTEQ found $\Delta \sigma_{W} = \pm 4\%$ at the Tevatron and $\pm 8\%$ at the LHC, taking only data errors into account. These surprisingly large errors are due to using older data and choosing a large tolerance, $\chi^{2} \sim \chi^{2}_{\text{min}} < 150$. In a new global analysis, which incorporates all the new precise data, MRST (2001) find $\Delta \sigma_{W} = \pm 2\%$ at both Tevatron and LHC energies, when only data errors are considered. However theory errors ($\Delta \alpha_{S}$ etc.) have to be included. Interestingly this latest analysis, which incorporates the new Tevatron jet data, gives an error of only $\pm 20\%$ on the high $x$ gluon density.

3 The description of $F_{2}$ data at low $x$

A good description of the $F_{2}^{ep}$ data in the low $x$ domain can be obtained in many different ways. First, there are empirical models, with very few input parameters (such as that of Haidt), which give excellent descriptions down to very low $Q^{2}$. Then there are dynamically-motivated parametrizations, such as the ‘saturation’ model of Golec-Biernat and Wüsthoff based on the $q\bar{q}$ dipole framework. Surprisingly, on a more comprehensive level, the NLO
DGLAP analyses continue to give satisfactory descriptions of the now very precise $F_2$ data, down to remarkably low $Q^2$ and $x$. However in these fits the gluon tends to go valence-like or negative at low $Q^2$, $Q^2 \sim 1 \text{ GeV}^2$, indicating that the approach ceases to be valid in this domain. This is also reflected in the anomalous behaviour of $F_L$ in this domain, perhaps indicating the need for $\ln(1/x)$ resummations. In fact, a good description of $F_2$ is also obtained in a unified approach which incorporates DGLAP and BFKL with higher order $\ln(1/x)$ contributions. Here a ‘flat’ gluon is input.

The message is clear. One low $x$ observable can be described in many ways. It is possible to trade effects of the different types of perturbative evolution with different choices of the input forms. Of course, the situation would be changed if another independent quantity, such as $F_L$, were precisely measured at low $x$.

4 BFKL and DIS

The resummation of the $\ln(1/x)$ contributions is now known to NLO. The resulting $x \to 0$ behaviour has the form $x^{-\omega}$ with

$$\omega = \omega_0 (1 - 6.5 \bar{\alpha}_S),$$

where $\omega_0 = \bar{\alpha}_S 4 \ln 2$ is the LO behaviour; as usual $\bar{\alpha}_S \equiv 3 \alpha_S / \pi$. At first it was thought that such large NLO corrections would mean that no stable small $x$ predictions could be made using the BFKL procedure. However it turns out that this is not the case. It is possible to identify higher-order terms and then to resum them. Indeed Ciafaloni et al. carry out an all-order $\ln(1/x)$ resummation of the following effects: (i) running $\alpha_S$, (ii) the non-singular DGLAP terms and (iii) the angular ordering and energy constraints. (Thorne has made a recent detailed study of effect (i)). The result is a stable $x^{-\omega}$ behaviour which is consistent with observations. In fact, prior to this, the fit of Kwiecinski et al. mentioned at the end of Section 3, was based on a unified equation which incorporates these all-order $\ln(1/x)$ contributions, where the imposition of a consistency (or kinematic) constraint plays a major role.

We may use a very simplified calculation to show how including effect (iii) tames the NLO behaviour of $[\text{4}]$. Recall that as we proceed along the BFKL gluon chain, we have ordering in the longitudinal momenta, $Y > Y'$, where $Y \equiv \ln(1/x)$. To allow for the energy/angular ordering constraints, the ordering takes the modified form

$$Y > Y' + \delta.$$  

(2)
Hence the BFKL equation for the unintegrated gluon distribution becomes

\[ f(Y) = f_0 + \int_{Y-\delta}^{Y} K f(Y') dY', \]  

(3)

or, in differential form

\[ \frac{\partial f(Y)}{\partial Y} = \omega_0 f(Y - \delta). \]  

(4)

Note that if \( \delta = 0 \) then \( f = f_0 \exp(\omega_0 Y) \sim x^{-\omega_0} \), as required. In general, the solution of (4) is \( f \sim \exp(\omega Y) \sim x^{-\omega} \), where

\[ \omega = \omega_0 e^{-\omega_0 \delta} = \omega_0 (1 - \omega_0 \delta + \ldots). \]  

(5)

Comparison with (1) reveals \( \delta = \frac{6.5}{4} \ln 2 = 2.3 \), so

\[ \omega = \omega_0 \exp(-2.3 \omega), \]  

(6)

which leads to a similar behaviour of \( \omega \) as a function of \( \bar{\alpha}_S \) to that found in Refs. 21, 19. This toy model illustrates how the inclusion of a summation of higher-order terms stabilizes the NLO result.

A more ‘phenomenological’ way of performing the resummation of higher-order contributions, and achieving the perturbative stability of the BFKL approach, has been proposed 22. This incorporates the one- and two-loop terms in the anomalous dimension, imposes momentum conservation and parametrizes the residual ambiguity in terms of a single parameter \( \omega_{\text{eff}} \) which specifies the \( x^{-\omega_{\text{eff}}} \) behaviour as \( x \to 0 \).

5 Parton shadowing

Let us start with the original GLR equation 23, written in terms of DGLAP evolution for the gluon density

\[ \frac{\partial (xg(x,Q^2))}{\partial Y \partial \ln(Q^2/\Lambda^2)} = \frac{N_C \alpha_S}{\pi} xg - \frac{\alpha_S^2}{R^2 Q^2} [xg]^2, \]  

(7)

where, as before, \( Y = \ln(1/x) \). The first term on the right-hand-side is the growth of \( g \) due to DGLAP evolution, and the second is the decrease due to gluon-gluon recombination. To gain insight into the origin of the shadowing term, note that the number, \( n \), of gluons per unit rapidity interval is \( xg(x,Q^2) \). Moreover the gluon-gluon cross section \( \hat{\sigma}(gg) \sim \frac{\alpha_S^2}{R^2 Q^2} \) and so

\[ \text{prob. of recomb.} \sim \frac{n^2 \hat{\sigma}}{\pi R^2} \sim \frac{\alpha_S^2}{R^2 Q^2} [xg]^2, \]  

(8)
where \(\pi R^2\) is the transverse area populated by gluons. The GLR equation effectively resums the ‘fan’ diagrams generated by the branching of QCD Pomerons. Recently there has been much activity in this area \(^{24}\), \(^{25}\), which has resulted in an improved knowledge of the triple-Pomeron vertex.

The structure of the triple-Pomeron vertex can be extracted from an equation \(^{23}\) for a quantity, \(N(r, b, Y)\), closely related to the cross section for the interaction of a \(q\bar{q}\) dipole of transverse size \(r\) with the proton target

\[
\sigma(r, Y) = 2 \int d^2 b \, N(r, b, Y). \quad (9)
\]

\(b\) is the impact parameter of the interaction. In the short-distance approximation \((r \ll b)\) the non-linear shadowing equation takes the simplified form

\[
\frac{\partial \tilde{N}(r, b, Y)}{\partial Y} = \frac{N_C \alpha_S}{\pi} \left\{ \text{K} \otimes \tilde{N} - \tilde{N}^2 \right\} \quad (10)
\]

in the large \(N_C\) limit, where \(K\) is the BFKL kernel and

\[
\tilde{N}(r, b, Y) = \int_0^\infty \frac{d\ell}{\ell} \, J_0(\ell r) \, N(\ell, b, Y). \quad (11)
\]

There have been two recent attempts to use \(^{15}\) to calculate the effect of gluon shadowing \(^{26}\), \(^{27}\). Both give similar results. Fig. 2 shows the results of Kimber et al. We see that shadowing is small in the domain accessible to experiments at HERA. However almost a factor of two suppression is anticipated for \(x \sim \text{few } \times 10^{-6}\). This domain may be accessible at the LHC for either prompt photon or high-\(Q_T\) Drell-Yan production at large rapidities, both of which can proceed via the subprocess \(gg \rightarrow \gamma q\). Another approach \(^{28}\) evaluates the corrections to \(\partial F_2/\partial \ln Q^2\) due to twist-4 gluon recombination.

If Fig. 2 shows that shadowing is small at HERA, does this rule out the saturation models, such as \(^{14}\)? Not necessarily; it is hard to distinguish the \(Q^2\) dependence with shadowing present from the pure DGLAP \(Q^2\) dependence, since the difference can be removed by an adjustment of the starting distributions \(^{29}\). However a distinction may be possible at the LHC. A much better way to identify shadowing is to study the \(A\) dependence of scattering on nuclei.

6 Unintegrated parton distributions

The natural framework with which to discuss DIS and related hard scattering processes is to use parton distributions \(f_a(x, k_t^2, \mu^2)\), unintegrated over the transverse momentum \(k_t\) of the parton, together with the \(k_t\) factorization
Figure 2. The effect of shadowing on the integrated gluon distribution $xg(x,Q^2)$ at $Q^2 = 4$ GeV$^2$, taken from Kimber et al. [26].

The unintegrated distributions have the advantage that they exactly correspond to the quantities which enter the Feynman diagrams and therefore allow for the true kinematics of the process even at LO. Here, for simplicity, we will discuss the gluon and so omit the subscript $a = g$. The distribution $f(x, k_t^2, \mu^2)$ depends on two hard scales — $k_t$ and the scale $\mu$ of the probe. The scale $\mu$ plays a dual role. On the one hand it acts as the factorization scale, while on the other hand it controls the angular ordering of the gluons emitted in the evolution. The two-scale distribution satisfies the CCFM equation [31], which embodies both DGLAP and BFKL evolution. In practice, it is complicated to solve the equation, and, up to now, it has only proved to be practical within Monte Carlo generators [32].

To gain insight, recall that both DGLAP and BFKL evolution are essentially equivalent to ordered evolution in the angles of the emitted gluons. In the DGLAP collinear approximation the angle increases due to the growth of $k_t$, while in BFKL the angle ($\theta \approx k_t/k_\ell$) grows due to the decreasing longitudinal momentum fraction as we proceed along the emission chain from the proton. The factorization scale $\mu$ separates the gluons associated with emission from different parts of the process, that is from the beam and target protons (in $pp$ collisions) and from the hard subprocess. For example, $\mu$ separates emissions from the beam (with polar angle $\theta \lesssim 90^\circ$) from those of the target (with $\theta \gtrsim 90^\circ$), and from the intermediate gluons associated with
the hard subprocess. This separation was proved by CCFM 31 and originates
from the destructive interference of the different emission amplitudes in the
angular boundary regions. Since the evolution process is essentially controlled
by one quantity, the emission angle, we should expect to be able to obtain
the unintegrated gluon distribution \( f(x, k_t^2, \mu^2) \) from a single-scale evolution
equation. Indeed it is possible to accomplish this and to follow an analytic
approach where the physical assumptions are more evident and where, in prin-
ciple, the NLO corrections can be included. The key observation is that the
\( \mu \) dependence enters only at the last step of the evolution 33.

To illustrate the last-step procedure, we start from the simplified case of
pure DGLAP evolution for \( G = xg \), where \( g \) is the conventional (integrated)
distribution
\[
\frac{\partial G(x, k_t^2)}{\partial \ln k_t^2} = \frac{\alpha_s}{2\pi} \left[ \int_x^{1-\Delta} P_{gg}(z) G \left( \frac{x}{z}, k_t^2 \right) dz - G(x, k_t^2) \int_0^{1-\Delta} z P_{gg}(z) dz \right].
\] (12)

Suppose that we were to omit the virtual contribution, then the unintegrated
density would be
\[
f(x, k_t^2) = \frac{\partial G(x, k_t^2)}{\partial \ln k_t^2} = \frac{\alpha_s}{2\pi} \int_x^{1-\Delta} P_{gg}(z) G \left( \frac{x}{z}, k_t^2 \right) dz. \tag{13}
\]
The virtual contributions do not change the \( k_t \) of the gluon and may be
resummed to give the survival probability \( T \) that the gluon remains untouched
in the evolution up to the factorization scale \( \mu \)
\[
T(k_t, \mu) = \exp \left( -\int_{k_t^2}^{\mu^2} \frac{dk_t'^2}{k_t'^2} \frac{\alpha_s}{2\pi} \int_0^{1-\Delta} z P_{gg}(z) dz \right), \tag{14}
\]
as in the Sudakov form factor. Thus the probability to find a gluon with
transverse momentum \( k_t \) (which initiates a hard subprocess with factorization
scale \( \mu \)) is
\[
f(x, k_t^2, \mu^2) = T(k_t, \mu) \left[ \frac{\alpha_s}{2\pi} \int_x^{1-\Delta} P_{gg}(z) g \left( \frac{x}{z}, k_t^2 \right) dz \right]. \tag{15}
\]
It is at this last step that the unintegrated distribution becomes dependent
on \( \mu \). Angular ordering requires that the cut-off \( \Delta = k_t/\left( \mu + k_t \right) \) in (15), and
analogously in (14). The last step causes the distribution to increasingly spill
over into the domain \( k_t > \mu \) as \( x \) decreases, as shown by the dotted curves in
Fig. 3.

To include \( \ln(1/x) \) effects, (12) is replaced by the unified equation of
Kwiecinski et al. 17, and a similar procedure followed 33. The continuous
The continuous curves show the $k_t^2$ dependence of the unintegrated gluon $f(x, k_t^2, \mu^2)$ for $x = 0.1$ and $0.0001$ at $\mu = 10$ GeV obtained from a unified BFKL/DGLAP study. The dotted curves are obtained from DGLAP evolution. The figure is taken from Kimber et al. [33].

curves in Fig. 3 are obtained, which are not very different from the previous DGLAP results. We conclude that the main physical effects come from angular ordering in the last step of the evolution, and not from $\ln(1/x)$ terms.

By imposing angular ordering in both the BFKL and DGLAP terms the integral up to $\mu^2$ of $f$ does not equal the integrated gluon. An evaluation of $f$ which imposes the equality, but does not have complete angular ordering, has also been made [33]. The difference is a NLO effect [34].

7 Generalized parton distributions

Two relevant examples, illustrating the need for generalized (or skewed) parton distributions, are shown in Fig. 4. The first shows either deeply virtual Compton scattering (DVCS) or the electroproduction of vector mesons at $t = 0$. The second is double-diffractive Higgs production in $pp$ collisions. These processes involve off-diagonal proton matrix elements $\langle p' | \ldots | p \rangle$, with longitudinal components of momentum transfer, $x' \neq x$. We have

$$\frac{d\sigma}{dt}(\gamma^* p \rightarrow \gamma p \text{ or } V p) \bigg|_{t=0} = \int dz \left[ \int \frac{d\ell_t^2}{\ell_t^4} \ldots f(x, x'; \ell_t^2, \mu^2) \right]^2,$$

(16)
\[
\mu^2 = z(1-z)Q^2 + k_t^2 + m_q^2,
\]
and
\[
\sigma(pp \to p + H + p) = \left[\ldots \int \frac{d\ell_t^2}{\ell_t^2} f\left(x_1, x_1'; \ell_t^2, \frac{M_H^2}{4}\right) f\left(x_2, x_2'; \ell_t^2, \frac{M_H^2}{4}\right)\right]^2.
\]

At small \(x\) and \(x' \ll x\), applicable to these processes, the generalized distributions are given by
\[
f(x, x'; \ell_t^2, \mu^2) = R f(x, \ell_t^2, \mu^2) = R \frac{\partial}{\partial \ln \ell_t^2} \left(\sqrt{T(\ell_t, \mu)} x g(x, \ell_t^2)\right),
\]
where the survival probability \(T\) is given by (14). For small \(x\) the ratio \(R\) of the skewed to the diagonal distribution is known. It leads, for example, to an enhancement of \(R^2 \simeq 2\) for \(\sigma(\gamma p \to \Upsilon p)\) at HERA which seems to be required by the data, and an enhancement of \(R^4 \simeq 2\) of \(\sigma(pp \to p + H + p)\) at the LHC.

Data for \(\gamma^* p \to \rho p\) now exist from high energies (H1, ZEUS), through intermediate energies (E665, NMC), down to low energies (HERMES). At high energies (\(x < 0.01\)) the gluon mechanism shown in Fig. 4(a) dominates, whereas at low energies (\(x > 0.1\)) quark exchange takes over. The calculation of Vanderhaeghen et al. gives a satisfactory description.

Good data are now starting to accumulate on the classic DVCS process, which allows generalized parton distributions to be studied without the complications associated with the vector meson wave function. Allowance must be made for the Bethe-Heitler process, but the first indications are that the theory expectations are in good agreement with the data.

The original aim for observing DVCS was to measure the generalized parton distributions, and then to determine the angular momentum distributions.
of the partons in the proton. However generalized distributions are interesting in their own right and it is an active area for theoretical study. There are many interesting observable asymmetries. The evolution of the distributions is known to NLO. In analogy to the Mellin moments of ordinary partons, we should consider the Gegenbauer moments of the generalized distributions. The inverse transform is known. There have been attempts to calculate the distributions in the non-perturbative region. An area of present activity is the preservation of gauge invariance, which requires a twist-3 contribution. The talk by Diehl gives an overview of some of these developments.

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