STRONG INTERACTIONS*

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The main topics covered in this summary talk are the R ratio in e^+e^- annihilation, jet cross sections at next-to-leading order, physics at low x, the strong coupling constant, the large transverse energy jet excess, the gluon density function and direct photon production.

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1. Introduction

During the week we have seen a wide range of talks packed with new results and usually followed by lively discussions. As usual, it is an impossible task to summarise them in a short talk, particularly, as many of the talks were themselves summaries. I therefore restrict myself to outlining a few of the main themes chosen according to my own personal prejudice.

Throughout, I will refer to the talks contained in these proceedings by simply giving the author's name in italics (with the relevant figures in that talk shown in parentheses) and relying on the individual talks to provide references to the original literature. The few references I make are for specific additional points.

Of course, I must omit many interesting topics. For example, the soft-to-hard transition being probed in Deep Inelastic Scattering at HERA is particularly interesting (*A. Zarnecki* (Fig. 5)) as is the improved agreement between QCD and experiment in diphoton production at the TEVATRON (*H. Piekarz* (Fig. 4)). For more information on these and other topics, I refer you to the bulk of the proceedings.

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2. $R(s)$

The ratio of the cross section for $e^+e^- \to$ hadrons with that for $e^+e^- \to \mu^+\mu^-$ is one of the 'gold standards' of strong interaction observables. It is an easy observable to experimentally measure, is amenable to a perturbative expansion and has been computed through to $\mathcal{O}(\alpha_s^3)$. At this order, the calculation is rather complex and involves a large number of three-loop diagrams (or four-loop counting at the level of squared matrix elements). In fact, the original calculation of the $\mathcal{O}(\alpha_s^3)$ contribution was in error [1], both in sign and magnitude, and corrected subsequently [2]. However, independent checks are vital and *K. Chetyrkin* presented a new calculation in an arbitrary gauge which confirms the results of [2]. In addition, the contribution from gluino loops is now known and the full expression in the $\overline{MS}$-scheme and $\mu = \sqrt{s}$ for n_f light quark flavours with charge Q_f and $n_{\tilde{g}}$ light gluino octets reads,

$$\begin{aligned} R(s) = {} & 3 \sum_f Q_f^2 \Bigg(1 + \frac{\alpha_s}{\pi} + \left(\frac{\alpha_s}{\pi}\right)^2 (1.96 - 0.115 n_f - 0.323 n_{\tilde{g}}) \\ & + \left(\frac{\alpha_s}{\pi}\right)^3 \Big(-6.64 - 1.20 n_f - 0.0052 n_f^2 - 2.85 n_{\tilde{g}} \\ & - 0.031 n_f n_{\tilde{g}} - 0.047 n_{\tilde{g}}^2 \Big) \Bigg) - \left(\sum_f Q_f \right)^2 \left(\frac{\alpha_s}{\pi}\right)^3 1.24 . \qquad (1) \end{aligned}$$

Once the light gluino contribution to the three loop β function is known, yet tighter constraints on the existence of light gluinos can be derived from the supersymmetric running of the strong coupling constant.

At relatively low $\sqrt{s}$ the effects of the quark mass may become important. *J. Kühn* reported on his recent work where the full m_b dependence of the $\mathcal{O}(\alpha_s^2)$ term in $R(s)$ has been completely evaluated. This may prove important in determining α_s just below the B meson threshold at either CLEO or one of the planned b-factories. Secondly, by improving the theoretical description of R, it may be possible to reduce the error on the electromagnetic coupling constant $\alpha(M_Z)$.

Given the theoretical and experimental accuracy on $R(s)$, it is one of the best observables for extracting the strong coupling constant. As discussed above, $R(s)$ is known to third order,

$$R(s) = 3 \sum_f Q_f^2 \left(1 + a + r_1 a^2 + r_2 a^3 + \mathcal{O}(a^4) \right), \qquad (2)$$

where $a = \alpha_s/\pi$. However, r_2 depends on the choice of scheme and a is a function of the renormalisation scale μ such that,

$$\frac{\partial a}{\partial \ln \mu} = -ba^2 \left(1 + c_1 a + c_2 a^2 + \mathcal{O}(a^3)\right). \tag{3}$$

The coefficient c_2 also depends on scheme and the extraction of α_s therefore also depends on both scheme and scale, the RS. However, as discussed by *K. Chetyrkin* and *P. Raczka*, it is possible to reduce this uncertainty by working with the related, but *RS independent*, Adler function,

$$D(q^2) = -12\pi^2 q^2 \frac{d\Pi(q^2)}{dq^2}. \tag{4}$$

Although this equation is valid for space-like momenta, R can be extracted by a contour integral in the complex s-plane,

$$R(s) = -\frac{1}{2\pi i} \int_C d\sigma \frac{D(\sigma)}{\sigma}. \tag{5}$$

Rather than applying the principle of minimal sensitivity or Padé approximants to estimate R directly, these methods are applied to D and the contour integral performed numerically. As shown by *P. Raczka*, this reduces (a) the RS dependence and (b) the difference between R calculated at NLO and NNLO. This is undoubtedly a good thing and needs to be carried out on the Z resonance.

3. NLO jet cross sections

In addition to inclusive quantities like R, perturbative QCD also describes semi-inclusive quantities like jet cross sections very well. While it is not possible (at present) to go beyond next-to-leading order, there have been many NLO calculations and for most $2 \to 2$ processes a Monte Carlo program is available. In fact, for many processes it is vital to go beyond leading order so that particle recombination and formation of jets can be modelled, the scale dependence reduced and the presence of large infrared logarithms identified.

The common problem in such calculations is how to remove the infrared divergences from the real radiation and to numerically combine it with those present in the virtual contribution. This can be illustrated with a toy integral.

$$I = \int_0^1 dx \frac{F(x)}{x} x^{-\varepsilon}, \tag{6}$$

where x represents the additional phase space integrations for the real radiation. $F(0)$ represents the lowest order matrix elements, while the real radiation matrix elements, $F(x)/x$, diverge as $x \to 0$ representing the infrared behaviour while ε is a regulator which is ultimately set equal to 0. There are two established techniques for evaluating I. In the **slicing** method, a small parameter δ is introduced so that,

$$\begin{aligned} I &= \int_\delta^1 dx \frac{F(x)}{x} x^{-\varepsilon} + \int_0^\delta dx \frac{F(x)}{x} x^{-\varepsilon} \\ &\sim \int_\delta^1 dx \frac{F(x)}{x} + F(0)\log(\delta) - \frac{F(0)}{\varepsilon} . \end{aligned} \tag{7}$$

For $x < \delta$, we have made the approximation that $F(x) \sim F(0)$ and can perform the second integral to make the singularity explicit as a pole in ε. This must then cancel against the poles from the virtual graph which are always of the form $F(0)/\varepsilon$. Numerical methods may be applied to the first two terms in equation (7) and this method has been generalised for all QCD processes. However, in order for the approximation to be good, δ must be taken small and there are large numerical cancellations between the two terms. Furthermore, one should check by varying δ that the physical cross section does not depend on the unphysical parameter δ.

On the other hand, in the **subtraction** technique, we avoid the introduction of the parameter δ by just adding and subtracting a term to the integrand,

$$\begin{aligned} I &= \int_\delta^1 dx \frac{F(x)}{x} x^{-\varepsilon} - \int_\delta^1 dx \frac{F(0)}{x} x^{-\varepsilon} + \int_\delta^1 dx \frac{F(0)}{x} x^{-\varepsilon} \\ &\equiv \int_\delta^1 dx \left(\frac{F(x) - F(0)}{x} \right) - \frac{F(0)}{\varepsilon} . \end{aligned} \tag{8}$$

As before, the $\varepsilon \to 0$ limit may be taken for the first term and numerical methods applied. Until recently, the subtraction term was generated on a case by case approach. However, *S. Catani* described how this approach can be generalised for any QCD process with no approximation in the kinematics. While there are still large cancellations in the numerator, the irritating presence of the cutoff δ has been avoided and this must be viewed as a significant improvement.

4. Low x physics

The rise in the F_2^{ep} structure function at low x is by now well known and has been thoroughly studied by the H1 and ZEUS collaborations. Rather amazingly, as shown by *A. Zarnecki* (Figs. 2 and 6), the rise persists even down to very low momentum transfer, $Q^2 \sim 1.5$ GeV^2. These observations lead to two important questions;

(1) what is the behaviour of F_2 as $x \to 0$,

(2) how does F_2 behave as $Q^2 \to 0$.

While the $Q^2 = 0$ photoproduction cross section rises slowly with s, the slope of F_2 is rather steeper and some new dynamics must enter in between $Q^2 = 0$ and $Q^2 = 1.5$ GeV^2. This is undoubtedly a very 'hot' topic, but, following the flavour of the talks presented here, I focus only on the first question. There are three contrasting approaches to the low x region. Here, the invariant mass of the hadronic final state $s = (1-x)Q^2/x$ becomes large and one must address the issue of whether or not to resum logarithms of the type $\log(x_0/x)$.

In the BFKL approach described by *L. Lipatov*, the cross section is dominated by reggeised gluon exchange in the t-channel. At lowest order, the cross section has the form,

$$\sigma \sim s^{\omega}, \qquad F_2 \sim x^{-\omega}, \tag{9}$$

with,

$$\omega = \frac{4\alpha_s N_c}{\pi} \log(2) \sim 0.5\,. \tag{10}$$

However, the scale at which the strong coupling is evaluated is not known so that the precise value of ω is uncertain. Furthermore, the Froissart bound appears to be violated. *L. Lipatov* reported on progress to resum the next-to-leading order logarithms which should help to pin down α_s. Whether or not this changes the $x^{-\omega}$ behaviour of F_2 is a fascinatingly open issue.

Ball and Forte [3] have proposed an experimental check of whether or not F_2 has a power like growth using the double leading log approximation. In this approximation logarithms of x_0/x and Q^2/Q_0^2 are treated equally. One of the predictions of this approximation is that $xg(x)$ grows faster than any power of $\log(1/x)$ but slower than any power of $1/x$,

$$xg \sim \exp\left(\log\left(\frac{1}{x}\right)\right). \tag{11}$$

Making the change of variables,

$$\sigma = \sqrt{\log\left(\frac{x_0}{x}\right)\log\left(\frac{Q^2}{Q_0^2}\right)}, \qquad \rho = \sqrt{\log\left(\frac{x_0}{x}\right)/\log\left(\frac{Q^2}{Q_0^2}\right)}. \tag{12}$$

Ball and Forte predict that the logarithm of a suitably scaled F_2 is linear in σ *iff* double asymptotic scaling behaviour is present. *A. Zarnecki* showed recent H1 data [4] that impressively confirms that this is so for $Q_0^2 = 2.5$ GeV^2 and $\rho \geq 2$. Clearly, if x gets small enough, the logarithms of x must dominate but it is nevertheless an interesting question as to where this simple approach breaks down.

Finally, although it is clear that the small x region must be dominated by logarithms of x, it is still not clear that either of the above approaches are necessary for currently accessible values of x. For example, applying the naive DGLAP evolution in Q^2 with a 'steep' input at some starting Q_0^2,

$$xg \sim x^A(1-x)^B(1+C\sqrt{x}+Dx), \tag{13}$$

with $A < 0$ may still describe the data well over the whole observable range of x and Q^2. As shown by *A. Zarnecki* (Figs. 1 and 2), the combined ZEUS and H1 data still shows the expected scaling violation. Unravelling the physics at low x and determining which dynamical effect really causes the rise in F_2 is one of the more important and interesting questions facing physicists at HERA.

5. The strong coupling constant

The generic strong interaction observable O can be described in the following way,

$$O = A + B\alpha_s + C\alpha_s^2 + \ldots + \frac{m^2}{Q^2} + \ldots, \tag{14}$$

where the first set of terms form part of a perturbative series and the second are non-perturbative (or higher twist) corrections. In extracting the strong coupling, usually the $1/Q^2$ terms are ignored and the polynomial in α_s directly solved. However, all extractions of α_s suffer from one or more of the following problems,

(a) large renormalisation scale dependence
(b) Q too small or large $1/Q^2$ corrections
(c) insufficient orders calculated
(d) confusion with electroweak radiative corrections
(e) confusion with parton density functions.

Until recently, measurements of $\alpha_s(M_z)$ have divided into two groups; those obtained at low Q^2 largely from Deep Inelastic Scattering and Lattice Gauge Theory favouring $\alpha_s(M_z) \sim 0.112$ and those at high Q^2 dominated by LEP with $\alpha_s(M_z) \sim 0.123$. The generally accepted world average lay between the two sets,

$$\alpha_s^{WA}(M_z) = 0.117 \pm 0.005\,, \tag{15}$$

with an estimated error encompassing the two extremes. Recently, some of the more extreme measurements have migrated towards the average and were reported here by *P. Burrows*. First, after extensive work on the energy calibration of their detector, the CCFR collaboration have obtained the preliminary value [5],

$$\alpha_s(M_z) = 0.119 \pm 0.0015 \pm 0.0035 \pm 0.004, \tag{16}$$

from Deep Inelastic Scattering measurements of xF_3 and F_2 at low and moderate Q^2. Second, it has proved possible to extract the strong coupling from Lattice Gauge Theory by computing the upsilon spectrum. Here the main developments have been the use of dynamical light fermions and a better understanding of how the lattice coupling is related to $\alpha_s^{\overline{MS}}$. Preliminary results presented this summer [6] give a slightly increased average value of,

$$\alpha_s(M_z) = 0.117 \pm 0.003\,. \tag{17}$$

The third and possibly best measurement comes from global electroweak fits of LEP observables. Recent improvements in the Z width and the Born cross section lead to the value presented here by *M. Pepe-Altarelli*,

$$\alpha_s(M_z) = 0.120 \pm 0.003 \pm 0.002\,. \tag{18}$$

The final world average quoted by *Burrows* is still around 0.118 with an error of possibly 5%.

On the theoretical front, one might hope for a better theoretical description of the observable O. Hadronic event shapes have been extensively used to extract α_s. However, although the NLO corrections have been known for a long time, a large scale uncertainty remained and the value of α_s extracted from each event shape differed greatly. A major improvement was the development of techniques for summing the leading and next-to-leading infrared logarithms, resulting in a measurement of

$$\alpha_s(M_z) = 0.122 \pm 0.007\,. \tag{19}$$

Work towards the α_s^4 NNLO corrections has been in progress for some time, and some results involving the pentagon one-loop contribution to $e^+e^- \to$ 4 partons are now available [7]. However, many technical problems remain, not least the evaluation of the two loop box graph with massless internal and external legs as well as the systematic isolation of singularities when two real external particles are unresolved. Nevertheless, the timescale is right for the Next Linear Collider and a 3% error on α_s measured from hadronic event shapes should be possible.

6. Single jet inclusive E_T distribution

The production of jets at 90° to the beam direction has been extensively studied by CDF and D0 collaborations at transverse energies up to almost half the beam energy and was discussed by *H. Piekarz* (Figs. 6 and 7). At moderate transverse energies $E_T \leq 200$ GeV, the experimental data shows good agreement with the NLO QCD prediction. However, at larger transverse energy CDF find an experimental excess for $E_T > 200$ GeV, while the D0 data is roughly in line with the theoretical expectations. In both cases, the statistical and systematic errors are large. Nevertheless, the CDF data shows a significant kink.

One can ask whether the theoretical prediction is reliable. In fact, the ratio $\sigma(\mathrm{NLO})/\sigma(\mathrm{LO})$ is roughly constant between 50–500 GeV and only falls off around $E_T \sim 600$ GeV as the jet rapidity cuts hit the phase space boundaries. For the specific choice of renormalisation and factorisation scale $\mu = 0.5E_T$, this ratio is 1.10 which indicates that that large unresummed logarithms are not present. Changing the input value of α_s alters the slope of the theoretical prediction. Although the parton level cross section increases with α_s, the parton luminosity is decreased at high E_T due to a faster DGLAP evolution. Alternatively, altering the choice of the renormalisation and factorisation scales alters the normalisation of the theoretical prediction. Either way, there is no structure visible in the $E_T \sim$ 100–400 GeV range.

Another important input is the choice of parton density functions. How uncertain are the parton distributions at large x? For $E_T \sim 300$ GeV, coresponding to a parton fraction of $x \sim 0.3$, roughly 80% of the cross section is generated by $q\bar{q}$ scattering while 20% comes from the qg subprocess. Therefore, a 20% increase in the total theoretical jet cross section could be generated by a 10% increase in the quark distributions. However, the quarks at large x are heavily constrained by measurements of F_2 by BCDMS and it is not possible to accommodate both the DIS and jet data simultaneously by changing the quark density functions [11]. Alternatively, a 100% increase in the gluon distribution could account for the observed excess. The canonical parameterisation of xg given by equation (13) does not have enough flexibility in the $x \sim 0.3$ region. The CTEQ collaboration [12] have found that with a parametric form,

$$xg \sim x^A(1-x)^B \exp(C + Dx), \tag{20}$$

it is possible to increase the predicted jet cross section and still be consistent with the constraints from direct photon production from WA70.

In fact the gluon density is still rather poorly constrained. At low x ($\leq 10^{-2}$) HERA is starting to play a valuable role in determining the gluon density as shown by *A. Zarnecki* (Fig. 3). At larger x, the constraint is from prompt photon data where the Compton scattering process $qg \to q\gamma$ is important. However, *M. Zielinski* reminded us that the agreement between many different experiments and NLO QCD is systematically awful [13]. In all cases, the experiment shows an excess over theory at the lower end of the probed p_T range. One possible explanation is an intrinsic smearing of the initial parton transverse energy due to multiple gluon emission. CTEQ have shown [13] that the CDF direct photon data can be described with an average k_T smearing of $\langle k_T \rangle \sim 3$ GeV while the E706 data presented here by *M. Zielinski* (Figs. 4 and 5) can be accommodated with $\langle k_T \rangle \sim 1.4$ GeV. Baer and Reno have now successfully incorporated this smearing into a NLO QCD Monte Carlo [14]. However, such smearing would also have to be present for other observables such as the single jet transverse energy distribution. There is also an uncertainty in the fragmentation functions $D_{q\to\gamma}$ and $D_{g\to\gamma}$ which are only experimentally relevant at large z where the photon is essentially isolated from other hadrons. Soft gluon emission from the parent quark can spoil the isolation and there is some evidence that resummation of logarithms of $(1-z)$ may be necessary. In addition, there is some confusion over the 'theological issue' of the power counting of α_s and the construction of fixed order calculations; is the fragmentation function $\mathcal{O}(\alpha)$ or $\mathcal{O}(\alpha/\alpha_s)$? In any event, this is an area where theoretical progress needs to be made, particularly in light of the plentiful high statistics data that will soon be available.

One final thought on the single jet inclusive excess is the dependence on the conesize. In both CDF and D0, the jets have been constructed so that calorimeter cells lying within $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2} = 0.7$ of each other are combined to form the jet. However, as E_T increases, α_s decreases and the perturbative size of the jet narrows. It would therefore not be unreasonable to reduce the experimental conesize. A hint that it may be important to vary the conesize with E_T (or that the out of cone corrections may not be well understood) comes from the new data at $\sqrt{s} = 630$ GeV where both CDF [8] and D0 [9] find an experimental depletion compared to NLO QCD. This is to be contrasted with the experience of UA2 (with a larger conesize of 1.3) who found good agreement with the NLO QCD cross section [10]. On the other hand, in the perturbative calculation, the dependence on the conesize comes from combining two partons together. This is effectively a leading order effect and it may be that the theoretical calculation is insufficiently accurate.

7. Summary

As we have seen in the past few days, the strong interaction is, in general, very successfully described by QCD. Recent technical achievements are very impressive, while the experiments are beginning to probe new and interesting areas. The situation can be described by the following time evolution equation for information concerning the strong interaction which depends on the theoretical input T, the experimental results E and the phenomenological coupling between the two p,

$$\frac{\partial I_s}{\partial t} \sim \left(\frac{p}{2\pi}\right) T \otimes E \gg 0. \tag{21}$$

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