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THEORY OF THE W/Z TRANSVERSE MOMENTUM DISTRIBUTION

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ABSTRACT

QCD predictions for the inclusive production of W and Z bosons with fixed transverse momentum Q_T in $p\bar{p}$ collisions at the Tevatron are reviewed in the light of recently completed calculations of the full $\mathcal{O}(\alpha_s^2)$ QCD radiative corrections.

1. INTRODUCTION

The Standard Model makes rather precise predictions for the production of W and Z bosons with invariant mass-squared $Q^2 = M_W^2$ or M_Z^2 as a function of transverse momentum Q_T at Tevatron energies, $\sqrt{S} \approx 2$ TeV. The total cross sections are $\sigma_W \approx 15$ nb and $\sigma_Z \approx 5$ nb. The typical momentum fractions of the annihilating partons in the underlying Drell-Yan mechanism are of order $x \approx \sqrt{\tau} = \sqrt{(Q^2/S)} \approx 0.04$. At these x values, gluons contribute roughly 30% of the events at large Q_T via the Compton process $qG \to W + X$: this is a dramatic qualitative difference from CERN Collider energies at which the dominant subprocess is $q\bar{q}$ annihilation. Thus the large Q_T distribution is a potentially useful probe of the gluon density in the proton. Since large values of Q_T probe short distances, the Q_T distribution is a good place to look for new physics. Possible signatures include anomalously large numbers of events with lepton + jet(s) + missing transverse energy ($\not{E}_T$) in the W distribution, and monojets or jets + $\not{E}_T$ in the Z distribution.

Perturbative QCD may be used to obtain precise testable predictions [1,2] for $d\sigma/dQ_T$ in the range 20 GeV $\lesssim Q_T \lesssim$ 250-300 GeV: the upper limit values are determined by Tevatron luminosity, and the lower limit is roughly the momentum at which large infrared logarithms of Q^2/Q_T^2 begin to spoil the convergence of the perturbation series. In the range 2 GeV $\lesssim Q_T \lesssim$ 20 GeV these infrared logarithms must be summed to all orders in α_s to yield a "Sudakov" form factor which causes the distribution to peak at $\approx$ 10 GeV. Theory seems to indicate [3] that this Sudakov region is cleanly separated from the region below $\approx$ 2 GeV where perturbatively intractable parton intrinsic momentum effects dominate.

2. CHOOSING RENORMALIZATION AND FACTORIZATION SCALES

The Lagrangian of QCD depends in principle on only one parameter $\Lambda \approx 0.2$ GeV if quark masses and the CP-violating θ parameter can be ignored in making predictions at high energies. In practise, the wave function of the proton cannot be calculated with sufficient accuracy to allow predictions to be made at high energies, and one must rely on the QCD-improved parton model formula

$$\frac{d\sigma}{dQ_T} = \sum_{ab} \int dx_a dx_b \, f_a(x_a, M^2) f_b(x_b, M^2) \frac{d\sigma^{ab}}{dQ_T^2}(\alpha_s(\mu^2), \mu^2, M^2) \,, \tag{1}$$

where σ^{ab} is the parton level distribution for partons with momentum fractions $x_{a,b}$ from which ultraviolet and collinear infrared divergences have been subtracted. This formula

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depends on the physical parameter Λ in an explicitly calculable way through the effective coupling

$$\frac{\alpha_s(\mu^2)}{2\pi} = \frac{6}{(33-2n_f)\ln(\frac{\mu^2}{\Lambda^2})} - \frac{\beta_2 \ln\ln(\frac{\mu^2}{\Lambda^2})}{\beta_1^3 \ln^2(\frac{\mu^2}{\Lambda^2})} , \tag{2}$$

as well as implicitly in an incalculable way through the parton densities $f_a(x, M^2)$. The renormalization scale μ is a remnant of the process of discarding ultraviolet divergent contributions "at scale μ" according to some arbitrary prescription (taken to be $\overline{\mathrm{MS}}$ renormalization in what follows). The factorization scale M is likewise a remnant of the process of discarding collinear singularities according to some arbitrary prescription (taken to be $\overline{\mathrm{MS}}$ factorization in what follows). In principle, the parameters μ and M are unphysical, unrelated to one another, and not specified by the theory in terms of any measured energy scale E_s (which can for example be chosen to be Q_T, $\sqrt{Q^2}$, $\sqrt{Q^2+Q_T^2}$ or $\sqrt{S}$). Indeed, proofs of renormalizability, and factorization theorems imply that

$$\mu\frac{\partial}{\partial\mu}\left(\frac{d\sigma}{dQ_T}\right)_{\text{exact}} = 0 , \qquad M\frac{\partial}{\partial M}\left(\frac{d\sigma}{dQ_T}\right)_{\text{exact}} = 0 . \tag{3}$$

In practise, this exact scale independence is spoiled if the perturbation series in α_s is truncated at any finite order. This unfortunate fact is signalled by the appearance in σ^{ab} of powers of $\ln(\mu^2/E_s^2)$ and $\ln(M^2/E_s^2)$ which are not precisely compensated by the μ dependence of α_s and the M dependence of $f(x, M^2)$. While the potentially destabilizing logarithms can be made small by choosing

$$\zeta_\mu \equiv \frac{\mu^2}{E_s^2} \approx 1 , \qquad \zeta_M \equiv \frac{M^2}{E_s^2} \approx 1 , \tag{4}$$

the theory makes absolutely no more precise determination of $\zeta_{\mu,M}$. This spurious scale dependence is of course precisely cancelled by higher (uncalculated) orders of perturbation theory. The logically optimal (but obviously unworkable) solution would be to choose scales which minimize the remainder of the perturbation series. Various strategies for guessing the unknown remainder have been suggested. Two such strategies will be discussed here. The first is (a simplified version of) Stevenson's "Principle of Minimal Sensitivity" [4] in which the unphysical scales are fixed at a local extremum or saddle point of $d\sigma/dQ_T$ in the variables ζ:

$$\zeta^{\mathrm{SP}}_{\mu,M}(Q_T): \ \zeta_\mu\frac{\partial}{\partial\zeta_\mu}\left(\frac{d\sigma}{dQ_T}\right)_{\mathrm{NLO}} = 0 , \quad \zeta_M\frac{\partial}{\partial\zeta_M}\left(\frac{d\sigma}{dQ_T}\right)_{\mathrm{NLO}} = 0 , \tag{5}$$

where SP stands for saddle point and the subscript NLO signifies that the parton level cross section is computed through next-to-leading order in perturbation theory and the parton densities have likewise been evolved using the next-to-leading order Altarelli-Parisi equations. Two remarks may be made concerning this choice of scales: First, there is the question of the existence and uniqueness of a saddle point. The answer to this question is not known, but in practise, there does usually seem to exist a fairly well defined saddle point at physically reasonable values of ζ. Second, even if a unique saddle point does exist, there is no guarantee that using the saddle point scales yields the best estimate of the uncalculated remainder of the series. The NLO cross section at the saddle point has the same derivatives with respect to ζ as the exact cross section, but not necessarily the same magnitude. A second strategy for choosing the unphysical scales ζ may be called the effective scales (ES)

prescription, which is a simplified version of Grunberg's "Method of Effective Charges" [5]. One would like to be able to choose the scales such that the NLO approximation is equal in magnitude to the exact distribution. In the absence of any information about the remainder of the series, we choose scales such that the leading order (LO) cross section is equal in magnitude to the cross section computed through next-to-leading order:

$$\zeta^{\rm ES}_{\mu,M}(Q_T): \quad \left(\frac{d\sigma}{dQ_T}\right)_{\rm LO} = \left(\frac{d\sigma}{dQ_T}\right)_{\rm NLO} . \tag{6}$$

This equation determines a line of possible effective scale choices in the $\zeta_\mu - \zeta_M$ plane.

The choice of scales was studied by Bawa and Stirling [6] for the non-singlet (i.e., valence quark) contributions to the Q_T distribution at $\sqrt{S} = 630$ GeV. They chose $E_s = Q_T$ and set $\zeta_M = \zeta_\mu$, and found that the one dimensional analogs of Eqs. 5 and 6 yielded $\zeta^{SP} \simeq \zeta^{ES} \simeq 0.16$ for $Q_T > 15$ GeV. The situation at $\sqrt{S} = 1.8$ TeV, where gluon initiated contributions cannot be neglected is more complicated. It was found in Refs. [1,2] that if ζ_M and ζ_μ were set equal, then the one dimensional version of Eq. 5 had no physically acceptable solution for $Q_T \lesssim 80$ GeV. We have since studied this problem [7] by allowing the two scales to vary independently of one another, as has been done for prompt photon production by Aurenche et al., [8]. We find that there is a well defined saddle point in the $\zeta_M - \zeta_\mu$ plane. Results of this analysis for W production and using the parton densities of Martin, Roberts and Stirling are shown in Fig. 1. The saddle point scales are seen to move away from the diagonal $\zeta_M = \zeta_\mu$ line for lower values of Q_T. It is interesting that there is an apparent clustering of ES choices at $\zeta_M \approx \zeta_\mu \approx .1$, and that the SP scales move toward this region for large Q_T.

What is the correct choice of scales? Fig. 2 illustrates what we believe is the answer to this question. In this figure are plotted the leading order (LO) and the full (NLO) distributions evaluated at the naive choice of scales $\zeta_M = \zeta_\mu = 1$, and at the saddle point scales. It is apparent that the LO distribution is quite sensitive to variations in scale. Inclusion of the radiative corrections considerably decreases this sensitivity and makes the choice of scales irrelevant! If experimental errors were to become smaller than the difference between the NLO curves in Fig. 2, the only rigorous way to improve the theoretical prediction would be to compute the $\mathcal{O}(\alpha_s^3)$ radiative corrections, thus further narrowing the difference and rendering a choice of scales unnecessary at this increased resolution.

3. OTHER SOURCES OF THEORETICAL UNCERTAINTY

The main emphasis of this review has been on renormalization and factorization scale ambiguities. In this section we briefly discuss two important additional theoretical issues.

Parton densities

The most significant contribution to the theoretical uncertainty at large Q_T arises from different choices of input parton densities. For consistency, next-to-leading order parton density parametrizations should be used in conjunction with the next-to-leading parton level cross sections of Refs. 1,2. Ref. 1 presents results using the parametrizations of Martin, Roberts and Stirling [9]. Cross sections derived using the MRSB and MRSE sets typically differ from one another by less than 20% for all $Q_T > 20$ GeV. Refs. 2 present results using the parametrizations of Diemoz et al. The different DFLM sets yield distributions that agree to within $\pm 15\%$. Taking both analyses into account, QCD with 5

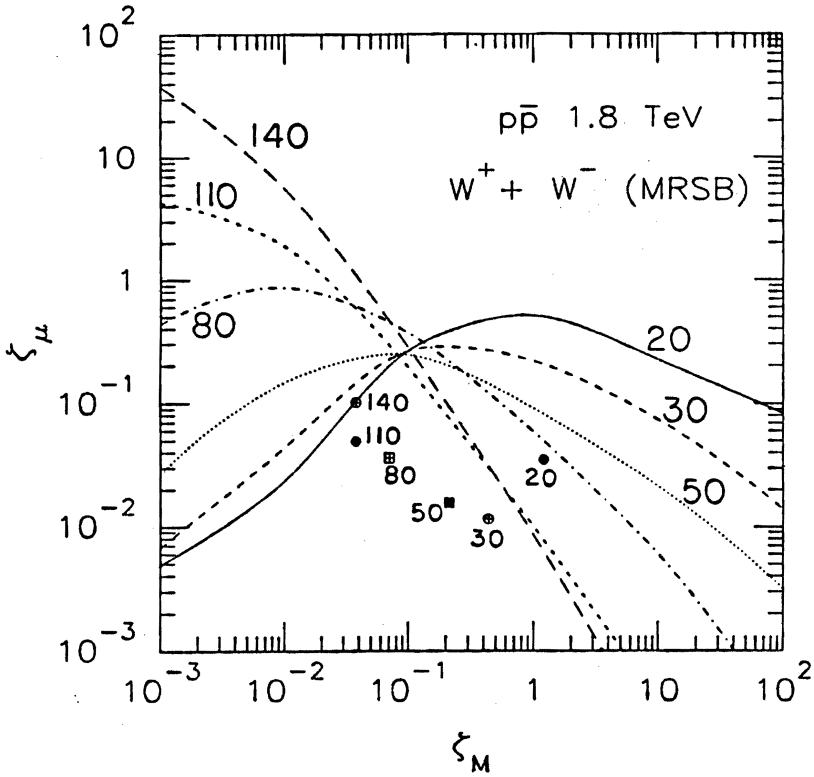

Fig. 1.

Numerical labels represent Q_T values in GeV. The points represent saddle point scales defined by Eq. 5, and the curves represent effective scales defined by Eq. 6.

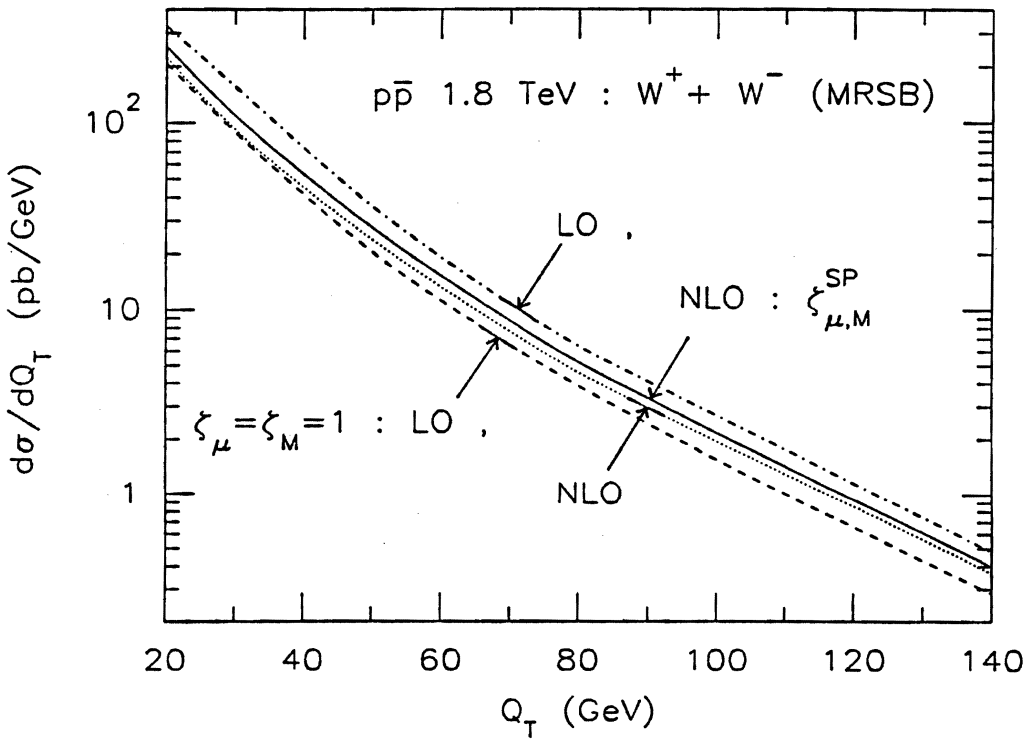

Fig. 2. Including $\mathcal{O}(\alpha_s^2)$ contributions (NLO) makes a choice of scales unnecessary.

light flavors yields absolute predictions for the Q_T distribution with $Q_T > 20$ GeV with a conservative uncertainty of no more than about 20-25%.

Sudakov logarithms

The predictions of QCD for $Q_T \lesssim 20$ GeV are less certain on account of the occurrence in the perturbation series of terms of the type $\alpha_s^n \ln^m(Q^2/Q_T^2)/Q_T^2$, $m \leq 2n-1$. These terms arise from gluon radiation off the initial state partons. They tend to suppress the emission of a W with very small Q_T, and in fact drive the cross section negative. This unphysical result must be dealt with by summing these logarithms to all orders. It has been shown, principally by Collins and Soper, that this can be done consistently [10]. Various perturbative coefficients required for this analysis have been computed through $\mathcal{O}(\alpha_s^2)$ for the non-singlet distribution [11], and a prescription for matching the exponentiated logarithms at small Q_T with the perturbative large Q_T tail has been given [3]. The principal result of a rather involved analysis is that the distribution should peak at $Q_T \approx 10$ GeV. However, a full analysis including gluon initiated events has not been performed.

4. CONCLUSIONS

QCD makes firm and reliable predictions for the W and Z transverse momentum distributions at large Q_T. Any deviation from these predictions would constitute a firm signal for new physics beyond the Standard Model. The inclusion of the full $\mathcal{O}(\alpha_s^2)$ QCD radiative corrections makes the choice of unphysical renormalization and factorization scales almost irrelevant. Accurate experimental measurement of the Q_T distribution can provide a means of probing the quark and especially the gluon densities in the proton. Further theoretical work needs to be done on summing Sudakov logarithms in the region of small $Q_T \lesssim 20$ GeV.

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REFERENCES

1. R.J. Gonsalves, J. Pawłowski and C.-F. Wai, Phys. Rev. D **40**, 2245 (1989).
2. P.B. Arnold and M.H. Reno, Nucl. Phys. **B319**, 37 (1989); P. Arnold, R.K. Ellis and M.H. Reno, Phys. Rev. D **40**, 912 (1989).
3. G. Altarelli et al., Nucl. Phys. **B246**, 12 (1984); Z. Phys. C **27**, 617 (1985).
4. P.M. Stevenson, Phys. Rev. D **23**, 2916 (1981); P.M. Stevenson and H.D. Politzer, Nucl. Phys. **B277**, 758 (1986).
5. G. Grunberg, Phys. Rev. D **29**, 2315 (1984).
6. A.C. Bawa and W.J. Stirling, Phys. Lett. **203B**, 172 (1988).
7. R.J. Gonsalves, J. Pawłowski and C.-F. Wai, UB-HET 90/3, submitted to Phys. Lett.
8. P. Aurenche, R. Baier, M. Fontannaz and D. Schiff, Nucl. Phys. **B286**, 509 (1987).
9. A.D. Martin, R.G. Roberts and W.J. Stirling, Phys. Rev. D **37**, 1161 (1988).
10. J.C. Collins, D.E. Soper and G. Sterman, Nucl. Phys. **B250**, 199 (1985).
11. C.T.H. Davies, W.J. Stirling and B.R. Webber, Nucl. Phys. **B256**, 413 (1985); J. Kodaira and L. Trentadue, Phys. Lett. **112B**, 66 (1982).