PHYSICAL REVIEW D VOLUME 25, NUMBER 3 1 FEBRUARY 1982

Identification of W bosons in $\bar{p}p$ collisions: A detailed study

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(Received 8 September 1981; revised manuscript received 28 December 1981)

We study the decay-lepton distribution of W bosons produced in high-energy $\bar{p}p$ collisions $\bar{p}p \to W^{\pm}X \to l^{\pm}X$. The accuracy that one can reach in calculating the shape of the lepton momentum distribution determines, along with the statistics of the experimental data, the accuracy of measuring M_W. Our study shows that in the experimentally relevant region, where the lepton's transverse momentum $p_T \simeq M_W/2$, the calculation is dominated by the behavior of the cross section for producing the W at small transverse momentum Q_T, e.g., $Q_T \simeq 5$ GeV for $\sqrt{s} = 540$ GeV. As $Q_T << M_W$, $O(\alpha_s)$ calculations are unreliable, in fact irrelevant. We calculate the Q_T distribution of the W to all orders in α_s using an impact-parameter formalism to sum the leading double logarithms. We present the lepton transverse-momentum distribution for $\bar{p}p$ collider energies of $\sqrt{s} = 540$ GeV, and 2 TeV. Further attention is devoted to the dependence of the signal on the center-of-mass angle θ, to the calculation of the prompt-lepton background, and to the relevance of identifying the missing energy associated with the neutrino from $W \to l\nu$.

It is widely believed[1] that the best chance to identify $W^{\pm}$ bosons in high-energy $\bar{p}p$ collisions will be through the Jacobian peak (or enhancement) in the transverse-momentum spectrum of the decay lepton in the neighborhood of $p_T \simeq M_W/2$. Although the observation of such an enhancement would be a clear signature of the W boson, there are good reasons to study the p_T spectrum in detail. Indeed, the accuracy that one can reach in calculating the shape of the distribution will determine, along with the experimental statistics, the accuracy of a measurement of M_W. Theoretical studies of this distribution are particularly important. It has recently been realized[2] that precise measurements of the $W^{\pm}$ and Z masses provide an opportunity to probe the gauge structure of the Weinberg-Salam model with $\bar{p}p$ colliders. Indeed, the one-loop corrections shift the well-known lowest-order predictions for M_W and M_Z by amounts of the order of 5%. Whereas measuring the mass of $Z \to l^+l^-$ is an experimental problem, measurement of the mass of $W \to l\nu$ is first a theoretical one. Knowledge of the momentum distribution of the observed lepton from the decay $W \to l\nu$ is required and calculations to $O(\alpha_s)$ in QCD have been performed.[3,4]

The crucial feature of such a calculation is the computation of the transverse-momentum (Q_T) distribution of the produced W boson, $\bar{p}p \to WX$. Our results will show that in the experimentally accessible region of phase space, where the lepton's momentum $p_T \simeq M_W/2$, the calculation is dominated by the behavior of the W production cross sections near $Q_T \simeq 0$. At $\sqrt{s} = 540$ GeV, W's with $Q_T \simeq 5$ GeV predominantly determine the behavior of the lepton decay distribution. $O(\alpha_s)$ calculations of the W transverse-momentum distribution are divergent here; they can only be trusted when $Q_T{}^2 \simeq M_W{}^2$. Although they were a qualitative guide to the important effects[3,4] of gluon emission in the production of weak bosons, they never addressed the problem where it experimentally counts, that is, for Q_T in the few-GeV region.

COMPUTATION OF THE W'S TRANSVERSE-MOMENTUM DISTRIBUTION

It is very easy to understand why an $O(\alpha_s)$ calculation inevitably breaks down when $Q_T{}^2 << M_W^2$. Consider a typical diagram [Fig. 1(a)]. It assumes that the antiquark $\bar{q}$ radiates a single gluon, not 2, 3,. . ., before annihilating with the quark q to form the W. This is a good assumption when

$Q_T{}^2 \simeq M_W{}^2$, but totally unrealistic when $Q_T{}^2 << M_W{}^2$. A reliable answer should allow multiple-gluon emission, i.e., sum all orders of perturbation theory in α_s. Progress[5–9] in the study of perturbative QCD in the region $Q_T{}^2 << M_W{}^2$ has fortunately made such a calculation possible. We proceed to sketch how this is done. In the limit $Q_T{}^2 << M_W{}^2$ the cross section for emitting a single gluon [diagram of Fig. 1(a)] is given by

$$\frac{Q_T{}^2}{\sigma_0}\frac{d\sigma}{dQ_T{}^2} = \left[\frac{\alpha_s}{\pi}\right] C_F \ln\left[\frac{M_W{}^2}{Q_T{}^2}\right] . \tag{1}$$

Here σ_0 is the (Q_T-integrated) Drell-Yan cross section for $\bar{q}q \to W$ and $C_F = \frac{4}{3}$ in QCD. The cross section for emitting two gluons [Fig. 1(b)] is given by

$$\frac{1}{\sigma_0}\frac{d\sigma}{dQ_T{}^2} = \frac{\pi}{2}\left[\frac{\alpha_s C_F}{\pi}\right]^2 \int \frac{d^2k_{T1}}{\pi k_{T1}{}^2} \int \frac{d^2k_{T2}}{\pi k_{T2}{}^2} \ln\left[\frac{k_{T1}{}^2}{M_W{}^2}\right] \ln\left[\frac{k_{T2}{}^2}{M_W{}^2}\right] \delta^2(\vec{k}_{T1} + \vec{k}_{T2} + \vec{Q}_T) . \tag{2}$$

In impact-parameter form Eq. (2) can be written

$$\frac{1}{\sigma_0}\frac{d\sigma}{dQ_T{}^2} = \frac{1}{8\pi}\left[\frac{\alpha_s C_F}{\pi}\right]^2 \int d^2b \int \frac{d^2k_{T1}}{\pi k_{T1}{}^2} \int \frac{d^2k_{T2}}{\pi k_{T2}{}^2} \ln\left[\frac{k_{T1}{}^2}{M_W{}^2}\right] \ln\left[\frac{k_{T2}{}^2}{M_W{}^2}\right] e^{i(\vec{k}_{T1} + \vec{k}_{T2} + \vec{Q}_T)\cdot\vec{b}} . \tag{3}$$

The k_{Ti} integrals now factorize and one can sum n-gluon emissions [Fig. 1(c)]. The result is

$$\frac{1}{\sigma_0}\frac{d\sigma}{dQ_T{}^2} = \frac{1}{4\pi}\int d^2b e^{i\vec{Q}_T\cdot\vec{b}} e^{\Delta(b)}$$
$$= \tfrac{1}{4}\int db^2 e^{\Delta(b)} J_0(Q_T b) \tag{4}$$

with

$$\Delta(b) = \frac{C_F}{\pi}\int \frac{d^2k_T}{\pi k_T{}^2}\alpha_s(k_T{}^2)\ln\frac{k_T{}^2}{M_W{}^2}[e^{i\vec{k}_T\cdot\vec{b}} - 1] \tag{5}$$

$$= \frac{C_F}{\pi}\int \frac{dk_T{}^2}{k_T{}^2}\alpha_s(k_T{}^2)\ln\frac{k_T{}^2}{M_W{}^2}[J_0(k_T b) - 1] . \tag{6}$$

In Eqs. (5) and (6) we allow α_s to run; the 1 in the square brackets takes care of the virtual-gluon diagrams, not shown in Fig. 1. This calculation correctly sums soft-gluon-emission terms of the form $[\alpha_s \ln^2(M_W{}^2/Q_T{}^2)]^n$ following the work of Dokshitser, D'Yakonov, and Troyan[5] and others.[6–8] The dominant contributions are coming from the region of phase space in which one gluon has transverse momentum of the order of Q_T and the others have negligible transverse momentum. However, when $Q_T \to 0$, this is no longer the case; then the main contributions are those in which two or more gluons have much larger transverse momenta that add vectorially to give the small Q_T. To take account of this, we have followed Parisi and Petronzio[6] and in Eq. (3) summed the leading logarithms in impact-parameter (b) space and then finally Fourier transformed back to Q_T space (see also Refs. 8 and 9). Thus for our result $Q_T{}^2 << M_W{}^2$, the transverse-momentum distribution of the W is

$$G(Q_T) \equiv \frac{M_W{}^2}{\sigma_0}\frac{d\sigma}{dQ_T{}^2} = \frac{M_W{}^2}{4}\int db^2 e^{\Delta(b)} J_0(bQ_T) , \tag{7}$$

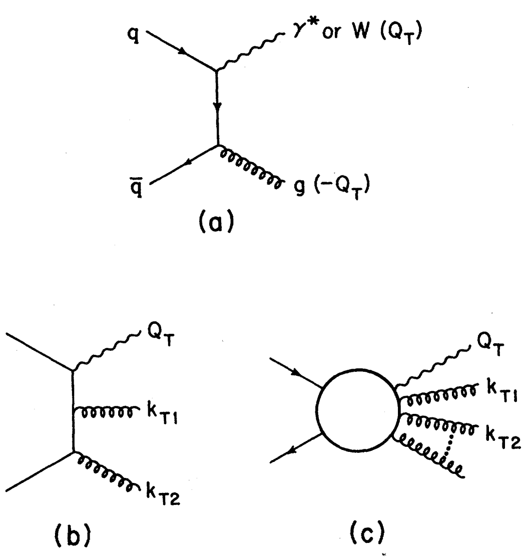

FIG. 1. (a) Example of $O(\alpha_s)$ graph contributing to the production of W's with transverse momentum $Q_T \neq 0$. (b) Same as Fig. 1(a) but $O(\alpha_s{}^2)$ emission of two gluons. (c) Emission of multiple gluons.

$$\Delta(b)=\frac{16}{33-2N_f}\int_{\Lambda^2}^{M_W{}^2}\frac{dk_T{}^2}{k_T{}^2}\frac{\ln(M_W{}^2/k_T{}^2)}{[\ln^2(k_T{}^2/\Lambda^2)+\pi^2]^{1/2}}\times[J_0(bk_T)-1]\ , \qquad (8)$$

where N_f is the number of quark flavors. Our particular form in the denominator of Eq. (8) results from the use of $|\alpha_s(k_T{}^2)|$ rather than α_s (as argued in Ref. 10) and makes $\Delta(b)$ insensitive to the choice of the lower integration limit. We will return to this point later. The transverse-momentum distribution of the W, calculated according to Eqs. (7) and (8), is shown in Fig. 2. Throughout the paper we choose $N_f=4$ in α_s $(k_T{}^2)$ for Eq. (8), since the major contribution to the integration comes typically from the region $k_T\simeq 5$ GeV. For comparison, we also show the result of an $O(\alpha_s)$ calculation which, as expected, is divergent in the $Q_T{}^2<<M_W{}^2$ regions shown in Fig. 2. It should be pointed out that there are ambiguities in the explicit form of Eq. (8) as we approach $Q_T{}^2\simeq M_W{}^2$, but they are of no concern to us, since in these regions the W production cross section is small, as we will explicitly show further on. There the $O(\alpha_s)$ calculations[3,4] give the right answer.

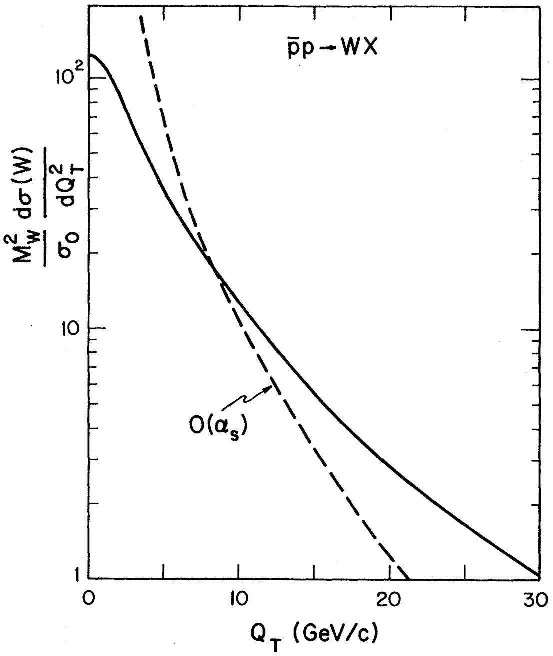

FIG. 2. The transverse-momentum (Q_T) distribution of the W, as defined in Eq. (7), calculated with $\Lambda=0.5$ GeV and $N_f=4$. The distribution for the Z is numerically very similar. Also shown (dashed line) is the $O(\alpha_s)$ result.

LEPTON DECAY SPECTRUM

The calculation of the lepton distribution from a known W distribution is straightforward.[11] Neglecting the lepton mass

$$p^0\frac{d\sigma}{d^3p}=\int\frac{d^3Q}{Q^0}[F(x_1,x_2)G(Q_t)]\times[\frac{B_l}{2\pi}\delta(Q\cdot p-\tfrac{1}{2}M_W{}^2)]\ , \qquad (9)$$

where the first term in square brackets is the Lorentz-invariant cross section for producing a W with momentum $Q(Q^2\equiv M_W{}^2)$, and the second term describes its two-body decay into $l\nu$ (with branching ratio B_l). The integration runs over all kinematic configurations Q which can yield a decay lepton of a given moemntum p. $G(Q_T)$ is given by Eqs. (7) and (8) and $F(x_1,x_2)$ represents the usual Drell-Yan cross section for producing a W from annihilation of a quark-antiquark pair with fractional longitudinal momenta x_1, x_2. For N_f quark flavors ($N_f/2$ doublets), the W^+ cross section is given by

$$F(x_1,x_2)=\frac{\pi\alpha}{3M_W{}^4\sin^2\theta_W}\left[d(x_1)u(x_2)\tfrac{3}{4}(1+\cos\hat{\theta})^2+\frac{N_f}{2}s(x_1)s(x_2)\tfrac{3}{4}(1-\cos\hat{\theta})^2+\left[\frac{N_f}{2}-1\right]s(x_1)s(x_2)\tfrac{3}{4}(1+\cos\hat{\theta})^2\right]\ , \qquad (10)$$

where u, d include valence and sea u, d contributions, and where $\hat{\theta}$ specifies the lepton direction with respect to the incident $\bar{p}$ direction in the rest frame of the W.[12] We assume that the sea distribution $s(x)$ is the same for all quarks. Equation (10) exhibits the $V-A$ structure of the $Wq\bar{q}$ coupling. It is useful to notice that

$$\sigma(p\bar{p}\to W^+X)(\theta)=\sigma(\bar{p}p\to W^-X)(\pi-\theta)\ . \qquad (11)$$

To transform the decay distributions from the $\bar{q}q$ to the $\bar{p}p$ center-of-mass frame, we note that

$$1+\cos\hat{\theta}=\frac{2p_T}{x_1\sqrt{s}}\frac{1+\cos\theta}{\sin\theta}\ , \tag{12}$$

$$1-\cos\hat{\theta}=\frac{2p_T}{x_2\sqrt{2}}\frac{1-\cos\theta}{\sin\theta}\ , \tag{13}$$

where θ is the angle at which the decay lepton is produced in the $\bar{p}p$ frame; it is also measured relative to the $\bar{p}$ direction. We use the quark structure functions of Owens and Reya[13] evolved in Q^2 to $M_W{}^2$. In a similar way we calculate the lepton distribution from Z^0 decay; $\bar{p}p\to Z^0X\to lX$. The corresponding form of $F(x_1, x_2)$ can be found, for example, in Ref. 14.

We present results for $\sin^2\theta_W=0.21$ [and $N_f=6$ in Eq. (10)]. That is, we use

$$\begin{aligned} &M_W=81\text{ GeV}\ ,\quad M_Z=91\text{ GeV}\ ,\\ &B_l(W\to l\nu)=0.08\ ,\quad B_l(Z\to l\bar{l})=0.03\ . \end{aligned} \tag{14}$$

RESULTS AND DISCUSSION

In Fig. 3 we show the lepton momentum distribution for $y=0$ ($\theta=90°$) at $\bar{p}p$ collider energy $\sqrt{s}=540$ GeV. Also shown for reference is a Drell-Yan calculation obtained assuming $Q_T=0$. As anticipated, our results are significantly different not only from the Drell-Yan calculation, but also from the previous $O(\alpha_s)$ calculations.[3,4] The enhanced broadening of the Jacobian peak reflects the width of the Q_T distribution of the parent W of Fig. 2. That is, the fact that in previous $O(\alpha_s)$ calculations some structure is retained at $M_W/2$ is a consequence of the narrower width of the Q_T distribution (dashed line in Fig. 2) before one sums all orders of perturbation theory. This is, of course, due to the divergence of the $O(\alpha_s)$ calculation and is not a consequence of QCD, graphically illustrating our arguments for the necessity of summing all orders in α_s.

We checked that any scales of order 1 GeV^2 do not appreciably change the distribution; for example, Γ_W, Λ, and the intrinsic transverse momentum of the colliding quarks. This latter effect can be easily explored in the impact-parameter formalism[6] by inserting a "smearing" factor $\exp(-b^2A/4)$ into the integrand of Eq. (7) corresponding to an intrinsic transverse-momentum (p_T) distribution $\exp(-p_T{}^2/A)$ and to an average value $\langle p_T{}^2\rangle=A$. The choice of structure functions does affect the normalization, but leaves the shape essentially unchanged. In Fig. 4 we show W, Z production at $\sqrt{s}=2000$ GeV. The order-of-magnitude increase is due to the dominant contributions coming from smaller x_1 or x_2 values ($\simeq 0.04$), where the sea quark distributions play a vital role [see also Fig.

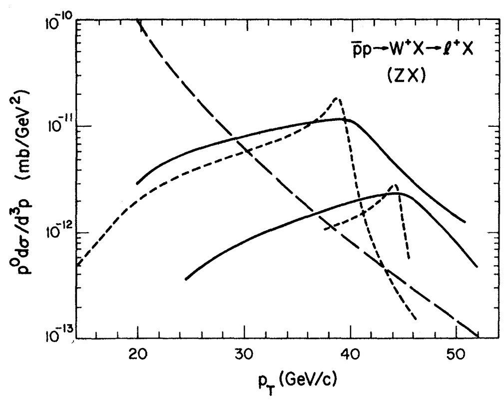

FIG. 3. The transverse-momentum distribution of decay leptons from W, Z bosons, Eq. (9), at $\theta=90°$ ($y_l=0$) and $\sqrt{s}=540$ GeV. For reference, our predictions (continuous curve) are compared with a Drell-Yan calculation with $Q_T=0$ (dotted curve). The dashed curve, taken from Ref. 15, shows an estimate of the leading background contribution which comes from the semileptonic decays of B and D mesons. The usual $K=2.3$ enhancement factor is not included in our predictions. (Ref. 18).

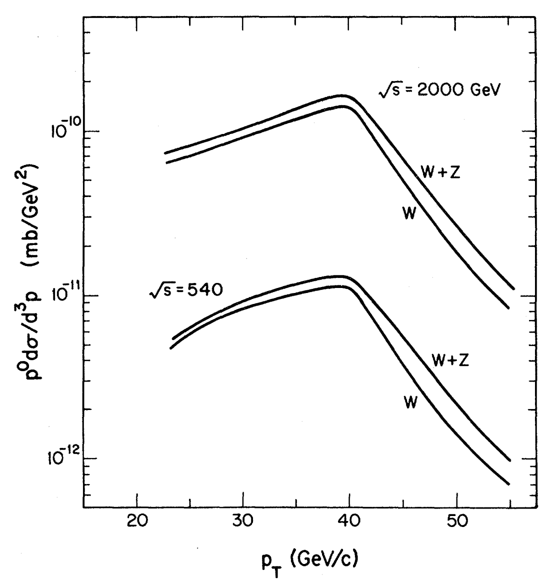

FIG. 4. The lepton spectrum from W, Z decays at $\theta=90°$ for $\bar{p}p$ collider energies of $\sqrt{s}=540$ GeV and 2 TeV.

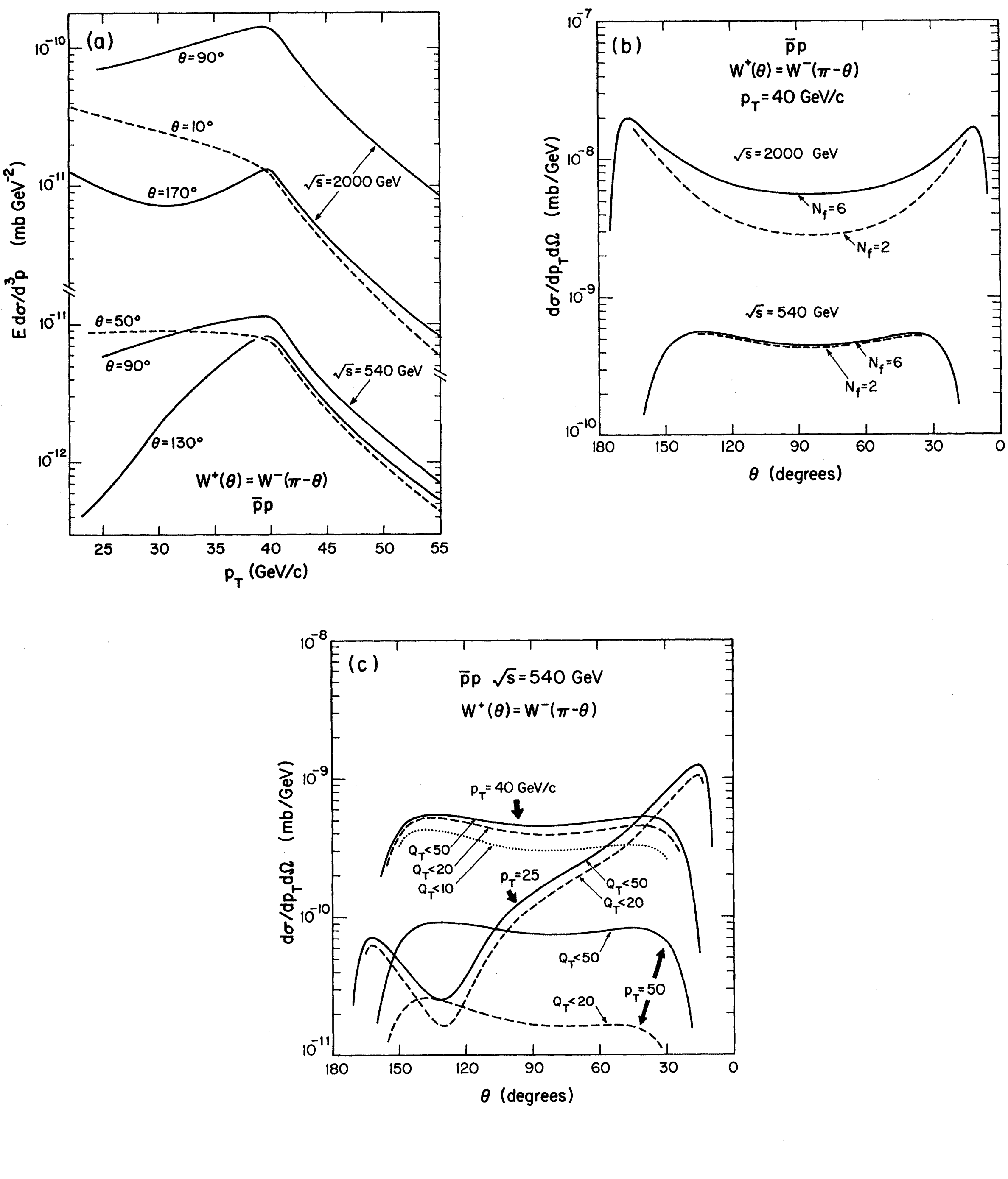

FIG. 5. (a) Transverse-momentum distribution of decay leptons from W^+ bosons for $\bar{p}p$ collisions with $\sqrt{s}=540$ GeV and $\sqrt{s}=2000$ GeV. The p_T distribution is shown for different values of the center-of-mass angle θ at which the lepton is produced relative to the $\bar{p}$ direction. The results for W^- are obtained from the same curves via Eq. (11). (b) Angular distribution of the decay lepton from W^+ produced at the maximum of the Jacobian peak shown in Fig 5(a), i.e., $p_T \simeq M_W/2 \simeq 40$ GeV. The results at $\sqrt{s}=2000$ GeV are sensitive to the number of flavors, i.e., N_f in Eq. (10). The curves show extreme choices $N_f=2,6$. (c) Angular distribution of the decay lepton from W^+ for various values of its transverse momentum p_T. The results are subdivided according to the contribution to the result from different ranges of the parent W's transverse momentum Q_T. The numerically dominant lepton cross sections are determined by the region of phase space where $Q_T < 20$ GeV, i.e., $Q_T << M_W$.

5(b)]. At $\sqrt{s}=540$ GeV the result depends mainly on the valence quark distributions. Further, we note that the tails of the calculated p_T distribution become sensitive to the region where $Q_T{}^2 \simeq M_W{}^2$, and are therefore unreliable.

In Figs. $5(a)-5(c)$ we exhibit the dependence of the lepton signal on the center-of-mass angle at which the lepton is detected. Figure 5(a) repeats the result of Fig. 3 on the Jacobian peak for different angles. Shown is the result for W^+; W^- is obtained from the same figure using Eq. (11). Figure 5(b) shows the dependence of the peak signal $p_T \simeq M_W/2 \simeq 40$ GeV on θ. Shown is the quantity

$$\frac{d\sigma}{dp_T d\Omega}=\frac{p_T}{\sin^2\theta}\left[E\frac{d\sigma}{d^3p}\right]. \tag{15}$$

We further illustrate our point regarding the dependence of the results at $\sqrt{s}=2000$ GeV on the degree of excitation of heavy flavors by varying N_f in Eq. (10) [but leaving the parameters of Eq. (14) unchanged]. $N_f=2$ ignores the contribution of all higher flavors, including strangeness. $N_f=6$ assumes a completely symmetric sea for all flavors, including t.

Figure 5(c) illustrates important experimental and theoretical points. Decay leptons within 30° of the beam constitute the dominant part of the total $W^\pm$ yield at $\sqrt{s}=540$ GeV. Or, as pointed out by Paige,[1] detectors ignoring small angles with respect to the beam pipe automatically sacrifice one order of magnitude of the $W^\pm$ rate. Figure 5(c) also proves after the fact the starting argument of this paper: small Q_T values of the W determine the shape of its lepton signature. As can be seen by comparing the solid and dashed curves, the $Q_T \simeq M_W$ region (say $Q_T > 20$ GeV) does not significantly affect the results of our calculation. The results are actually dominated by a Q_T value of 4 GeV. The small Q_T approximations implicit in Eqs. (5), (6), and (10) are therefore safe. Conversely, however, the calculations presented here are unreliable when $Q_T \simeq M_W$. The results only become sensitive to this region of Q_T when the lepton p_T itself is required to be large [see the curves labled $p_T=50$ in Fig. 5(c)]. The $O(\alpha_s)$ calculation is the correct approach to this region of phase space, as pointed out in the introduction. The latter case is of course difficult to access experimentally. Drell-Yan data at $\sqrt{s}\simeq 27$ GeV could be used as a laboratory to guess the range of validity of the $O(\alpha_s)$ calculation and the calculation described in this paper. Whereas the latter is able to accommodate[6] the data over the full Q, Q_T range, the $O(\alpha_s)$ result only describes the data for $Q_T > 3-4$ GeV for $Q \simeq 7-9)$ GeV, which translates in $Q/\sqrt{s}$ values in excess of 0.3.

THE PROMPT-LEPTON BACKGROUND AND THE IMPORTANCE OF IDENTIFYING MISSING ν ENERGY

In Fig. 3 we show the prompt-lepton background[15] under the Jacobian peak of W. It shows the favorable single-to-background ratio for the lepton signature of the W (see Halzen and Scott[2]) as opposed to the hadron-jet signature. The prompt-lepton background is made up of multiple sources[15]: single leptons from the Drell-Yan production of virtual photons, internal conversion of real photons with large p_T, decay of ψ's, but the dominant source of background is the production at large p_T of pairs of b quarks (c quarks to a lesser extend[15]), followed by their semileptonic decay. We would like to comment that this background calculation is much less reliable than that of the signal: e.g., in the case of b's, one uses $O(\alpha_s)$ perturbative QCD for producing the quarks, one has to guess the D functions to form B states, and one has to further guess a semileptonic decay distribution. One real worry here is the possibility of substantial diffractive production of b states. This is not included in the calculation. The crucial question is whether such b's are produced with increased transverse momentum at high energy. Reference 16 is an example of a scenario where this indeed happens. It is important to cover one's bets by rejecting these backgrounds by detecting the energy flow opposite the direct lepton.[17] Indeed, opposite a lepton from b decay one finds a hadron jet; opposite a lepton from W decay one finds an undetected ν, i.e., a gap in the energy flow.

Our calculations show, however, that it is impossible to enhance the Jacobian peak itself by doing $e-\nu$ coincidences. Indeed, the ν energy would have to be measured to better than a few GeV, as this is the type of Q_T imbalance that determines its shape.

Note added. After completion of this paper, similar work[19] by Paige and Protopopescu was brought to our attention.

ACKNOWLEDGMENTS

We are grateful to Max Dechantsreiter, John Dowell, Mike Pennington, Roberto Petronzio, Roger Phillips, Paul Stevenson, and Bryan Webber for help, advice, and many discussions. This research was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, in part by the Department of Energy under contract DE-AC02-76ERO0881, and in part by the Science and Engineering Research Council of Great Britain.

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