Single-electron/muon signature for heavy leptons and W gauginos at the CERN $p\bar{p}$ collider

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The signatures of several hypothetical fermions in the electron-plus-missing-transverse-energy channel at the CERN $p\bar{p}$ collider are studied. A new sequential lepton and a new lepton with $V+A$ coupling to the W boson, i.e., a "mirror" fermion, are studied in detail for several values of the lepton mass. The backgrounds of W decay through $e+\bar{\nu}_e$ and $\tau+\bar{\nu}_\tau \rightarrow e\bar{\nu}_e\nu_\tau\bar{\nu}_\tau$ will dominate a sequential lepton; however, a mirror lepton will be distinguishable at forward angles and moderate p_T for a lepton mass up to 50 GeV. The possibility of observing the supersymmetric decay of the W through W gaugino plus photino is also discussed. Subsequent W gaugino decay into photino + electron (or muon) + neutrino offers a signature differing only slightly from heavy-lepton decay due to the inherent mixing of $V\pm A$ couplings as determined by detailed mass mixing in the theory. If microvertex electronics and/or triggering on hadronic decay modes of the τ effects the isolation of the $W\rightarrow\tau\nu_\tau$ background (and by universality, the $W\rightarrow e\nu_e$ background), then the single-electron signature discussed here can reveal the existence of any new fermion produced in W decay.

I. INTRODUCTION

The discoveries of the W and Z bosons[1] are likely to open a window on new physics through their decays. The decay of the W is of special interest because of its nondiagonal nature. It may decay into a heavy new fermion approaching the mass of the W if the mass of the new partner is light. A new lepton may be expected to have a light neutrino as a partner, so we look to rare decays of the W as a likely place to find a new lepton.

In supersymmetric theories, the W may decay into W gaugino ($\tilde{\omega}$) and photino if these new particles are sufficiently light. The photino is expected to be very light and long-lived in many such theories. The decays of the $\tilde{\omega}$ may be such that an electron and missing transverse energy are the signatures of the event. Unfortunately, many of the details of the supersymmetric decays are model dependent. Nevertheless, there is a real possibility that we may see supersymmetry in the decay of the W.

This paper is a continuation of previous work[2] which presented a covariant formalism for the decay of a vector particle into a single observed particle. We applied the formalism to the case of W decay through a new heavy fermion, to electron or muon and neutrinos. Analytic formulas were presented for the invariant functions describing the W decay for arbitrary vector and axial-vector coupling of the new fermion. The cases of $V+A$ and $V-A$ were studied in the rest frame of a polarized W. In that frame, for either coupling, there exists a kinematic region where a new lepton would dominate the background from decay through the τ.

Here we continue our study by calculating the electron spectrum for the CERN $p\bar{p}$ collider.[3,4] We find that a sequential-lepton signal would be masked by the decays of the W through $e+\bar{\nu}_e$ and $\tau+\bar{\nu}_\tau$. However, a mirror lepton with $V+A$ coupling[5] would appear above background if its mass is below about 50 GeV. We will also consider the possibility of observing the W decay into W gaugino and photino. Microvertex detectors offer the possibility of isolating the $W\rightarrow\tau\nu_\tau$ signal from the data, and recently we have learned[6] of the possibility of isolating the τ signal by selecting exclusive hadronic decay modes of the τ. Knowledge of the $W\rightarrow\tau\nu_\tau$ distribution determines the direct $W\rightarrow e\nu_e$ distribution by universality, and determines the $W\rightarrow\tau\rightarrow e$ distribution when $W\rightarrow\tau\nu_\tau$ is convoluted with the standard four-fermion weak decay formula for the τ. Thus, one may hope for the eventual isolation of the $W\rightarrow\tau\rightarrow e$ and $W\rightarrow e$ backgrounds from any new sources of electron-plus-missing-transverse-momentum events. If this hope is realized, than any new fermion resulting from W decay may be observable. A quantitative search for the new fermion could be made using the formulas of this paper and Ref. 2.

In Sec. II of this paper, we recall the W decay formalism presented in Ref. 2, and extend that formalism to W's produced by $p\bar{p}$ annihilation. We also describe our choice for parton distributions and how we deal with the fact that W bosons are produced with considerable transverse momentum. Section III recalls our results for the electron angle and energy distributions in the W rest frame, and contains our numerical predictions for the same distributions in the $p\bar{p}$ laboratory frame. The graphs presented there show how the new lepton will manifest itself. Section IV contains a discussion of the possible observation of supersymmetry through the decay $W\rightarrow\tilde{\omega}+\tilde{\gamma}$. Finally, our conclusions are summarized in Sec. V. An appendix presents the totally differential cross section for the production of a W followed by its cascade decay through a new heavy fermion, with arbitrary vector and axial-vector coupling for that fermion.

II. COVARIANT FORMALISM

We recall the formalism presented in Ref. 2 for the momentum spectrum of a single particle chosen from the decay products of a spin-one particle. Denote the momentum of the decaying particle by Q^μ and that of the observed final particle, labeled f, by p^μ. Define the covariant tensor

$$W_{\mu\nu}=\frac{d^3p}{(2\pi)^3 2p_0}\int dP_{\rm LI}\,\langle 0\,|\,J_W^{\nu\dagger}(Q)\,|f(p),F\rangle \times\langle f(p),F\,|\,J_W^{\mu}(Q)\,|0\rangle\;, \qquad (1)$$

where J_W is the current coupling to the weakly decaying spin-one particle, F labels the final-state particles excluding f, and

$$dP_{\rm LI}=(2\pi)^4\delta^4\left[Q^\mu-p^\mu-\sum_F p_F^\mu\right]\prod_F\frac{d^3p_F}{(2\pi)^3 2p_F^0} \qquad (2)$$

is the Lorentz-invariant phase space for the set F. If particle polarizations are not measured, a spin sum is implied in Eq. (1), and we have the covariant tensor expansion:

$$W^{\mu\nu}=-g^{\mu\nu}W_1+p^\mu p^\nu W_2+i\epsilon^{\mu\nu\alpha\beta}p_\alpha Q_\beta W_3 +Q^\mu Q^\nu W_4+p^{\{\mu}Q^{\nu\}}W_5+ip^{[\mu}Q^{\nu]}W_6. \qquad (3)$$

The notation $a^{[\mu}b^{\nu]}=a^\mu b^\nu-a^\nu b^\mu$ and $a^{\{\mu}b^{\nu\}}=a^\mu b^\nu+a^\nu b^\mu$ is employed. The properties of the Lorentz-invariant functions, W_i are discussed in Ref. 2, and analytic formulas are presented there for W cascade decay through a heavy fermion, with subsequent decay to light (massless) fermions. We allowed an arbitrary combination of vector and axial-vector coupling between the new fermion, its partner, and the W boson. For completeness we present in the Appendix the squared matrix element for the cascade decay. We also present there the squared matrix element for production of a W by quark-antiquark annihilation followed by cascade decay. These formulas are valuable for analyzing signatures involving detection of more than one final-state particle or jet.

Here we apply the calculations of Ref. 2 to a realistic experimental setting by introducing and integrating over parton distributions and W transverse momenta, as appropriate for the CERN $p\bar{p}$ collider. The Feynman diagram for the hadronic production and subsequent decay of a W^- through a heavy new fermion, and our conventions, are given in Fig. 1. The antiproton momentum defines the $+z$ axis. For parton distributions we choose those of Owens and Reya,[7] scaled up to M_W^2. For the W transverse momentum distribution we fit the published data[8] on W production with the form

$$G(Q_T)=\frac{a^2}{2\pi}e^{-aQ_T}, \qquad (4)$$

$G(Q_T)$ is normalized such that its integral over $\mathbf{Q}_T$ is unity. The result of the fit is $a=0.35\ {\rm GeV}^{-1}$.

Using standard parton-model assumptions, we find

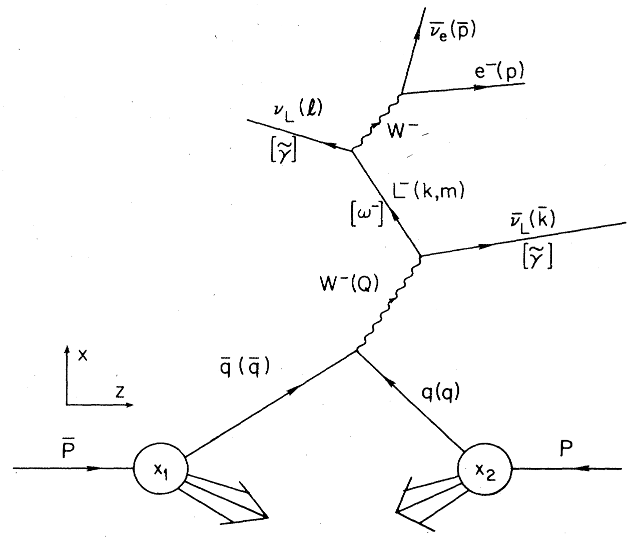

FIG. 1. The Feynman diagram for the production and decay of a W^- through a new heavy fermion. Particle labels in brackets identify a possible supersymmetric decay mode. The quark and antiquark may also originate from the antiproton and proton, respectively.

$$d\sigma=\frac{G_F M_W^2}{\sqrt{2}}\,\frac{1}{3}\int\frac{dQ^2}{s}G(Q_T)\frac{d^3Q}{Q_0} \times\frac{1}{(Q^2-M_W^2)^2+M_W^2\Gamma_W^2} \times\sum_{i,j}[f_{\bar{q}_i}(x_1)f_{q_j}(x_2)D_{++} +f_{q_i}(x_1)f_{\bar{q}_j}(x_2)D_{--}]\;, \qquad (5)$$

where s is the proton-plus-antiproton center-of-mass energy squared. The indices i and j run over all flavor pairs which may annihilate to give a W^-. Weak-mixing angles are neglected. Specifically, if there are N_f flavors active at the W mass scale, then

$$\sum_{i,j}f_{\bar{q}_i}(x_1)f_{q_j}(x_2)=u(x_1)d(x_2) +\left[\frac{N_f}{2}-1\right]s(x_1)s(x_2)$$

and

$$\sum_{i,j}f_{q_i}(x_1)f_{\bar{q}_j}(x_2)=\frac{N_f}{2}\,s(x_1)s(x_2)\;,$$

where the quark densities $u(x)$ and $d(x)$, are defined, as usual, with respect to a proton, and include both the valence and sea contributions. The sea is given by $s(x)$. We assume an SU(4)-symmetric sea by setting $N_f=4$. For W^+ production, replace $u(x_1)d(x_2)$ by $d(x_1)u(x_2)$

and change the sign of W_3. The fractional momenta, x_1, x_2, refer to the antiproton and proton, respectively. They take the values $(Q_0 \pm Q_z)/\sqrt{s}$, respectively. In Ref. 4, it was emphasized that one must not put the W on mass shell for the direct decay. Accordingly, we do not use this approximation for the direct decay into $e+\bar{\nu}_e$. We do, however, put the W on mass shell for the cascade decays by replacing $[(Q^2-M_W^2)^2+M_W^2\Gamma_W^2]^{-1}$ in Eq. (5) with $\delta(Q^2-M_W^2)\ \pi/(M_W\Gamma_W)$. When enough data accumulate to warrant a careful comparison with theoretical formulas, the rigorous expression of Eq. (5) can be used. The $D_{\pm\pm}$ are the density matrix elements for decay of completely polarized W's. They are obtained by contracting $W_{\mu\nu}$ of Eq. (3) with the appropriate tensor $\epsilon_\pm^\mu \epsilon_\pm^{\nu *}$ constructed from W polarization vectors.

Blindly convoluting the no-q_T parton result with a Q_T spectrum for the W will yield negative values for the cross section at large values of Q_T. This unphysical behavior can be traced to a violation of the invariant constraint, $Q\cdot\epsilon=0$, for on-shell W's. [In the naive parton model, $\epsilon_\pm^\mu=(0,1,\pm i,0)/\sqrt{2}$ so $Q\cdot\epsilon\neq 0$ if $Q_T\neq 0$.] Accordingly, for consistency of the model, we must explicitly arrange to maintain the constraint $Q\cdot\epsilon=0$. From helicity conservation arguments, one expects that to order $(q^2/Q^2,\bar{q}^2/Q^2)$ the W will have helicity ± 1 in the frame where the $q\bar{q}$ producing it are colinear. Finding the colinear frame amounts to solving QCD dynamics;[9] hence we approximate the colinear frame with the $Q_T=0$ frame obtained by a purely transverse boost from the laboratory frame. For simplicity, let the W transverse momentum be in the x direction, the beam axis in the z direction. Boosting back to the laboratory frame, we find

$$\epsilon_\pm^\mu=\frac{1}{\sqrt{2}}(\gamma\beta,\gamma,\pm i,0) \tag{6}$$

with

$$\gamma=\frac{Q_0}{(Q^2+Q_z^2)^{1/2}} \tag{7}$$

and

$$\gamma\beta=\frac{Q_T}{(Q^2+Q_z^2)^{1/2}} . \tag{8}$$

Let the electron momentum be determined by the polar angles (θ,η) with respect to the z axis. Contracting $\epsilon_\pm^\mu \epsilon_\pm^{*\nu}$ with the decay tensor, $W_{\mu\nu}$, we then get from Eqs. (3), (6), (7), and (8),

$$\begin{aligned} D_{\pm\pm}=W_1+\frac{p_T^2}{2}\left[\sin^2\eta+\frac{(Q_0\cos\eta-Q_T/\sin\theta)^2}{Q^2+Q_z^2}\right]W_2 \\ \mp[Q_0(Q_0\cos\theta-Q_z)p_0+Q_T(Q_z p_T\cos\eta-Q_T p_z)]\frac{W_3}{(Q^2+Q_z^2)^{1/2}} . \end{aligned} \tag{9}$$

The terms included by using the boosted polarization vectors for the W differ from the unboosted by magnitudes of order $(\beta=Q_T/Q_0)$, which need not be small.

III. SINGLE AND DOUBLE DIFFERENTIAL CROSS SECTIONS

In Ref. 2, we detailed the single differential cross sections as a function of p_T or angle in the W rest frame. Let us recall those results. For the τ, the p_T distribution peaks very close to zero, whereas for the 20, 40, and 60 GeV $V-A$ leptons, the peaks are at $\sim$ 3, 9, and 18 GeV, respectively. For $V+A$ leptons, the average p_T is smaller, and the peaks are at $\sim$ 2, 7, and 14 GeV, respectively.

There are three significant differences between the W rest frame and the laboratory frame. First, in the laboratory frame, the W has longitudinal momentum because the annihilating quark and antiquark have unequal momentum fractions. This longitudinal boost would not affect the electron transverse-momentum spectrum, but would affect energy and angle distributions. The second difference comes from initial state gluon bremsstrahlung and primordial parton q_T. These result in the W having transverse momentum, and a spin direction slightly off the beam axis. Thus we expect that the p_T distributions we calculated in Ref. 2 will be smeared. The third difference is that the direct electron energy is no longer fixed at $M_W/2$. Rather, it ranges over the interval $(\sqrt{s}/2)\ [M_W^2/s,1]$ and presents a serious background to cascade-decay signatures. [The direct-electron minimum energy is 6 GeV (5 GeV) for $\sqrt{s}$ = 540 GeV (630 GeV), probably too low to provide a useful experimental cut.] Fortunately, the conclusion remains that the cascade decays of interest here peak at relatively small values of p_T, whereas the direct decay develops its characteristic Jacobian peak at $M_W/2$. Thus, one may expect an enhanced signal to background ratio at smaller p_T.

In Ref. 2 we also showed the angular distribution of decay electrons in the rest frame of a polarized W. We found that for $V-A$ leptons, the new-particle decay dominates the τ decay only close to the direction of the W spin. However, for $V+A$ leptons, with 20 or 40 GeV masses, there is a considerable angular range where the new physics dominates the τ background. In either case, we must look in the direction of W spin. In the laboratory frame, the W polarization is in the direction of the antiproton momentum, but it is reduced by the contribution of the sea partons.

Motivated by the above considerations, we choose to display cross sections as a function of $\cos(\theta)$ for three values of electron p_T, 5, 10 and 20 GeV. We display curves for $V\pm A$ for lepton masses of 20, 40, and 60 GeV.

The two backgrounds $W \rightarrow e\bar{\nu}_e$ and $W \rightarrow \tau\bar{\nu}_\tau \rightarrow e\bar{\nu}_e \nu_\tau \bar{\nu}_\tau$ are summed. All distributions are normalized by the direct cross section to reduce uncertainties coming from nonperturbative QCD and from the W width. We assume the existence of a 40 GeV top quark. This increases the W width, and the heavy lepton width for $m_L > 45$ GeV. On the $V-A$ plots we have also displayed the τ contribution separately, in order to more clearly show the trend as the lepton mass changes.

In Figs, 2, 3, and 4 we show the cross sections for the background, τ, and heavy-lepton channels for the three p_T values considered. Apparently, a fourth-generation sequential lepton would be very difficult to see in this channel. For $p_T = 10$ and 20 GeV, the lepton would be very well masked by the direct and τ decays. For $p_T = 5$ GeV, up to $m = 40$ GeV, a new lepton would provide an excess of some fraction of the background, still not a very promising possibility.

The difficulty of seeing a new sequential lepton through its leptonic decay was discussed in Ref. 4. There additional analysis indicated that the hadronic decays of L give a promising signature when appropriate cuts are introduced to reduce the background due to $q+\bar{q} \rightarrow g+W \rightarrow g+\tau\nu_\tau$, where g denotes a gluon, and to heavy quark production followed by semileptonic decay.

On the other hand, a $V+A$ lepton shows a very promising signal. If $m = 20$ GeV, the lepton contribution exceeds the background for $\cos(\theta) \geq 0.25$, at p_T of 5 GeV. At p_T of 10 GeV, we see that for $m = 20$ GeV, signal exceeds background for $\cos(\theta) \geq 0.65$ and for $m = 40$ GeV, signal exceeds background for $\cos(\theta) \geq 0.75$. At p_T of 20 GeV, the direct and τ decays dominate any electrons from new leptons.

We reiterate the potential significance of deducing the $W \rightarrow \tau\nu_\tau$ cross section from microvertex techniques, or by triggering on exclusive hadronic decay modes of the τ. Such a measurement then allows, in principle, subtraction of the $W \rightarrow e$ and $W \rightarrow \tau \rightarrow e$ backgrounds, leaving a clean electron plus missing transverse momentum signal for a new sequential or mirror lepton, or for a gaugino from supersymmetric models, to which we now turn.

IV. SUPERSYMMETRIC DECAYS OF THE W

One particularly exciting class of decays of the W boson is that including the superpartners of ordinary particles. Supersymmetry,[10] in addition to its mathematical beauty, is interesting as a possible solution to the gauge hierarchy problem.[11] There has been considerable work recently trying to uncover the best ways to find evidence for supersymmetry. For recent reviews see Ref. 12.

Some time ago, Weinberg[13] noted that in a large class of supersymmetric models, those based on broken supergravity,[14] there will be light gauge fermions into which the gauge bosons may decay. A number of authors have investigated the branching ratios of W and Z into gauge fermions.[15] We should, in fact, be more precise, and note that the mass eigenstates of these models are not the gauge fermions themselves, but mixtures of the gauge fermions and the neutral and charged fermionic partners of the Higgs multiplets required for gauge symmetry breaking. Unfortunately, the details of the mixing, and hence the branching ratios, depend on several unknown parameters. Let us briefly summarize the situation.

The low-energy sector of a supersymmetric model is characterized by a set of superfields, a superpotential which describes the interactions, and a set of additional interactions which softly violate supersymmetry and depend on the details of the complete theory.[16] Minimally, a model will entail vector superfields for each of the SU(3)$\times$SU(2)$\times$U(1) gauge bosons, chiral superfields for each generation of quarks and leptons, and a pair of chiral Higgs superfields. As in the ordinary case, our ignorance is concentrated in the Higgs sector. The superpotentials for models which break SU(2)$\times$U(1) at the tree level contain as a minimum the terms:[17]

$$\mu \hat{H}\hat{H}' + \lambda \hat{U}\hat{H}\hat{H}' + \lambda' \hat{U}^3 , \tag{10}$$

where $\hat{U}$ is a gauge singlet chiral superfield, and $\hat{H}$ and $\hat{H}'$ are the two Higgs doublet superfields. Finally, the soft-supersymmetry-breaking terms include direct gaugino mass terms as well as mass terms for the various scalars. These terms play a role in determining the vacuum expectation values of the neutral scalar components of the Higgs doublets. $\langle H^0 \rangle$ and $\langle H'^0 \rangle$ are crucial parameters in the theory since they determine the SU(2)$\times$U(1) breaking, and hence the $W^\pm$ and Z masses. Thus, $|\langle H^0 \rangle|^2 + |\langle H'^0 \rangle|^2$ is determined by experiment and only $\langle H^0 \rangle / \langle H'^0 \rangle$ remains an undetermined parameter.

The decays in which we are most interested involve the coupling of the W boson to the gauge fermions. We must determine mass eigenstates since the gauge fermions may mix with the Higgs fermions. The contributions to the mass matrix come from the $\hat{H}\hat{H}'$ terms in the superpotential, the gauge-fermion—Higgs-boson—Higgs-fermion gauge terms dictated by supersymmetry (with the neutral Higgs fields replaced by their vacuum expectation values), and soft-supersymmetry-breaking gaugino mass terms. Thus, the eigenstates depend upon four parameters: $\mu + \lambda\langle U \rangle$, $\langle H^0 \rangle / \langle H'^0 \rangle$, and two gaugino mass terms.

Having diagonalized the mass matrix, one may reexpress the trilinear coupling to the W boson in terms of the mass eigenstates. This, of course, has been done in the literature.[15] The result is a set of couplings which are an admixture of V and A. The admixture depends upon the four undetermined parameters just identified, and must be considered as completely arbitrary. Thus the angular distribution of electrons from the mode (see Fig. 1) $W \rightarrow \tilde{\gamma} + \tilde{\omega} \rightarrow \tilde{\gamma} + \tilde{\gamma} + W_{\text{virtual}} \rightarrow \tilde{\gamma} + \tilde{\gamma} + e + \nu$ may resemble the $V-A$ or $V+A$ curves we have exhibited in Figs. 2—4, or any shape in between. The rate for this mode is expected to be somewhat suppressed relative to the heavy-lepton channel by mixing angles arising from mass diagonalization. (If the gluino is light and scalar quarks not too heavy the decay $\tilde{\omega} \rightarrow q\,\bar{q}\,\tilde{g}$ will further suppress the $W \rightarrow \tilde{\omega} \rightarrow e$ channel considerably.)

A more complete calculation of the $W \rightarrow \tilde{\gamma}\tilde{\omega} \rightarrow \tilde{\gamma}\,\tilde{\gamma}\,e\,\nu$ mode should include antisymmetrization of the final-state photinos, include the amplitude for $\tilde{\omega}$ decay through a virtual scalar electron in addition

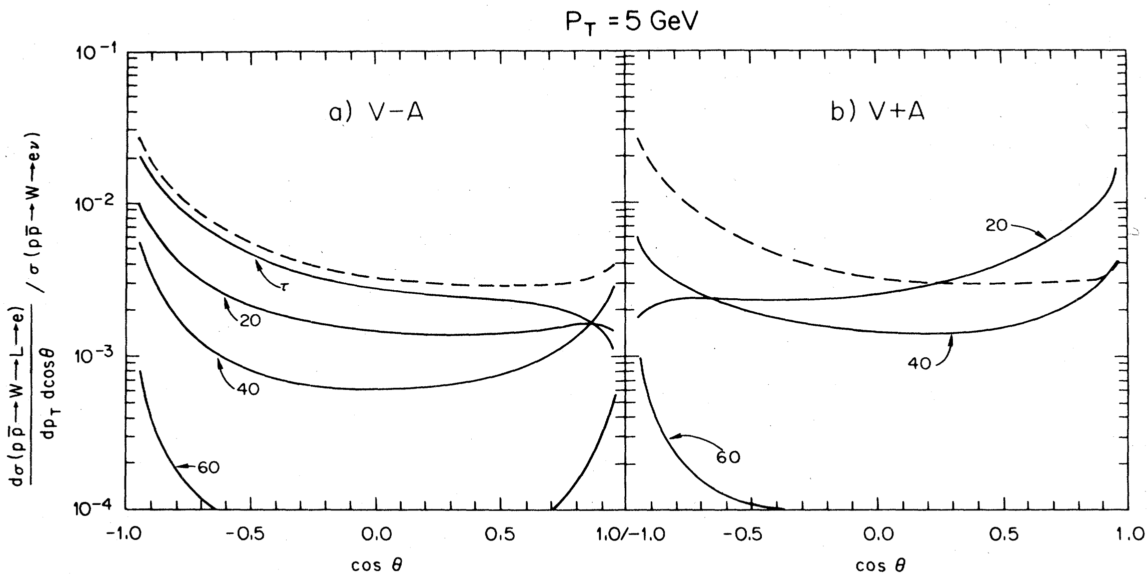

FIG. 2. The cross section, as a function of electron (or muon) angle with respect to the antiproton momentum. The cross section is for a fixed value of $p_T=5$ GeV, and $\sqrt{s}=540$ GeV. It is normalized to the cross section for the production and direct decay of a W. Curves are shown for the background, i.e., direct decay plus decay through τ, and various masses for the heavy lepton. The lepton coupling is $V-A$ in (a), and $V+A$ in (b).

to virtual W, and allow for a possibly large photino mass.[18] We do not include these refinements here. If the scalar-electron mass is large compared to the W mass, neglect of the second amplitude is warranted. If the photino mass is small on the scale of the W gaugino, neglect of its mass is also warranted. The appendixes of Ref. 2 contain the necessary formulas for the inclusion of a photino mass.

There are other possible decay modes available to the W gauginos depending upon the supersymmetric mass spectrum. For each gauge coupling of the W, there is a corresponding coupling of the $\tilde{\omega}$ with one of the other particles replaced by its superpartner. Thus, the $\tilde{\omega}$ couples to W + photino, quark + scalar quark, electron + scalar neutrino, scalar lepton + neutrino, as well as additional couplings to Higgs bosons and fermions and other gauge particles. If one of the scalar leptons or scalar quarks is lighter than the $\tilde{\omega}$, there can be a two-body decay of the $\tilde{\omega}$ in addition to the three-body decay discussed above. Many of these decay modes lead to a final

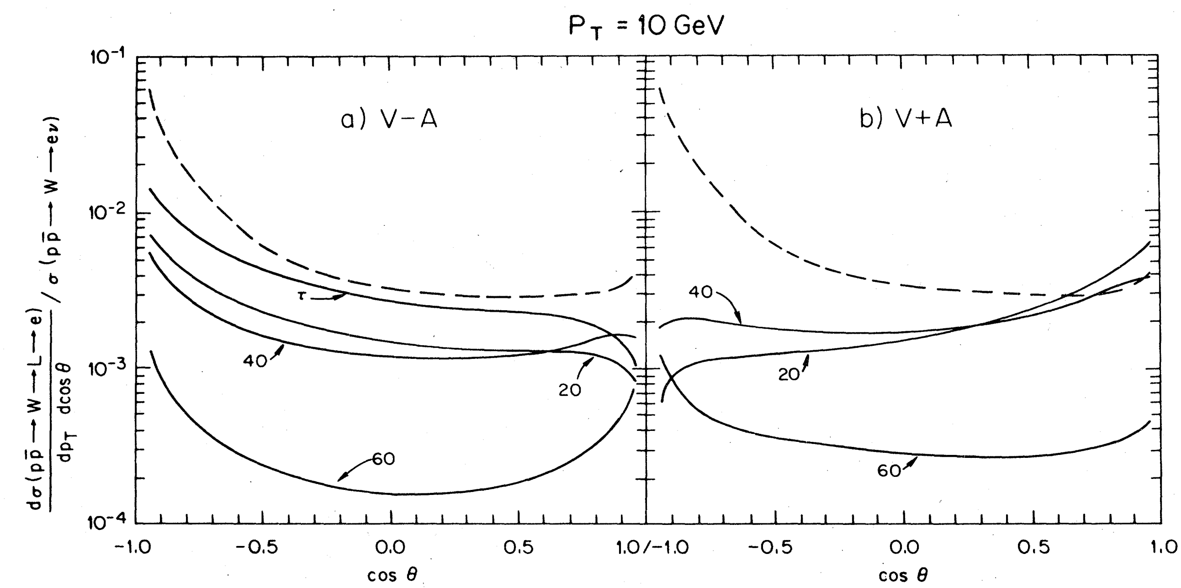

FIG. 3. Same as for Fig. 2, except that $p_T=10$ GeV.

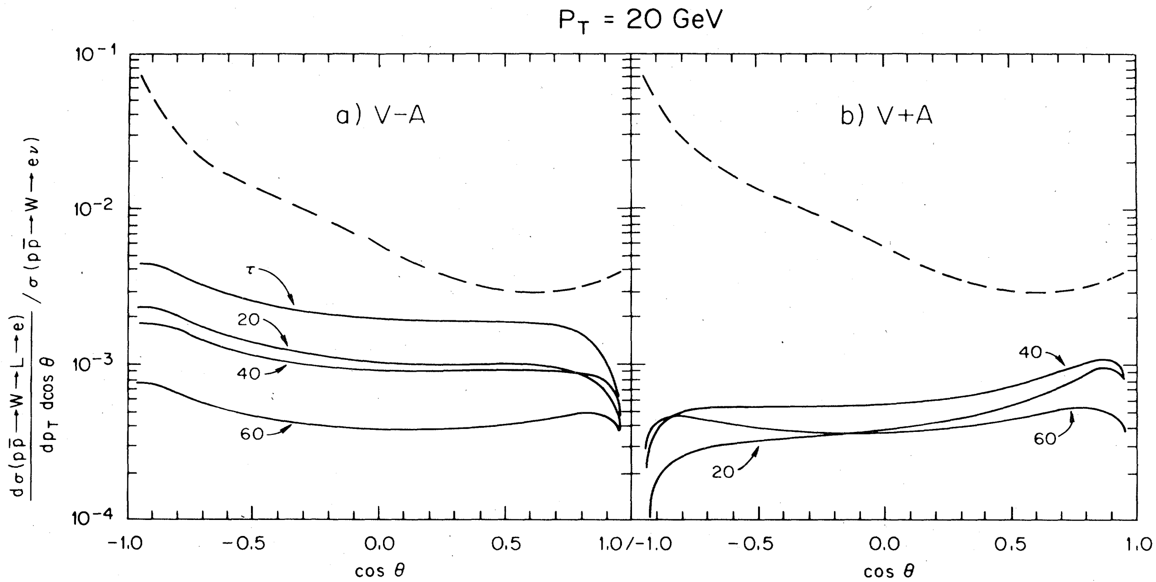

FIG. 4. Same as for Fig. 2, except that $p_T = 20$ GeV.

state consisting of a charged lepton and unobserved neutrals, again yielding the same signature, electron plus missing transverse energy, as the heavy-lepton decay discussed previously.

In addition to W and Z decays into supersymmetric fermions, there is the possibility of decay into new scalars, specifically, decay into scalar electron and scalar neutrino[19] or scalar quark and scalar antiquark. This, of course, requires that these two scalar leptons or scalar quarks be sufficiently light. Since only one of the superpartners will be stable in most supersymmetry models, whether the W decays into fermions or bosons, we will be looking for the decay products of those particles, not the new particles themselves. This situation is again analogous to the decay of the W into a new lepton. The similarity of signatures resulting from W cascade decay through a new heavy lepton (charged or neutral[20]) or through supersymmetric particles has been noted by many authors. We wish to emphasize that a complete treatment of the decay chain, such as that afforded by the formalism of Ref. 2, will be necessary to accurately predict experimental signatures.

There has been some prior work in which the decay chain has been followed to the final electron. Barbieri, Cabibbo, Maiani, and Petrarca[19] consider the decay of the W into scalar electron and scalar neutrino. Dicus, Nandi, Repko, and Tata[3] consider the decay into W gaugino and photino. Backgrounds as a function of electron energy are considered along with the decay spectrum. Their treatment of the angular distribution does not take proper account of the initial W polarization, as our formalism does. Barnett, Lackner and Haber[19,21] and Baer, Ellis, Nanopoulos, and Tata[22] calculate the W decay into scalar electron and scalar neutrino, or treat the decay into W gaugino and photino, but concentrate on the possible two-body decay of the W gaugino into electron and scalar neutrino or scalar electron and neutrino.

Given the model dependence of the supersymmetric decay scenarios, it seems that considerable care will be required to calculate the final electron signals from supersymmetric decays. If an excess of electrons is seen it will be necessary to decide if they are coming from a new lepton, or from a supersymmetric decay. It is quite likely that only by looking at additional decay channels and other signatures will one be able to decide.

V. CONCLUSIONS AND SUMMARY

We have applied the formalism developed in Ref. 2 to the decay of W bosons as produced in the CERN collider, into new matter. Taking a conservative viewpoint, we conclude that it would be very difficult to see a new sequential $V-A$ lepton in the electron plus missing energy channel which was used to discover the W. The backgrounds of direct decay and τ cascade decay are larger than the signal for all ranges of electron p_T and angle. Thus a very reliable background subtraction would be required.

On the other hand, for a new $V+A$ lepton, the direct and τ cascade decay backgrounds will only be a fraction of the signal at forward angles and low p_T. Thus, a $V+A$ lepton could be seen if its mass is less than about 50 GeV.

If the $W \to \tau\nu$ (and by universality, $W \to e\nu$) backgrounds can be eliminated by microvertex techniques or by inferring the τ distribution from measurement of exclusive hadronic decay modes of the τ, then electron plus missing transverse momentum becomes an excellent signature for revealing a new heavy lepton.

Finally, the possibility of seeing new decays of the W and Z into superpartners of ordinary particles is real, and very exciting. In particular, the $W \to \tilde{\gamma}\,\tilde{\omega} \to \tilde{\gamma}\,\tilde{\gamma}\,\nu\,e$ channel will yield an electron distribution intermediate between the $V+A$ and $V-A$ heavy-lepton results. This is because mass mixing in the gaugino sector causes an admixture of $V\pm A$ in the $W\,\tilde{\gamma}\,\tilde{\omega}$ coupling.

(*Note added*: After this research was completed, we received a copy of Report No. LAPP-TH-104 by P. Aurenche and R. Kinnunen which also investigates the possibility of observing a new heavy lepton in $p\bar{p}$ collisions.)

ACKNOWLEDGMENTS

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APPENDIX: TOTALLY DIFFERENTIAL CROSS SECTION

Since the amplitude we calculate may also be used to describe cascade decay modes of the W where more than a single final particle or jet is detectable, we display the completely differential cross section in this appendix. Particle names and momenta are identified in Fig. 1.

The amplitude for the cascade decay is given by

$$M^\mu = i\left[\frac{G_F M_W^2}{\sqrt{2}}\right]^{3/2} \times \frac{J^\mu}{[k^2-m^2+im\Gamma_L][(p+\bar{p})^2-M_W^2+iM_W\Gamma_W]} \tag{A1}$$

with

$$J^\mu = [\bar{u}_{\nu_L}(l)\gamma^\lambda(v+a\gamma^5)(\not{k}+m)\gamma^\mu(v+a\gamma^5)v_{\bar{\nu}_L}(\bar{k})] \times[\bar{u}_e(p)\gamma_\lambda(1-\gamma^5)v_{\bar{\nu}_e}(\bar{p})]\ . \tag{A2}$$

Summing over final-state spins, we define

$$J_{\mu\nu} = \sum_{\text{spins}} J_\mu J_\nu^\dagger\ , \tag{A3}$$

$$\begin{aligned} J^{\mu\nu} = {} & 256 g_L^4\, p\cdot l\, \mathrm{Tr}(1-\gamma^5)(2k\cdot\bar{p}\,\not{k}-k^2\bar{\not{p}})\gamma^\mu \bar{\not{k}}\gamma^\nu \\ & +256 g_R^4\, \bar{p}\cdot l\, \mathrm{Tr}(1+\gamma^5)(2k\cdot p\,\not{k}-k^2\not{p})\gamma^\mu\bar{\not{k}}\gamma^\nu \\ & +256 m^2 g_L^2 g_R^2\{p\cdot l\,\mathrm{Tr}(1+\gamma^5)\bar{\not{p}}\gamma^\mu\bar{\not{k}}\gamma^\nu \\ & +\bar{p}\cdot l\,\mathrm{Tr}(1-\gamma^5)\not{p}\gamma^\mu\bar{\not{k}}\gamma^\nu\}\ , \end{aligned} \tag{A4}$$

where $g_L=\frac{1}{2}(v-a)$, and $g_R=\frac{1}{2}(v+a)$. The third term in Eq. (A4) assumes m_{ν_L} is negligible. The first two terms make no assumptions on final-state masses. $J^{\mu\nu}$ is related to the $W^{\mu\nu}$ in the text by

$$W^{\mu\nu} = \left[\frac{G_F M_W^2}{\sqrt{2}}\right]^3 \frac{d^3p}{(2\pi)^3 2p_0}\int \frac{d^3\bar{p}}{(2\pi)^3 2\bar{p}_0}\frac{d^3\bar{k}}{(2\pi)^3 2\bar{k}_0}\frac{d^3 l}{(2\pi)^3 2l_0} \times \frac{J^{\mu\nu}(2\pi)^4\delta^4(Q-p-\bar{p}-\bar{k}-l)}{[(k^2-m^2)^2+(m\Gamma_L)^2]\{[(p+\bar{p})^2-M_W^2]^2+(M_W\Gamma_W)^2\}}\ . \tag{A5}$$

The analytic expression for $W^{\mu\nu}$ is available in Ref. 2.

If the W boson is produced through the annihilation of a quark and antiquark, the amplitude for the process is given by the contraction of M^μ with a production vector P_μ:

$$P_\mu = -\left[\frac{G_F M_W^2}{\sqrt{2}}\right]^{1/2}\frac{1}{Q^2-M_W^2+i\Gamma_W M_W}\,\bar{v}(\bar{q})\gamma^\mu(1-\gamma^5)u(q)\ . \tag{A6}$$

Terms of order m_q/M_W are neglected, and $Q^\mu=(q+\bar{q})^\mu$ is the W four-momentum. Concentrating for the moment upon the factor involving the spinors, we define

$$P_{\mu\nu} = \frac{1}{12}\sum_{\text{spins}}[\bar{v}(\bar{q})\gamma_\mu(1-\gamma^5)u(q)]\,[\bar{u}(q)\gamma_\nu(1-\gamma^5)v(\bar{q})]\ . \tag{A7}$$

The factor of $\frac{1}{12}$ accounts for the average over initial spins and colors. We have,

$$\begin{aligned} P_{\mu\nu}J^{\mu\nu} = \frac{2^{13}}{3}[& g_L^4\, l\cdot p\,\bar{k}\cdot q\,(2k\cdot\bar{p}\,k\cdot\bar{q}-k^2\bar{p}\cdot\bar{q})+g_R^4\, l\cdot\bar{p}\,\bar{k}\cdot\bar{q}\,(2k\cdot p\,k\cdot q-k^2 p\cdot q) \\ & +m^2 g_L^2 g_R^2\,(p\cdot l\,\bar{p}\cdot q\,\bar{k}\cdot\bar{q}+\bar{p}\cdot l\,\bar{k}\cdot q\,p\cdot\bar{q})]\ . \end{aligned} \tag{A8}$$

The absolute matrix element squared, averaged over initial spins and colors, and summed over final spins is given by

$$|M|^2 = \left[\frac{G_F M_W^2}{\sqrt{2}}\right]^4 \frac{J^{\mu\nu}P_{\mu\nu}}{\{[(q+\bar{q})^2-M_W^2]^2+\Gamma_W^2 M_W^2\}[(k^2-m^2)^2+\Gamma_L^2 m^2]\{[(p+\bar{p})^2-M_W^2]^2+\Gamma_W^2 M_W^2\}}\ . \tag{A9}$$

This result for a $V-A$ coupling, i.e., $g_L = 1$, and $g_R = 0$, has been given in Ref. 4. Finally, dividing Eq. (A9) by the parton flux, and integrating over parton momentum probabilities gives the totally differential cross section

$$\frac{2\bar{k}^0}{d^3\bar{k}}\frac{2l^0}{d^3l}\frac{2\bar{p}^0}{d^3\bar{p}}\frac{2p^0}{d^3p}\,d\sigma=\frac{1}{(2\pi)^8}\int\frac{dx_1dx_2}{2Q^2}\sum_{i,j}[f_{\bar{q}_i}(x_1)\,f_{q_j}(x_2)+f_{q_i}(x_1)\,f_{\bar{q}_j}(x_2)]\,|M|^2\delta^4(Q-p-\bar{p}-\bar{k}-l)\ . \tag{A10}$$

To introduce W transverse momentum, we make the replacement

$$dx_1dx_2=\frac{dQ^2}{s}\frac{dQ_z}{Q_0}\to\frac{dQ^2}{s}\frac{d^3Q}{Q_0}\,G(Q_T)\ ,$$

where s is the square of the proton-antiproton center-of-mass energy and $G(Q_T)$ is the normalized W transverse-momentum distribution.

The analog of Eqs. (A4) and (A8) for the case where the heavy fermion (labeled L in Fig. 1) is an antiparticle rather than a particle is obtained by interchanging g_L and g_R, and identifying l with the $\bar{\nu}_L$ momentum and $\bar{k}$ with the ν_L momentum. For any decay mode, p remains the particle momentum, and $\bar{p}$ remains the antiparticle momentum. The analogs for the single-particle expressions (3) and (A5) have been given in Ref. 2: If the final antiparticle labeled by $\bar{\nu}_e$ in Fig. 1 is detected, rather than the particle labeled by e^-, the sign of W_3 is reversed. If the heavy fermion L and the detected particle have opposite fermion number, an interchange of g_L and g_R is required.

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