Signatures of heavy-neutrino production at the CERN collider

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(Received 7 July 1986)

We discuss the possibility of discovering at the CERN collider a new sequential neutrino or a mirror neutrino (which can mix with the light neutrinos). Three classes of event signatures are examined: purely leptonic decay channels, monojets with missing p_T, and monojets with isolated electrons and very little missing p_T. This last signal is studied in detail; a set of experimental cuts are suggested and a number of possible backgrounds are estimated. At least 1 pb^{-1} of data is needed before nontrivial limits on heavy-neutrino parameters can be achieved.

I. INTRODUCTION

Three families of quarks and leptons are observed experimentally, but the existence of more massive families is not excluded by any experiments or theoretical prejudices. Indeed, a variety of superstring,[1] composite,[2] supersymmetric,[3] grand unified, and family unified theories[4–6] suggest the existence of new fermions beyond the presently known three families. Current data from e^+e^- colliders imply that a new charged lepton or quark must be heavier than about 22.5 GeV (Ref. 7). The CERN collider may be able to increase these bounds somewhat. Some limits also have been obtained for new neutral leptons, which depend on their masses and mixing angles. The existence of massive neutral leptons are severely restricted for masses below 10 GeV (Refs. 7–9). On the theoretical side, arguments have been made which put an upper bound on new-fermion masses. Such arguments are based on computing the effective potential (including radiative corrections) and requiring that the spontaneous-symmetry-breaking vacuum be the true minimum.[10] Other theoretical limits on fermion masses are derived using the infrared fixed-point structure of the $SU(3)\times SU(2)\times U(1)$ renormalization-group equations for Yukawa couplings.[11] Finally, the experimental observation that $\rho = m_W{}^2/(m_Z{}^2\cos^2\theta_W) \approx 1$ limits the mass splittings of new quark and lepton isodoublet partners.[12] All the bounds described above are quite weak for the mass range we consider here, and they leave a good deal of room for the existence of new families.

Cosmology provides the only significant restriction on extra families: standard big-bang nucleosynthesis calculations[13] imply a conservative upper bound of four essentially massless neutrinos (assuming typical weak-interaction couplings). With reasonable values for the fractional abundances of deuterium and ^{4}He, this bound may be lowered to three. Therefore, we will require that additional neutral leptons beyond the three known light neutrinos are massive. Such neutral leptons will henceforth be denoted by N.

Neutrinos may obtain masses from several sources. The vacuum expectation value of a weak isotriplet gives Majorana masses to neutrino doublets. In such models, one must be careful not to upset the relation $\rho \approx 1$. In grand unified theories with both weak doublet and singlet neutrinos, the Gell-Mann–Ramond–Slansky–Yanagida[14] mechanism induces ultralight Majorana masses ($\lesssim 100$ eV), which arise from a mass matrix consisting of Dirac masses for the doublet neutrinos and superheavy Majorana masses ($\gtrsim 10^{10}$ GeV) for singlet neutrinos. If we forbid this mechanism to operate by introducing a horizontal or family symmetry (which would prohibit the generation of a Majorana mass term), then a neutrino with a typical Dirac mass would remain.

In the O(18) family unified theory, the neutrinos from extra families have Dirac masses.[5] This is a consequence of a horizontal discrete symmetry in the low-energy effective theory. In the context of the low-energy theory, to determine mass bounds on new particles, a renormalization-group analysis of the evolution of the Yukawa couplings of extra families was performed. The neutrino Yukawa couplings for the predicted five new families have been examined at the weak scale. All five neutrinos are expected to have Dirac masses below 40 GeV.

While the O(18) and isotriplet scalar methods for giving neutrinos mass are not the most general, it is plausible that if there exist new families with heavy neutrinos, they may be in a range kinematically accessible to W and Z decays. In this paper we will discuss a few experimental signals of heavy neutrinos. We will concentrate on the low-energy framework suggested by the O(18) theory of family unification because of its concrete predictions of mass limits and its rich neutrino structure. In the following discussion, we will assume that all additional charged leptons are more massive than the heavy neutrinos.

This paper is organized as follows. First, we elaborate on our model of neutrino mixing. Second, we describe a few of the signals for heavy-neutrino production at the CERN collider. We concentrate on the process $p\bar{p} \to l^\pm N \to l^\pm + \mathrm{jet(s)}$, where $l^\pm$ is an ordinary charged lepton which is isolated from the jet. The main background comes from $p\bar{p}\to W^{\pm}+(g$ or $q)$ followed by $W^{\pm}\to l^{\pm}\nu$. This sample of events is also of interest in the top-quark search which focuses on $p\bar{p}\to l^{\pm}+$jet(s); it is important to see how to distinguish between these various processes. Other backgrounds from the semileptonic decays of c and b quarks are briefly mentioned. Finally, the prospects for observing heavy neutrinos at the CERN collider are evaluated.

II. MODEL OF NEUTRINO MIXING

Theories with massive neutrinos typically violate electron, muon, and tau number separately in their charged-current couplings; this is the analogue of Cabibbo mixing in the quark sector. If the neutrinos that mix belong to different SU(2) representations, neutral-current couplings are flavor off-diagonal as well. The low-energy effective theory based on O(18) has both types of couplings.

The O(18) theory of family unification predicts five additional massive neutrino species at the weak scale. Four of the new species belong to mirror families, so they have opposite $SU(2)\times U(1)$ quantum numbers to the ordinary neutrinos. One of the new families is of the usual $V-A$ type. The full pattern of mixing among neutrinos is rather complicated[5] so we shall make a few simplifying assumptions. We will consider here the addition of a *single* neutrino species for two special cases: (1) a sequential neutrino with a mass eigenstate that contains a small component of an ultralight neutrino ν_l and (2) a mirror neutrino with a small component of ν_l. In the first case, neutral-current couplings will be flavor diagonal, whereas in the second case, couplings not allowed by the Glashow-Iliopoulos-Maiani mechanism appear in the neutrino sector. Furthermore, in both cases, we will also assume that there is no mixing among the charged leptons. (Mirror neutrinos have also been treated in Ref. 15, although those models are somewhat different from the one we consider here.)

In terms of the left-hand components of the mass eigenstates ν_l and N, we parametrize the mixing by

$$\begin{bmatrix}\nu_l'\\ N'\end{bmatrix}=\begin{bmatrix}1 & \epsilon\\ -\epsilon & 1\end{bmatrix}\begin{bmatrix}\nu_l\\ N\end{bmatrix}, \tag{1}$$

where primes indicate weak-interaction eigenstates. Then, the charged-current couplings of the neutrinos with the charged lepton l are

$$\mathscr{L}_{\rm CC}=\frac{g}{2\sqrt{2}}\bar{l}\gamma_\mu(1-\gamma_5)(\nu_l+\epsilon N)W^\mu+\text{H.c.} \tag{2}$$

When a mirror neutrino mixes with an ordinary neutrino, flavor—off-diagonal Z couplings occur. To order ϵ, the neutral-current couplings between mass eigenstates ν_l and N are

$$\mathscr{L}_{\rm NC}=\frac{g}{4\cos\theta_W}(\bar{\nu}_l\bar{N})\gamma_\mu(1-\gamma_5)\begin{bmatrix}1 & 2\epsilon\\ 2\epsilon & -1\end{bmatrix}\begin{bmatrix}\nu_l\\ N\end{bmatrix}Z^\mu . \tag{3}$$

Neutral-current decays of N to ν_l have a Cabibbo-type suppression factor of $4\epsilon^2$. We have calculated the branching fractions for possible N decays in the two special cases described above. Our results are displayed in Table I. (We have accounted for the charged-current interference with the neutral-current amplitude in the branching ratios for the mirror neutrino.)

Universality in charged weak currents restricts $|\epsilon_e|^2<0.043$ and $|\epsilon_\mu|^2<0.008$ for heavy-neutrino mixing with the electron or muon neutrino, respectively.[16] Experimental and theoretical uncertainties in the tau-lepton lifetime are such that ν_τ-N mixing is not as well constrained: $|\epsilon_\tau|^2\lesssim 0.14$ (Ref. 16). Despite the potentially enhanced signals of heavy neutrinos that couple to ν_τ, we do not consider that possibility here. As we shall see below, many characteristic decay signatures are tagged by the charged lepton whose neutrino partner mixes with N. If that charged lepton also decays within the detector, the signal is less easily distinguished from the backgrounds. Therefore, we will assume below that $l=e$ or $l=\mu$.

III. SIGNATURES OF NEW NEUTRINOS AT THE CERN COLLIDER

Sufficiently light neutrinos with very small couplings to ordinary charged leptons will be observable by their finite decay lengths. For example, sequential neutrinos which are pair produced in Z decays with $m_N=40$ GeV and $|\epsilon|\simeq 10^{-4}$ have mean decay lengths of approximately 1 mm (Ref. 17). In practice, we take $|\epsilon|\simeq 10^{-1}$—10^{-2}, motivated by the suggestion that ϵ should be related to the Cabibbo angle, as it is in the O(18) framework. For the range of neutrino masses from 10 to 40 GeV, and for such relatively large values of ϵ, the mean decay lengths are less than 1 mm. Gaps resulting from such decays may eventually be observable by using vertex detectors. We will proceed here under the assumption that heavy neutrino decays would appear as prompt decays in the UA1 and UA2 detectors.

Monojet-event signatures can result from Z decays into heavy mirror neutrinos. Both $Z\to N\bar{N}$ and $Z\to\nu_l\bar{N}+\bar{\nu}_lN$ may contribute to a signal of jets and missing transverse momentum (denoted by $\not{p}_T$). Consider the hadronic decay of a heavy neutrino: $N\to q\bar{q}+$lepton. Naively, one would expect that such a decay should be

TABLE I. Approximate branching fraction for a sequential and mirror neutrino with a mass in the range of 15–40 GeV and mixing with ν_e. All other members of the new generation containing N as well as the top quark are assumed to be kinematically inaccessible. The masses of fermions in the final state are taken to be negligible relative to the N mass.

Sequential neutrino		Mirror neutrino	
Channel	Branching fraction	Channel	Branching fraction
$e(\bar{e}\nu_e+\bar{\mu}\nu_\mu)$	0.22	$\nu_e(\bar{e}e+\bar{\mu}\mu)+e(\bar{\mu}\nu_\mu)$	0.08
$e(\bar{\tau}\nu_\tau)$	0.11	$\nu_e(\bar{\tau}\tau)+e(\bar{\tau}\nu_\tau)$	0.07
$e(\bar{q}q')$	0.67	$\nu_e(\bar{\nu}\nu)$	0.13
		$\nu_e(\bar{q}q)+e(\bar{q}q')$	0.72

characterized by two separated hadronic jets. However, one must take into account the jet selection criteria employed by the various experiments. For example, the UA1 Collaboration regards two energy clusters as one jet if they satisfy $(\Delta\eta)^2+(\Delta\phi)^2\leq 1$, where $\Delta\eta$ and $\Delta\phi$ are the difference between the pseudorapidity and the azimuthal angle of the two clusters, respectively. In addition, UA1 also requires that the transverse energy of a given cluster (after possible coalescing if the above condition is satisfied) be greater than 12 GeV, in order that it be counted as a separate jet. Thus, the hadronic decay of the N will often be observed in the experiment às containing only one hadronic jet. Using the branching ratios of Table I and assuming unit probability for N to decay into a single jet, we estimate the monojet cross section due to the presence of a heavy mirror neutrino to be

$$\frac{\sigma(Z\to\text{monojet}+\not p_T)}{\sigma(Z\to\nu_e\bar{\nu}_e)}\simeq 0.19(1-4|\epsilon|^2)F_{NN}(m_Z)+5.76|\epsilon|^2F_{N\nu}(m_Z)$$
$$\lesssim 0.22\ (m_N=20\ \text{GeV})$$
$$\lesssim 0.12\ (m_N=40\ \text{GeV})\ , \tag{4}$$

where $F_{NN}(m_Z)$ and $F_{N\nu}(m_Z)$ are the kinematic suppression factors

$$F_{NN}(m_Z)=\left[1-\frac{m_N{}^2}{m_Z{}^2}\right]\left[1-\frac{4m_N{}^2}{m_Z{}^2}\right]^{1/2},$$
$$F_{N\nu}(m_Z)=\left[1+\frac{m_N{}^2}{2m_Z{}^2}\right]\left[1-\frac{m_N{}^2}{m_Z{}^2}\right]^{2}. \tag{5}$$

A mixing angle $|\epsilon|^2<10^{-2}$ was used in Eq. (4). For reference, using $m_N=40$ GeV, $m_Z=94$ GeV, and $m_W=83$ GeV, we find $F_{NN}(m_Z)\simeq 0.43$, $F_{N\nu}(m_Z)\simeq 0.73$, $F_{NN}(m_W)\simeq 0.20$, and $F_{N\nu}(m_W)\simeq 0.66$. The monojets from Z decays cannot come from a sequential neutrino, since neutral-current decays of N are required. The monojet signal and backgrounds have been discussed in more detail by Hall, Kim, and Nelson.[18] [Hall, Kim, and Nelson[18] point out that the upper bounds in Eq. (4) for N mixing with ν_τ are somewhat larger. We find that the limits for N-ν_τ mixing corresponding to Eq. (4) are 0.83 and 0.63 for $m_N=20$ and 40 GeV, respectively.]

Purely leptonic Z decays would also provide a clear signal for heavy neutrinos. With both heavy neutrinos decaying into muons and/or electrons plus ultralight neutrinos, we find

$$\frac{\sigma(Z\to N\bar{N}\to l^+l^-l^+l^-\nu_l\bar{\nu}_l)}{\sigma(Z\to\nu_e\bar{\nu}_e)}\simeq 6.4\times10^{-3}F_{NN}(m_Z)\ \text{(mirror)}$$
$$\simeq 4.8\times10^{-2}F_{NN}(m_Z)\ \text{(sequential)}\ . \tag{6}$$

With the identification of four charged leptons, little missing transverse energy, and a reconstructed mass of approximately m_Z, this process is relatively background free.

Another striking signal with little missing transverse energy appears in W decays. Charged-gauge-boson decays $W\to lN$ followed by the subsequent hadronic N decay will be identified by an isolated charged lepton $l^\pm$ with $p_T\sim m_W/2$. Indeed, the UA1 and UA2 top-quark searches are carried out by using a hard-isolated-lepton trigger. Our estimate of the total cross section is

$$\frac{\sigma(W\to lN\to l+\text{jets})}{\sigma(W\to e\nu_e)}\simeq\begin{cases}0.67F_{N\nu}(m_W)|\epsilon|^2 & \text{(sequential)}\ ,\\ 0.72F_{N\nu}(m_W)|\epsilon|^2 & \text{(mirror)}\ .\end{cases} \tag{7}$$

One important feature of this process is that very often only one jet will appear in the final state after the UA1 jet criteria are applied. In particular, we shall focus on these monojets plus hard lepton events as being a good signal for new physics. In contrast, events with top quarks are characterized by two or more jets. For the remainder of this paper we will study those events which result in a single jet and an isolated hard lepton. For definiteness, we take $l=e$.

IV. ANALYSIS OF MONOJET EVENTS WITH A HARD ISOLATED ELECTRON

To determine the cross section for $p\bar{p}\to W\to eN$ followed by subsequent N decay into a single jet, we consider the parton subprocess $u\bar{d}\to W^+\to\bar{e}l_1q_1\bar{q}_2$. When N decays via charged currents, $l_1=e$, whereas for neutral-current decays of N, $l_1=\nu_e$. In either case, keeping general the vector and axial-vector couplings of the gauge boson V mediating the N decay, we find for the squared matrix element summed over initial and final spins and colors

$$|\mathcal{M}|^2=9\times2^{12}\times\frac{G_F{}^4m_W{}^8}{x_V}\frac{\pi}{\Gamma_Wm_W}\delta(xys-m_W{}^2)\frac{\pi}{\Gamma_Lm_N}\delta(N^2-m_N{}^2)$$
$$\times(s_{q_1\bar{q}_2}-M_V{}^2)^{-2}[(v_V{}^2+a_V{}^2)\mathcal{F}-2v_Va_V\mathcal{G}]\,|U_W|^2|U_V|^2\ , \tag{8}$$

where

$$\mathscr{F}=(N\cdot q_1 l_1\cdot\bar{q}_2+N\cdot\bar{q}_2 l_1\cdot q_1)N\cdot\bar{d}\,\bar{e}\cdot u-\tfrac{1}{2}m_N{}^2\bar{e}\cdot u(l_1\cdot q_1\bar{q}_2\cdot\bar{d}+l_1\cdot\bar{q}_2 q_1\cdot\bar{d})\ ,$$
$$\mathscr{G}=(N\cdot\bar{q}_2 l_1\cdot q_1-N\cdot q_1 l_1\cdot\bar{q}_2)N\cdot\bar{d}\,\bar{e}\cdot u-\tfrac{1}{2}m_N{}^2\bar{e}\cdot u(l_1\cdot q_1\bar{q}_2\cdot\bar{d}-l_1\cdot\bar{q}_2\bar{d}\cdot q_1)\ . \tag{9}$$

The four-momentum of each particle is denoted by the corresponding particle symbol, and

$$x_W=1,\ \ U_W=\epsilon,\ \ v_W=1,\ \ v_Z=2T_3-4Q\sin^2\theta_W\ ,$$
$$x_Z=4\cos^4\theta_W,\ \ U_Z=2\epsilon,\ \ a_W=-1,\ \ a_Z=-2T_3\ . \tag{10}$$

The heavy neutrino has $V-A$ couplings, as indicated in Eq. (3), so here, Q refers to the charge of the quark in the final state and T_3 is $\pm\frac{1}{2}$, the weak isospin of the quark. The matrix element for $\bar{u}d\to W^-\to e\bar{l}_1\bar{q}_1 q_2$ is obtained from Eqs. (8)–(10) by making the following substitutions: $l_1\to e$, $\bar{e}\to\bar{l}_1$, $u\to-\bar{q}_1$, $\bar{d}\to-q_2$, $q_1\to-\bar{u}$, and $\bar{q}_2\to-d$. In this calculation, we have taken both W and N to be on mass shell and have used the narrow-width approximation. For $V=W$, we have the matrix element relevant to the case of a heavy sequential neutrino. Mirror neutrinos decay via both W- and Z-mediated processes. The matrix element squared for the parton subprocess is converted into a differential cross section in the usual manner

$$d\sigma=\int dx\,dy\frac{1}{(2\pi)^8}\frac{1}{9\times 4}\sum|\mathscr{M}|^2\delta^4\left[\sum\text{momenta}\right]\frac{K}{2xys}f_u(x)f_{\bar{d}}(y)\frac{d^3\bar{e}}{2E_{\bar{e}}}\frac{d^3l_1}{2E_l}\frac{d^3q_1}{2E_1}\frac{d^3q_2}{2E_2}\ . \tag{11}$$

The functions f_u and $f_{\bar{d}}$ are the parton distribution functions, both valence and sea, for charge $\frac{2}{3}$ and $-\frac{1}{3}$ quarks, respectively. Higher-order perturbative corrections tend to enhance the theoretical rate somewhat; we have taken this into account by inserting a K factor of $K=1.5$ in Eq. (11).

The quark distribution functions used in the calculation are those of Eichten, Hinchliffe, Lane, and Quigg.[19] Phase-space integrals in the differential cross section are performed using a Monte Carlo integration technique. (See Ref. 20 for a detailed discussion of our techniques.) Events are generated, appropriately weighted, and passed through cuts designed to simulate those of the UA1 detector. Final-state quarks are interpreted as hadronic energy clusters with four-momenta q_1 and $\bar{q}_2$, respectively.

The two clusters are coalesced into a single jet when $\Delta R\equiv[(\Delta\phi)^2+(\Delta\eta^2)]^{1/2}\le 1$. We select events in which the two quarks are considered to form a single jet, or in which one quark passes the transverse-momentum cut, $p_T^{\rm jet}>12$ GeV, but the second quark does not. For neutrino decays via charged currents, we further require that the electron resulting from the decay of the N meet one of the following requirements: (1) $p_T^e<15$ GeV, (2) $|\eta_e|>2.5$, or (3) $\Delta R_{e,\rm jet}<0.4$. For (1) and (2), the electron is either lost or indistinguishable from a charged pion. For (3), the electron is embedded in the jet. The electron from the W decay must not satisfy any of the three conditions above.

The neutrino or missing charged lepton in the N decay will have small transverse energy, so we impose a cut on the missing transverse energy

$$E_T^{\rm miss}<15\ \text{GeV or } E_T^{\rm miss}<4\sigma\ , \tag{12}$$

where $\sigma=0.7\sqrt{E_T}$ and E_T is the total scalar transverse energy of the event (obtained experimentally by summing the transverse energies deposited in each calorimeter cell). The quantity E_T is not well understood theoretically, since a significant part of it is due to soft nonperturbative processes. We estimate E_T to be

$$E_T=E_T^e+\sum E_{T_i}+E_r\ , \tag{13}$$

where E_T^e is the transverse energy of the primary electron, E_{T_i} is summed over the observable decay products of the N, and E_r is the residual transverse energy coming from the underlying "background" of the event. (We take E_r to be distributed with a mean of about 40 GeV; for further details, see Ref. 20.)

We have imposed two further cuts to reduce the standard-model background coming from $p\bar{p}\to W+(g$ or $q)$. First, we require that $p_T^{\rm jet}>20$ GeV. Second, we require the transverse mass of our monojet-plus-electron sample to lie in the range 60 GeV $<m_T<$ 100 GeV. Most of the $p\bar{p}\to W+(g$ or $q)\to e+$jet background is eliminat-

TABLE II. The total cross sections (after cuts) are given for $p\bar{p}\to e+$jet arising from the production of a heavy neutrino N with mass as shown above (assuming a mixing parameter of $\epsilon_e=0.1$). For comparison, we also give the total cross section (after cuts) for two important background processes: the production of a $W^\pm$ plus a jet (computed to lowest order in QCD), and $W^\pm\to tb$ (assuming $m_t=40$ GeV). The cross sections above are obtained by imposing a set of UA1-inspired cuts, as discussed in the text.

Mass (GeV)	$\sigma_{\rm cut}$ (pb)
15	1.3
25	0.9
40	0.4
$(W^\pm\to\nu e)+(g$ or $q)$	0.9
$W^\pm\to tb$	0.7

ed, as we shall see below.

Our results are presented in Figs. 1(a)–1(c). Figure 1(a) illustrates the differential cross section versus the transverse momentum of the primary electron, for the case of a sequential heavy neutrino of masses 15, 25, and 40 GeV with $\sqrt{s}=630$ GeV. Figures 1(b) and 1(c) depict the differential cross section versus the jet transverse momentum and the invariant mass of the electron-jet system, respectively. For sequential neutrinos, the plots are nearly identical so we do not exhibit them here. The total cross sections for the three different masses, including cuts, are shown in Table II.

These cross sections are rather small, but we find that the backgrounds from $p\bar{p}\to W+g$ or $W+q$ to be even smaller. Geer and Stirling[21] have calculated that $\sigma(p\bar{p}\to W+\text{jet})\simeq 0.26$ nb for $\sqrt{s}=630$ GeV, $|\eta^{\text{jet}}|$

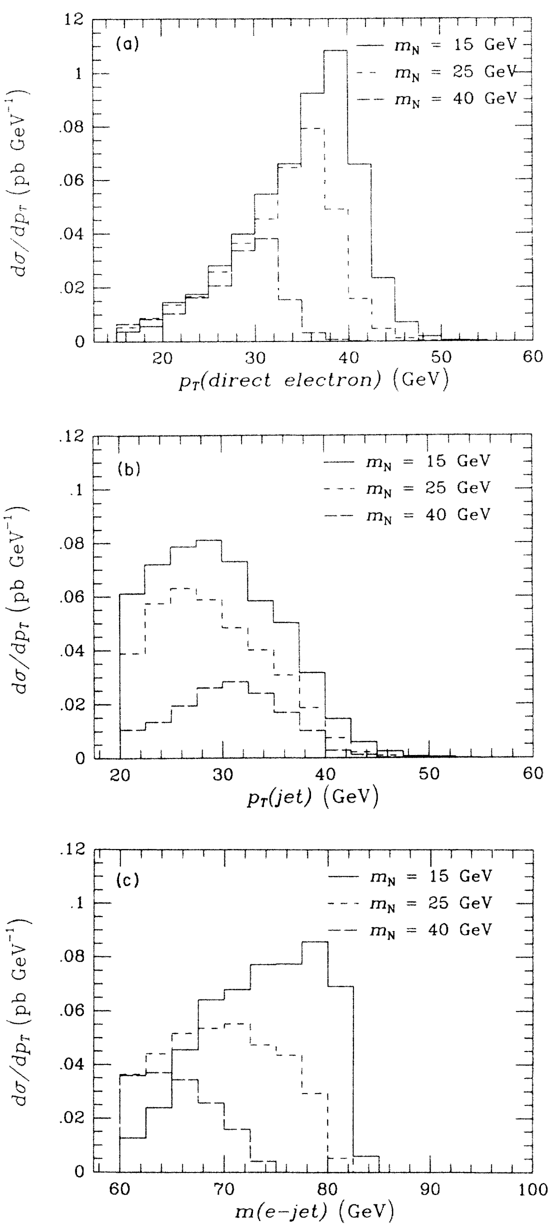

FIG. 1. Differential cross section versus (a) transverse momentum of the primary electron, (b) transverse momentum of the jet, and (c) invariant mass of the electron-jet system, for $p\bar{p}\to W^\pm\to e^\pm+$jet with $m_N=15$, 25, and 40 GeV. (Cross sections for both charges have been added.) All cuts, as described in the text have been imposed. The neutrino N has both charged- and neutral-current decays. We have taken the neutrino mixing parameter to be $\epsilon=0.1$.

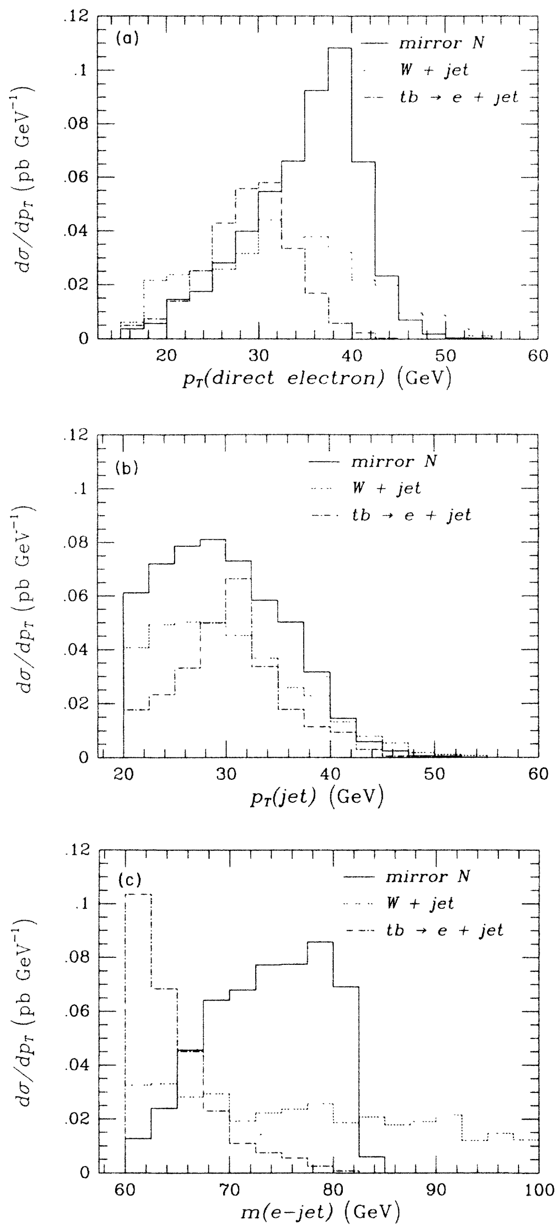

FIG. 2. Differential cross section versus (a) transverse momentum of the primary electron, (b) transverse momentum of the jet, and (c) invariant mass of the electron-jet system, for $p\bar{p}\to e^\pm+$jet for a mirror neutrino with $m_N=15$ GeV, compared with backgrounds from $p\bar{p}\to W^\pm+\text{jet}\to e^\pm+$jet and from $W^\pm\to t\bar{b}(\bar{t}b)\to e^\pm+$jet. All cuts, as described in the text, have been imposed. We have taken the neutrino mixing parameter to be $\epsilon=0.1$.

<1.75, and $p_T^{\text{jet}}>20$ GeV. Including the branching ratio for $W\rightarrow e\nu$, $B_{e\nu}\simeq 0.09$, the total cross section reduces to 0.023 nb. For these cuts with our jet-finding algorithm, we find a total cross section $\sigma\simeq 0.016$ nb. The more stringent cuts outlined above further suppress the background to $\sigma\simeq 0.9$ pb. This is just below the signal of a heavy neutrino with a mass of 15 GeV and a mixing parameter given by $|\epsilon|^2\simeq 10^{-2}$. We illustrate the differential cross sections from these backgrounds together with the signal from a neutrino with $m_N=15$ GeV in Figs. 2(a)–2(c).

Other processes that contribute to the single-electron-plus-jet background include $Z\rightarrow\tau\bar{\tau}$, $Z\rightarrow q\bar{q}$, and $W\rightarrow q\bar{q}'$ where q and $\bar{q}'$ are heavy quarks. For Z decays, the higher invariant mass of the Z can presumably be used to discriminate these sources of an isolated lepton plus a single jet. For both the Z and W decays, the electron comes from a sequential decay, so typically will not satisfy both p_T^e and isolation cuts.

As an example of the heavy-flavor background that produces an isolated electron, we consider specifically $W\rightarrow t\bar{b}$. Our parton model estimate of its contribution to the electron plus jet signal is below the level of the signal itself when our cuts are imposed. As can be seen from Fig. 2(a), the transverse-momentum distribution of the charged lepton is peaked lower than that of the relatively light (15 GeV) neutrino, as is the electron-jet invariant mass. A stronger cut on p_T^e and $m_{e\text{-jet}}$ would effectively eliminate this background.

In addition, Drell-Yan heavy-quark production followed by a semileptonic decay of one of the quarks could mimic the signal. Typically, the transverse momenta of the charged leptons are low.[22] Requiring the invariant mass of the Drell-Yan produced heavy quarks to match that of the W further reduces this contribution, as does the requirement that the lepton be isolated.

V. CONCLUSIONS AND DISCUSSION

Many of the present experimental bounds on heavy neutrino mass and mixing angle rely on a measurable decay length. Others measure the lepton end-point spectra of meson decays. Typically, neither are sensitive to the range of neutrino masses and mixing angles considered here: m_N in the 10–40-GeV mass range and $\epsilon\simeq 0.1$. In order to constrain this region of parameter space, one can make use of distinctive decay signatures, as described in this paper.

From the purely leptonic decay modes of pair-produced heavy neutrinos, signals are expected at the level of 10^{-1}–10^{-2} of the cross section for Z decays into e^+e^-. Limits on the neutrino mass come only from the phase-space suppression. Since the cross section for pair-produced neutrinos is independent of the mixing angle, this signature is not useful for placing bounds on mixing angles.

More interesting for constraints on mixing angles are the flavor nondiagonal Z and W decays. Our focus here is on the W decay into an ordinary charged lepton and a heavy neutrino that subsequently decays hadronically. We have calculated the cross section for $W\rightarrow eN$ to produce an isolated charged lepton balanced by a monojet that reconstructs (roughly) to the W mass. By making use of a Monte Carlo technique which approximates the UA1 jet-finding algorithm and cuts, we find that for $\epsilon=0.1$, the heavy-neutrino signal would appear at the picobarn level. As far as we know, no published data exists at present for such events; *a careful investigation of the collider data for event signatures of this type would be most desirable.* With several hundred decays of W into an electron and its neutrino partner, one may begin to see this lepton-flavor-nonconserving process. Both the QCD background and backgrounds from top-quark production are comparable to or below the signal for $M_N\lesssim 25$ GeV. Both mirror and sequential neutrinos can yield such a signal. Mirror neutrinos, however, may also exhibit monojet-plus-missing-transverse-momentum signatures when produced by Z decays, whereas sequential neutrinos cannot.

Absolutely massless neutrinos are a special feature of the Glashow-Weinberg-Salam model. If we go beyond the standard model and attempt to explain why quarks and leptons belong to families and why the families repeat, we are led to theories which contain massive neutrinos. Although the ordinary neutrinos are ultralight, additional neutrinos need not be. Massive neutrinos appear naturally in O(18), which predicts[5] five new families of neutrinos with masses below 40 GeV. With the integrated luminosity of the CERN collider approaching 1 pb^{-1}, experiments are just reaching the edge of sensitivity to this proposed new physics. More dedicated running time is needed to confirm or exclude sequential and mirror neutrinos at this energy scale.

ACKNOWLEDGMENTS

We would like to thank H. Baer for useful discussions. We are grateful for the hospitality of the SLAC Theory group where some of this work was done. This work was supported in part by the Department of Energy. Fermilab is operated by Universities Research Association Inc. under contract with the United States Department of Energy.

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