Supercollider signatures and correlations of heavy neutrinos

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Forthcoming proton supercolliders can discover a heavy neutrino lying hidden in the 10^2-10^3 GeV mass range, one that may harbinger a new set of heavy particles. For an unstable neutrino, the optimal search strategy will be to study a final state of charged dileptons and jets without missing p_T. Dilepton charge signs *or* angular correlations can discriminate between Dirac and Majorana heavy neutrinos.

This work is concerned with very heavy neutrinos which have practically no perceptible effects in cosmology, astrophysics, or low-energy nuclear and particle physics. These can only be studied at high-energy accelerators through their coupling to the weak neutral Z boson. Current restrictions on nonstandard Z decay from CERN LEP 1 data require such a neutrino (provided its coupling to the Z is not reduced by an unnaturally small mixing angle) to be heavier than about $\frac{1}{2}M_Z$. LEP 2 will be able to probe more massive neutrinos with masses $\lesssim 100$ GeV. Still heavier ones can only be sought in the proposed supercolliders where such a quest will be an exciting venture. Here we discuss the means of discovering and studying a heavy neutrino in the 10^2-10^3 GeV range via pair production in pp collisions at supercollider energies of tens of TeV in the center of mass. If such a neutrino exists, along with extra heavy quarks, the prospects are shown to be excellent.

Much theoretical speculation abounds on very heavy neutrinos in scenarios going beyond the standard model such as those with a fourth-generation [1], left-right symmetry [2], or superstring-generated E_6 grand unification [3]. In many of these it is theoretically reasonable to expect the occurrence of a heavy neutrino in the region of hundreds of GeV. Moreover, such a neutrino should naturally couple to the Z, either directly or through a mixing mechanism. On the other hand, it is likely to be associated with other heavy particles arising from the new physics that is responsible for the generation of the heavy neutrino mass. To be definite, we work within a scenario with at least one extra heavy-quark doublet (U,D), though an extra singlet quark would suffice.

The question of the Dirac or Majorana character of a heavy neutrino is important. Let N_l be such a neutrino of mass M_N which, after production, decays (via mixing of N_l to ν_l) into a light charged lepton l (e, μ, or τ) and a real or virtual charged weak gauge boson within the detector. We confine our attention to its inclusive semileptonic decay. One then has a distinct signature, namely, dileptons and jets (mostly four in number) without missing p_T. We use superscripts D and M to refer to the Dirac and Majorana cases, respectively. (a) In the Dirac case $N_l^D \neq \bar{N}_l^D$ and the pair-produced neutrinos are $(N_l^D, \bar{N}_l^D)$. These yield unlike-sign charged dileptons via $N_l^D \to lX$, $\bar{N}_l^D \to \bar{l}\,\bar{X}$, where X consists mainly of two jets. (b) In the Majorana case $N_l^M = \bar{N}_l^M$ and the produced neutrinos make the identical pair (N_l^M, N_l^M) with consequent decays $N_l^M \to lX$, $N_l^M \to \bar{l}\,\bar{X}$ with equal total rates [4] yielding both like- and unlike-sign charged dileptons. Knowledge of the sign of l would make the distinction [4] between cases (a) and (b) relatively straightforward. However, the sign determination of a very fast l—as produced in a supercollider—may not prove easy. A study of the angular correlation of the dilepton pair can, nonetheless, discriminate between the two cases.

Signatures and correlations of dileptons among the decay products of pair-produced heavy neutrinos in fermion-antifermion annihilation have already been discussed [5] and can be used for the quark annihilation subprocess $q\bar{q} \to Z^* \to N_l \bar{N}_l$, where Z^* is an off-shell Z. However, as shown by Willenbrock and Dicus [6], gluon fusion into an off-shell Z^* mediated by a heavy-quark loop, is generally more important for heavy-lepton pair production in pp collision at supercollider energies. There are two main reasons for this. The quark annihilation cross section falls off as $\hat{s}^{-1}$ for large $\hat{s}$, where $\hat{s}$ is the subprocess incident total four-momentum squared. In contrast, the gluon-fusion cross section does not decrease as much with $\hat{s}$. Yang's theorem, applied to the ggZ^* vertex, generates [7] a factor proportional to $\hat{s}-M_Z^2$ in the latter amplitude canceling the $(\hat{s}-M_Z^2)^{-1}$ factor of the Z propagator. Additionally, between gluon and quark luminosities, the former far exceed the latter at supercollider energies. We shall, therefore, largely concentrate on the $gg \to Z^* \to N_l \bar{N}_l$ process comparing with quark annihilation only at the production level.

The key element in the amplitude under consideration is the ggZ^* vertex $v_Q^{\mu\nu\rho}$ mediated by a heavy-quark (Q) triangular loop. The indices μ,ν correspond to the two gluons and ρ to the Z^*. The two diagrams to the lowest order lead to [8]

$$v_Q^{\mu\nu\rho}(q_1,q_2)=(-1)_Q\frac{\alpha_S}{2\pi}e(\sin\theta_W\cos\theta_W)^{-1}\hat{s}^{-1}\epsilon^{\mu\nu\alpha\beta}q_{1\alpha}q_{2\beta}(q_1+q_2)^\rho\left[1+2m_Q^2\hat{s}^{-1}\int_0^1\frac{dx}{x}\ln[1-\hat{s}m_Q^{-2}x(1-x)+i\epsilon]\right]. \tag{1}$$

In (1) q_1,q_2 are gluon four-momenta, e is the electron charge, α_S is the QCD fine-structure constant, m_Q is the mass of the heavy quark Q in the triangular loop, θ_W is the Weinberg angle, and $\hat{s}=(q_1+q_2)^2$. Moreover, the sign $(-1)_Q$, originating from the quantum number T_{3LQ} of $SU(2)_L$, is positive or negative for Q being an up-type or down-type quark, respectively. The first term within the large parentheses in (1) drops out when summed over each quark doublet and the second term makes an appreciable contribution only for generation 3 or higher. We take the $ZN_l\bar{N}_l$ vertex to be the standard $Z\nu\bar{\nu}$ coupling times an unknown mixing parameter $\sin\phi$. Then, to the lowest nontrivial order, the subprocess $g_a(q_1,\epsilon_1)g_b(q_2,\epsilon_2)\rightarrow Z^*\rightarrow N_l(p_3,S)\bar{N}_l(p_4,\bar{S})$ has the differential cross section

$$\begin{aligned}\frac{d\hat{\sigma}^D}{d\Omega_N}=\frac{1}{4}\frac{d\hat{\sigma}^M}{d\Omega_N}&=\frac{1}{2048}\left[\frac{\alpha_S}{\pi}\right]^2\alpha_{\rm EM}^2[\sin^4(2\theta_W)]^{-1}M_N^2M_Z^{-4}\hat{s}^{-2}\sin^2\phi(1-4M_N^2\hat{s}^{-1})^{1/2}(1-\zeta\cdot\bar{\zeta})\\&\times\left|[1+iM_Z\Gamma_Z(\hat{s}-M_Z^2)^{-1}]^{-1}\sum_Q(-)_Q m_Q^4\int_0^1\frac{dx}{x}\ln[1-\hat{s}m_Q^{-2}x(1-x)+i\epsilon]\right|^2\\&\equiv f(\hat{s})(1-\zeta\cdot\zeta)\ ,\end{aligned} \tag{2}$$

where ζ and $\bar{\zeta}$ represent the spin vectors of N and $\bar{N}$ in their respective *rest frames* and the function $f(\hat{s})$ stands explicitly defined. Furthermore, for total cross sections $\hat{\sigma}^D$ equals $\frac{1}{2}\hat{\sigma}^M$ because of two final identical particles in the Majorana case.

The components of the neutrino momentum and spin four-vectors can be written in the subprocess c.m. frame in terms of the velocity β as

$$p_3=(E,0,0,\beta E),\quad p_4=(E,0,0,-\beta E)\ ,$$
$$S=[\beta(1-\beta^2)^{-1/2}\zeta_z,\zeta_x,\zeta_y,(1-\beta^2)^{-1/2}\zeta_z]\ ,$$
$$\bar{S}=[-\beta(1-\beta^2)^{-1/2}\bar{\zeta}_z,\bar{\zeta}_x,\bar{\zeta}_y,(1-\beta^2)^{-1/2}\bar{\zeta}_z]\ .$$

The subprocess total cross sections are related to the rapidity distribution in a pp collision with the square of c.m. energy s via

$$d\sigma^{D,M}/dy=\int d\tau F_G(\sqrt{\tau}e^y)F_G(\sqrt{\tau}e^{-y})\hat{\sigma}^{D,M}(\tau s)\ .$$

Here $\tau=\hat{s}/s$, the τ integration goes from $4m_N^2/s$ to $e^{-2|y|}$, $F_G(x)$ is the distribution of gluons carrying a fraction x of the proton momentum, and y is the rapidity of the subprocess c.m. frame with respect to the lab frame. Also, σ equals $\int dy\, d\sigma/dy$ with integration limits $\pm\ln(\sqrt{s}/2M_N)$. We can present the corresponding results numerically. Figure 1 shows the quark annihilation and gluon-fusion contributions to the total cross section for the process $pp\rightarrow N_l\bar{N}_l+$soft particles (Dirac case, $\phi=45°$), as explained in the caption. For $M_N<200$ GeV at the Superconducting Super Collider (SSC) and $M_N<300$ GeV at the CERN Large Hadron Collider (LHC), the signals are large but predominantly from quark annihilation. On the other hand, for 500 GeV $<M_N<1.5$ TeV at the SSC and for 700 GeV $<M_N<1.2$ TeV at the LHC (higher luminosity for the latter can compensate for smaller rates), there is clear domination by the gluon fusion mechanism. Beyond the upper limits the rates become too small. In the inter-

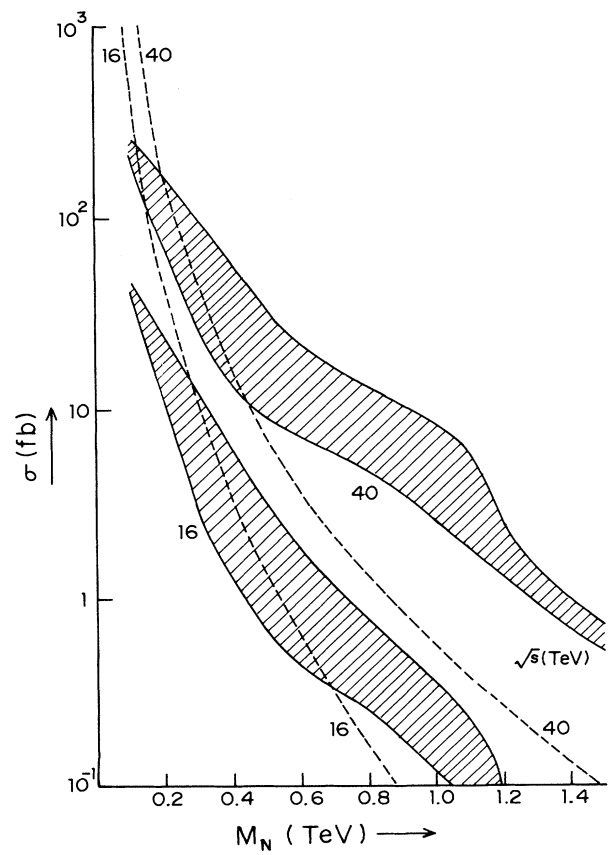

FIG. 1. Illustrative plots of the total cross sections (in femtobarns) of the process $pp\rightarrow N_l\bar{N}_l+\cdots$ due to gluon fusion (shaded bands corresponding to a variation in the top mass from 130 to 170 GeV and in the mass M_U of the U from 400 to 1200 GeV with $M_D=M_U-100$ GeV) and quark annihilation (dashed curves) as a function of the heavy neutrino mass at $\sqrt{s}=16$ TeV (LHC) and 40 TeV (SSC).

mediate region, the two processes are comparable.

We now consider the inclusive semileptonic decays of the produced $N_l, \bar{N}_l$ produced exclusively by gluon fusion and focus on the angular correlation [9] of the charged dilepton system. Following Ref. [9], we assign the respective polarization vectors [10] $\mathbf{w}_-$ and $\mathbf{w}_+$ to N_l and $\bar{N}_l$. Take the charged leptons emanating from the corresponding semileptonic decays to be l and $\bar{l}$; neglect their masses but assign three-momenta $\mathbf{q}$ and $\bar{\mathbf{q}}$ to them, respectively (in the Majorana case they could also be both l's or $\bar{l}$'s but the analysis below still goes through). Let us, in fact, be more definite and specify the heavy neutrino decays as $N_l \to lW$, $\bar{N}_l \to \bar{l}\,\bar{W}$ and $W, \bar{W} \to$hadrons with left-handed couplings. It is convenient to make the definitions

$$C_N = \frac{M_N}{128} \alpha_{\rm EM} B \sin^{-2}\theta_W \sin^{-2}\chi (1 - M_W^2 M_N^{-2}) ,$$
$$\alpha_N = 1 + M_N^2 M_W^{-2} - 2M_W^2 M_N^{-2} , \quad (3)$$
$$\beta_N = 4(1 - \tfrac{1}{2} M_N^2 M_W^{-2}) ,$$

where χ is the mixing angle through which $N_l, \bar{N}_l$ couple to the W and B is the branching ratio for the decay $W \to$hadrons, being nearly 2/3. Now the differential decay rates can be given as

$$\frac{d\Gamma_N}{d\Omega_l} = C_N(\alpha_N - \beta_N M_N^{-1} \mathbf{q}\cdot\mathbf{w}_-) , \quad (4a)$$

$$\frac{d\Gamma_{\bar{N}}}{d\Omega_{\bar{l}}} = C_N(\alpha_N + \beta_N M_N^{-1} \bar{\mathbf{q}}\cdot\mathbf{w}_+) . \quad (4b)$$

The sign difference in the polarization-dependent terms of (4a) and (4b) arises from the interrelation between the helicity of the decaying neutrino and the momentum of the charged lepton forced by the parity-violating part of the coupling. The total decay rates $\Gamma(N_l \to lX)$ and $\Gamma(\bar{N}_l \to \bar{l}\,\bar{X})$ are both equal to $4\pi C_N \alpha_N$.

Given the decay distributions (4) and the heavy neutrino production rates (2), one can use the formalism given in Sec. IV of Ref. [9] [in particular, Eq. (4.26)] to obtain the combined angular distribution of the decay products for a fixed production angle, i.e., $d^3\sigma / d\Omega_N d\Omega_l d\Omega_{\tilde{l}}$ where the dilepton pair is generically denoted as $l, \tilde{l}$. In the Dirac case $\tilde{l}$ is always $\bar{l}$ and one has

$$\frac{d^3\hat{\sigma}^D}{d\Omega_N\, d\Omega_l\, d\Omega_{\tilde{l}}} = f(\hat{s})[\Gamma_{\rm tot}^D]^{-2} C_N^2 (\alpha_N^2 + \beta_N^2 \mathbf{q}\cdot\bar{\mathbf{q}} M_N^{-2}) . \quad (5)$$

Here $\Gamma_{\rm tot}^D$ is the total width of N_l^D (or $\bar{N}_l^D$). The right-hand side (RHS) of (5) clearly shows a strong dependence on the angle between the detected leptons. The triple differential distribution is a maximum when the latter are parallel and a minimum when they are antiparallel.

In the Majorana case, there are three possibilities: $l = \tilde{l} = l^-$, $l = \tilde{l} = l^+$, and $l \neq \tilde{l} = \bar{l}$. The corresponding combined angular distributions are

$$\frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^-}\, d\Omega_{l^-}} = \frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^+}\, d\Omega_{l^+}} = 4f(\hat{s})(\Gamma_{\rm tot}^M)^{-2} C_N^2 (\alpha_N^2 - \beta_N^2 \mathbf{q}\cdot\bar{\mathbf{q}} M_N^{-2}) , \quad (6a)$$

$$\frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^+}\, d\Omega_{l^-}} = 8[\Gamma_{\rm tot}^D / \Gamma_{\rm tot}^M]^2 \frac{d^3\hat{\sigma}^D}{d\Omega_N\, d\Omega_l\, d\Omega_{\tilde{l}}} . \quad (6b)$$

In (6) $\Gamma_{\rm tot}^M$ is the total decay width of the Majorana neutrino N_l. (If the field ψ_{N_l} couples only to the lepton field ψ_l and the charged-weak-boson field W_μ, then $\Gamma_{\rm tot}^M = 2\Gamma_{\rm tot}^D$.) If the lepton momenta are detected but their charge signs are left undetermined, the data would yield

$$\frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_l\, d\Omega_{\tilde{l}}} = \frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^+}\, d\Omega_{l^+}} + \frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^-}\, d\Omega_{l^-}} + \frac{d^3\hat{\sigma}^M}{d\Omega_N\, d\Omega_{l^+}\, d\Omega_{l^-}} = 16 f(\hat{s})[\Gamma_{\rm tot}^M]^{-2} C_N^2 \alpha_N^2 . \quad (7)$$

The RHS of (7) shows no dependence at all on the opening angle of the dileptons in marked contrast with (5). This difference can be experimentally utilized to distinguish between a Dirac and a Majorana heavy neutrino. This distinction is independent of the unknown mixing angles ϕ and χ. *We emphasize the novelty and significance of this result which arises out of the specific structure of the ggZ* vertex. In contrast, in the quark annihilation case, such a simple method of distinction is not possible* [5].

A left-handed $WN_l\bar{l}$ coupling was explicitly chosen for the above discussion. In principle, the heavy neutrino could have only a right-handed coupling through some heavier new gauge boson W'. The entire analysis would still go through *mutatis mutandis.* A new semiweak coupling constant squared over 4π would need to replace $\alpha_{\rm EM}$ and the signs of $\mathbf{q}$- or $\bar{\mathbf{q}}$-dependent terms in (4), would have to flip. The production cross section (2) would only change its overall factor $|\alpha_{\rm EM}[1 + iM_Z\Gamma_Z(\hat{s} - M_Z^2)^{-1}]^{-1}|^2$ in that there would be an extra term within the modulus involving the new coupling and the mass and width of the Z' corresponding to the W'. (The last remark also applies for any extra heavy neutral vector boson of superstring-motivated E_6 models.) The decay analysis would change through the off-shell effect of the W', there being an extra integral left over from the three-body phase space. *Nevertheless, the general statement about the distinct dilepton angular correlations in the Dirac and Majorana cases would still hold.*

We have not gone into a detailed discussion here of the kinematic cuts that would be needed on the data to keep the signal of a lepton-pair plus four jets without missing p_T significantly above backgrounds. Such backgrounds could arise in several ways. One example is special configurations in heavy-quark pair production where the heavy quarks decay semileptonically and the remnant lighter quarks themselves decay hadronically, jets merge and missing p_T's cancel out. Our signal is sizable by supercollider standards and we believe that all such backgrounds can be [11] kept below the signal by suitable lepton isolation and p_T-conservation criteria at least at the SSC.

In summary, we have pointed out how a massive neutrino in the 10^2-10^3 GeV mass range can be probed in forthcoming *pp* supercolliders, our focus being on the dominance of the gluon fusion mechanism for 0.5 TeV $< M_N <$ 1.5 TeV at the SSC and for 0.7 TeV $< M_N <$ 1.2 TeV at the LHC. Characteristic ways of differentiating between Dirac and Majorana heavy neutrinos have been highlighted. Produced in a pair, the former will yield only unlike-sign dileptons while the latter can lead to both like- and unlike-sign ones. The corresponding dilepton angular correlations are very different, being pronounced for the former but nonexistent for the latter so long as gluon fusion dominates. Both the LHC and the SSC are capable of discovering a heavy neutrino in the proposed mass range. But the isolation and study of the gluon fusion mechanism with its unique feature will be easier at the SSC because of larger cross sections. Evidently, heavy neutrino search needs to be given high priority in plans for supercollider experiments.

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[8] Given a real and positive $\alpha, \frac{1}{2}\int_0^1(dx/x)\ln[1-\alpha x(1-x)+i\epsilon]$ equals $-(\arcsin\sqrt{\alpha}/2)^2$ and $(\mathrm{arccosh}\sqrt{\alpha}/2)^2-\pi^2/4+i\pi\,\mathrm{arccosh}(\sqrt{\alpha}/2)$ for $0\leq\alpha\leq 4$ and $4\leq\alpha$, respectively.
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