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# Heavy-Majorana-neutrino production

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Heavy-Majorana-neutrino production and detection at colliders is examined. Majorana-neutrino decay is shown to be simply calculable from Dirac-neutrino results. Majorana-neutrino production $f\bar{f} \to Z \to NN$ is calculated and compared to Dirac-neutrino production. Near the threshold, the Dirac cross section rises much more quickly. Far above the threshold, the two types of neutrinos have the same spin-independent cross sections while the spin-dependent cross sections are very different. This difference could be observed in the correlations between the angle and energy dependence of the heavy-neutrino decay products.

Massive neutrinos can come in two different types: as Dirac or Majorana particles. Dirac fermions have distinct particle and antiparticle degrees of freedom while Majorana fermions make no such distinction and have half as many degrees of freedom.[1] Thus fermions with conserved charges such as color, electric charge, baryon number, or lepton number must be of the Dirac type, while fermions without conserved charges may be of either type. The distinction between Dirac and Majorana neutrinos vanishes if the mass vanishes because then both types are equivalent to two-component Weyl fermions since the standard weak interactions couple only to left-handed states. Experimental results on the currently known three generations[2,3] of neutrinos constrain the masses to be small but allow the neutrinos to be of either the Dirac or Majorana type.

New neutrinos could have large masses and be of either type. There are extensive experimental limits on heavy, fourth-generation, neutrinos from colliders assuming that the neutrinos are of the Dirac type.[4–6] These limits do not directly apply to Majorana neutrinos since production and decay properties are different for the two types. In this paper we discuss some of the relevant issues for extending these limits to neutrinos of the Majorana type.

## I. MAJORANA-NEUTRINO PRODUCTION

One of the best possibilities for producing heavy neutrinos is through a real or virtual $Z$, $f+\bar{f} \to Z \to N+\bar{N}$. Thus the CERN collider LEP and the SLAC Linear Collider will be good places to search for heavy neutrinos. However the cross section for this process will depend on whether the neutrinos are Dirac or Majorana particles. A Majorana-neutrino pair can coherently interfere with each other since particle and antiparticle are identical while a Dirac pair will not interfere. The Feynman diagrams for the Majorana-neutrino production process are shown in Fig. 1. For an electron (or quark) coupling of $\bar{f}\gamma^\mu(P_L-2x)f$ ($x=\sin^2\theta_W$ times 1, $\frac{1}{3}$, or $\frac{2}{3}$ for $f=e$, $d$, or $u$), $P_L=(1-\gamma_5)/2$, and a fourth-generation neutrino of mass $M$, the differential cross section is

$$\begin{aligned}\frac{d\sigma^M}{dt}=\frac{G_F^2}{4\pi\mathscr{s}^2}|R(\mathscr{s})|^2\big((1-4x)M[p\cdot(\bar{s}-s)\bar{p}\cdot(q-\bar{q})-\bar{p}\cdot(\bar{s}-s)p\cdot(q-\bar{q})]\\ +(1-4x+8x^2)\{(p\cdot\bar{q})(\bar{p}\cdot q)+(p\cdot q)(\bar{p}\cdot\bar{q})-M^2(p\cdot\bar{p})+(q\cdot\bar{q}-M^2)[(\bar{s}\cdot p)(s\cdot\bar{p})+(\bar{s}\cdot\bar{p})(s\cdot p)]\\ -(\bar{s}\cdot q)[(\bar{q}\cdot p)(s\cdot\bar{p})+(\bar{q}\cdot\bar{p})(s\cdot p)]+(\bar{s}\cdot q)(p\cdot\bar{p})(\bar{q}\cdot s)-(\bar{s}\cdot s)(p\cdot\bar{p})(\bar{q}\cdot q)\\ +(\bar{s}\cdot s)[(p\cdot\bar{q})(\bar{p}\cdot q)+(\bar{p}\cdot\bar{q})(p\cdot q)]-(\bar{q}\cdot s)[(\bar{s}\cdot p)(q\cdot\bar{p})+(\bar{s}\cdot\bar{p})(q\cdot p)]\}\big)\ , \qquad (1)\end{aligned}$$

where the spin and momentum vectors are defined in Fig. 1, $M$ is the neutrino mass, $\mathscr{s}=(p+\bar{p})^2$, and

$$R(\mathscr{s})=\frac{M_Z^2}{\mathscr{s}-M_Z^2+iM_Z\Gamma_Z}\ .$$

This expression includes an explicit statistical factor of $\frac{1}{2}$ because of the two identical particles in the final state. The initial $f$'s are assumed to be unpolarized and their mass is neglected.

It was assumed above that $N$ was a fourth-generation neutrino, but Eq. (1) can easily be adapted to other models with a heavy Majorana neutrino. Equation (1) is independent of the vector neutrino coupling since this term cancels out in the interference of the Majorana neutrinos,

$$\bar{u}(q)\gamma^\mu P_L v(\bar{q})-\bar{u}(\bar{q})\gamma^\mu P_L v(q)=-\bar{u}(q)\gamma^\mu\gamma_5 v(\bar{q})\ .$$

The Majorana-neutrino vector coupling cancels out for any combination of left- and right-handed couplings. For

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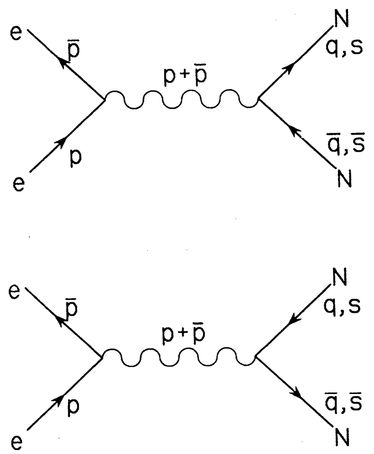


FIG. 1. Feynman diagrams and notation used for calculation of $e^+e^- \rightarrow NN$.

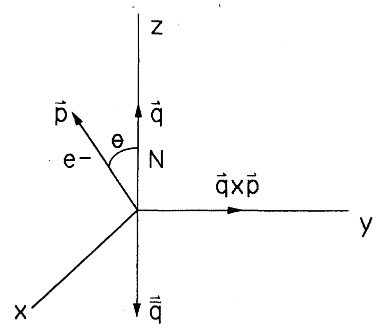


FIG. 2. Coordinate system for the final-state neutrinos.

production of neutrinos coupled with less than full strength, like a heavy "right-handed" neutrino in a "seesaw" model,[7] Eq. (1) is simply multiplied by a small mixing factor, $\sin^4\beta$. Equation (1) is also easily adapted to describe neutrino production through new, extra $Z$ bosons,[8,9] which might be observable at the Superconducting Super Collider.

In order to understand Eq. (1) let us evaluate it in the center-of-mass reference frame with the coordinate system shown in Fig. 2. The differential cross section for Majorana-neutrino production $e^+e^- \rightarrow NN$ takes the form

$$\frac{d\sigma^M}{d\Omega}=\frac{G_F^2}{128\pi^2}|R(s)|^2 s\beta^3\{-(1-4x)2\cos\theta(S_z+\bar{S}_z)+(1-4x+8x^2)[(1+\cos^2\theta)(1+S_z\bar{S}_z)+\sin^2\theta(S_y\bar{S}_y-S_x\bar{S}_x)]\}\ , \tag{2}$$

where **S** and $\bar{\mathbf{S}}$ are unit spin vectors defined in the rest frame of the neutrinos with momentum **q** and $\bar{\mathbf{q}}$, respectively, and $\beta=[1-(2M/s)^2]^{1/2}$ is the velocity of a final-state neutrino. For comparison the Dirac-neutrino production cross section $e^+e^- \rightarrow N\bar{N}$ is

$$\begin{aligned}\frac{d\sigma^D}{d\Omega}=\frac{G_F^2}{128\pi^2}|R(s)|^2 s\beta\{&(1-4x+8x^2)(1+\beta^2C_\theta^2)+(1-4x)2\beta C_\theta\\ &-(S_z+\bar{S}_z)[(1-4x+8x^2)\beta(1+C_\theta^2)+(1-4x)C_\theta(1+\beta^2)]\\ &-(S_x+\bar{S}_x)[(1-4x+8x^2)\beta C_\theta+(1-4x)]S_\theta(1-\beta^2)\\ &+S_z\bar{S}_z[(1-4x+8x^2)(\beta^2+C_\theta^2)+(1-4x)C_\theta 2\beta]\\ &+(S_z\bar{S}_x+S_x\bar{S}_z)[(1-4x+8x^2)C_\theta+(1-4x)\beta]S_\theta(1-\beta^2)+S_x\bar{S}_x(1-4x+8x^2)S_\theta^2(1-\beta^2)\}\ .\end{aligned} \tag{3}$$

Here **S** ($\bar{\mathbf{S}}$) is the spin of the neutrino (antineutrino), $C_\theta=\cos\theta$ is the angle between the electron and the neutrino three-momentum.

Equations (2) and (3) appear quite different; however, as stated in the introductory paragraph there should be no difference between a Dirac or Majorana neutrino when $M=0$. In the limit of $\beta\rightarrow 1$, the rate for the Dirac neutrino, Eq. (3), smoothly approaches

$$\left.\frac{d\sigma^D}{d\Omega}\right|_0=\frac{G_F^2}{128\pi^2}|R(s)|^2 s[(1-4x+8x^2)(1+C_\theta^2)+(1-4x)2C_\theta](1-S_z)(1-\bar{S}_z)\ . \tag{4}$$

For Majorana neutrinos with $M=0$, we again take $\beta=1$ but now the dependence on $S_y$, $\bar{S}_y$, $S_x$, and $\bar{S}_x$, must be explicitly dropped. These quantities represent transverse polarizations which are unphysical degrees of freedom for a massless fermion for which only $S_z(\bar{S}_z)=\pm 1$ are allowed. The differential cross section takes the form

$$\begin{aligned}\left.\frac{d\sigma^M}{d\Omega}\right|_{M=0}=\frac{G_F^2}{256\pi^2}|R(s)|^2 s\{&[(1-4x+8x^2)(1+C_\theta^2)+(1-4x)2C_\theta](1-S_z)(1-\bar{S}_z)\\ &+[(1-4x+8x^2)(1+C_\theta^2)-(1-4x)2C_\theta](1+S_z)(1+\bar{S}_z)\}\ .\end{aligned} \tag{5}$$

The first term is exactly one half of the Dirac-neutrino result. It gives a nonzero probability for emitting a left-handed neutrino in the direction of the three-momentum $+z$ and a right-handed "anti"neutrino in the direction $-z$. The second term gives a nonzero probability for emitting a right-handed "anti"neutrino in the direction $+z$ and a left-handed neutrino in the direction $-z$. The coordinate system for the Dirac-neutrino expression always has the left-handed neutrino propagating along the $+z$ direction; thus, in order to compare expressions we must flip $z \to -z$ for the last term in the Majorana-neutrino expression, Eq. (5). Then we see that Eqs. (4) and (5) are equal so that the differential cross sections for Dirac and Majorana neutrinos are equivalent when $M=0$.

In the comparison above, the Dirac expression smoothly approaches the massless limit far above the production threshold while for the Majorana expression there is a definite distinction between $M=0$ and $\beta \to 1$. The transverse degrees of freedom gradually decouple as $\beta \to 1$ in the Dirac case because the coupling is left handed, $\gamma^\mu(1-\gamma_5)$, but for Majorana neutrinos the corresponding coupling is $-2\gamma^\mu\gamma_5$, due to the interference of the identical particles, so the coupling is not to a single helicity. Combining Eqs. (2) and (5), the differential cross section for Majorana neutrinos can be written as

$$\frac{d\sigma^M}{d\Omega} = \beta^3 \left[ \left. \frac{d\sigma}{d\Omega} \right|_{M=0} + \frac{G_F^2}{128\pi^2} |R(s)|^2 s [(1-4x+8x^2)\sin^2\theta (S_y \bar{S}_y - S_x \bar{S}_x)] \right] . \quad (6)$$

The first term above smoothly approaches the Dirac expression far above threshold, as shown previously, while the second term does not. The second term gives a finite probability for emitting Majorana neutrinos with transverse polarizations, even far above threshold. The second term above disappears for massless neutrinos or if one does not observe the spin of both of the neutrinos. If only one neutrino is observed, the spin of the other neutrino is summed, $\sum_a S_\mu^a = 0$, and the final term above vanishes. Thus the spin-independent Majorana and Dirac cross sections will approach each other far above threshold, $\beta \to 1$. However, in heavy-neutrino searches, the neutrino lifetime is typically assumed to be short enough so that both neutrinos decay inside the detector; hence, the second term in Eq. (6) will be physically relevant. Since the neutrino decays are parity violating, the production spin dependence will show as correlations[10,11] in the angle and energy dependence between the decay products of the neutrino pair. Thus there are measurable differences between Dirac- and Majorana-neutrino production, even far above threshold.

Near the threshold, there will be large differences between the production of Dirac or Majorana neutrinos for the spin-independent cross sections as well as the spin-dependent differences discussed above. From Eqs. (3) and (6) we see that the cross section is proportional to $\beta^3$ for Majorana neutrinos[12,13] and to $\beta$ for Dirac neutrinos.[14,15] Thus the Majorana-neutrino cross section rises much more slowly above threshold than the Dirac-neutrino cross section.

## II. MAJORANA-NEUTRINO DECAY

Heavy-neutrino decay occurs through the mixing of the heavy neutrino with the light neutrino states and has been discussed previously by other authors.[16] A Majorana neutrino can decay as a Dirac neutrino or as a Dirac antineutrino. There is no interference between these two types of decays due to the Majorana nature of the initial neutrino because the final states are distinct. For instance, the Majorana neutrino can decay analogous to a Dirac neutrino to $e^+e^-\nu_e$ or it can decay analogous to the Dirac antineutrino to $e^+e^-\bar{\nu}_e$ and (neglecting a possible small Majorana mass for the $\nu_e$) because the $\nu_e$ and the $\bar{\nu}_e$ are distinguishable there will be no interference between the two possible final states. Thus to quantitatively describe heavy-Majorana-neutrino decay, one merely sums incoherently the expressions describing heavy Dirac neutrino and antineutrino decays. The predicted total lifetime for a Majorana neutrino will be simply $\frac{1}{2}$ the corresponding Dirac expression since for every Dirac-neutrino decay there is a corresponding Dirac-antineutrino decay of equal width, both of which are allowed for the Majorana neutrino.

If only one of the produced neutrinos is observed, it would be difficult to decide whether it was a Majorana or Dirac particle. However, if both of the produced neutrinos are observed, then there can be clear and dramatic differences between the two types. Assuming the neutrino couples dominantly to a single light charged lepton of flavor $L$, $N \to L + W^*$, then the Dirac neutrinos will always decay into oppositely charged $L$'s while the Majorana neutrinos will decay like that half of the time but the other half into like-sign $L$'s. Like-sign charged leptons could provide a clean and unambiguous signal for production of Majorana neutrinos if the $L$ were an electron or muon. However, if the $L$ were a $\tau$, it would decay before reaching the detector, and heavy-neutrino production would be more difficult to discern.

## III. MAJORANA- AND DIRAC-NEUTRINO DECAY CORRELATIONS

The correlations between neutrino decay products can be useful for separating the neutrino signal from the background. Also, correlations can provide important independent checks to confirm the presence of neutrino production. Qualitatively, it is easy to understand the type of effects that correlations in the produced neutrino's spin can give rise to. For example, consider the decay of a heavy Majorana neutrino $N$ into a charged lepton $L$ and where the virtual $W$ decays into two light fermions. An $L^-$ ($L^+$) is emitted preferentially in the opposite (same) direction as the heavy neutrino's spin. Thus the positive $S_y\bar{S}_y$ term in Eq. (6) causes like- (unlike-) sign $L$'s to lie preferentially on the same (opposite) side of the $x$-$z$ plane while the negative $S_x\bar{S}_x$ causes like- (unlike-) sign $L$'s to lie preferentially on opposite (same) sides of the $y$-$z$ plane. Quantitatively, it is often difficult to analytically describe the physically relevant correlations and one is forced to use numerical methods. Here we will give some of the angle and energy correlations which can be calculated analytically[17] in order to illustrate the differences between Dirac and Majorana neutrinos.

Define $X_1$ and $X_2$ to be two of the decay products of the produced neutrinos where each $X$ comes from a different neutrino. The rate for the weak decay of $N \to X + \text{anything}$ is described by

$$\frac{d\Gamma}{d\Omega} = A + B\hat{\mathbf{s}}\cdot\hat{\mathbf{k}} \tag{7}$$

in the rest frame of $N$, where $\hat{\mathbf{s}}$ is the spin of $N$ and $\hat{\mathbf{k}}$ is the unit momentum vector of $X$. The first correlation we consider between $X_1$ and $X_2$ is in the angle between the two. We define

$$z = \cos\theta_1\cos\theta_2 + \sin\theta_1\sin\theta_2\cos(\phi_1 - \phi_2) \ . \tag{8}$$

This is the angle between $X_1$ and $X_2$ in the rest frames of $N$, which coincides with the angle measured in the laboratory only near threshold. Then folding in Eqs. (2) and (7) yields

$$\frac{1}{\sigma}\frac{d\sigma}{dz} = \frac{1}{2}\left[1 + \frac{1}{3}\left[\frac{B_1B_2}{A_1A_2}\right]z\right] \tag{9}$$

for the angular correlation between $X_1$ and $X_2$ produced from the decay of Majorana neutrinos. For Dirac neutrinos, Eqs. (3) and (7) yield an expression which turns out to be identical to Eq. (9). It is reasonable that the Dirac and Majorana expressions are identical far above threshold. There the differences in the spin correlations, as exemplified in Eq. (6), come in as two terms of identical size, opposite sign, and transverse to the heavy-neutrino momentum. The two terms cancel out when only the overall angle between two decay products is considered. In order for the two terms to give a nonzero contribution, a correlation must be considered with a preferred direction transverse to the heavy-neutrino axis.

Next we consider the double correlation of $X_1$ and $X_2$ with respect to the beam axis. We define

$$z_i = \cos\theta\cos\theta_i + \sin\theta\sin\theta_i\cos\phi_i \tag{10}$$

which is the angle between $X_i$ and the electron in the rest frame of the $N$ (we stress again that this only corresponds to the physical angle near threshold). Then the double angular correlation for Majorana neutrinos can be written as

$$\frac{1}{\sigma}\frac{d^2\sigma}{dz_1dz_2} = \frac{1}{4}\left[1 - \frac{1}{2}\frac{1-4x}{1-4x+8x^2}\left[\frac{B_1}{A_1}z_1 + \frac{B_2}{A_2}z_2\right]\right] . \tag{11}$$

For Dirac neutrinos, the double angular correlation takes the form

$$\frac{1}{\sigma}\frac{d^2\sigma}{dz_1dz_2} = \frac{1}{4}\left[1 - \frac{3-\beta^2}{3+\beta^2}\frac{1-4x}{1-4x+8x^2}\left[\frac{B_1}{A_1}z_1 + \frac{B_2}{A_2}z_2\right] + \left[\frac{1-\frac{7}{15}\beta^2}{1+\frac{1}{3}\beta^2}\right]\left[\frac{B_1B_2}{A_1A_2}\right]z_1z_2\right] . \tag{12}$$

As $\beta \to 1$, the terms linear in $z_i$ in Eqs. (11) and (12) approach each other while the bilinear terms do not. This difference in the $z_1z_2$ term comes from the $S_y\bar{S}_y - S_x\bar{S}_x$ term for Majorana neutrinos [Eq. (6)]. The coefficient of the linear terms is quite small due to the factor of $[1-4\sin^2\theta_W]$; however, the coefficient of the Dirac bilinear term is much larger, ranging from 1 near threshold to $\frac{2}{5}$ far above threshold. Since the bilinear term of the Majorana neutrino vanishes, the difference between the two angular correlations is substantial.

In addition to the angular correlations given above, moments of the energy distribution can be calculated. To be specific, we consider $X_1$ and $X_2$ to be charged leptons with energies $E_1$ and $E_2$, respectively, in the laboratory frame; and where each virtual $W$ goes into two light fermions. For Dirac neutrinos, the average energy of the oppositely charged leptons can be written as

$$\begin{aligned} \langle E_1 \rangle = \langle E_2 \rangle &= \frac{7}{20}\frac{M_N}{\sqrt{1-\beta^2}}\left[1 + \frac{1}{7}\frac{\frac{4}{3}\beta^2}{1+\frac{1}{3}\beta^2}\right] , \\ \langle E_1E_2 \rangle &= \left[\frac{7}{20}\frac{M_N}{\sqrt{1-\beta^2}}\right]^2\left[1 + \frac{2}{7}\frac{\frac{4}{3}\beta^2}{1+\frac{1}{3}\beta^2} + \frac{1}{7^2}\frac{\beta^2(\beta^2+\frac{1}{3})}{1+\frac{1}{3}\beta^2}\right] , \end{aligned} \tag{13}$$

where the mass of the final-state fermions has been neglected. For $\beta \to 1$, the Dirac neutrinos are produced in left-handed helicity states and the above expressions reproduce the well-known results for weak decay (as in muon decay). For Majorana neutrinos, the average energy of the charged leptons can be written as

$$\begin{aligned} \langle E_1 \rangle = \langle E_2 \rangle &= \frac{7}{20}\frac{M_n}{\sqrt{1-\beta^2}} , \\ \langle E_1E_2 \rangle &= \left[\frac{7}{20}\frac{M_n}{\sqrt{1-\beta^2}}\right]^2\left[1 + \epsilon\frac{1}{7^2}\beta^2\right] , \end{aligned} \tag{14}$$

where $\epsilon = -1$ ($+1$) for like- (opposite-) sign charged $L$'s. These expressions are quite different from the Dirac results, Eq. (13), far above threshold. However, unlike the differences in the double angular correlations, this difference does not come from the $S_y\bar{S}_y - S_x\bar{S}_x$ term in Eq. (6) which does not contribute at all to the average energy of the $L$'s since $x$ and $y$ are transverse to the heavy-neutrino momentum. Instead the difference occurs because massive Majorana neutrinos are not emitted in left-handed helicity states, even far above threshold.

In summary, we have examined the relevant issues for constraining heavy Majorana neutrinos from collider experiments.[18] The production cross section has been calculated and compared to that for Dirac neutrinos. Pair production of the two types of neutrinos is completely different near the threshold, while far above the threshold the spin-independent cross sections are the same but the spin-dependent cross sections are not. This difference is observable in the angle and energy correlations of the neutrino decay products. Majorana-neutrino decay is discussed and shown to be calculable from previous results for Dirac neutrinos. Majorana neutrinos can have a strong like-sign-lepton signal which would be readily observable.

## ACKNOWLEDGMENTS

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