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Fourth-Generation Leptons at LEP2

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ABSTRACT

From non-observation at LEP1, a lower mass limit of 45 GeV has been established on any additional sequential fermion beyond the three generations. Precision measurements have further constrained the number of such additional generations, either degenerate or nearly so, to be at most one. LEP2, an energy-upgraded version of LEP1, would provide greater mass-reach in the search for such particles. We study the pair-production of fourth-generation leptons for various LEP2 energy options. We find that in most cases such particles could be discovered/ruled out up to the kinematic limit.

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1 Introduction

Despite the vindication of the Standard Model (SM) at LEP1, a few questions remain unanswered, not the least of which concerns the twin issues of fermion masses and family replication. While, as yet, there exists no understanding of either issue, there have been some speculations to the effect that these might be interrelated [1, 2, 3]. Furthermore, the existence of such particles are likely to have a significant impact even on low-energy observables, *e.g.* particularly in the context of CP violation in the B-system [4]. The issue of exploring the existence of new fermions at LEP2 thus turns out to be imperative.

Possible new quarks or leptons can be subdivided into two categories : sequential (*i.e.* with gauge quantum numbers identical to the SM fermions) or exotic [3]. The latter class would include all such fermions that have no analogue in the SM.

The most model-independent bounds on fermions beyond the SM can be inferred from their non-observation at LEP1. Unless their couplings to the Z are highly suppressed, for all such fermions [5]: $m_F > 45$ GeV. At the Tevatron, on the other hand, quark pair-production far outstrips that for leptons. Assuming that a heavy quark decays within the detector via its neutral current or charged current mixing with the standard quarks, one obtains a stronger bound $m_Q > 85$ GeV [5, 6], which is likely to improve with new data. As lepton production at the Tevatron proceeds through Drell–Yan processes or through weak-gauge boson fusion, the corresponding limits are expected to be weaker. This analysis has not yet been reported though. Over and above these direct bounds, the bound on the oblique parameter S (equivalently ϵ_3) from the precision electroweak measurement restricts the number of additional degenerate chiral generations to just one [7]. It is thus quite clear that any addition to the fermionic sector of the SM cannot be too arbitrary and hence all search strategies should be devised keeping the existing bounds in mind. In this article we examine the possibility of discovering new fermions at the forthcoming upgrade of the LEP machine at CERN. As it is supposed to operate at $\sqrt{s} = 140, 176, 192$ and (hopefully) 205 GeV, heavy quark pair-production is interesting only for the last two energy options. Single production in association with a light quark is ruled out (at an observable level) from the severe bounds on flavour-changing neutral currents (FCNC) involving light quarks. Hence, in the present work we concentrate on heavy leptons only.

Now, the search pattern for exotic particles [9] necessarily differs from that for sequential fermions. Although the mass reach is much more for the former, as they may be singly produced, the production cross section suffers a suppression from the low-energy bounds on lepton mixing [8]. Pair production of the exotics has been considered in ref. [9] and we shall thus confine ourselves to sequential leptons.

Since we restrict our analysis to sequential fermions, it is clear that there are no tree-level FCNCs in the theory. The production of a heavy charged lepton pair ($E\bar{E}$) at LEP2 then proceeds primarily through a Drell–Yan-like mechanism (*viz.*exchange of s-channel γ/Z). For $N\bar{N}$ production, the γ channel is obviously absent. There is an additional diagram, though, involving a t-channel W exchange. However since this contribution to the amplitude involves the eN mixing, which is already restricted to be very small ($\sin^2\theta_{eN} \sim \sin^2\theta_{\nu N} \lesssim 0.005$), the

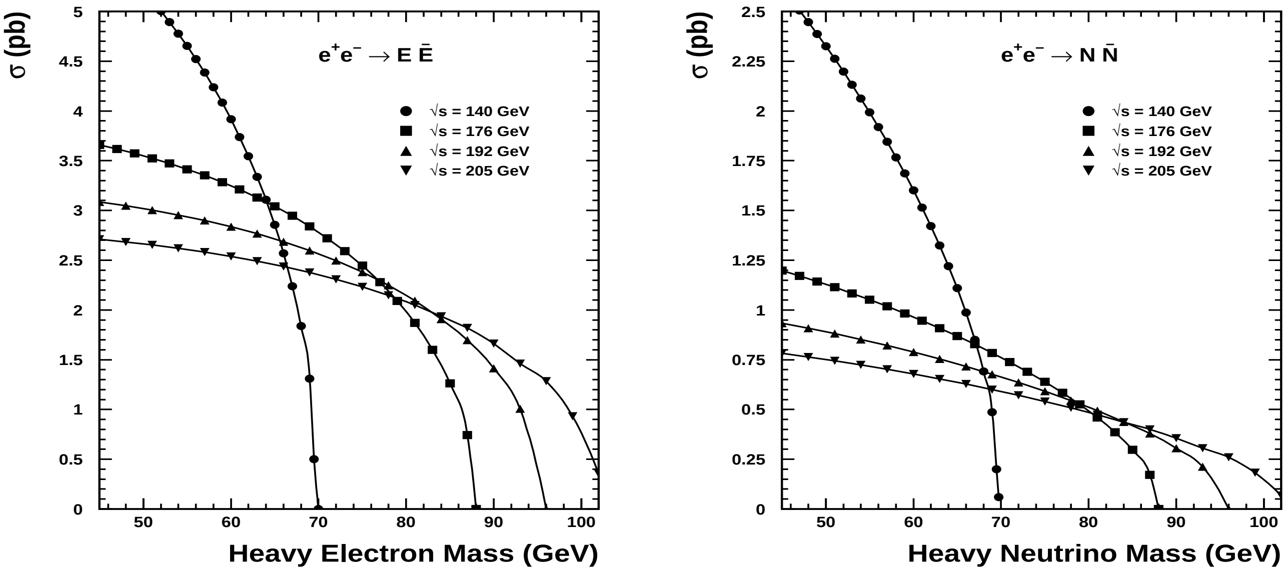

Figure 1: *The total cross section (at various c.m. energies) for E^+E^- and $N\bar{N}$ production at LEP2.*

t-channel exchanges can be essentially neglected. The production cross-section, can then be calculated in a straightforward manner and is given in the appendix. In Fig. 1, we display the total cross section for the pair-production of both N and E as a function of their masses[1]. That the $E\bar{E}$ cross section is typically larger than the corresponding $N\bar{N}$ cross section is due to the presence of the extra (γ–exchange) diagram for the former. As can be easily seen, the cross-section is quite large almost up to the kinematic limit of LEP2. Were all the produced leptons available for detection, this would be a very good theatre for discovery. It is thus meaningful to consider the potential backgrounds and evaluating prospect of detecting such fermions at LEP2.

Since an absolutely stable heavy lepton (charged or neutral) is disfavoured on astrophysical/cosmological grounds [12], any such new particle must necessarily decay. Although a very small width can evade this bound, such particles would be effectively stable as far as the collider experiments are concerned. While a stable N would go undetected in the LEP experiments, it might just be possible to detect a stable E through a calorimetric measurement. We shall, however, assume that these particles, if produced, decay within the detector volume. Since tree-level FCNC's are absent in the simplest fourth generation scenario, all decays proceed through the charged current interaction. Current bounds on possible lepton mixing suggests the heavy lepton L $(= E, N)$ would, if allowed to kinematically, tend to decay predominantly into its isospin partner. Rather than look for such cascade decays, it makes more sense to concentrate instead on the lighter of E, N, especially since we shall show that both such particles

[1]Here, as in most of the following discussion, we neglect the effect of initial state radiation (ISR). As $N\bar{N}$ production proceeds mainly through an s-channel Z-exchange, the effect of ISR is expected to be severe on account of the tendency to return to the Z-pole. However, from the quantitative analysis of ref. [10], we see that the effect is actually not severe enough to change our qualitative predictions. For charged leptons (E), the effect is even smaller.

can be discovered for almost all of the parameter space. This particle may then decay only into fermions of the first three generations. Decay within the detector volume is assured as long as the mixing angle is θ_{lN} ($\theta_{\nu E}$) $\gtrsim 10^{-6}$.

The rest of the article is planned as follows: in section 2 we discuss the different decay modes of $L(E, N)$, the kinematic cuts implemented in our analysis, the SM backgrounds; we then highlight the post-cut efficiencies. We draw our conclusion in section 3. In the appendix, we list all the relevant formulae.

2 The Signal

According to our criterion, the heavy charged lepton E would decay into one of the lighter neutrinos and a W (on- or off-shell), *i.e.*

$$E^- \to \nu_i W^{-(*)} \qquad \text{and} \qquad E^+ \to \bar{\nu}_i W^{+(*)} \ , \tag{1}$$

where $i = 1, 2, 3$. The individual mixing matrix elements are of no consequence as long as at least one of them $\gtrsim 10^{-6}$. When considering decays of E into neutrinos of different flavour, care must be taken though that the choice of parameters does not lead to unacceptably large leptonic FCNC's [8]. However, since such decay modes do not afford us any particular experimental advantage, we shall assume that E decays to only one specific light neutrino flavour. The W's may then go into either hadronic or leptonic channels. The final state is thus one of

1. $\nu_i\bar{\nu}_i + q_a\bar{q}_b q_c\bar{q}_d$,
2. $\nu_i\bar{\nu}_i + q_a\bar{q}_b(l_k\bar{\nu}_k + \bar{l}_k\nu_k)$,
3. $\nu_i\bar{\nu}_i + l_k\bar{\nu}_k\bar{l}_n\nu_n$.

Of the three modes, the last one is of little use. The signal would be overwhelmed by the background from $e^+e^- \to W^+W^-$ with each W decaying leptonically. The second option is somewhat better. We shall, however, confine ourselves to the first and the most promising channel with the signal

$$e^+e^- \to E^+E^- \to 4 \text{ jets} + \not{p}_T \tag{2}$$

where $\not{p}_T$ represents the missing transverse momentum.

We now turn to the case where the fourth-generation neutrino N is lighter than its charged counterpart. Unlike the charged particle, N could have a Majorana mass too. Apart from leading to interesting possibilities in low-energy neutrino phenomenology [13], such an eventuality leads to the beautifully clean signal of like-sign dileptons with hadronic activity but without any missing momentum. The backgrounds are twofold : (i) cascade decays of a heavy quark pair (say $b\bar{b} \to ce^-\bar{\nu}\bar{c}q_i\bar{q}_j \to e^-e^-\nu\nu+$ jets) or (ii) effects like B–$\overline{B}$ mixing. Although both these backgrounds ostensibly lead to missing momentum, these are still relevant as the neutrinos may be slow, or the transverse component ($\not{p}_T$) of the total missing momentum may be small. These can however be easily eliminated by a combination of isolation cuts and imposition of a upper bound on missing momentum. Having argued for the ease of detection in the Majorana N case, we shall desist from discussing it any further and will rather concentrate on the case where N is a Dirac particle. The primary decay vertices then are of the form

$$N \to l_i^- W^{+(*)} \qquad \text{and} \qquad \bar{N} \to l_j^+ W^{-(*)} \ , \tag{3}$$

where $i, j = 1, 2, 3$. Some of the most striking signals would emanate in the case $i \neq j$ in eq. (3). However, such a scenario will induce leptonic FCNC's at the one-loop level. The non-observation of such effects can be translated to considerably strong constraints [8] in such scenarios. We assume henceforth that N, like E, has charged-current coupling with only *one* of the light flavours. Even with this simplification, two further possibilities exist: (i) N decays mainly into e or μ. (ii) N decays mainly into τ. Since τ's decay within the detector volume, the signal profile changes in an essential manner. Consequently, these two cases must be discussed separately. We shall, for the present, concentrate on the first case as τ-identification is tricky, especially in the presence of hadrons.

Quite analogous to the case of E, the W's in eq. (3) may go into either hadronic or leptonic channels. The final states are of the following types:

1. $l_i \bar{l}_i + q_a \bar{q}_b q_c \bar{q}_d$,
2. $l_i \bar{l}_i + q_a \bar{q}_b (l_k \bar{\nu}_k + \bar{l}_k \nu_k)$,
3. $l_i \bar{l}_i + l_k \bar{\nu}_k \bar{l}_n \nu_n$.

The first two modes are better suited for discovery (the first especially, as it affords easy mass reconstruction). Although the last of the three can also be useful, we shall not consider it here.

Having identified our final states, it now remains to calculate the effective transition matrix element. We do this within the narrow width approximation (albeit keeping track of the spin-spin correlations) and this derivation is presented in the appendix. It turns out that, on account of the peculiar kinematics of the system and our choice of observables, the spin correlations are not particularly significant.

At this stage it must be pointed out that the above-mentioned signals are not necessarily the only viable ones. In fact, quite a few studies [10, 11] have focussed on an $N\bar{N}$ production signal of isolated leptons accompanied by hadronic activity. While the effective signal strength is higher for such configurations, the backgrounds are larger too, and consequently special kinematic cuts (for example, removal of the ZZ or $Z\gamma^*$ backgrounds) are necessary. We, on the other hand, have chosen final states such that the corresponding SM backgrounds are easily reduced to rather innocuous levels.

2.1 Kinematic cuts

A multitude of cuts are necessary both for experimental reasons and to suppress all possible SM backgrounds. To list:

1. Each quark in eq. (2) should be sufficiently away from the beam pipe and should carry sufficient energy and transverse momentum. We require

$$E_j > 10\,\text{GeV}, \qquad p_{Tj} > 5\,\text{GeV}, \qquad 10^\circ < \theta_j < 170^\circ. \tag{4}$$

2. Since we do only a parton-level simulation, we must ensure that the angular separation between any two quarks must be large enough for them to lead to recognisably different jets. To be conservative, we require, for any two jets,

$$\theta_{jj} > 30^\circ. \tag{5}$$

   This cut also eliminates potentially large contributions from collinear singularities in the SM matrix elements.

3. For final states with charged leptons, we require

$$p_{Tl} > 5\,\text{GeV}, \qquad 10^\circ < \theta_l < 170^\circ. \tag{6}$$

   The angular cut also eliminates collinear singularities for the case $l = e$. On the other hand, for the $E\bar{E}$ case, where the final state contains neutrinos, we demand that the missing transverse momentum[2] should be sufficiently large as to be observable :

$$\not{p}_T > 10\,\text{GeV}. \tag{7}$$

4. We further impose the following isolation cuts (whenever applies depending on the final states)

$$\theta_{jl}, \theta_{ll}, \theta_{j\not{p}} > 20^\circ. \tag{8}$$

   Apart from facilitating proper detection, this also serves to reduce the background contribution where the lepton (or neutrino) arises from a heavy quark decay. In addition, the cut on θ_{ll} removes electromagnetic collinear singularities in the SM matrix elements.

2.2 Backgrounds

As our final states typically comprise of six fermions, power counting (in terms of the coupling constants) would naively dictate that the expected SM backgrounds be small. A possible counterargument could be that the large number of diagrams contributing to such processes

[2] Note that missing longitudinal momentum is not necessarily a good signal, as it can possibly be faked by ISR.

might enhance the effects. In fact, it is precisely this (*viz.*, a large number of graphs) that makes an accurate calculation of the SM backgrounds an extremely difficult task. There does not exist any computational tool that calculates the full $2 \to 6$ matrix elements in an acceptable time-frame. One should also bear in mind the possibility that the backgrounds may be enhanced on account of the presence of resonant processes and/or collinear emissions.

However, very reasonable estimates may be made by classifying the possible Feynman diagrams in terms of well-identifiable "subprocesses", *viz.*$e^+e^- \to e^+e^- W^{(*)}W^{(*)}$ with the $W^{(*)}$'s decaying hadronically, or $e^+e^- \to \nu\bar{\nu}q\bar{q}$ followed by $q\bar{q} \to 4$ jets. Such piecemeal calculations are expected to give an estimate not drastically different from the full $2 \to 6$ calculation[3].

These reduced processes were calculated with the help of the helicity amplitude package MadGraph [14] and some of them were counterchecked against an independent analysis performed with GRACE [15]. With the adoption of the kinematic cuts listed in the next subsection, each of these individual contributions is reduced to well below 1 fb, and it may thus be safely assumed that the SM backgrounds to our signals are finally of no consequence.

2.3 The Efficiencies

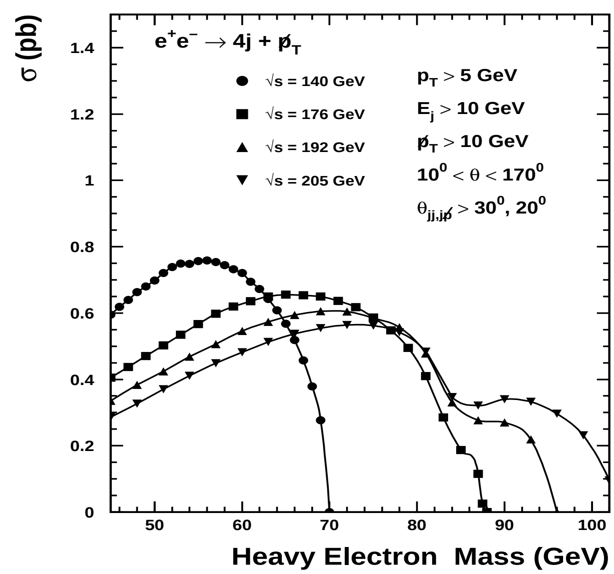

Figure 2: *The effective cross section (at various c.m. energies) for E^+E^- production after imposition of the cuts of section 2.1.*

While the cuts imposed above serve to eliminate all the SM backgrounds, they are not particularly severe on the signal. In Fig. 2, we show the effective signal strength (*i.e.* after the imposition of the above-mentioned cuts) for E^+E^- production with both $W^{(*)}$s decaying

[3]The caveat is that such an analysis ignores those diagrams that cannot be broken down in terms of simple 'subprocesses'. But then, these are truly higher-order in the coupling constants and the corresponding contributions are small.

hadronically. On the other hand, Fig. 3 shows the same for two different decay modes of the $N\bar{N}$ pair. The post-cut efficiencies (without folding any detector effects) depend on m_L ($L = E, N$) and can be deduced by comparing these figures with those in Fig. 1. Since the phase-space volume available to the decay products from $L(E, N)$ grows as m_L, it is obvious that for low m_L, the twin requirements of minimum jet (lepton) energy and the isolation cuts would substantially reduce the efficiency. However, this is not a cause for worry as the production cross-sections are rather large. On the other hand, for $m_L \gtrsim m_W$, L prefers a two-body decay (into an on-shell W) rather than a genuine three-body decay. Consequently, for m_L just above m_W the primary lepton is left with very little momentum and such configurations are eliminated primarily by the cuts of eq. (6). This leads to the dips in Figs. 2 and 3. As can easily be

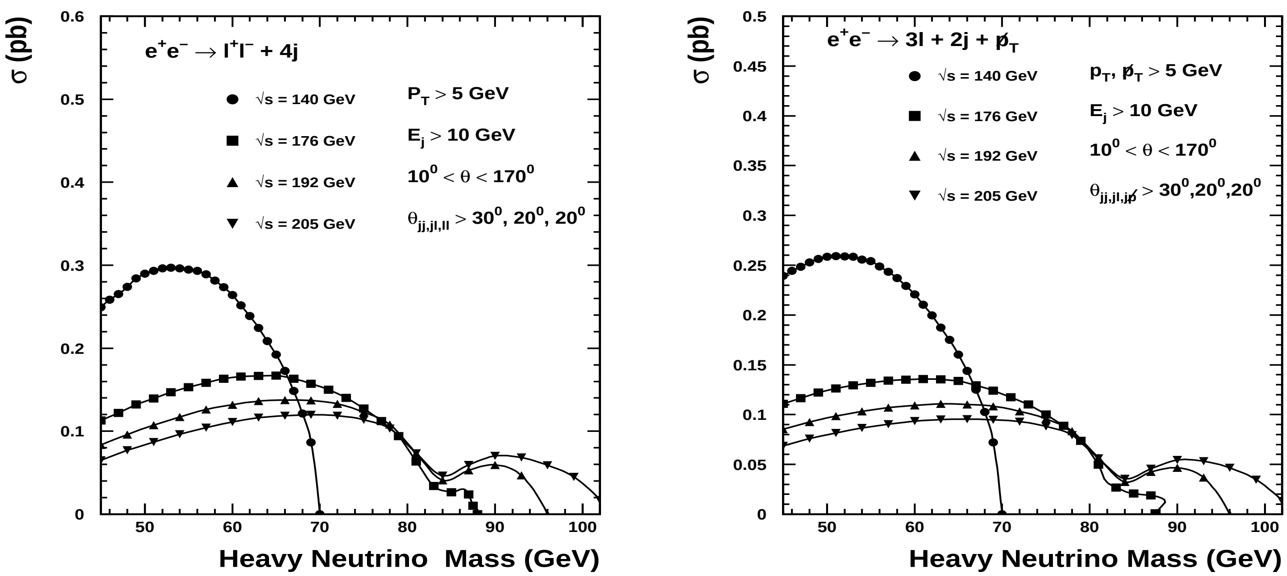

Figure 3: *The effective cross section (at various c.m. energies) for pair-produced N's decaying into two particular channels. The cuts of section 2.1 have been imposed.*

ascertained, even after the imposition of such strong cuts, the effective cross sections are large enough to warrant a discovery claim right up to the kinematic limit for the expected luminosity of 300 pb^{-1}. On the other hand, if an integrated luminosity of 10 pb^{-1} is achieved for the run at $\sqrt{s} = 140$ GeV, it might be possible to rule out $m_E \lesssim 60$ GeV.

Although the cross sections are smaller for $N\bar{N}$, than for E^+E^-, there are some advantages associated with the first, not the least of which being mass reconstruction and a window to the angular distribution (see Fig. 4a). The latter can differentiate between heavy neutrinos with different gauge quantum numbers through their polar angle distributions. Although the imposition of the cuts tends to distort the distribution, to a certain extent the characteristic signatures are still preserved. In Fig. 4b, we concentrate on the case of the $(l^+l^- + 4j)$ signal and show the distribution and function of two relevant angles : (i) the opening angle of the two leptons (θ_{ll}), and (ii) their azimuthal separation ($\Delta\phi_{ll}$). The "peaks" are characteristic, though they tend to get flattened as $m_N/\sqrt{s}$ increases.

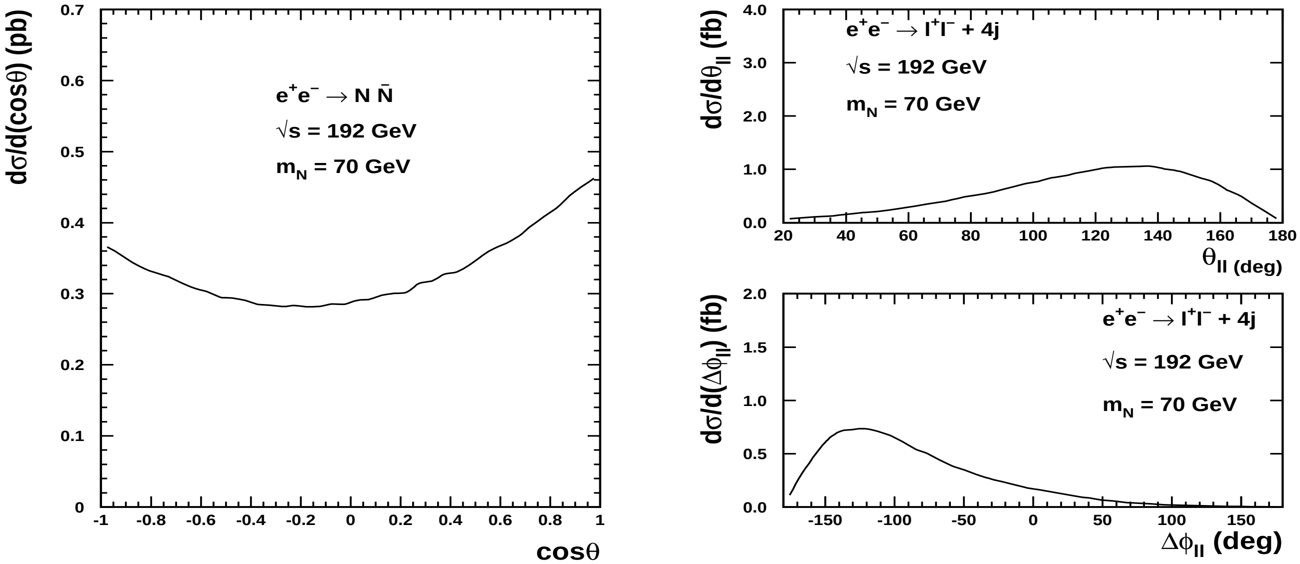

Figure 4: *Angular distributions for $N\bar{N}$ production. (a) The variation with the production angle* before *the cuts are imposed. (b) The angular distributions (*after *imposition of cuts of section 2.1): (i) the opening angle and (ii) the difference of azimuthal angles) for the leptons from the primary decay vertices of N and $\bar{N}$.*

3 Conclusions

We have examined the possibility of observing a fourth-generation lepton in the forthcoming runs at LEP2. The prognosis for quarks is not favourable. Only the runs at $\sqrt{s} = 192$ and 205 GeV may explore mass ranges not ruled out as yet, and that window too might already be closed from the upcoming Tevatron data. On the other hand, both charged and neutral leptons can be discovered (or ruled out) up to the kinematic limit. Although we attempt only a parton-level simulation, our phase-space cuts are conservative enough and, therefore, the conclusions should not change significantly even on the inclusion of hadronization effects. Initial-state radiation, which has been neglected here, is perhaps more likely to have a discernible effect, particularly for $N\bar{N}$ production, but still would not change the conclusions significantly, as independent studies have indicated. For a heavy neutrino, mass reconstruction works very well and angular distributions can be used to determine the quantum numbers. If the neutrino were to have a significant amount of Majorana mass, a like-sign dilepton pair would be a very distinct signal.

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Appendix

In this appendix we present the relevant formulae for pair-production and subsequent decay of *polarized* heavy fermions at LEP2. Since the W-exchange diagrams can be neglected, the production is essentially given by s-channel diagrams mediated by γ, Z exchange. The amplitude for the process $e^-(p_1)e^+(p_2) \longrightarrow L(p_3,s_3)\bar{L}(p_4,s_4)$ (where $L = E, N$) can then be expressed as

$$\mathcal{M} = e^2 \sum_{i=\gamma,Z} \mathcal{P}_i \; \bar{u}(p_3,s_3)\,\gamma_\mu \left(v_i^{(L)} + a_i^{(L)}\gamma_5\right)\, v(p_4,s_4)\;\, \bar{v}(p_2)\,\gamma_\mu \left(v_i^{(e)} + a_i^{(e)}\gamma_5\right)\, u(p_1)\;, \tag{9}$$

where

$$\mathcal{P}_i = (s - m_i^2 + i\Gamma_i m_i)^{-1} \tag{10}$$

and

$$\begin{aligned} v_\gamma^{(f)} &= Q_e & a_\gamma^{(f)} &= 0 \\ v_Z^{(f)} &= \frac{t_{3L}^{(f)} - 2Q_f \sin^2\theta_W}{2\sin\theta_W \cos\theta_W} & a_Z^{(f)} &= \frac{-t_{3L}^f}{2\sin\theta_W \cos\theta_W} \end{aligned} \tag{11}$$

for $f = e, L$. For convenience, we define

$$\begin{aligned} \mathcal{K}_1 &= 2\sum_{i,j} \mathcal{P}_i\mathcal{P}_j^* \left(v_i^{(e)}v_j^{(e)} + a_i^{(e)}a_j^{(e)}\right)\left(v_i^{(L)}v_j^{(L)} + a_i^{(L)}a_j^{(L)}\right) \\ \mathcal{K}_2 &= 2\sum_{i,j} \mathcal{P}_i\mathcal{P}_j^* \left(v_i^{(e)}v_j^{(e)} + a_i^{(e)}a_j^{(e)}\right)\left(v_i^{(L)}v_j^{(L)} - a_i^{(L)}a_j^{(L)}\right) \\ \mathcal{K}_3 &= 2\sum_{i,j} \mathcal{P}_i\mathcal{P}_j^* \left(v_i^{(e)}a_j^{(e)} + a_i^{(e)}v_j^{(e)}\right)\left(v_i^{(L)}a_j^{(L)} + a_i^{(L)}v_j^{(L)}\right)\;. \end{aligned} \tag{12}$$

The production matrix element squared can then be expressed as

$$|\mathcal{M}_{\rm prod}|^2 = e^4 \sum_{i,j} \left[\mathcal{A} + 4\mathcal{B}_{\mu\nu} s_3^\mu s_4^\nu\right], \tag{13}$$

where

$$\begin{aligned} \mathcal{A} &= \mathcal{K}_1\left[(m_L^2 - u)^2 + (m_L^2 - t)^2\right] + \mathcal{K}_3 s(t-u) + 2\mathcal{K}_2 m_L^2 s \\ \mathcal{B}^{\mu\nu} &= m_L^2\mathcal{K}_1\,(p_1^\mu p_2^\nu + p_1^\nu p_2^\mu) + m_L^2\mathcal{K}_3\,(p_1^\nu p_2^\mu - p_1^\mu p_2^\nu) \\ &+ \frac{\mathcal{K}_2}{2}\Big[(m_L^4 - ut)g^{\mu\nu} + (s - 2m_L^2)\,(p_1^\mu p_2^\nu + p_1^\nu p_2^\mu) + s p_3^\nu p_4^\mu \\ &\quad - p_4^\mu\left\{(m_L^2 - u)p_1^\nu + (m_L^2 - t)p_2^\nu\right\} - p_3^\nu\left\{(m_L^2 - t)p_1^\mu + (m_L^2 - u)p_2^\mu\right\}\Big]\,. \end{aligned} \tag{14}$$

The decay vertex of L is of the form $\bar{l}LW^{(*)}$, where $W^{(*)}$ is either real or virtual decaying as $W^{(*)} \to f_1(q_1)\bar{f}_2(q_2)$, l is a light neutrino for $L = E$ and a light charged lepton for $L = N$. The decay matrix element squared can then be parametrized as

$$|\mathcal{M}_{L\ \rm decay}|^2 = C\,(1 + D_\mu s_3^\mu)\,, \tag{15}$$

where

$$
\begin{aligned}
C &= 2g^4 \sin^2\theta_{lL} \, \frac{(p_3 \cdot q_2)\ (p_l \cdot q_1)}{\left[(p_3 - p_l)^2 - m_W^2\right]^2 + \Gamma_W^2 m_W^2} \\
D_\mu &= \frac{m_L}{p_3 \cdot q_2}\ q_{1\mu},
\end{aligned}
\tag{16}
$$

and $\sin\theta_{lL}$ is the relevant mixing angle. Similar statements can be made for $\bar{L}$ as well.

Within the narrow-width approximation, eqs. (13 & 16) can then be convoluted to give the effective matrix element squared:

$$
\begin{aligned}
|\mathcal{M}|^2_{\rm eff}\left(e^+e^- \to L\bar{L} \to 6 \text{ fermions}\right) &= \Gamma_L^{-2} \sum_{s_3,s_4} |\mathcal{M}_{\rm prod}|^2\ |\mathcal{M}_{L\ \rm decay}|^2\ \left|\mathcal{M}_{\bar{L}\ \rm decay}\right|^2 \\
&= 4e^4 \frac{C\bar{C}}{\Gamma_L^2}\left[\mathcal{A} + \frac{4}{9}\mathcal{B}_{\mu\nu} D_\sigma \bar{D}_\lambda \left(g^{\mu\sigma} - \frac{p_3^\mu p_3^\sigma}{m_L^2}\right)\left(g^{\nu\lambda} - \frac{p_4^\nu p_4^\lambda}{m_L^2}\right)\right]
\end{aligned}
\tag{17}
$$

Integrating over the appropriate phase space (including the three δ-functions originating from energy-momentum conservation), gives us the effective cross-section. The expression in eq. (17) includes all spin–spin correlations. It is easy to see that the factor C/Γ_L, on integration, is the same as the branching fraction of L into the specific channel being considered. Since, by definition, all decay modes for L involve the factor $\sin^2\theta_{lL}$, the branching fraction is independent of this quantity. Thus, any mixing element at the decay vertex is irrelevant as long as it is large enough for L to decay within the detector.

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