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THE DECAY OF THE SCALAR NEUTRINO*

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ABSTRACT

One major problem for supersymmetry has been the lack of any experimental motivation. Although scalar neutrinos are usually predicted to be among the lighter new particles, their presence has been expected to be hidden because of their decay into unobserved neutrals. We calculate the decays of the scalar neutrino and show that there may be a substantial rate into charged particles. These decay modes lead to very distinctive signatures for supersymmetry in e^+e^- physics and in $W^\pm$, Z^0 decays.

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1. Introduction

Supersymmetric theories have been of interest recently, because they may provide a solution to the gauge hierarchy problem[1] by explaining why the scale of electroweak interactions (300 GeV) is many orders of magnitude less than the grand unification scale (or the Planck mass). In supersymmetric theories, for each presently known fermion there exists a scalar partner. If supersymmetry is relevant to the solution of the hierarchy problem, then the masses of these scalars cannot be much larger than the weak interaction scale of 300 GeV. In fact, they could even be lighter. Thus, one way to test supersymmetric models experimentally is to search for new scalar particles. However, it has been a major disappointment that absolutely no experimental evidence for or against supersymmetry has been found.[2] Therefore, it is essential to seek further means for determining whether supersymmetry is relevant and which classes of models are indicated.

Experimental searches at PEP and PETRA have already set lower limits of approximately 17 GeV for the masses of charged scalar leptons and scalar quarks.[3] Since it has been assumed that the scalar neutrino decays only into invisible particles (neutrino plus photino or Goldstino), no limits at all exist for the mass of the scalar neutrino. The purpose of this paper is to show that this assumption is not always true, opening new ways for experimentalists to search for scalar neutrinos.

There are two aspects to the search for scalar neutrinos: production mechanisms and decay signatures. In certain favorable circumstances, the production of scalar neutrinos could be large.[1] For example, certain supersymmetric models predict a light wino[6–7] which could significantly enhance the process $e^+e^- \rightarrow$

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[1]Production of the scalar neutrino has been considered in the literature for e^+e^- machines[4] and e^-p machines.[4,5]

$\nu_s \bar{\nu}_s$ via wino exchange. Furthermore, at a Z^o factory, assuming $M_{\nu_s} < \frac{1}{2} M_Z$, we find:

$$\frac{\Gamma(Z^0 \to \nu_s \bar{\nu}_s)}{\Gamma(Z^0 \to \nu \bar{\nu})} = \frac{1}{2}\left(1 - \frac{4M_{\nu_s}^2}{M_Z^2}\right)^{3/2} \tag{1}$$

which corresponds to a branching ratio for each $\nu_s \bar{\nu}_s$ pair of about three percent provided the phase space suppression is not too severe. Finally, at hadron-hadron colliders which are energetic enough to produce the W^+, we expect the decay $W^\pm \to e_s^\pm \nu_s$ to occur if it is kinematically allowed. The signatures for the above processes are described later. In this paper, we concentrate on the decay modes of the scalar neutrinos. Details of the calculation which follow and further discussion of production mechanisms will be provided in a forthcoming paper.[8]

2. Two-Body Decays of the Scalar Neutrino.

In order to perform an experiment capable of detecting scalar neutrinos, it is crucial that the branching ratios of ν_s into charged particles be appreciable. Naively, one would not expect this to be the case; a first guess would be that the dominant channel is $\nu_s \to \nu + \tilde{\gamma}$ (we assume in this paper that the photino is light; e. g. $M_{\tilde{\gamma}} \leq 5 GeV$).[2]

A light $\tilde{\gamma}$ is expected to behave more or less like a neutrino.[10] In particular, we assume that it will neither stop nor decay in any experimental apparatus. (Otherwise, positive evidence for supersymmetry will be available long before the discovery of the ν_s !) If $\nu + \tilde{\gamma}$ is the dominant decay mode, there is little hope that the ν_s is observable. Hence, it is important to do a careful computation of its two-body decay rate. In particular, if four-body final states (with charged particles) can be appreciable relative to this rate, then the observation of the ν_s might be possible.

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[2] In some supersymmetric models,[9] one finds that the proper mass eigenstate is not the $\tilde{\gamma}$, but rather is mostly the supersymmetric partner of the $U(1)_Y$ gauge boson. However, in such models, one usually finds at least one nearly massless new supersymmetric partner (such as a higgsino) which might play an analogous role in ν_s decays. We will not pursue this alternative further.

First, we discuss the two-body decays of the ν_s. We discuss here only the scalar electron neutrino, although we will generalize our remarks to other flavors in the final section. We assume that $M_{\nu_s} < M_\omega$ (ω is a wino), so that $\nu_s \to \omega^+ e^-$ is prohibited.[3] One possible decay mode is $\nu_s \to \nu \tilde{G}$ where $\tilde{G}$ is the massless Goldstino which would result from the spontaneous breaking of supersymmetry. Current algebra arguments[11] allow one to compute the couplings of $\tilde{G}$; the strength of the coupling varies inversely with the scale M_s at which supersymmetry breaks. The result for the lifetime is well known:[11]

$$\tau = \frac{8\pi \not{h} M_s^4}{M_{\nu_s}^5} = 1.65 \times 10^{-23} \times M_s^4 M_{\nu_s}^{-5} \text{ seconds} \tag{2}$$

where all masses are given in GeV. In the early days of supersymmetric model building,[1] it was hoped that M_s would be roughly on the order of 1 TeV. At present, there are many reasons to suspect that this is not true.[12] In fact, it seems more likely that $M_s \sim (M_P M_W)^{\frac{1}{2}} \sim 10^{10} GeV$ (where M_P is either the grand unification scale[9,13] or the Planck mass).[4] For such a large value of M_s, the decay rate (Eq. (2)) becomes utterly negligible. Even for smaller values of M_s the $\nu \tilde{G}$ branching ratio could be negligible. For example, for $M_{\nu_s} = 20 GeV$, the $\nu \tilde{G}$ mode becomes dominant only for $M_s \leq 3 TeV$.

We therefore turn to the expected dominant decay mode: $\nu_s \to \nu + \tilde{\gamma}$. Because there is no bare $\nu_s \nu \tilde{\gamma}$ vertex in supersymmetric models, this process must occur by one-loop graph. Hence, in order to perform a calculation, we must specify a model. We wish to do this, by choosing a procedure which should produce a result that will be fairly model-independent. Our method is as follows: we first construct an unbroken supersymmetric version of the $SU(2) \times U(1)$ model of electroweak interactions. We will introduce the minimal number of fields necessary to break the gauge symmetry down to $U(1)_{EM}$. Second, we will add soft

[3] Of course, if this were not the case, the decay $\nu_s \to \omega^+ e^-$ would occur at the tree level and hence would be the dominant mode and be very visible.

[4] Such a relation is suggested in the context of supergravity models.[7,14] In these models, $\tilde{G}$ is absorbed into the gravitino which then becomes massive with mass on the order of the weak interaction scale. But, even if the decay $\nu_s \to \nu + gravitino$ is energetically allowed, its rate would be totally negligible.

supersymmetry breaking terms (following the rules of Girardello and Grisaru[15]). We will add only those terms necessary for a sensible calculation - specifically, we will only add mass terms for the scalar quarks and leptons. We will not, for example, add Majorana mass terms to the gauge fermions. The resulting model is not meant to be totally realistic but is only needed to provide a framework for a sensible calculation. At the end, we will comment on how the results would change if additional soft supersymmetric breaking terms were to be added. It is interesting to note that recent work in supergravity shows that the effective low energy theory of a spontaneously broken supergravity coupled to matter fields is simply a globally supersymmetric theory broken by various soft terms.[7,14,16]

We now briefly specify the model that we use. Fayet[17] first wrote down a supersymmetric $SU(2) \times U(1)$ model containing two $SU(2)$ doublet chiral superfields and one $SU(2) \times U(1)$ singlet chiral superfield along with the usual gauge supermultiplets.[5] We will make use of that model but unlike the authors of Ref. 17, 18, we will interpret the chiral superfields mentioned as the Higgs boson multiplets. We add to the model the necessary quark and lepton supermultiplets (i. e. $SU(2)$ doublets corresponding to the left-handed fermions and $SU(2)$ singlets for the right-handed fermions). For simplicity, we neglect the masses of all quarks and leptons. Hence we banish the interactions of the quark and lepton supermultiplet with the Higgs supermultiplet in the superpotential. Thus, the quark and lepton supermultiplets interact only through the gauge supermultiplets. The resulting particle spectrum is given in Table 1.

Some consequences of the model are as follows: the particles arrange themselves into new supermultiplets: e. g. $(H^{\pm}, \omega_1^-, \omega_2^+, W^{\pm})$ with mass M_W, $(H^0$, z_1, z_2, $Z^0)$ with mass M_Z and $(\tilde{\gamma}, \gamma)$ with zero mass. In particular, the winos ω_1^- and ω_2^+ are four-component Dirac spinors which are made up of fermions from the gauge and Higgs supermultiplets. For convenience, we have chosen to make use of Majorana fermions z_1 and z_2 as opposed to combining them into one Dirac spinor. Further details regarding this model are discussed in Ref. 8. Two

[5]This model has recently been examined in detail in Ref. 18.

aspects of the model are important to note. First, the ω_1^- couples only to $\nu_s e^-$, whereas the ω_2^+ couples only to νe_{sL}^+[6] (and similarly for other weak doublets). Second, both z_1 and z_2 couple to $\nu_s \nu$, $e_{sL}^- e^+$, and $e_{sR}^- e^+$ in such a way as to allow for $e_{sL}^- e^+ \to e^- e_{sR}^+$ by s-channel z-exchange.[7] Finally, all leptons, quarks and their scalar partners are massless. As stated previously, we arbitrarily give masses to the scalar partners (to be determined by experiment); this softly breaks the supersymmetry.

All the graphs which contribute to $\nu_s \to \nu + \tilde{\gamma}$ are given in Fig. 1. Most of the graphs are divergent, but renormalizability of the theory requires that the sum of all the graphs be finite. The details of the calculations are very instructive and are given in Ref. 8. We found that the divergences do cancel and obtained

$$\Gamma(\nu_s \to \nu + \tilde{\gamma}) = \frac{M_{\nu_s} \alpha^3}{128\,\pi^2 \sin^4 \theta_W} \Big[F(r_1, r_2)\Big]^2 \tag{3}$$

where $r_1 \equiv \frac{M_{\nu_s}^2}{M_W^2}$ and $r_2 \equiv \frac{M_{e_sL}^2}{M_W^2}$. The function $F(r_1, r_2)$ is given explicitly as follows:[8]

$$\begin{aligned} F(r_1, r_2) \; = \; & \frac{2(1-r_1)}{r_1} Li_2(r_1) \; - \; \frac{2r_2}{r_1}\left(\frac{1}{4}\log^2 r_2 - \log^2\left(\frac{1-r_1+r_2+\lambda}{2\sqrt{r_2}}\right)\right) \\ & + \left(\frac{2-r_1-r_2}{r_1}\right)\left(Li_2(1-r_2) - Li_2\left(\frac{1-r_2+r_1-\lambda}{2}\right)\right. \\ & \left. - Li_2\left(\frac{1-r_2+r_1+\lambda}{2}\right)\right) \end{aligned} \tag{4}$$

where $\lambda \equiv [(1-r_1-r_2)^2 - 4r_1 r_2]^{1/2}$. The notation of the dilogarithm is the one used by Lewin:[21] $Li_2(z) = -\int_0^z dx \frac{\log(1-x)}{x}$. It is useful to note the following

[6] This was also noted in Ref. 19. Note that if we add Majorana mass terms for the gauge fermions, this will no longer be the case.

[7] One would interpret $e_{sL}^- e^+ \to e^- e_{sR}^+$ as a reaction which violated fermion number unless one of the bosons was given a non-zero "fermion number" (see Ref. 20). Note that the z_1 and z_2 couplings are such that $e_{sL}^- e^+ \to e^- e_{sL}^+$ does not occur.

[8] Due to the properties of the dilogarithm, $Im F(r_1, r_2) = 0$ if $0 \leq r_1 \leq 1$ independent of r_2.

limits of Eq. (4). For $r_1 = r_2 \equiv r$,

$$F(r,r) = -\frac{2(1-r)}{r}\log r \log(1-r) + \frac{2(1-2r)}{r}\left(\frac{1}{4}\log^2 r - \log^2\left(\frac{1+\sqrt{1-4r}}{2\sqrt{r}}\right)\right) \tag{5}$$

and for $r_1 = 0$,

$$F(0,r_2) = -\frac{r_2}{1-r_2}\left(1 + \frac{r_2 \log r_2}{1-r_2}\right) \tag{6}$$

In the supersymmetric limit where r_1, $r_2 \to 0$, then $F(r_1,r_2) \to 0$. This must be true, since in this limit supersymmetry relates $F(r_1,r_2)$ to the electromagnetic form factor of the neutrino at $q^2 = 0$ which is well known to vanish.[22]

Plugging in the numbers shows that the ν_s lifetime in this mode is

$$\tau = \frac{\hbar}{\Gamma} = 1.13 \times 10^{-16}\left(\frac{1}{M_{\nu_s}(GeV)}\right)\frac{1}{[F(r_1,r_2)]^2}\ \text{seconds} \tag{7}$$

$F(r_1,r_2)$ tends to be a number between 0.1 and 1.0 for interesting masses. It is important to note that the decay rate Γ does *not* decrease with increasing $M_{e_{sL}}$ (this is most easily seen in Eq. (6)). As we shall show, the four-body decays become very suppressed as $M_{e_{sL}}$ is increased. Therefore, the two-body decay mode will certainly dominate unless there are fairly light supersymmetric particles which are relevant to this decay. An exception to the rule just stated would occur if $M_{e_{sL}} < M_{\nu_s}$. In such a case, the three-body decay $\nu_s \to e_s^- e^+ \nu$ (which is a tree level process) would clearly be the most important decay mode if sufficient phase space is available. However, most models which predict scalar lepton masses favor a lighter ν_s (for a counterexample, see Ref.23). Details on three-body decay rates will be presented in Ref. 8.

3. Four-Body Decays of The Scalar Neutrino

We show in Fig. 2 all possible diagrams leading to four-body final states within the model we are using. Unlike the case of $\nu_s \to \nu + \tilde{\gamma}$, many different four-body channels are accessible. To get a first estimate of the rates, let us look at the diagram which is expected to dominate, namely $\nu_s \to e^- u \bar{d} \tilde{g} + e^- c \bar{s} \tilde{g}$.

There are a number of enhancements worth mentioning. First, because the gluino $\tilde{g}$ is being produced, there is a color factor of 4 for the squared amplitude. Second, the gluino couples with the strength of the strong QCD coupling constant (as opposed to the electromagnetic coupling of the photino $\tilde{\gamma}$). Thus, unless the scalar quarks are significantly heavier than the charged scalar leptons or the gluino mass is much more than 5 GeV, it is clear that this process will be the dominant four-body decay mode. The neutral current processes $\nu_s \rightarrow \nu q \bar{q} \tilde{g}$ are somewhat lower in rate but are still important. We may obtain an analytic expression for the dominant rate in the limit that $M_{\nu_s} \ll M_{u_{sL}}, M_W$ and all final-state particles are taken as massless:

$$\Gamma(\nu_s \rightarrow e^- u \bar{d} \tilde{g}) = \frac{\alpha^2 \alpha_s}{11520\, \pi^2 \sin^4 \theta_W} \frac{M_{\nu_s}^9}{M_W^4 M_{u_{sL}}^4} \tag{8}$$

We may compare with this with Eq. (3) and (6). For example, if $M_{u_{sL}} = M_{e_{sL}} = M_W \gg M_{\nu_s}$ then,

$$\frac{\Gamma(\nu_s \rightarrow e^- u \bar{d} \tilde{g})}{\Gamma(\nu_s \rightarrow \nu \tilde{\gamma})} \approx \frac{2}{45}\left(\frac{\alpha_s}{\alpha}\right)\left(\frac{M_{\nu_s}}{M_W}\right)^8 \tag{9}$$

We see that if the ν_s is light and if all other relevant supersymmetric particles have the mass of the W, then the four-body decays will be negligible. On the other hand, from Eq. (8) we see that if $M_{u_{sL}}$ is significantly smaller than M_W, the four-body rates could become appreciable. Similarly, if the ω_1 is significantly lighter than the W, a similar conclusion would follow. We have computed all the diagrams in Fig. 2 explicitly, and have calculated the phase space numerically (with some analytic checks).

We present some of our results in Fig. 3 where $M_{\nu_s} = M_{e_{sL}}$ is varied. For convenience, all masses of scalar quarks and charged scalar leptons were set equal to $M_{e_{sL}}$. The four-body final states which include the gluino could have significant branching ratios totaling about 40% if the parameters of the supersymmetric model are in the approximate ranges discussed above. One should note, however, that the branching ratio for all four-body modes drops sharply if $M_{e_{sL}}$ ($M_{u_{sL}}$) is increased with M_{ν_s} held fixed (as indicated by Eq. 8). Our results show that the branching ratio drops by an order of magnitude as $M_{e_{sL}}$ is increased from 20 to 30 GeV (for M_{ν_s}= 20 GeV). Therefore, if the scalar electron (quark) is significantly heavier than the scalar neutrino, the four-body modes become negligible.

These results suggest a technique for finding the scalar neutrino. Many $\nu_s \bar{\nu}_s$ events will have a striking signature. One of the scalar neutrinos decays into two invisible neutral particles carrying away a large amount of energy and transverse momentum. The other scalar neutrino decays into charged particles. If $M_{\nu_s} \approx M_{e_{sL}}$ then Fig. 3 indicates that this signature occurs almost half the time. If the scalar neutrinos are produced clearly above threshold, the charged decay particles are all in one hemisphere of the center of mass system. Therefore, it is a very impressive signature for an e^+e^- machine. At a Z^0 factory the pair production rate of (light) scalar neutrinos is approximately 3%. Furthermore, for the dominant charged-current process the only missing energy on the side of the charged particles will result from the decay of the gluino into a photino. The total invariant mass of the charged-particle side will be nearly equal to the mass of the ν_s.

One additional test of this picture would be at the hadron-hadron colliders which run at energies large enough for W production. If the ν_s and e_s masses are not too large, then it will be important to consider the decay $W^{\pm} \rightarrow e^{\pm}_{sL} + \nu_s$ (and analogous decays involving other generations of scalar leptons). The branching ratio is

$$\frac{\Gamma(W^- \rightarrow e^-_{sL} \bar{\nu}_s)}{\Gamma(W^- \rightarrow e^- \bar{\nu})} = \frac{1}{2}\lambda^3 \quad . \tag{10}$$

where λ defined below Eq. (4) is a kinematical factor. Thus using the same numbers as in the previous example, we would expect a branching ratio $BR(W^- \rightarrow e^-_{sL} \bar{\nu}_s) \approx 4\%$ which would lead to two classes of signals. If the $\bar{\nu}_s$ decays into neutrals, one would see an unbalanced electron jet (corresponding to the decay $e^-_{sL} \rightarrow e^- \tilde{\gamma}$). The electron would tend to be much softer than those from $W^- \rightarrow e^- \bar{\nu}$ decay. If the $\bar{\nu}_s$ decays into charged particles (presumably dominated by quarks and a gluino), then one would see an electron jet on one side and a hadron jet on the other side. Thus, we hope that further studies of W production at the CERN $p\bar{p}$ collider will be able to put significant constraints on the masses of the ν_s and e_s or else discover them.

4. Discussion and Conclusions

We have shown that there exists a range of the supersymmetric parameters in which four-body charged-particle modes of the ν_s are competitive with the unobservable $\nu\tilde{\gamma}$ mode. One may wonder how dependent our results were on the model which we used. In particular, in our model the ω_1^- and ω_2^+ (the winos) were degenerate in mass with the W. If we were to add Majorana mass terms to the gauge fermions, the masses of the ω_1^- and ω_2^+ would split on either side of the W. As noted in Refs. 6 and 7, certain supersymmetric models suggest that one of the winos could be significantly lighter than the W. Without making any detailed calculations, we can make the following observations. In the case of the four-body decays, a lighter ω_1 will certainly enhance diagram 2(e). Equations (4)-(6) suggest that a lighter wino is likely to have a relatively small effect on the two-body decay rate. The end result is that the four-body modes are enhanced further. This just means there is a larger range of supersymmetric parameters in which ν_s decays are observable.

We have concentrated on the decay of the scalar electron neutrino. It is important to consider whether any of our conclusions would change for the decay of the scalar muon (or tau) neutrino. The answer depends on various masses and mixing angles in the supersymmetric model which is used. For example, a diagram analogous to Fig. 2(e) could lead to $\nu_{\mu s} \to \mu^- e^+ \nu_{es}$ if the mass difference between $\nu_{\mu s}$ and ν_{es} were large enough. Furthermore, mixing between $\nu_{\mu s}$ and ν_{es} (analogous to Cabibbo mixing in the quark sector) could lead to unusual decay signatures. However, there are strong phenomenological constraints on the mass differences between successive generations of scalar quarks and leptons.[24] The absence of flavor-changing neutral-currents typically leads to the requirement that mass-squared differences between (at least) the first two generations be very small. Many supersymmetric models predict that $\Delta M_{sl}^2 = \Delta M_l^2$ (up to radiative corrections) and similarly for the scalar quarks, which is consistent with the above requirement. This would imply that the three generations of scalar neutrinos are nearly degenerate in mass. If such is the case, all of the results of this paper remain true for the decay of the scalar muon (or tau) neutrino. Furthermore, mixing effects would be negligible, because in the limit that the scalar neutrinos are exactly degenerate in mass, electron, muon and tau numbers are separately conserved global quantum numbers.

If supersymmetry is relevant for the solution of the hierarchy problem, then it follows that there are new particles likely to be discovered either at current accelerators or at machines now under consideration. It is important to search for all possible types of new particles that the theory predicts. In this paper, we have concentrated on the properties of scalar neutrinos. It is entirely possible that the relevant parameters are such that the dominant decay mode of the ν_s is into the unobservable channel $\nu\,\tilde{\gamma}$. On the other hand, there exists a large range of parameters in which the ν_s will decay appreciably into charged particles. In this case, those parameters suggest that production rates of $\nu_s\,\bar{\nu}_s$ could be significant. This would be a fortuitous situation and allow for the discovery of the scalar neutrino. Perhaps, the best chance is to observe at a Z^0 factory the decay $Z^0 \rightarrow \nu_s\,\bar{\nu}_s$ where one of scalar neutrino decays into unobservable neutrals and the other one into charged modes. Such an observation would be an important step in confirming the supersymmetric picture.

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Table 1

SYMBOL	*NAME*	*PARTNER*	*MASS*
$\tilde{\gamma}$	photino	photon	0
ω_2^+	wino	W^+, H^+	M_W
ω_1^-	wino	W^-, H^-	M_W
z_1, z_2	zino	Z^0, H^0	M_Z
$\tilde{h}$	higgsino	h_i^0	M_h
ν_s	scalar neutrino	ν	0
l_{sL}, l_{sR}	scalar leptons	l	0
q_{sL}, q_{sR}	scalar quarks	q	0

Particle Spectrum of a Supersymmetric $SU(2) \times U(1)$ Model. The supersymmetry is broken softly by adding explicit mass terms for the scalar quarks and leptons. If no Majorana mass terms for the gauge fermions are added, the photino remains massless, ω_2^+ couples to $e_s^+\nu$ but not $e^+\nu_s$, and ω_1^- couples to $e^-\bar{\nu}_s$ but not $e_s^-\bar{\nu}$. Note that neutral Higgs particles h_i and the higgsino $\tilde{h}$ are not needed in this paper.

FIGURE CAPTIONS

1. One-loop diagrams for $\nu_s \rightarrow \nu\,\tilde{\gamma}$. These are the contributing graphs in a supersymmetric $SU(2) \times U(1)$ model with the supersymmetry broken only by explicit mass terms for the scalar quarks and leptons. Note that in (e), the loop consists of either $W^- \omega_2^+$ or $W^+ \omega_1^-$. Similarly in (f) with H replacing W. In (g) the loop consists of either $e_{sL}^- e^+$ or $e_{sR}^+ e^-$. In the model we use, graph (g) vanishes exactly.

2. Four-body decays of the scalar neutrino. See caption to Fig. 1. Note that we use the symbol u and d for all up-type and down-type quarks, etc. For convenience, the Cabibbo angle is neglected.

3. Branching ratio of four-body modes of the scalar neutrino. We label the various modes as follows:

   (a) $e^- u\,\bar{d}\,\tilde{g} + e^- c\,\bar{s}\,\tilde{g}$;

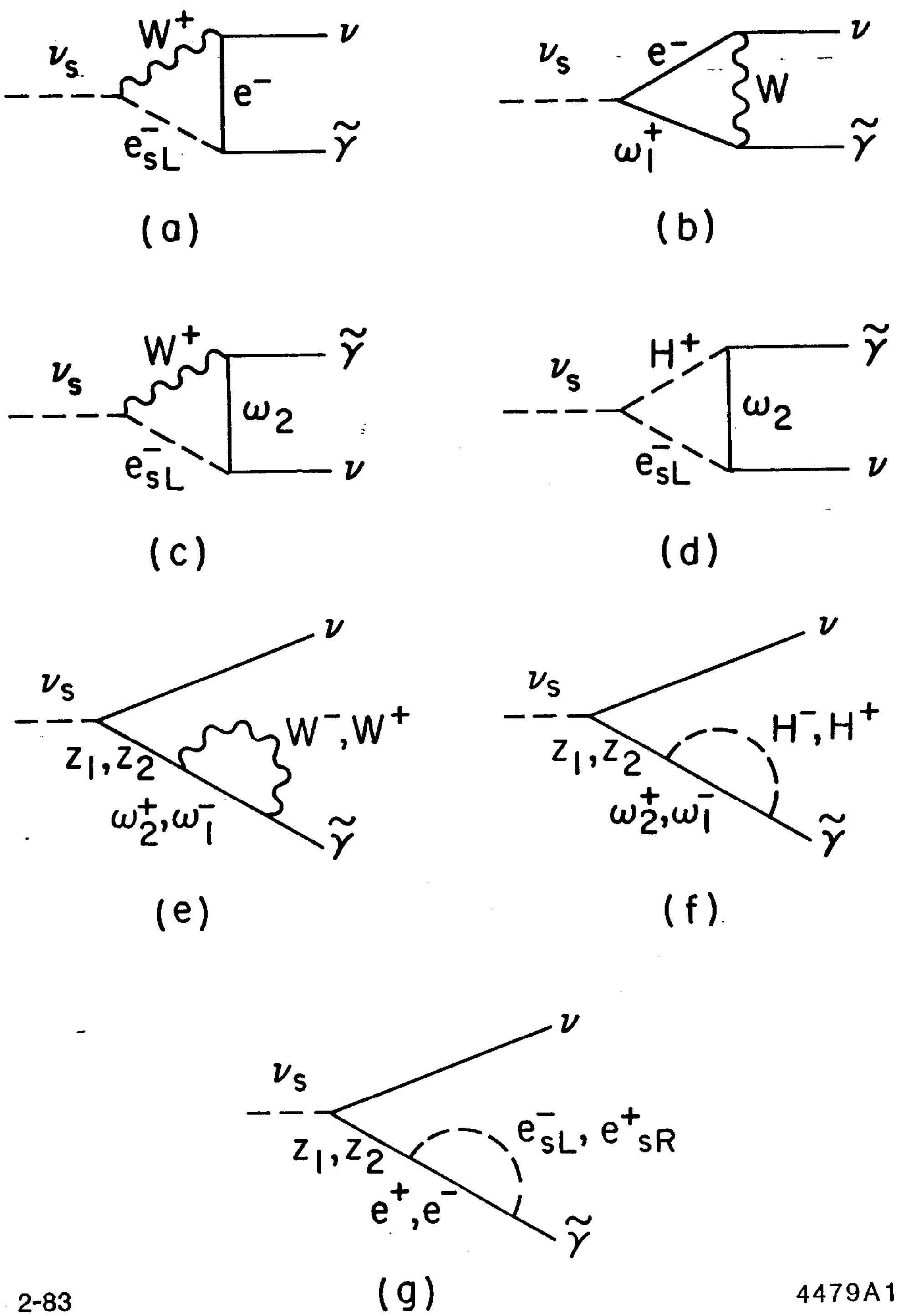
   (b) $\sum_i \nu\,\bar{q}_i\,q_i\,\tilde{g}$;

   (c) $e^- u\,\bar{d}\,\tilde{\gamma} + e^- c\,\bar{s}\,\tilde{\gamma}$;

   (d) $\nu\,e^+ e^-\,\tilde{\gamma}$;

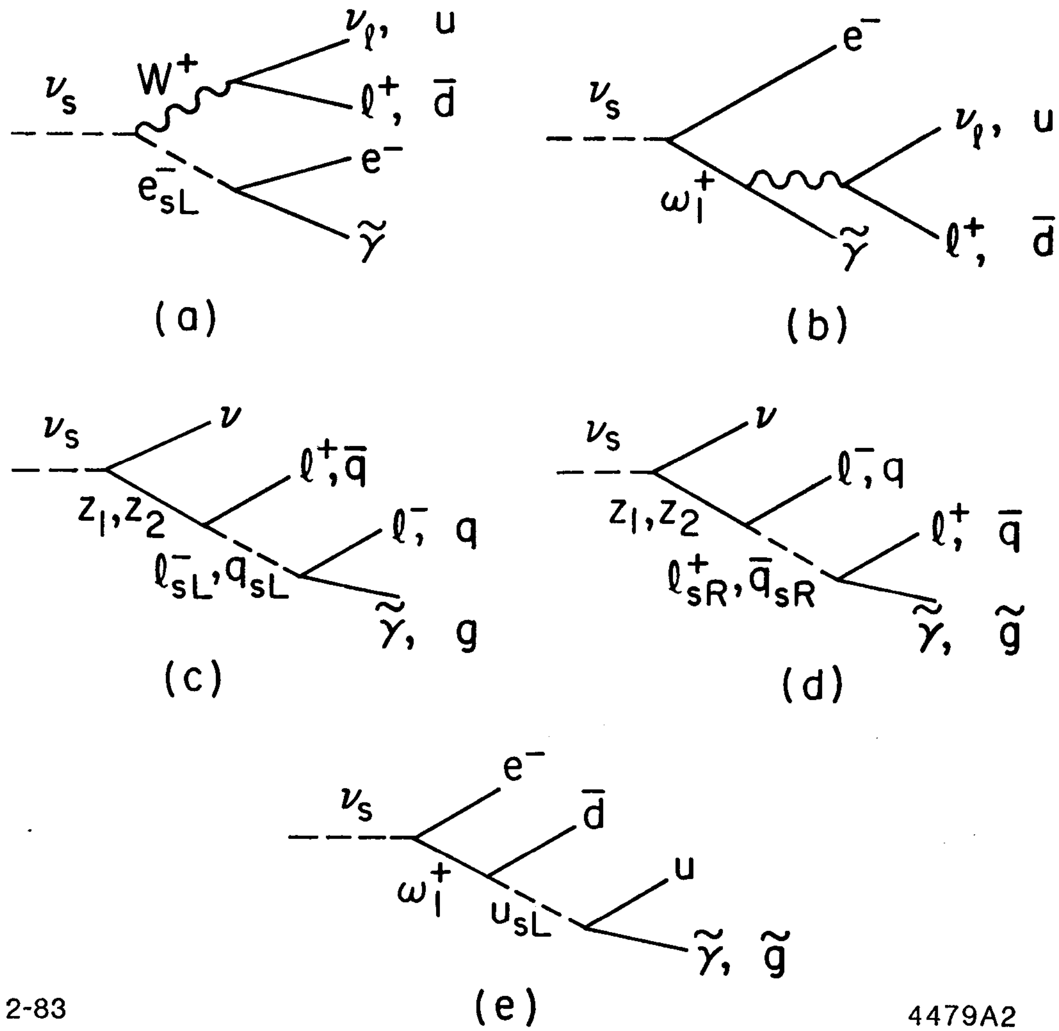
   (e) $\nu\,\mu^+ \mu^-\,\tilde{\gamma}$ or $\nu\,\tau^+ \tau^-\,\tilde{\gamma}$;

   The rates for $\nu\mu^+ e^-\,\tilde{\gamma}$, $\nu\tau^+ e^-\,\tilde{\gamma}$ and $\sum_i \nu\,\bar{q}_i\,q_i\,\tilde{\gamma}$ each occur approximately at the level of (d). For convenience the Cabibbo angle is neglected. We use a value of $\alpha_s = 0.24$. We have assumed that charged scalar lepton and scalar quark masses are equal.



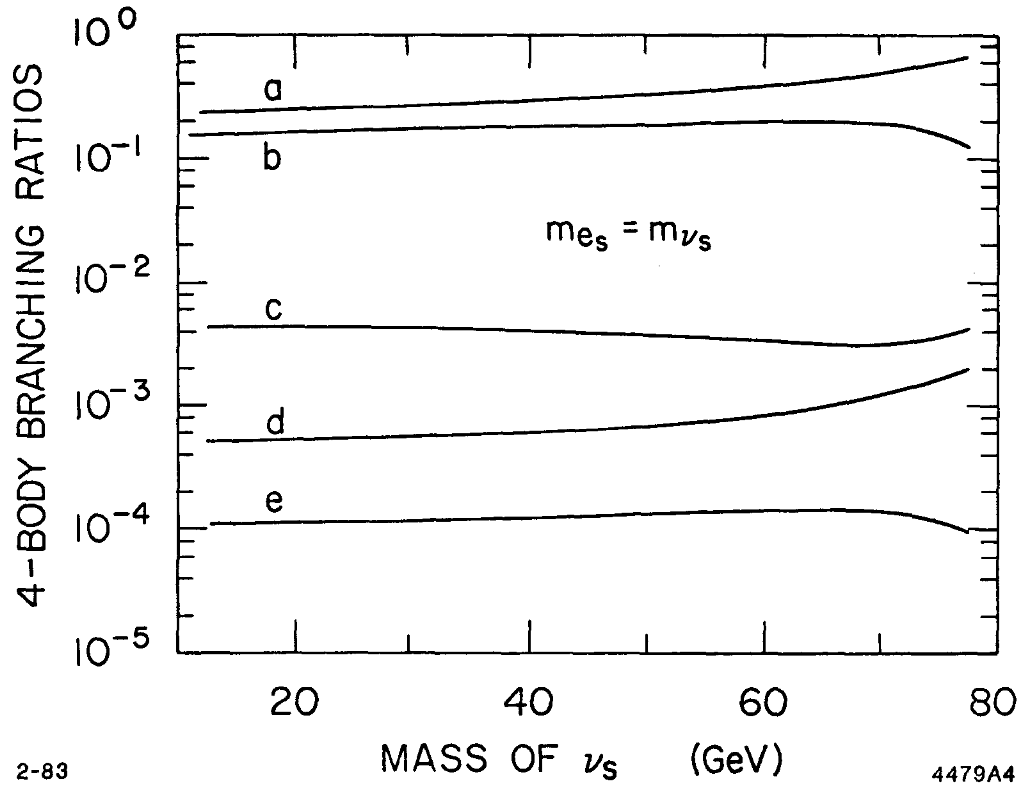
νs
W+
e−
esL−
ν
γ̃
(a)
νs
e−
W
ω1+
ν
γ̃
(b)
νs
W+
ω2
esL−
γ̃
ν
(c)
νs
H+
ω2
esL−
γ̃
ν
(d)
νs
z1,z2
W−,W+
ω2+,ω1−
ν
γ̃
(e)
νs
z1,z2
H−,H+
ω2+,ω1−
ν
γ̃
(f)
νs
z1,z2
esL−, e+sR
e+,e−
ν
γ̃
(g)
2-83
4479A1

Fig. 1



ν_s
W+
ν_ℓ, u
ℓ+, d̄
e−
e−_sL
γ̃
(a)
ν_s
e−
ν_ℓ, u
ω+_1
γ̃
ℓ+, d̄
(b)
ν_s
ν
ℓ+, q̄
Z1, Z2
ℓ−, q
ℓ−_sL, q_sL
γ̃, g
(c)
ν_s
ν
ℓ−, q
Z1, Z2
ℓ+, q̄
ℓ+_sR, q̄_sR
γ̃, g̃
(d)
ν_s
e−
d̄
u
ω+_1
u_sL
γ̃, g̃
(e)
2-83
4479A2

Fig. 2



4-BODY BRANCHING RATIOS
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
a
b
$m_{e_s} = m_{\nu_s}$
c
d
e
20
40
60
80
MASS OF ν_s (GeV)
2-83
4479A4

Fig. 3